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AND CONTROLLED
NUCLEAR FUSION RESEARCH
1990

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AND CONTROLLED NUCLEAR FUSION RESEARCH
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In three volumes

VOLUME 1

INTERNATIONAL ATOMIC ENERGY AGENCY
VIENNA, 1991
FOREWORD

The 1990 International Atomic Energy Agency Conference on Plasma Physics and Controlled Nuclear Fusion Research was characterized by reports of steady technical progress in research on both magnetic and inertial confinement fusion, leading towards the long term goal of producing commercial energy from controlled fusion power generators. Also, major results were reported from completion of the Conceptual Design Activities of the International Thermonuclear Experimental Reactor (ITER) project, which has been conducted since 1988 under the auspices of the IAEA.

This conference, the thirteenth in a series held biennially, was organized in cooperation with the United States Department of Energy, to whom the IAEA wishes to express its gratitude. Over 640 participants and observers from thirty countries and two international organizations attended the conference.

Over two hundred technical papers were presented in thirty technical sessions, including eight poster sessions. There were contributions on tokamak experiments; inertial confinement; non-tokamak confinement systems; magnetic confinement theory and modelling; plasma heating and current drive; ITER; technology and reactor concepts; and the economic, safety and environmental aspects of fusion. The opening session of the conference included an address by Admiral James D. Watkins, Secretary, United States Department of Energy; a round table discussion entitled Why Fusion?; a summary talk on the ITER project; and the traditional Artsimovich Memorial Lecture.

These proceedings, which include all the technical papers and five conference summaries, are published as a supplement to the IAEA journal *Nuclear Fusion*.

The IAEA contributes to international collaboration and exchange of information in the fields of plasma physics and controlled nuclear fusion by organizing these biennial conferences and by sponsoring technical committee meetings, workshops, consultants meetings and advisory groups on relevant topics. Through these and other activities, the IAEA hopes to contribute significantly towards bringing forward the day whose night will be lit by fusion generated power.
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This text was compiled before the unification of Germany in October 1990. Therefore the names German Democratic Republic and Federal Republic of Germany have been retained.

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ARTSIMOVICH MEMORIAL LECTURE
AND
TOKAMAK EXPERIMENTS

(Session A)

Chairmen

Session A-I  R. PARKER  United States of America
Session A-II  J.D. CALLEN  United States of America
Session A-III  H. KISHIMOTO  Japan
Session A-IV  A. GIBSON  CEC
Session A-V  Y. HUO  China
Session A-VI  K. RAZUMOVA  Union of Soviet Socialist Republics

Session A-VII  (Posters)
Mr. Chairman, Ladies and Gentlemen,

I want to thank the International Atomic Energy Agency for asking me to give this lecture. It is my great honour to be here and to speak to you who share a strong interest and enthusiasm for developing plasma physics and technology with the common goal of attaining thermonuclear energy for the welfare of mankind.

I think that fusion energy is essentially different, in many respects, from traditional energy sources. Fusion is really a very difficult technology, to capture the 'sun' within a man-made cage.

In this lecture, allow me to consider only magnetic confinement fusion. In this case, we have to keep ultra-high temperature plasmas stationary within a limited volume; but energetic plasma is fairly violent, and we have to know its nature thoroughly before we can confine it well. Even in common internal combustion engines, physical and chemical processes taking place within combustion cylinders are also extremely complicated and rapidly changing in time, and we are far from understanding the detailed mechanism of combustion. Consequently the control of chemical products in exhaust gases, for example, is still not easy. However, the phenomena and the relevant materials are now within our reach and can be investigated experimentally well enough to obtain the data necessary for the design of practical engines.

Now, when we recollect the history of fusion research and development, we first imagined a high temperature plasma to be like a gaseous octopus or jellyfish that has no bones or shell and is difficult to catch, especially when it becomes violent at high temperatures. Our only hope, if there was any, was that the electromagnetic field might possibly play a role in the long range order in a plasma and produce a sort of macroscopic order. Our vague idea was first to arrange the electromagnetic field and other boundary conditions to make the confinement macroscopically stable — this step corresponds to providing a hard shell to confine the octopus. The next step was to improve plasma confinement by means of reducing various losses of plasma particles and energy which were thought to occur owing to certain instabilities. We found that the shells were sometimes not strong enough; for instance, a shell could automatically break or change form. Further new difficulties appeared one after another when we approached the regions of breakeven conditions. We tried to find clues to solve the problems through the combined use of plasma physics, diagnostics and technologies. However, the non-linear behaviour of plasmas was often beyond what we had anticipated and, when we tried to push a projecting head under the shells, several hidden legs were likely to appear.
The realization of nuclear fusion encountered difficulties not only because it is an unprecedented technology but also because it is a qualitatively new attempt to confine the totality of a slippery object in a reasonably quiet and stationary manner. We have overcome many of these difficulties successively. What were the secrets of successful cases of our efforts and how should we learn from them to go ahead to arrive at the final goal? I want to consider these questions with respect to two key areas, namely plasma physics and international collaboration.

I hope you will forgive me for basing my observations frequently on research made in Japan or by Japanese researchers. This is only because I have many acquaintances in Japan who are willing to tell me their research stories.

I think you will immediately remember Academician Artsimovich for his outstanding contribution in these two key areas. In 1968, at the IAEA international conference held at Novosibirsk in the USSR, remarkable results were reported indicating that electron temperatures up to 2 keV as estimated by diamagnetic measurements had been obtained in the tokamak T-3. The T-3 programme had been under the strong leadership of Academician Artsimovich. In response to his call, the Culham Laboratory sent the ‘Culham Mission’, consisting of Dr. Peacock and his group, to confirm the results by using the newly developed laser scattering diagnostics. In this way, the T-3 results confirmed by this new technique received worldwide recognition. This is a good example of success attained through scientists’ enthusiasm and by international collaboration. You know, international collaboration was rather exceptional in those days. It is worth noting that international collaboration has been very effective in increasing the reliability of research results, and the 1968 IAEA conference provided a chance to promote international collaboration.

Secondly, let me tell you about another, more recent example. It is the case of efforts towards achieving higher beta in tokamaks. Computational codes of MHD theory were developed in many countries, each in an original way; they were compared with one another and gradually a common understanding was achieved. Around 1980, JAERI’s small tokamak JFT-2, with a circular plasma cross-section, gave a value of beta of about 3% for the first time.

Theory also predicted that the upper bound of beta could be much increased when plasma cross-sections were elongated. This prediction was beautifully verified by the DIII-D experiment with US–Japanese collaboration. The values of beta were about 5% in 1982 and have recently been above 9% in DIII-D, with a D shaped cross-section. This is also a story of success in which physics researchers started from personal, original ideas and competitive research work, followed by exchange of information among them, leading to an international common understanding, and finally to major results obtained by joint experiments through international collaboration. I think this is a model of productive R&D.

Thirdly, I want to tell you how physics was useful in finding reliable linkages among results obtained by facilities of different sizes and types. Let us observe, as an example, the case of current drive, which is thought to be one of the most useful means to attain stationary tokamaks.
In 1966, Dr. Shoichi Yoshikawa and others at the Princeton Plasma Physics Laboratory showed the possibility of driving current by means of electromagnetic wave injection using the C-stellarator. After preliminary studies, the theory of current drive by means of lower hybrid waves was developed and experiments started in tokamaks in many laboratories, namely WT-2, JFT-2, JIPP-II and many other devices throughout the world. A recent result in this area is the realization of large current (2 MA) and long pulse drive (over one hour) at JT-60 and TRIAM-1M (Kyushu University), respectively.

Current drive by means of neutral beam injection also began with theoretical work by Dr. Tiiho Ohkawa (General Atomics), and experiments were performed in JET and TFTR, among others. Current drive for 2.5 s was demonstrated using only NBI drive in the DIII-D facility. Driving efficiency was also studied extensively and was shown to agree with theoretical estimates. In this way, you see how various ideas of current drive were conceived and proved using many tokamaks.

Fourthly, let me consider how the interaction of physics and technology is important. The success of lower hybrid current drive experiments in TRIAM-1M was brought about by a clever combination of technical and physical ideas. In this case, plasma position was determined using real time sensing, and was controlled and appropriately swept in order to disperse heat load to the plasma facing wall. In this way experimenters succeeded in maintaining a quiet and pure plasma for more than one hour by relatively small (30 kA) current driven by LH waves. Thus we see that progress in plasma physics will reduce the technical burden imposed upon facilities and vice versa.

Now, although plasma physics has made remarkable progress, it is not yet so powerful that the design of plasmas of next phase facilities or reactors can be done relying only upon physics without extrapolation of empirical scaling laws. In order, however, to make design really reliable, it is clearly desirable that the design be done on the basis of natural laws as far as possible. However, the situation is not so easy. First, since many degrees of freedom are still alive in high temperature plasmas, even the most advanced diagnostics are not adequate to measure the physical quantities of plasma with sufficient spatial and temporal precision to completely characterize the plasma. Secondly, although computational physics has made tremendous progress, the plasma is non-linear in principle and we cannot solve the basic equations exactly. What we can do is to find solutions on a case by case basis by giving the physical quantities (e.g. initial and boundary conditions) necessary to characterize the plasma under consideration.

Such being the case, what we have to do from the physics side may be, among other things, to uncover and clarify important factors hidden behind various plasma phenomena, say, self-organization, for example pinches or bootstrap current, disruption, etc., by using the method of analysis and simulation. Most important in this case is the power of insight to point out the most essential factors hidden behind the phenomena. Such insight is originally based upon personal intuition, but it can be polished by training face to face with reality. In this sense, mutual contact and
collaboration among theoretical and experimental physicists and engineers are essential.

Central problems to be attacked by this kind of collaboration are as follows: the characterization of H mode and other improved modes, elucidation of various kinds of anomalous transport, further understanding of plasma–wall interaction, establishment of a method of particle and heat removal from the reactor core, etc., and finally one of the most fundamental problems of plasma physics — long pulse stationary operation with good confinement.

I think this is just the occasion to refer to the present activities of the epoch making international collaboration, ITER. Everybody concerned recognizes the important progress made in the confinement database, namely the construction of an internationally evaluated database and efforts to improve it by new experiments and analyses done at various research institutions. Moreover, strong correlations have become evident among a number of variables — confinement, divertor, density limit, beta limit, etc. — each of which had been studied separately. Significant progress has been made and effective procedures established to examine overall consistency. I would like here to stress the importance of individual, original ideas, enthusiasm, collaboration and a consistent physical and engineering approach; these are indispensable to attaining a reliable design and possible future construction.

Finally I want to consider what policy is needed in order to succeed in attaining the goal of fusion energy.

Firstly, independent tests and verifications are extremely important in research work, and no results are believable unless similar results can be reproduced in analogous facilities. You know well that even the existence of four large tokamaks — JET, JT-60, TFTR and T-15 — was not always enough for the needs of ITER design work. How to meet this need for a diversity of machines will be a very important question if we pursue the line of a single large international machine.

Secondly, the question of flexibility is also serious and cannot be bypassed. Namely, when we follow the main line of a standard tokamak project like ITER, I think we have to encourage the birth of various new ideas and always be prepared to respond to them; for instance, the idea of a stationary reactor based upon positive use of bootstrap current, of a microscopically more stable tokamak with low aspect ratio, or of a non-tokamak reactor, which might become more desirable for commercial purposes, and so on. Thus we face a serious problem of how to divide limited human and material resources in a balanced way among a variety of research items. This problem needs to be discussed and a solution has to be found that is understandable to all the people concerned.

Next, I want to discuss the importance of different points of view. You know that the free activity of researchers is essential in producing new ideas which might have an epoch making effect on the development of science and technology. Therefore, when many people always work together towards the construction of a single large machine, we have to take special care not to kill a free and critical atmosphere.
I would like to remark that, generally speaking, joint international work towards the common goals of mankind gives us a strong belief in its being one of the surest pathways to attaining global understanding. I have personally shared, as a co-chairperson of the US–Japan Co-ordinating Committee, in the important experience of continued collaboration between the USA and Japan in the past twelve years. The collaboration has been successful, as you see, for instance, from the good results of Doublet; but our path has never been easy. Besides the benefit of cost and risk sharing, we shall obtain a precious treasure called mutual understanding, which only arrives after patient efforts to overcome such cultural sources of friction as differences in budgetary systems, industry participation, planning ideas and the innumerable trifles of daily life. I am sure people working together on ITER also share the same impression.

When the mutual understanding obtained by such international collaborative programmes as ITER spreads in many other fields, for example global environmental studies, west to east and north to south, I can truly hope we shall find a glorious entrance to a peaceful world.
RECENT TFTR RESULTS


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\textsuperscript{9} Canadian Fusion Fuels Technology Project, Toronto, Ontario, Canada.
Abstract

RECENT TFTR RESULTS.

TFTR experiments have emphasized the optimization of high performance plasmas as well as studies of transport in high temperature plasmas. The recent installation of carbon composite tiles on the main bumper limiter has allowed operation with up to 32 MW of neutral beam injection without degradation of plasma performance by large bursts of carbon impurities ("carbon blooms"). Plasma parameters have been extended to $T_i(0) \sim 35$ keV, $T_e(0) \sim 12$ keV, $n_e(0) \sim 1.2 \times 10^{20}$ m$^{-3}$, producing D-D reaction rates of $8.8 \times 10^{16}$ reactions per second. The fusion parameter $n_e(0)T_i(0)$ in supersonic plasmas is an increasing function of heating power up to an MHD stability limit, reaching values of $\sim 4.4 \times 10^{20}$ m$^{-3}$ sec keV. Peak-concentration-profile hot-ion plasmas with the edge characteristics of the H-mode have been produced in a circular cross-section limiter configuration with $n_e(0)T_i(0)$ values characteristic of supershots, namely up to four times those projected for standard H-modes with broad density profiles. Reduced transport is also observed in the core of high-density ICRF-heated plasmas when the density profile is peaked. At the highest performance, the central plasma pressure in TFTR reaches reactor level values of 6.5 atmospheres. These regimes, MHD instabilities with $m/n = 1/1, 2/1, 3/2$ and $4/3$ are often observed concurrent with a degradation in performance. High $\beta_p$ plasmas with $\delta_{pe} \approx 1.6$ and $\beta/(\Omega_a^2)$ $\approx 4.7$ (\%) have demonstrated confinement enhancement over the low-mode confinement time with $\tau_e/\tau_L \sim 3.5$ and a bootstrap current of about 65% of the total plasma current. The best TFTR supershots in deuterium plasmas have $Q_{\text{obs}} = 1.9 \times 10^3$, corresponding to an equivalent $Q_{\text{DT}} \sim 0.31$ if the plasma and beam parameters are exactly the same in a 50/50 D/T plasma. Performance enhancements in D-T due to higher heating power, higher density, lower $Z_{\text{eff}}$, ion mass effect and optimum programming of the D-T mix in the beam extrapolate to D-T plasma performance near the breakeven regime, $Q_{\text{MHD}} \sim 0.5 - 0.7$, that would produce $\sim 15 - 25$ MW of fusion power. Lower $Q$ ($\sim 0.3$) plasmas at higher temperatures are expected to produce alpha particle betas suitable for testing collective alpha instability theories.

1. INTRODUCTION

During the past two years, the TFTR program has concentrated on:

(1) optimization of high performance plasma regimes at modest (1.6 - 2 MA) plasma currents. These regimes utilize peaked-electron-density profiles to enhance deuterium fusion power;

(2) detailed measurements of local plasma transport and fluctuations in a variety of plasma regimes; and

(3) preparation for D-T experiments, including the achievement of deuterium plasma parameters that extrapolate to the D-T breakeven regime, theoretical studies of alpha physics and improvements in the D-T configuration.

These activities have important consequences for optimizing the plasma regimes of BPX, ITER and a reactor of the ARIES type. The peaked-density-profile regimes enhance fusion performance ($P_{\text{fusion}}/P_{\text{loss}}$) by increasing the central density and ion temperature as well as the global energy confinement time. More fundamental understanding arising from transport studies in TFTR may yield new techniques for reproducing the favorable confinement properties of these regimes. Bootstrap and beam-driven-current studies are of major interest for ITER and ARIES operating scenarios, while the recent high $\beta_p$ experiments may provide the basis for a new regime of operation near the boundary of the first regime of MHD stability.
The D-T program will provide information to BPX, ITER and ARIES on confinement of D-T plasmas, initial studies of alpha particle effects and a demonstration of significant fusion power.

2. HARDWARE CAPABILITIES

The TFTR hardware capabilities are summarized in Table I and are compared to the values specified in the 1975 Final Conceptual Design of TFTR [1]. TFTR has a circular cross-section plasma with minor radius variable from 0.4 m to 0.96 m and major radius variable from 2.1 m to 3.1 m. The toroidal field system is routinely operated at the full rating of 5.2 T and plasma currents of up to 3 MA have been achieved. The neutral beam system has injected a total power of up to 32 MW. The ICRF power systems have been operated up to their design value, with 6 MW at 47 MHz and 3 MW at 40 - 80 MHz. Two antennas, each with two current straps and carbon Faraday shields, have coupled a total of 6.3 MW to the plasma with an RF power density of 10 MWm^{-2} on one antenna. The pulse length of the auxiliary heating systems (NB and ICRF) is 2 seconds, which is ~ 10 times the typical confinement time at high temperature.

<table>
<thead>
<tr>
<th>TABLE I. TFTR HARDWARE CAPABILITIES</th>
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<tbody>
<tr>
<td></td>
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<tr>
<td>a(m)</td>
</tr>
<tr>
<td>R(m)</td>
</tr>
<tr>
<td>ELONGATION</td>
</tr>
<tr>
<td>B\textsubscript{T}(T) AT 2.48 m</td>
</tr>
<tr>
<td>I\textsubscript{p}(MA)</td>
</tr>
<tr>
<td>P\textsubscript{NB}(MW)</td>
</tr>
<tr>
<td></td>
</tr>
<tr>
<td>NEUTRAL BEAM</td>
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<tr>
<td>PULSE LENGTH (sec)</td>
</tr>
<tr>
<td>LIMITER</td>
</tr>
<tr>
<td>P\textsubscript{ICRF}(MW)</td>
</tr>
</tbody>
</table>

* INCREASED IN 1980
† INCREASED IN 1988
The limiter system for TFTR consists of an axisymmetric inner-wall bumper limiter and two carbon/carbon (C/C) composite poloidal arcs on the outer cross-section to protect the ICRF antennas. In 1988-1989, the bumper limiter consisted of 2000 0.1 m x 0.1 m POCO graphite tiles. During high power operation, small cracks or high spots on the limiter formed local hot-spots (> 1700 °C) which would evaporate significant quantities of carbon into the plasma, thereby dramatically reducing the neutron production in a high performance supershot. A typical case is illustrated in Fig. 1, where 22 MW of neutral beam power produced a "carbon bloom" at ~ 0.6 sec. It is speculated that electrons deposited their energy in a very short scrape-off layer during the disruption of a previous high energy discharge and thereby initiated cracks which propagated through the graphite tiles. In early 1990, the graphite tiles in the high heat-flux regions of the inner-wall bumper limiter were replaced with C/C fiber reinforced tiles. The new C/C limiter system has allowed operation up to 32 MW for 1 sec or 25 MW for 2 sec without carbon blooms limiting performance.

TFTR has two pellet injectors. An eight barrel D2 pellet injector, developed and built by Oak Ridge National Laboratory, produces pellets of 3, 3.5 and 4 mm diameter at speeds of ~ 1.5 km sec\(^{-1}\). A two barrel lithium or carbon pellet injector developed, built and operated by Massachusetts Institute of Technology, uses 2 mm pellets at 0.8 km sec\(^{-1}\) for fueling and diagnostic experiments.
3. TFTR PLASMA REGIMES AND PERFORMANCE

The TFTR program has emphasized optimization and study of enhanced-confinement regimes with peaked-density-profiles such as the supershot, peaked-density-profile H-mode, and high-density pellet-injected plasmas. In addition, high $\beta_p$ regimes have been produced at low plasma current.

The ratio of fusion power ($P_{\text{fusion}}$) to power lost from the plasma ($P_{\text{loss}}$) is a useful measure of fusion plasma performance. For a thermonuclear plasma with average ion temperatures in the 8 - 23 keV range, $P_{\text{fusion}}/P_{\text{loss}} = n_i(0)T_i(0)\tau_S$ where $\tau_S = W_p/P_{\text{heat}}$, $W_p$ being the plasma energy and $P_{\text{heat}}$ the non-fusion external heating power. This formulation gives the proper description for both steady-state and transient cases with large $dW_p/dt$. Additionally, it can be shown that $\tau_S = W_p/P_{\text{heat}}$ is the correct confinement time to use on the Lawson diagram with transient plasmas. TFTR was designed to enhance the fusion power output using beam-target and beam-beam reactions so that $P_{\text{fusion}}/P_{\text{loss}} = M n_i(0)T_i(0)\tau_S$ where $M$ is the fusion reactivity multiplier due to beam-target and beam-beam reactions. It is customary to express the confinement time in terms of an enhancement factor ($H$) over the original L-mode expression [2] for the confinement time. Therefore:

$$P_{\text{alpha}}/P_{\text{loss}} \sim M \left( \frac{n_D}{n_e} \right) \gamma_n \gamma_T \left[ T_i/T_e + T_i \right] H^2 I_p^2 \left( \frac{R}{a} \right)^{2.5} a^{-0.24}$$  \hspace{1cm} (1)

where $\gamma_n$ is the density profile peaking factor and $\gamma_T$ is the temperature profile peaking factor. The principal goal of TFTR experiments is to increase $M$, $\gamma_n$, $\gamma_T$, $T_i/T_e$ and $H$ simultaneously to obtain high performance at modest plasma currents.

<table>
<thead>
<tr>
<th>TABLE II. TFTR PLASMA PARAMETER RANGE</th>
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<tbody>
<tr>
<td>DESIGN</td>
</tr>
<tr>
<td>--------------------------------------------------</td>
</tr>
<tr>
<td>$T_i(0)$ (keV)</td>
</tr>
<tr>
<td>$&lt;T_i&gt;$ (keV)</td>
</tr>
<tr>
<td>$T_e(0)$ (keV)</td>
</tr>
<tr>
<td>$&lt;T_e&gt;$ (keV)</td>
</tr>
<tr>
<td>$n_e(0) \ (10^{20} \text{m}^{-3})$</td>
</tr>
<tr>
<td>$n_e \ (10^{20} \text{m}^{-3} \text{sec})$</td>
</tr>
<tr>
<td>$n_e \ (10^{20} \text{m}^{-3} \text{sec} \text{keV})$</td>
</tr>
<tr>
<td>CENTRAL PRESSURE (ATMOSPHERE)</td>
</tr>
</tbody>
</table>

* INCREASED IN 1980

# WITH DEUTERIUM PELLETS
The plasma parameters achieved in TFTR are summarized in Table II. Central ion temperatures of 35 keV and particle-average ion temperatures of 12.4 keV have been attained. Similarly, central electron temperatures of 11 keV and particle-average electron temperatures of 5.2 keV have been obtained in supershot plasmas with central electron densities up to 9.5 x 10^19 m^-3. A slight major radius compression (C = 1.07) has been used to increase these values to T_e = 12 keV, and n_e(0) = 1.2 x 10^20 m^-3. In these high density cases, the total plasma pressure is up to 6.5 atmospheres with only 10% of the pressure due to unthermalized beam ions. Pellet injection has been used to produce central electron densities of 5 x 10^20 m^-3 during 14 MW of neutral beam heating, while n_e(0)τ_e = 1.5 x 10^20 m^-3 sec has been obtained in ohmically-heated plasmas. During supershots, Lawson triple products of n_e(0)τ_eT_i(0) = 4.4 x 10^20 m^-3 sec keV are produced, exceeding the value expected in 1975 for TFTR. In 1980, the performance goal was increased to n_e(0)τ_e = 4 x 10^19 m^-3 sec at T_i(0) = 10 keV. Although these n_e(0)τ_e values have not been realized in high-temperature plasmas, TFTR has achieved the corresponding Lawson triple product at 27 keV ion temperature, which has the same QDT value.

The TEXTOR boronization technique has been employed on TFTR [3]. The benefit has been a reduction in oxygen impurities, which has greatly accelerated both startup after a major vacuum opening (one week instead of four weeks) and recovery after a major disruption. However, boronization has not reduced the carbon impurity level in supershots. The TFTR C/C limiter was conditioned for supershot experiments by running ~ 50 helium discharges at 1.4 MA which remove hydrogen and deuterium from the limiter. This produces a pumping limiter with a low recycling coefficient, R ~ 0.5, and a capacity for pumping the particle input from about 50 supershots (100 Torr-liters of injected deuterium).

The impurity content of TFTR plasmas is dominated by carbon, except soon after vacuum openings or disruptions when oxygen can become significant. Metallic impurities (Fe, Ni) can become important when the plasma minor radius is increased from 0.8 m to 0.96 m and the distance from the plasma to the metal walls is reduced. High power (30 MW) supershots have Z_{eff} = 2 - 3 at <n_e> ~ 4 x 10^19 m^-3: this level of impurities imposes a severe limit on the performance of supershots. The quality of supershots is correlated with the carbon level prior to beam injection. The carbon content can be reduced by by injecting Li pellets ~ 1 sec prior to the neutral beams [4] or intentionally generating a carbon bloom on the preceding shot by moving the plasma to a new portion of the limiter.

### 3.1 SUPERSHOT REGIME

In TFTR the highest performance plasmas are supershots with peaked-density profiles (Fig. 2), which have performance, n_e(0)T_i(0)τ_e, enhanced by a factor of ~ 20 over a comparable L-mode plasma, or factor of ~ 5 over the standard H-mode with a broad density profile. The total enhancement is multiplicative since the peaked-density profile increased confinement (x 3), central electron density and T_i simultaneously in supershot discharges. The hot-ion aspect of this regime was discovered on PLT [5] in 1978 and has been extended by TFTR at ion temperatures up to ~ 35 keV. Over the last four years, the supershot characteristics have been produced in plasmas with currents up to 2 MA and neutral beam power levels up to 32 MW.
The enhanced confinement of supershots is related to the peaking of the density profile, \( n_e(0)/\langle n_e \rangle \), as shown in Fig. 3. In plasmas with constant beam power, the enhancement over L-mode rises to \( \sim 3 \) as \( n_e(0)/\langle n_e \rangle \) increases to 3. The confinement enhancement has been further increased to \( \sim 3.5 \) by using very high power injection to build up the density profile and then reducing the beam power during the pulse. An important feature of the supershot regime is that the confinement time does not decrease with added heating power [6] in contrast with L-mode [2] and H-mode plasmas [7] where \( \tau_E \sim P_{\text{heat}}^{-1/2} \). This feature is also evident in the local transport coefficients for supershots and L-modes [8], and suggests that the basic mechanism causing transport is substantially modified in supershots relative to L-mode plasmas (see Section 4). The absence of confinement degradation with heating power has the important consequence that \( n_e(0)\tau_E T_i(0) \) increases rapidly with power at fixed plasma current in supershots (Fig. 4) until the beta limit is reached. Higher plasma currents allow the injection of higher power before this limit takes place. At the highest \( n_e T \) values, the central plasma pressure is near the MHD stability limit. Indeed, MHD-like internal modes with \( m/n = 1/1, 2/1, 3/2 \) and \( 4/3 \) are often observed concurrently with a broadening of the density profile and confinement degradation. The MHD activity can be affected by programming the plasma current evolution, conditioning the wall and limiters and modifying the current profile using the neutral beam and bootstrap driven current. As the plasma current is increased to stabilize pressure-driven instabilities, sawteeth instabilities are more difficult to suppress and degrade performance at \( \sim 2 \) MA. ICRH has been able to stabilize sawteeth and will be used to further extend the operating regime of supershot discharges [9].

### 3.2 PEAKED-PROFILE LIMITER H-MODE REGIME

The enhanced central confinement of the supershot has been combined with the enhanced edge confinement of a limiter H-mode to produce a confinement regime superior to the standard flat-profile divertor H-mode. The properties of the limiter...
FIG. 3. Enhancement of the energy confinement time over the prediction of L-mode scaling as a function of density profile peaking for a variety of supershots.

FIG. 4. Lawson triple product versus beam power for supershots.
FIG. 5. Characteristics of a limiter H-mode plasma in TFTR, with $I_p = 0.8$ MA, $B_T = 4$ T, and $P_{NB} = 11$ MW (balanced). The plasma parameters were $n_e(0) \approx 4.7 \times 10^{19}$ m$^{-3}$, $T_e(0) \approx 22$ keV and $\tau_e \approx 0.12$ sec.

H-mode plasma are shown in Fig. 5. After 0.4 sec of beam heating, the edge Ha drops while the edge electron density and temperature increase, as is characteristic of H-modes. The plasma core retains the properties of the supershot and confinement time is enhanced 2.5 times relative to the L-mode, with density peaking of 2.3 and $T_i/T_e \sim 2$ resulting in $n_e(0)\tau_e T_i(0)$ about four times that for a "standard" H-mode. The $n_e(0)\tau_e T_i(0)$ value is not higher than that in the supershot from which the H-mode evolved, but the favorable confinement is retained despite a moderate drop in the peaking factor, which arises when the edge density increases.

3.3 HIGH-DENSITY PELLET INJECTED REGIME

Since central particle fueling by neutral beams is not feasible in near-ignition reactor plasmas, different techniques are required to produce peaked-density profiles. Peaked-profile high-density plasma experiments were done on ALCATOR C, where confinement in ohmic plasmas was found to be near the neoclassical value [10] and enhanced by a factor of $\sim 2$ above ohmic L-mode. Subsequent experiments on TFTR have used injection of deuterium pellets into ohmic plasmas to obtain $n_e(0) \sim 3 \times 10^{20}$ m$^{-3}$ and $T_i(0) \sim 1.4$ keV with confinement times of $\sim 0.6$ sec which give $n_e(0)\tau_e \sim 1.5 \times 10^{20}$ m$^{-3}$ sec - about half the minimum value required for ignition. The confinement time in the TFTR plasmas was about twice the L-mode value [2]. Similar high-density plasmas heated with $\sim 2.5$ MW of ICRH using hydrogen minority had a central enhanced confinement region with transport coefficients reduced below
FIG. 6. Confinement and transport during the density decay following pellet injection at 3.0 sec. Hydrogen minority ICRF heating is applied at 3.05 sec. At 3.4 sec, $n_e(0) = 10^{20} \text{ m}^{-3}$, $T_e \sim 2.4 \text{ keV}$ in a 1.4 MA deuterium plasma. $\chi_{\text{eff}}$ is average of the electron and ion thermal conductivity.

those of L-mode [4], and global confinement $\sim 1.5$ L-mode (Fig. 6) when the density profile was peaked. The studies of this regime will be extended with the planned increase of ICRF heating from 7 MW to 12.5 MW, thereby allowing tests of confinement with higher temperature which will be more relevant to the BPX.

3.4 HIGH POLOIDAL BETA REGIME

High poloidal beta plasmas are of interest since they offer reactor concepts with the potential advantages of lower plasma current, significant bootstrap current drive and the possibility of improved confinement due to changes in the magnetic configuration. Plasmas with $\beta_B \sim 1.6$, near the equilibrium limit, and $\beta/(U_{\text{aB}}) \sim 4.7$ ($\% \text{ mT/MA}$) have been produced [11] in high temperature plasmas ($T_i \sim 20 \text{ keV}$, $T_e \sim 8 \text{ keV}$) with confinement times significantly enhanced ($\tau_E/\tau_L \sim 3.5$) using a current ramp-down during the ohmic phase of the discharge to modify the current profile prior to neutral beam injection. The range of parameters produced relative to previous supershot results is
FIG. 7. Summary of $\epsilon\beta_p$ achieved in TFTR as a function of $q^* = (5a^2B_p/R_0J_p)(1 + x^2 \times (1 + 28^2 - 1.25^2))/2$. Plasmas became diverted when $\epsilon\beta_p > 1.25$ and an apparent equilibrium limit was found at $\epsilon\beta_p \sim 1.6$.

FIG. 8. Comparison of the total $\chi_i$ including ion conduction and convection at $r = a/3$ versus $T_i (r = a/3)$ for L-mode plasmas and supershots.
shown in Fig. 7. At $\varepsilon \beta_p \geq 1.2$, the plasma cross-section becomes highly oblate ($\kappa \sim 0.7$) and a natural poloidal divertor forms at small major radius which is sustained for several energy confinement times. The central $q$ value, measured using polarimetry of the emission from an injected lithium pellet [12], was found to be $q_0 = 1.6 \pm 0.5$ in 275 kA plasmas ($9 \leq q \leq 12$). Stability calculations indicate that, while these plasmas are near the $\varepsilon \beta_p$ equilibrium limit, the plasmas are largely in the first stability regime with the outer part of the plasma near the predicted stability boundary for high-$n$ ballooning modes. As indicated in Fig. 7, the stored energy achieved in these plasmas is roughly a factor of two larger than previous supershots at a given plasma current over the full range of current studied (275 kA - 850 kA).

Previous analyses have shown that the measured surface voltage in TFTR is in agreement with the theory for beam driven and bootstrap current [13]. In the most collisionless high-$\beta_p$ plasmas, transport analysis shows that the bootstrap current accounts for about 65% of the plasma current for several hundred milliseconds during the heating pulse.

4. LOCAL TRANSPORT STUDIES

Local transport coefficients $X_i$, $X_e$ and $X_\phi$ have been determined using radial power balance for the plasma regimes described in Section 3. In L-mode plasmas, it is found that $X_i \sim X_\phi \sim (2 - 4) X_e$ throughout the radial profile for different plasma parameters, and the $X$-values scale inversely with $I_p$ as would be expected from the scaling of global confinement. As the electron temperature was varied [14] by increasing the neutral beam heating power at fixed density, the transport coefficients $X_i$, $X_e$ and $D$ varied as $T_e^x$ where $x = 1.5$ to 2.5. The relative magnitudes of the various $X$-values in the L-mode, especially the similarity of the momentum diffusivity $X_\phi$ with the thermal diffusivities, is in general agreement with the predictions of $E_xB$ transport due to drift wave turbulence, and would seem to be at variance with transport models based on magnetic fluctuations. Microwave scattering experiments [15] indicate that fluctuations with drift wave characteristics ($\omega \sim \omega^*, k_p \sim 1$) are present with amplitudes comparable to mixing length estimates, $\delta n/n \sim 1/(k\ell_n)$.

Since the global confinement in supershots does not degrade with power like L-mode confinement, it is instructive to compare the transport in the core of supershot and L-mode plasmas [16]. The ion channel is dominant for L-mode and all but the best supershots. The total $X_i$ (i.e., including both conduction and convection) at $r = a/3$ increases rapidly with $T_i$ ($r = a/3$) for L-mode plasmas, but decreases rapidly with temperature in the high temperature supershot regime as shown in Fig. 8. This feature is evidently a key factor in the performance of high $T_i$ supershots.

5. DEUTERIUM FUSION EXPERIMENTS

The present supershots produce significant D-D fusion reactions which are used as a measure of plasma performance. TFTR has an extensive neutron detection system: six detectors for flux determination, a ten channel radial profile system, one detector for fast neutron fluctuations, four detectors for D-T neutron flux measurements and a six-position foil-activation system. The D-D neutron detectors were calibrated in situ by placing a Cf$^{252}$ neutron source or a D-D and D-T neutron generator inside the TFTR vacuum vessel at many locations to provide an accurate determination of the multiple scattering of neutrons between the plasma source and the detectors. The accuracy of
these measurements is typically ~15%. The total D-D neutron rate is shown in Fig. 9 as a function of beam heating power for a variety of plasma conditions. The upper boundary of the points varies as \( P_{\text{heat}}^{1.8} \) where \( P_{\text{heat}} \) is the heating power. The maximum D-D fusion reaction rate (2 reactions per neutron) has increased to \( \sim 8.8 \times 10^{16} \text{ sec}^{-1} \). The maximum \( Q_{\text{DD}} = P_{\text{fusion}}/P_{\text{heat}} \) is \( 1.9 \times 10^{-3} \). The D-D fusion power for the best cases is \( \sim 50 \text{ kW} \) producing a yield of \( \sim 50 \text{ kJ} \) per pulse with a total of \( \sim 50 \text{ MJ} \) of D-D fusion energy having been produced during the lifetime of the TFTR experiments.

6. D-T OBJECTIVES

The original purpose of the TFTR D-T experiments was "The generation of reactor-like fusion power densities under conditions of approximate energy breakeven, in D-T operation", as described in the 1975 Final Conceptual Design Report. During the intervening years, the scope of the planned D-T experiments has expanded from a very brief demonstration of the use of D-T to broader studies of D-T physics issues. The present, expanded objectives are to:

- Study confinement of D-T plasmas, especially the isotope effect in tritium plasmas, and plasma heating with neutral beams and ICRF,
- Determine effects of alpha particles, such as single particle confinement, alpha driven instabilities, and first indications of alpha heating,
- Demonstrate D-T technical capability, such as tritium handling, recovery of tritium inventory in tokamak in-vessel components and operation of an activated machine, and
- Demonstrate D-T power production at fusion reactor power densities of \( \sim 1 \text{ MWm}^{-3} \) at approximate breakeven.
The initial D-T experiments will study the heating and confinement scaling of a deuterium-tritium plasma. A more critical issue for the success of subsequent burning-plasma experiments is the investigation of possible instabilities driven by the finite pressure of the alpha particles, such as the fishbone, the toroidal Alfvén eigenmode (TAE) [17-19], high-n ballooning modes and alpha-driven sawteeth. The theoretical thresholds for these instabilities depend on $\beta_G$ and $\nabla \beta_G$. Some aspects of these instabilities can be simulated using neutral beam injection to produce sufficient energetic particle beta using particles with parallel velocity exceeding the Alfvén velocity. A simulation experiment to excite the TAE modes via neutral beam injection has been carried out on TFTR, using up to 14 MW of nearly balanced deuterium beams with energy 70 - 110 kV injected into plasmas with $B = 1 - 1.2$ T, $q(a) \sim 3 - 4$, $<n_e> \sim 2.5 \times 10^{19}$ m$^{-3}$, and $<\beta> \sim 1\%$. Bursts of magnetic fluctuations near the TAE frequency appeared in the Mirnov coil signal when the beam ion velocity exceeded 70% of the Alfvén velocity. A mode with $n = 3$ has been identified. Growth of this mode was accompanied by a drop in neutron emission and a burst in $H_\alpha$ emission, indicative of expulsion of energetic ions from the plasma core. Very recent data from beam-emission spectroscopy has shown that density fluctuations with the same frequency exist within the plasma core ($r \sim a/2$). Detailed analysis to determine the radial mode structure needs to be completed. The $\beta_G$ expected for various regimes in TFTR is compared with theoretical estimates for alpha-driven instabilities in Table III and this comparison indicates that the TFTR D-T plasmas with the parameters anticipated will be able to test the theoretical predictions for the thresholds of essentially all of the major modes of instability, as well as the effect of the instabilities on the alphas and plasma.

### Table III. Estimates of Theoretical Thresholds for Alpha Collective Effects in TFTR Supershoot Plasmas

<table>
<thead>
<tr>
<th>INSTABILITY</th>
<th>$\beta_G (@ q = 1)$</th>
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<tbody>
<tr>
<td>FISHBONES</td>
<td>~ 1%</td>
</tr>
<tr>
<td>SAWTOOTH STABILIZATION</td>
<td>~ 0.5%</td>
</tr>
<tr>
<td>TOROIDAL ALFVEN (GLOBAL MODEL)</td>
<td>~ 0.3%</td>
</tr>
<tr>
<td>HIGH-n BALLOONING</td>
<td>~ 0.1%</td>
</tr>
<tr>
<td>TOROIDAL ALFVEN (LOCAL MODEL)</td>
<td>~ 0.02%</td>
</tr>
</tbody>
</table>

### 7. Projection of D-T Plasma Performance

The D-T performance of TFTR has been estimated using codes which have been benchmarked against many D-D plasmas. The calculated stored plasma energy and neutron production rate for D-D supershots are consistent with the measured values to
within their uncertainties. The observed dependences of the D-D neutron rate and QDD on beam power and total plasma energy are reproduced by the codes. In the D-T simulations, the measured temperature and electron density profiles were scaled uniformly to account for changes in beam power and expected confinement. The normalization of the temperature profiles is determined by matching the expected TE while maintaining $X_\| = X_\perp$ at $r = a/2$ (as in the best discharges). The power balance at $r = a/2$ assumes a convective multiplier of 1.5 and includes alpha particle heating. The fusion gain is taken to be $Q_{DT} = P_{\text{fusion}}/P_{\text{heat}}$ where $P_{\text{heat}}$ is the total auxiliary heating power. No reduction in the auxiliary heating power was made to correct for shinethrough, fast-ion orbit losses or $dW_p/dt$.

The performance of equivalent D-T discharges was simulated by injecting equal power $D^0$ and $T^0$ beams into a 50/50 D/T plasma using the measured temperatures and electron density profiles, $Z_{\text{eff}}$, beam power, voltage, and species mix of a comparable D-D shot (Table IV). Equivalent $Q_{DT}$ values up to $\sim 0.31(\sim 185 \text{ QDD})$ were obtained in this way. In this case the thermonuclear reactions contributed 42% while beam-target and beam-beam were 51% and 7% respectively. Thus, for this mode of operation, $M = 2.4$.

Significantly more fusion power and alpha production is expected if the best discharge conditions are extrapolated to anticipated D-T conditions at full beam power. In these extrapolated cases, the total neutral beam power was increased to 35 MW while the

<table>
<thead>
<tr>
<th>TABLE IV: PROJECTIONS OF D-T PERFORMANCE</th>
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<tbody>
<tr>
<td>SHOT 53848</td>
</tr>
<tr>
<td>Pb (MW)</td>
</tr>
<tr>
<td>$\bar{n}_e$ ($10^{19}$m$^{-3}$)</td>
</tr>
<tr>
<td>$Z_{\text{eff}}$</td>
</tr>
<tr>
<td>$t\Gamma$ (sec)</td>
</tr>
<tr>
<td>$T_{e0}$ (kev)</td>
</tr>
<tr>
<td>$T_{i0}$ (kev)</td>
</tr>
<tr>
<td>beam-beam (%)</td>
</tr>
<tr>
<td>beam-target (%)</td>
</tr>
<tr>
<td>Thermonuclear (%)</td>
</tr>
<tr>
<td>$&lt;\beta_{\alpha}&gt;$ (core) (%)</td>
</tr>
<tr>
<td>$Q_{DT}$</td>
</tr>
<tr>
<td>$P_{\text{fusion}}$ (MW)</td>
</tr>
<tr>
<td>$P_{\alpha}$ (MW)</td>
</tr>
</tbody>
</table>
beam voltage was increased to 120 keV and it was assumed that $\tau_\text{E} \sim A_1^{0.5}$. The electron density was increased by 25%, according to present scaling with power, and $Z_{\text{eff}}$ was reduced following the observed trend with density. For the full power case with roughly 50/50 D$^0$/T$^0$ beams into a roughly 50/50 D/T plasma, the extrapolated $Q^{*}\text{DT}$ is in the range 0.5 (300 QDD). Thermonuclear reactions contribute 50% of the rate, while beam-target and beam-beam are 42% and 8% respectively, as shown in Table IV. These plasmas have central $\beta_\alpha$ in the range of $\sim 0.7\%$ which is above the theoretical estimate of 0.3% for the global TAE threshold of 0.3%.

The largest $Q$ and fusion power are obtained using the original TFTR D-T scenario which maximizes beam-target reactions by injecting 35 MW of D$^0$ beams into a tritium plasma. The target tritium plasma could be preheated by ICRF or $\sim 10$ MW of tritium neutral beams (in the latter case the D$^0$ beams are at 25 MW). The fusion reactions are predominantly beam-target reactions and were projected in 1975 to give $Q$ in the breakeven regime of 0.3 to 1.0. Present projections for this mode of operation, based on the best D-D supershots, indicate that $Q^{*}\text{DT} = 0.4 - 0.7$ could be achieved. In this case the beam-target reactions contribute 89% while thermonuclear reactions are 11% of the total rate. These will be transient cases, since the plasma composition will evolve to that of the beams, but the high-$Q$ conditions are expected to last for times of order $\tau_\text{E}$. The central $\beta_\alpha$ value for this case is $\sim 0.9\%$ (Table IV).

The addition of 12 MW of ICRF heating is expected to stabilize sawtooth instabilities and, through heating the electrons will increase $\beta_\alpha$, thereby expanding the range of study for alpha-driven instabilities.

8. FUTURE PLANS

Experiments over the next two years will concentrate on improving the performance of supershots, especially at higher density, using ICRF and current profile control along with improved wall conditioning techniques. The ICRF heating power will be doubled to 12.5 MW and will be used for sawtooth and MHD-mode control, as well as to extend the high-density peaked-profile regime in preparation for D-T experiments to a BPX-relevant high-temperature plasma scenario. The transport studies will continue with emphasis on correlating fluctuations with transport and theoretical models. In addition, activities will increase to prepare for D-T experiments in 1993-1994.

9. SUMMARY

TFTR has developed plasma regimes with deuterium plasma parameters that exceed the TFTR goals for temperature and $n_\text{e}(0)/\tau_\text{E}T_1(0)$. These deuterium plasma conditions project to D-T plasmas in the breakeven regime with fusion powers of 15 - 25 MW and alpha particle parameters suitable for testing collective alpha instability theories. The TFTR enhanced performance plasma regimes and D-T results will provide important data for the design and operation of BPX, ITER and ARIES.
ACKNOWLEDGMENTS

We wish to acknowledge the dedicated efforts of the TFTR staff in the support of the experiments reported here. The contributions of collaborators from many other institutions have been critical to the success of these experiments.

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REFERENCES


DISCUSSION

G. GRIEGER: I have a question concerning the supershots and the low ion heat conductivity observed under such conditions. Is anything known about radial electric fields or poloidal rotation when the effect occurs?
D.M. MEADE: The supershot performance is best when the toroidal rotation is near zero, that is, close to balanced co- and counter-injection. Poloidal rotation, as determined from the frequency shift of edge fluctuations, is small compared to the electron diamagnetic drift velocity.

A. GIBSON: The highest triple product value of $4.4 \times 10^{20} \text{m}^{-3} \cdot \text{s} \cdot \text{keV}$ which you quoted for TFTR was given in terms of $n_e$. What was the $n_D/n_e$ ratio for this case?

D.M. MEADE: The ratio is $\sim 0.75$ for the best supershots.

A.K. SEN: The peaked density profiles in supershot regimes necessarily indicate lower $\eta_i$ and $\eta_e$. Do you observe lower levels of electrostatic fluctuations under those conditions?

D.M. MEADE: Electrostatic fluctuations in the plasma interior have been examined on TFTR with microwave scattering and beam emission spectroscopy. Microwave scattering is sensitive to $2 < k_\theta < 10 \text{ cm}^{-1}$ and has found $k$ spectra $\sim k^{-3}$ for both L mode and supershot plasmas. The estimate of $\delta n/n$ requires a number of simplifying assumptions, but roughly indicates that $\delta n/n \sim \langle k \rangle L_n^{-1}$ and, within a factor of 2–4, is the same for L mode and supershots. Beam emission spectroscopic estimates are sensitive to $0 < k_\perp < 1.5 \text{ cm}^{-1}$ and show some internal, poloidally coherent oscillation in supershots with $\delta n/n = 1.0$–1.5%. In L mode discharges, fluctuations are below the limit of sensitivity ($\delta n/n = 0.5\%$) except at the edge.

S.I. ITOH: What sort of $Z_{\text{eff}}$ scaling (dilution) did you find with density? (You clearly used some kind of scaling for the projection of TFTR performance.) And what about the time derivative of $Z_{\text{eff}}$, $dZ_{\text{eff}}/dt$? This is a key parameter for long pulse operation, in addition to $dP_{\text{rad}}/dt$.

D.M. MEADE: $Z_{\text{eff}}$ decreases with increasing density in TFTR supershots and this scaling was used for the D–T projections. $Z_{\text{eff}}$ reaches equilibrium rapidly, and $dZ/dt = 0$ for high performance supershots.
RECENT JET RESULTS
AND FUTURE PROSPECTS

JET TEAM*
(Presented by P.H. Rebut)

JET Joint Undertaking,
Abingdon, Oxfordshire,
United Kingdom

Abstract

RECENT JET RESULTS AND FUTURE PROSPECTS.

The latest results of JET plasmas in transient and steady states are presented. Substantial improvements in plasma purity and corresponding reductions in plasma dilution have resulted from the use of beryllium as the first wall material facing the hot plasma. As a consequence, plasmas with a fusion triple product \( n_0(T) \tau_T \tau_E \) in the range \((8 - 9) \times 10^{20} \text{ m}^{-3} \cdot \text{s} \cdot \text{keV}\) have been achieved (within a factor of 8 of that required in a fusion reactor), albeit under transient conditions. The general JET performance has also improved, allowing the parameters of a reactor plasma to be individually achieved in JET. In view of their importance for reactors, the JET results are presented with particular emphasis on their significance for the formulation of a plasma model for the Next Step. However, impurity influxes limit the attainment of better parameters and prevent the realisation of steady state conditions at high heating powers. To address this problem of impurity control, and those of plasma fuelling and helium ash exhaust, a New Phase is planned for JET with an axi-symmetric pumped divertor configuration that will allow operating conditions close to those of a reactor. The divertor configuration should demonstrate a concept of impurity control and determine the size and geometry needed to fulfil this concept in a reactor. It should identify appropriate materials for plasma facing components and define the operational domain for the Next Step.

1. INTRODUCTION

The objective of JET is to obtain and study plasmas in conditions and dimensions approaching those needed in a thermonuclear reactor [1,2]. The present paper concentrates on progress towards this objective during the last two years of JET operation, since the 1988 IAEA Conference [3]. The transient and steady state behaviour of JET plasmas is presented in view of their importance for reactor plasmas. JET operation and the consequences of using a beryllium "first wall" are discussed and the best fusion performance and general plasma behaviour in JET are reported. The underlying results are presented with particular emphasis on their significance for the formulation of a plasma model for a Next Step tokamak. In view of the importance of dilution and exhaust for ignition [4] and the need for adequate impurity control and understanding of the scrape-off layer plasma, a New Phase is planned for JET with a divertor configuration and this is also discussed.

* See Appendix I.
2. JET OPERATION

Since 1988, further additions and technical enhancements to JET have been made:

• reinforcement of the vacuum vessel to withstand radial and vertical instabilities and permit operation at full power up to 7MA in the material limiter configuration and up to 6MA in the X-point configuration (inconel rings fitted above and below the mid-plane at the in-board walls to stiffen and strengthen the vessel);

• belt limiter tiles, ion cyclotron resonance heating (ICRH) antenna screens and lower X-point dump plates were changed to beryllium and beryllium evaporation has also been used;

• one neutral beam injection (NBI) box was modified to operate at 140kV to provide better penetration at higher density. The other NBI box remained at 80kV. This gave a total power of 18MW (instead of 21MW with both boxes at 80kV);

Table I
JET Parameters

<table>
<thead>
<tr>
<th>Parameters</th>
<th>Design Values</th>
<th>Achieved Values</th>
</tr>
</thead>
<tbody>
<tr>
<td>Plasma Major Radius (R₀)</td>
<td>2.96m</td>
<td>2.5-3.4m</td>
</tr>
<tr>
<td>Plasma Minor Radius</td>
<td></td>
<td></td>
</tr>
<tr>
<td>horizontal (a)</td>
<td>1.25m</td>
<td>0.8-1.2m</td>
</tr>
<tr>
<td>Plasma Minor Radius</td>
<td></td>
<td></td>
</tr>
<tr>
<td>vertical (b)</td>
<td>2.1m</td>
<td>0.8-2.1m</td>
</tr>
<tr>
<td>Toroidal Field at R₀</td>
<td>3.45T</td>
<td>3.45T</td>
</tr>
<tr>
<td>Plasma Current:</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Limiter Mode</td>
<td>4.8MA</td>
<td>7.1MA</td>
</tr>
<tr>
<td>Single null X-point</td>
<td>not foreseen</td>
<td>5.1MA</td>
</tr>
<tr>
<td>Double null X-point</td>
<td>not foreseen</td>
<td>4.5MA</td>
</tr>
<tr>
<td>Neutral Beam (NB) Power</td>
<td></td>
<td></td>
</tr>
<tr>
<td>(80kV, D)</td>
<td>20MW</td>
<td>21MW</td>
</tr>
<tr>
<td>(140kV, D)</td>
<td>15MW</td>
<td>8MW</td>
</tr>
<tr>
<td>(one box converted, so far)</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Ion Cyclotron Resonance Heating (ICRH) Power to Plasma</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>15MW</td>
<td>22MW</td>
</tr>
</tbody>
</table>
each ICRH generator power unit was upgraded to 2MW, offering the potential of 32MW total power source and ~24MW coupled to the plasma through eight antennas;

- a prototype lower hybrid current drive (LHCD) system has been installed, offering the potential of 4MW at 3.7GHz;

In addition, a faulty toroidal field coil has been removed and replaced successfully with a spare coil.

JET is now about midway through its experimental programme. The technical design specification of JET has been achieved in all parameters and exceeded in several cases (see Table I). The plasma current of 7MA in the limiter configuration [5] and the current duration of up to 30s at 3MA are world records and are over twice the values achieved in any other fusion experiment. 5.1MA and 4.5MA are also world records in the single-null and double-null divertor configurations, respectively [6]. NBI heating has been brought up to full power (~21MW) and ICRH power has been increased to ~22MW in the plasma. In combination, these systems have delivered 35MW to the plasma.

3. THE USE OF BERYLLIUM IN JET

Over the last two years impurities and density control have been the main obstacles to the improvement of JET performance. Carbon first-wall components had been developed so that they were mechanically able to withstand the power loads encountered. However, the interaction of the plasma with these components, even under quiescent conditions, caused unacceptable dilution of the plasma fuel. In addition, imperfections in the positioning of the components led to localised heating at high power, and the following problems occurred:

- the production of impurities increased with the input power to the plasma;
- at high power, the heat load on the tiles caused a plasma evolution which exhibited a catastrophic behaviour - the so-called "carbon catastrophe". Increased plasma dilution, increased power radiated, reduced neutral beam penetration and a threefold fall of fusion yield resulted from the carbon influx;
- for lower input power with long duration, problems were also encountered. Without fuelling, deuterium was pumped by the carbon and replaced by impurities, resulting in severe dilution of the plasma;
- the maximum density achieved without the occurrence of plasma disruptions appeared to be limited by edge radiation.

The situation has been redressed by the progressive introduction of beryllium "first-wall" components since 1989 [7]. First, beryllium was evaporated as a thin layer on the carbon walls and limiters; then, as the material for the limiter tiles; and finally, as the material for the lower X-point target tiles and the open screens of the ICRH antennas.

With a carbon first-wall, the main impurities were carbon (2-10%) and oxygen (1-2%). With beryllium evaporated inside the vessel, oxygen was reduced by factors >20, and carbon by >2. Although beryllium increased, carbon remained the dominant impurity for this phase. With beryllium limiters, the carbon concentration was reduced by a further factor of 10, but beryllium levels increased by ~10, and became the dominant impurity. Due to the virtual
FIG. 1. (a) Dilution factor, \( n_{\text{b}}/n_{\text{e}} \), and (b) the effective charge, \( Z_{\text{eff}} \), as functions of power per particle \( P_{f}/\langle n_{\text{e}} \rangle \) for carbon limiter tiles, beryllium gettering and beryllium limiter tiles.
elimination of oxygen and replacement of carbon by beryllium, impurity influxes were reduced significantly, in line with model calculations [8] which take account of impurity self-sputtering. In addition, nickel was eliminated from the plasma when the nickel screens for the ICRH antennas were replaced by beryllium.

During 1989, plasma dilution and the effective plasma charge, $Z_{\text{eff}}$, were reduced significantly in ohmic plasmas and with strong additional heating. Fig. 1(a) shows the dilution factor, $n_D/n_e$, as a function of input power per particle, $P_I/n_e$. The corresponding values of $Z_{\text{eff}}$ are shown in Fig. 1(b). With moderate power, it was not possible to maintain $n_D/n_e$ much above 0.6 with carbon, but values greater than 0.8 were routinely achieved with beryllium. Furthermore, high power operation was possible only with beryllium.

Impurity radiation was also reduced and operation with beryllium gettering allowed improved density control (due to high wall pumping of both deuterium and helium). On the longer timescale (minutes to hours), very little deuterium was retained compared with a carbon first-wall; >80% of the neutral gas admitted to JET is recovered, compared with ~50% with a carbon first-wall. This has important advantages for the tritium phase of JET operation.

4. JET PERFORMANCE

4.1. Fusion performance

With carbon X-point target plates, the length of the H-mode has been extended (up to 5.3s) either by sweeping the X-point (both in the radial and vertical directions) to reduce the X-point tile temperature, or by using strong gas puffing in the divertor region. This, together with the better plasma purity achieved with a beryllium first-wall, resulted in increased ion temperatures ($T_i(0)$ in the range 20-30keV) and improved plasma performance, with the fusion triple product ($n_D(0)\tau_E T_i(0)$) increasing significantly. Such improved fusion performance could otherwise have been achieved only with a substantial increase in energy confinement.

In a particular case, the central ion temperature reached 22keV, the energy confinement time, $\tau_E$, was 1.1s, with a record fusion triple product ($n_D(0)\tau_E T_i(0)$) of $8-9\times10^{20}$m$^{-3}$sk eV. The neutron yield for this discharge was also amongst the highest achieved on JET at $3.5\times10^{16}$ns$^{-1}$, with $Q_{\text{DD}}=2.4\times10^{-3}$. A full D-T simulation of the pulse showed that 12MW of fusion power would have been obtained transiently with 16MW of NBI power, giving an equivalent fusion amplification factor $Q_{DT} \sim 0.8$, reaching near breakeven conditions and within a factor of 8 of that required by a reactor. Similar results were also obtained at medium temperatures, with $T_e\sim T_i\sim 10$keV.

The overall fusion triple product as a function of central ion temperature is shown in Fig. 2 for a number of tokamaks.

4.2. General behaviour

Reduced impurity levels allowed prolonged operation at higher densities and improved the general JET performance, as follows:

- the pumping of deuterium with a beryllium first-wall was more efficient than with a carbon first-wall and provided improved density control. This
permitted low density and high temperature (up to 30keV) operation for times >1s;
• the density limit increased [9], and a peak density of $4 \times 10^{20} \text{m}^{-3}$ was achieved with pellet fuelling. The density is limited principally by fuelling and not by disruptions, as was found with carbon limiters;
• sawtooth free periods exceeding 5s were achieved, but the stabilisation mechanism is still not yet clear [10]. The central electron temperature appears to saturate at about 12keV, even though the central heating power to the electrons can be higher than that to the ions;
- H-modes were established with ICRH alone and for periods >1s. With beryllium antenna screens, H-modes were established with either monopole or dipole phasing [11]. The confinement characteristics of ICRH H-modes were similar to those with NBI alone;
- $\beta$-values up to the Troyon limit were obtained in double-null X-point plasmas [9].

Thus, the parameters of a reactor plasma have been achieved individually in JET.

However, the best fusion performance was obtained in a transient state and could not be sustained in steady state. Ultimately, the influx of impurities caused a degradation in plasma parameters. Furthermore, a severe carbon influx ("carbon catastrophe") was still a problem for inner wall and X-point operation and is a serious limitation in H-mode studies.

5. UNDERLYING RESULTS AND THEIR SIGNIFICANCE

The underlying JET results are presented with particular emphasis on their significance for the formulation of a plasma model for the Next Step.

\[ n_e (10^{19} \text{m}^{-3}) \]

![Electron density profiles for different fuelling and heating methods.](image)

**FIG. 3.** Electron density profiles for different fuelling and heating methods.
5.1. Density limit

With a carbon first-wall, the plasma density was limited, in general, when the radiated power reached 100% of the input power. This led to the growth of MHD instabilities and ended in a major disruption. The density limit was dependent on plasma purity and power to the plasma.

With a beryllium first-wall, the maximum operating density increased significantly by a factor of 1.6-2. A record central density of $4 \times 10^{20} \text{m}^{-3}$ was achieved by strongly peaking the density profile using a sequence of 4mm solid deuterium pellets injected at intervals throughout the current rise phase of an X-point discharge. Furthermore, the nature of the density limit changed and the frequency of disruptions at the density limit was much reduced. Disruptions did not usually occur, and the limit was associated rather with the formation of a poloidally asymmetric, but toroidally symmetric radiating structure (a "MARFE"), which limits the plasma density to within the stable operating...
domain. These results constitute a substantial enhancement of JET's operating capability.

Heating and fuelling were varied systematically, using both gas and pellet fuelling. With deep pellet fuelling and either NBI or ICRH, peaked profiles were obtained (Fig. 3). Just before a density limit MARFE occurred, pellet fuelled discharges reached the same edge density as gas fuelled discharges, but the central densities were considerably higher. The central density depends, therefore, on the fuelling method used. The profiles are similar near the edge, but are remarkably flat with gas fuelling.

These observations suggest that the edge density may be correlated with the density limit and is found to increase approximately as the square root of power (Fig. 4). This endorses the view that the density limit is determined by a power balance at the plasma edge and the cause of disruptions is related to radiation near the q=2 surface. Thus, under beryllium conditions, when the radiation is low, or confined to the outermost edge, there are no density limit disruptions.

\[ n_e(D) \times 10^{19} \text{ m}^{-3} \]

\[ \tau_{\text{decays}} = 1.8 \text{ s} \]

\[ 3 \text{MA/3.1T} \]

\[ \text{Pellet Time (s)} \]

\[ \text{FIG. 5. Decay of central density following pellet injection into discharges with (a) ohmic heating only and (b) ~8 MW ICRH.} \]
FIG. 6. (a) Nickel density profiles from X-ray tomography at times following laser ablation and (b) normalised particle fluxes versus normalised density gradients for nickel impurities at different plasma radii.
FIG. 7.  (a) Temporal evolution of Ni XXVI emission in (a) the L-phase and (b) the H-phase of two similar discharges with ~9 MW of additional heating.

5.2. Density profiles and transport

Of significance also are the density profiles obtained with edge fuelling, which tend to be flatter with the lower $Z_{\text{eff}}$ achieved with beryllium, in contrast to those obtained with carbon, which tended to be more peaked, even with edge fuelling. The occurrence of flat density profiles suggests that there is no need for an anomalous inward particle pinch, except perhaps on impurities. This observation poses important questions related to particle transport, and in particular, the transport and exhaust of helium ash products.

The relaxation of the peaked density profiles achieved with pellet injection allows an estimate of particle transport. For a 4mm pellet injected into 3MA/3.1T plasma, the decay of the central electron density is shown in Fig. 5. Following injection, the decay constants are 1.8s for the ohmically heated discharge and 1s when ~8MW ICRH is applied. The global energy confinement times are in a similar ratio. It is therefore reasonable to assume that particle and energy transport are linked. Furthermore, modelling studies of similar discharges suggest that the diffusion coefficient is lower in the central plasma than further out and that there is no need for a large anomalous inward particle pinch in the central plasma [12].
Impurity transport studies have been possible from the measurement of emissivity profiles by the soft X-ray cameras following the injection of laser-ablated, high-Z impurities. The evolution of the nickel density profile has been determined (Fig. 6(a)) and so the particle fluxes can be plotted as a function of the density gradient at different radii. Fig. 6(b) shows that fluxes and gradients are linearly related through a diffusion coefficient which increases with radius. There is no evidence of a significant pinch term for the impurities.

This measurement also provides evidence of better confinement in H-modes. The temporal evolution of NiXXVI emission is shown in Fig. 7 for the L- and H-phases of two similar discharges with ~9MW of additional heating. In contrast to the decaying signal of the L-phase, the signal rises rapidly to a steady value which persists to the end of the H-phase. This shows that impurities have considerably longer confinement times in the H-phase and endorses the view that an edge transport barrier exists, which could be destroyed (for example, by ELMs) on transition from the H- to the L-phase.

5.3. Temperature

High ion temperatures have been obtained at the low densities possible with a beryllium first-wall and with the better penetration afforded by NBI at 140kV.

FIG. 8. Central ion ($T_i$) and electron ($T_e$) temperatures as functions of power per particle $(P/n)_{ei}$ to either species.
FIG. 9. Temporal evolution of electron temperature perturbations (normalised to the central electron temperature before the collapse of a sawtooth) at different radii for 3 MA/2.8 T discharges with (a) ohmic heating only and (b) 9.5 MW ICRH. Dashed lines are from model calculations using $\chi_{\text{HP}} = 3.2 \text{ m}^2 \cdot \text{s}^{-1}$. 
Record ion temperatures were achieved of up to 18keV in limiter plasmas and up to 30keV in X-point plasmas (with powers up to 17MW). In this mode, the ion temperature profile is sharply peaked and the electron temperature is significantly lower than the ion temperature, by a factor of 2-3. The central ion temperature (as shown in Fig. 8) increases approximately linearly with power per particle up to the highest temperatures, indicating that ion thermal losses are anomalous, but ion confinement degrades little with input power. On the other hand, the central electron temperature saturates at ~12keV, even though with ICRH the central heating power to the electrons can be higher than that to the ions. Electron thermal transport is also anomalous and electron confinement degrades strongly with increased heating power. This suggests that electrons are primarily responsible for confinement degradation.

At higher densities with combined NBI and ICRH, central ion and electron temperatures were both above 11keV in a 3MA plasma for power input of 33MW (21 MW NBI and 12MW ICRH).

Extensive studies have also been performed in the 'monster-sawtooth' regime [10] in which sawtooth oscillations have been suppressed for up to 5s by central ICRH. Peaked temperature profiles (with both central ion and electron temperatures above 10keV) were maintained for several seconds. In an equivalent D-T mixture, this would result in a significant enhancement in the time-averaged neutron rate compared with a sawtoothing discharge. This does not mean, however, that ion losses are necessarily small.

5.4. Electron heat pulse propagation

The propagation of temperature perturbations (determined from the electron cyclotron polychromator) and density perturbations (determined from the multichannel reflectometer) following the collapse of a sawtooth provide good measurements of energy and particle transport. The decay of the temperature perturbation at different radii in a 3MA/3.1T ohmically heated discharge is shown in Fig. 9(a). This decay can be modelled with a heat pulse diffusivity, $\chi_{HP} \approx 3.2m^2s^{-1}$, which should be compared with $\chi_e \approx 1m^2s^{-1}$, obtained from power balance considerations. The results in an L-mode plasma, heated with 9.5MW of ICRH, are shown in Fig.9(b) and indicate that, although $\chi_e \approx 2m^2s^{-1}$, the same $\chi_{HP} \approx 3.2m^2s^{-1}$ can be used in the simulation to fit the data. It is also found that, within experimental uncertainties, the same $\chi_{HP}$ can be used also for H-regime plasmas and does not depend on heating power.

The propagation characteristics of the density perturbation indicate that the density pulse is slower than the temperature pulse and that the density pulse is comprised of both an outward and inward propagating perturbation (the latter resulting from the earlier interaction of the temperature pulse with the limiters) [12].

Furthermore, simultaneous measurements of the temperature and density perturbations indicate that the particle pulse diffusion coefficient, $D_{DP} \approx D_e \times \chi_{HP}$.

5.5. Global energy confinement

With a carbon first-wall, the energy confinement time improves with increasing current and degrades with increasing heating power, independent of the heating
FIG. 10. Global energy confinement time ($\tau_E$) during the H-mode as a function of net input power for different plasma currents and first wall materials.

Method. With a beryllium first-wall, energy confinement times and their dependences are effectively unchanged: energy confinement does not appear to be affected by the impurity mix (carbon or beryllium in deuterium plasmas).

In the X-point configuration, high power H-modes (up to 25MW) have been studied. In comparison with limiter plasmas, confinement is a factor ~2 better, but the dependences with current and heating power are similar (Fig. 10).

With a carbon first-wall, H-modes with ICRH alone were not obtained. Beryllium evaporation on the nickel antenna screens led to lower impurity production and H-modes were successfully obtained with ICRH alone. With beryllium antenna screens, the threshold for the H-mode was reduced somewhat for dipole phasing, and ICRH H-modes were also obtained with monopole phasing. However, this required feedback control of the plasma position to allow for the movement of the plasma boundary during the L-H transition. In all cases, H-mode confinement with ICRH alone was similar to that with NBI, that is independent of the heating method.

A particular discharge (Fig. 11) exhibited an H-mode and pellet enhanced performance [11]. A 4mm pellet was injected into a 3MA/2.8T double null X-point plasma, heated with 9MW of ICRH and 2.5MW of NBI (which served primarily as a diagnostic for the measurement of the ion temperature). The
stored plasma energy increased to 8MJ (still increasing at a rate \( \sim 4\text{MW} \) when the period of pellet enhanced confinement ceased). The energy confinement time reached \( \sim 1.0\text{s} \) and with central electron density \( \sim 8 \times 10^{19} \text{m}^{-3} \), \( Z_{\text{eff}} \sim 1 \) and central electron and ion temperatures \( \sim 10\text{keV} \), the fusion triple product was \( \sim 8 \times 10^{20} \text{m}^{-3}\text{skeV} \). The neutron yield was \( 10^{16} \text{s}^{-1} \). This phase terminated as the central density decayed, although the stored energy remained high (again favouring the existence of an edge confinement barrier) until the end of the H-mode. Subsequently, plasma temperatures recovered, with the bulk deuterium ions being effectively heated in this ICRH scheme which used a high concentration of hydrogen minority ions.
5.6. Beta limits

Experiments have explored the plasma pressure (indicated by the $\beta$-value) that can be sustained in JET and investigated the plasma behaviour near the expected $\beta$-limit in a double-null H-mode configuration, at high density and temperature and low magnetic field ($B_t = 1T$). $\beta$ values up to $\sim 5.5\%$ were obtained, close to the Troyon limit $\beta_t(\%) = 2.8I_p(MA)/B_t(T)a(m)$, where $I_p$ is the plasma current and $a$ is the plasma minor radius [9]. Significantly, the JET limit does not appear to be disruptive at present power levels. Rather, a range of MHD instabilities occurs, limiting the maximum $\beta$-value without causing a disruption. The behaviour near both the density and $\beta$-limits may be interpreted in terms of resonant instabilities which have the magnetic topology of an island.

5.7. Alpha-particle simulations

The behaviour of alpha-particles has been simulated in JET by studying energetic particles such as 1MeV tritons, and $^3$He and H minority ions accelerated to a few MeV by ICRH [11]. The energetic population has up to 50% of the stored energy of the plasma and possesses all the characteristics of alpha-particles in an ignited plasma, except that in the JET experiments, the ratio of the perpendicular to parallel pressure was above three, while in a reactor plasma the distribution will be approximately isotropic. The mean energy of the minority species was about 1MeV, and the relative concentration of the $^3$He ions to the electron density was 1-2%, which is comparable to the relative concentration of alpha-particles in a reactor (7%). Under conditions with little MHD activity, no evidence of non-classical loss or deleterious behaviour of minority ions was observed, even though the ratio of the fast ion slowing down time to the energy confinement time in JET is greater than that expected in a reactor.

Fusion reactivity measurements were undertaken on the $D$-$^3$He reaction when minority $^3$He ions were accelerated to energies in the MeV range using ICRH. During a 5s monster sawtooth produced by 10MW of ICRH, a reaction rate of $4 \times 10^{16}s^{-1}$ was achieved, which corresponded to 100kW of fusion power and $Q=1\%$ was reached. This was carried out with a beryllium first-wall and benefitted from an improved fuel concentration $n_p/n_e$ of up to 0.7. Comparison of the measurements with theoretical predictions suggests nearly classical trapping and thermalisation of the energetic particles.

6. PLASMA MODELLING

6.1. Formulation of a plasma model

Any model used to predict the performance of a Next Step tokamak must be consistent with the foregoing data. In particular, it must explain:

- the resilience of the electron temperature to heating;
- the heat and density pulse propagation studies;
• no intrinsic degradation of ion confinement with ion heating;
• the density decay after pellet injection; and
• the similar behaviour of particle and heat transport.

It is possible to formulate a transport model, based on one phenomenon and consistent with JET data and with physics constraints. Specifically, above a critical threshold, \((\nabla T_e)_c\), in the electron temperature gradient, the transport is anomalous and greater than the underlying neoclassical transport. The electrons are primarily responsible for the anomalous transport, but ion heat and particle transport are also anomalous. The general expressions for the conductive heat fluxes and the anomalous transport coefficients are:

\[
Q_e = -n_e\chi_e \nabla T_e = -n_e\chi_{an,e} (\nabla T_e - (\nabla T_e)_c) \\
Q_i = -n_i\chi_i \nabla T_i \quad \text{and} \quad \chi_i = 2\chi_e \frac{Z_i}{\sqrt{1+Z_{eff}}} \sqrt{\frac{T_e}{T_i}} \\
D_i = (0.3+0.5) \chi_i
\]

The critical electron temperature gradient model of Rebut et al [13] specifies possible dependences for \(\chi\) and \(D\) and this is explored further in [12].

6.2. To model a Next Step tokamak

The fuelling, impurity control and exhaust capability of a Next Step tokamak will be dependent on whether deuterium and impurities (including helium) accumulate in the plasma centre. The production and transport of helium ash towards the plasma edge (where it must be exhausted) will depend on the relative importance of energy and particle confinement, the effect of sawteeth, the effect of the edge transport barrier in the H-mode and the behaviour of the scrape-off-layer plasma.

Although the transport model of Section 6.1 applies in both L- and H-regimes, and ensures that particle and energy transport will follow each other in the transition from the L- to H-regime (since \(\chi\) and \(D\) are linked in the model), it is necessary to include several additional elements in order to complete the model for a Next Step tokamak. For example, the spontaneous improvement in edge confinement has yet to be modelled. The reduction in MHD activity observed experimentally suggests the presence of some other instability at the edge of L-mode plasmas, where the effect of atomic physics on MHD might be important. This is apparently easier to suppress in an X-point configuration with high edge magnetic or rotational shear. Furthermore, an understanding of the scrape-off-layer (SOL) plasma is also needed.

7. IMPURITY CONTROL AND THE NEW PHASE PLANNED FOR JET

Achieving control of the impurity influx into the plasma is a pre-requisite to the construction of a tokamak reactor. In the case of high \(Z\) impurities, radiation
losses may prevent reaching ignition. The presence of low Z impurities, in addition to helium produced by nuclear reactions, dilutes the concentration of reacting ions and therefore reduces the alpha-particle power. Under present conditions, the lifetime of the plasma facing components would be severely limited. Understanding the SOL plasma is needed because of the importance of dilution and exhaust for ignition.

So far, JET has concentrated on passive methods of impurity control. Studies of active methods of impurity control are a natural development of the JET programme and, accordingly, a New Phase for JET is planned to start in 1992 [4], with first results becoming available in 1993 and continuing to the end of 1996.

The aim of the New Phase is to demonstrate, prior to the introduction of tritium, effective methods of impurity control in operating conditions close to those of the Next Step, with a stationary plasma of 'thermonuclear grade' in an axisymmetric pumped divertor configuration. Successful impurity control would lead also to an increase in alpha-particle power by more than a factor of two.

Specifically, the New Phase should demonstrate:
• control of impurities generated at the divertor target plates;
• decrease of the heat load on the target plates;
• control of plasma density;
• the exhaust capability;
• a realistic model of particle transport.

7.1. Key concepts of the JET pumped divertor

Since sputtering of the target plates cannot be suppressed, the impurities produced must be retained close to the target plates for effective impurity control. This can be achieved by friction with a strong plasma flow, directed along the divertor channel plasma (DCP) towards the target plates. If large enough, the frictional force between the plasma and the impurities should prevent impurity migration towards the X-point. The plasma flow will be generated by a combination of gas puffing, the injection of low speed pellets and the recycling and recirculation of some of the flow at the target plates towards the X-point. The connection length along the magnetic field line between the X-point and the target plates should be sufficiently long to achieve effective screening of impurities.

Rapid sweeping of the target plates to limit the localised heat load, to limit erosion and to affect redeposition is an important feature. Methods of ensuring that a substantial fraction of input power can be radiated in a controlled way in the DCP remain key elements.

7.2. Modelling the edge plasma

The plasma behaviour in the SOL and the DCP can be qualitatively understood. Impurity retention in the divertor is determined from the steady state momentum equation for impurity ions, which for the simplest, realistic case gives the
impurity density, $n_z$, decaying exponentially with distance from the target on a scale length, $\lambda_z$, given by:

$$\lambda_z^{-1} = \lambda_F^{-1} - \lambda_e^{-1},$$

with

$$\lambda_F^{-1} = \frac{m_z v_z}{T \tau_z}$$

and

$$\lambda_e^{-1} = \alpha_z \frac{1}{T} \frac{dT}{dx}$$

The temperature gradient scale length is given by the heat transport equation with electron heat conductivity parallel to the magnetic field ($\kappa = k_0 T^{5/2}$) being dominant and dependent primarily on the input power. To ensure impurity control, the frictional force must exceed the sum of the pressure and thermal forces, that is:

$$\lambda_F^{-1} > \lambda_e^{-1}.$$ 

To increase the hydrogen-impurity friction and radiated power in front of the target plates, a low temperature, high density plasma is needed. This will also limit erosion and therefore increase the lifetime of the target plates.
To solve the full set of classical fluid equations for the conservation of particles, momentum and energy in the SOL and DCP, a numerical 1-1/2 D transport model has been developed [14]. Monte Carlo methods are used for neutral deuterium and impurities in the flux surface geometry of the pumped divertor. Erosion at the target plates is determined by a model of sputtering tested against JET experimental data [8]. The calculations show that impurities can be retained near the target plates for plasma flows, typically \( \~10^{23}\text{s}^{-1} \) near the X-point. The steady state distributions of \( Z_{\text{eff}} \) (with beryllium impurities), for conditions in the SOL and DCP, with and without flow, are shown in Fig. 12(a). These results are obtained for an electron density \( \~10^{20}\text{m}^{-3} \) at the target plates. At target densities approaching \( 10^{21}\text{m}^{-3} \), the reduction of erosion and the plasma flow associated with high recycling at the target plates ensures impurity control. Furthermore, the calculations show that the ion temperature in the SOL can be substantially larger than the electron temperature (Fig. 12(b)). In present JET discharges, probe measurements indicate that, at low density, the electron temperature at the target plates is lower than the ion temperature, determined from broadening of the \( \text{H}_\alpha \) emission and power balance considerations [7].

**FIG. 13.** Cross-section of JET showing the plasma and main elements of the pumped divertor planned for the New Phase of JET.
7.3. The pumped divertor configuration

The aims of the New Phase can be realised with the internal multi-coil configuration shown in Fig. 13. The design allows a large plasma volume at 6MA and the operational flexibility to modify the magnetic configuration in the vicinity of the X-point independent of the plasma current and separately on the high and low field sides. In contrast to the normal configuration for a divertor, all divertor coils carry current flowing in the same direction as the plasma current.

Water-cooled, hypervapotron elements, made of copper and covered by beryllium, will be used for the high heat flux components of the target plates, and these are expected to accommodate power fluxes up to $15\text{MWm}^{-2}$ at the copper-beryllium interface. A pumping chamber is introduced in the vicinity of the target plates to provide control of the main plasma density. Pumping is achieved by a cryo-pump to avoid excessive hydrogen retention and to be compatible with the tritium phase.

With this configuration, single null X-point operation should be possible for performance and impurity control studies. Plasmas should be obtained at 6MA for 10s, a volume ~93m$^{-3}$ and a connection length from the X-point to the target plates of 3m, and at 5MA for 10s, a volume ~80m$^{-3}$ and a connection length ~10m should be achieved. In addition, it should be possible to run 3MA double null X-point plasmas for up to 20s at 3.4T and for up to 1 minute at 2.1T.

7.4. JET programme in the New Phase

A schedule for the JET programme incorporating the New Phase is shown in Table II. The earliest date to have a pumped divertor in JET is 1992. Further optimisation would likely be necessary about 18 months later, in the light of new experimental results.

By the end of 1994, all information on particle transport, exhaust and fuelling, first wall requirements and enhanced confinement regimes needed to construct a Next Step tokamak, should be available. Final tests with tritium, including alpha-particle heating studies, could be performed in the two years following, leading to completion of the JET programme by the end of 1996. During the tritium phase of operation, and even later, tests on prototype elements for a Next Step tokamak could be undertaken on the JET site, using the test facilities, tritium plant and power supplies.

8. CONCLUSIONS

In summary, the paper sets out the main JET results and future prospects. JET has successfully achieved and contained plasmas of thermonuclear grade. Individually, parameters required for a fusion reactor have been obtained, and simultaneously the fusion product $(n_D(0)\tau_E T_i(0))$ has reached $8-9\times10^{20}\text{m}^{-3}\text{skeV}$, for both medium (~10keV) and high (>20keV) central temperatures, and is within a factor of 8 of that required in a fusion reactor;
neutron yield has increased to $\sim 3.5 \times 10^{16} \text{ns}^{-1}$, corresponding to an equivalent $Q_{\text{DT}} \sim 0.8$. However, these values were in a transient state and could not be sustained in steady state.

A clearer picture of energy and particle transport also starts to emerge. The resilience of the electron temperature suggests that the electrons are primarily responsible for confinement degradation. The occurrence of flat density profiles suggests that there is no need for an inward pinch term, except perhaps for impurities. Particle and energy transport exhibit similar behaviour. The critical electron temperature gradient model is one such model consistent with these observations.

In a reactor, the density limit should not be a problem; high densities and flat profiles are likely; an advanced divertor concept for impurity control is required. These conditions of high plasma density at the separatrix are unfavourable for the methods of non-inductive current drive envisaged at present.

A New Phase is planned for JET with an axi-symmetric pumped divertor configuration to operate with a stationary plasma (10s-1 minute) of thermonuclear grade. In this New Phase, JET should be able to:

- demonstrate a concept of impurity control;
- determine the geometry needed to fulfil this concept in a reactor;
- identify appropriate materials for plasma facing components;
- define the operational domain for the Next Step.

To ensure the success of a Next Step ignition device, it is imperative that all aspects of plasma behaviour, impurity control and plasma exhaust be included in the model used to define the size of the device and its toroidal field, plasma current and operating conditions.
Appendix I

THE JET TEAM

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DISCUSSION

R.J. GOLDSTON: It is important to distinguish the density limit in the L-mode from the density limit in the H-mode. Could you tell us the maximum $n_n R/B_T$ achieved in JET, in a high confinement H-mode? It would be good to know the values with gas fuelling and with pellets, separately.

P.H. REBUT: The density limit in the H-mode is generally an $H \to L$ transition and corresponds to a power balance between the heating and the radiation. As the density profile is generally flat, this occurs at a lower mean density.

C.E. SINGER: It is difficult to obtain a plasma for which the collisional fluid theory is valid in the scrape-off layer, particularly for the high energy electrons which carry most of the energy. Was the fluid theory valid for the scrape-off simulations you referred to?

P.H. REBUT: The heat conductivity has to be limited by non-collisional transport and this is not a problem in the simulation.

B. COPPI: There are problems with using the product $n_D(0) T_D(0) \tau_E$ to measure the rate of progress towards ignition conditions. ‘Penalty’ factors should in fact be introduced when $Z_{\text{eff}}$ exceeds unity and the ratio $\tau_E/\tau_\alpha$ (energy confinement time/alpha particle slowing-down time) is less than unity.

P.H. REBUT: This is true, but on JET the dilution factor and the ratio $\tau_E/\tau_\alpha$ do not bring a penalty outside the error bar. But as far as the ratio $T_i/T_e$ is concerned, we have $T_i = T_e$, as in a reactor.

J.C. HOSEA: In the H-mode, you showed a trapping of nickel injected impurity, but earlier you stated that particle confinement follows energy confinement. Are you concluding that the impurity transport differs from the majority ion (deuterium) transport? I am presuming that energy confinement only increases by a factor of approximately 2-3 in the transition from the L- to the H-mode.

P.H. REBUT: The presence of a confinement barrier at the edge of the plasma affects both the energy and the particles. With the same relative improvement in $\chi$ and $D$, the effect on particle transport is much larger, as it is controlled mainly by the recirculation at the edge. The neoclassical effects on high-Z impurities are also important in an H-mode plasma.

G. FUSSMANN: I could not follow your argument that there is no need for an anomalous inward drift for the electrons. You showed two $n_e$ profiles, one for carbon and the other for beryllium as the dominant impurity, which differed considerably in shape. What, in your opinion, is the physical mechanism producing the peaked $n_e$ profile?

P.H. REBUT: This is a complex argument. An inward pinch for the impurities, that is, a cross-term, also has an effect on the electrons through the radial electric field due to quasi-neutrality.
RECENT EXPERIMENTS IN JT-60

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Abstract

RECENT EXPERIMENTS IN JT-60.

The emphasis of the latest JT-60 experiments has been placed on confinement improvement with profile control, and steady state operation research. Improved confinement and fusion product accompanied by centrally peaked pressure profiles was demonstrated by profile control with pellet injection, LHCD and ICRF heating. Extensive research on LHCD has made it possible to find a limiter H mode with a nearly steady state ELM free phase of up to 3.3 s. The main issues for steady state operation research were non-inductive current drive and particle and heat control with a divertor. LHCD has been achieved with a maximum \( \frac{T}{\rho} \) of \( 3.4 \times 10^{19} \text{ m}^{-2} \cdot \text{A}^{-1} \cdot \text{W}^{-1} \), which is very close to the efficiency necessary for ITER steady state operation. The ratio of neoclassical bootstrap current to the total current was found to increase in proportion to \( \beta_p \) and reached 80\% at \( \beta_p \sim 3 \). In L mode divertor discharges, \( \alpha \) particle production with a helium NB was associated with a high helium pressure in the divertor region, enough for the helium exhaust in future devices with a reasonable pump speed. Strong remote radiative cooling of up to 50\% of the total input power and a dense, cold divertor plasma were successfully demonstrated in 23 MW NB heated divertor discharges. JT-60 has been shut down since November 1989 for modification to the JT-60 Upgrade. In early 1991, JT-60 Upgrade will initiate a new regime of experiments including: (1) a confinement study in a new parameter regime of high \( I_p \), \( B_T \) and aspect ratio; (2) non-inductive current drive by a negative ion source neutral beam, bootstrap current and LHCD; and (3) energetic particle physics relating to burning plasmas.

1. INTRODUCTION

JT-60 is a large tokamak designed to study: (a) the scaling of plasma behaviour as parameters approach the reactor range; (b) plasma–wall interaction in these conditions; (c) plasma heating; and (d) steady state operation. JT-60 started up in April 1985 and has since been progressively upgraded [1, 2]. This paper provides an overview of the results since the last IAEA conference (Nice, 1988), namely the experimental results in 1989, JT-60 has been shut down since November 1989 for modification to the JT-60 Upgrade. The experiment will resume in early 1991. This paper also discusses the future prospects of JT-60 Upgrade experiments.
2. STATUS OF THE DEVICE

After the Nice conference, there were two major additions to the operation area in 1989 that have had a significant impact on expanding the performance capability of the tokamak: (1) a four barrel pneumatic pellet injector with a speed of 2.3 km/s and sizes of 3 mm and 4 mm; and (2) a new multijunction type lower hybrid launcher with sharp power spectra of $\Delta N_e \sim 0.5$ and high directivity. These additions have given us extensive capabilities in density profile control and very sharp profile control of the production of energetic electron current. Associated with this additional equipment, the plasma operation regime has been extended to investigate two major issues: improvements in plasma confinement with profile control, and steady state operation.

In JT-60 Upgrade, the plasma current will be increased up to 6 MA with a lower single null divertor. The poloidal field coils and the vacuum vessel are being replaced for this modification. NB heating power of up to 40 MW and a total of \( \sim 20 \) MW of LH and ICRF will be available using deuterium as the working gas. The major radius and the toroidal field will be 3.4 m and 4.2 T. Plasma parameters are expected to further approach the reactor range. Furthermore, the addition of a 500 keV, 10 MW negative ion source neutral beam in 1994, which will be primarily used for NBCD, will allow us to investigate physics relating to burning plasmas, where the global plasma behaviour will be greatly affected by energetic particles.

3. IMPROVED CONFINEMENT REGIME WITH PROFILE CONTROL

Improved confinement and fusion product \( (n_e(0)T_e(0)) \) were investigated in hydrogen discharges. One of the most effective ways to enhance the fusion product is to produce peaked plasma pressure profiles, which are usually accompanied by sawtooth stabilization and improved confinement inside the \( q = 1 \) surface. During the profile control experiments in 1989, pellet injection, LHCD and higher harmonic ICRF were found to stabilize the sawtooth and to enhance the fusion product.

3.1. Pellet injection

Improved energy confinement has been obtained by injection of hydrogen pellets for auxiliary heating power of up to 25 MW. Figure 1 compares the time evolution of the SXR signals for four pellet injection discharges with \( q_a \sim 2.3 \) (\( I_p = 2.8-3.1 \) MA). The pellet penetration becomes deeper from the top to the bottom of the figure. With increasing $\beta_p$ (0.07 to 0.35) or density gradient inside the \( q = 1 \) surface, the sawtooth period increases and the release of the central kinetic energy due to the sawtooth decreases. The strongly peaked electron density profile with \( n_e(0)/\langle n_e \rangle < 3 \) and \( n_e(0) \) up to \( 2.7 \times 10^{20} \) m$^{-3}$ was sustained within 0.5–1 s after a series of pellet injections. Central plasma pressure reaches 2 atm, and achieved
values of $n_e(0)\tau E_T(0)$ ($\sim 1.2 \times 10^{20}$ m$^{-3}$·s·keV) at $I_p = 3.1$ MA are two times larger than those obtained with gas puffing. Improved energy confinement for the pellet fuelled plasma is mainly due to the peaked density and pressure profiles inside the $q = 1$ rational surface, as shown in Fig. 2. The pressure gradient of these discharges locally reaches the marginal value for the infinite $n$ ballooning mode inside the $q = 1$ surface. The poloidal beta within the $q = 1$ surface, $\beta_p^\perp$, has not reached the ideal ballooning limit. The analysis shows that the attained maximum $\beta_p^\perp$ is consistent with the internal $n = 1$ kink mode beta limit [3].

FIG. 1. Time evolution of SXR signals (chords at $r = 0$, 42 and 56 cm) and the electron density profile for pellet fuelled discharges with $q_a \sim 2.3$ ($I_p = 2.8$-3.1 MA). The pellet penetration depth becomes greater from the top to the bottom of the figure.
3.2. Sawtooth stabilization with LHCD

Sawtooth stabilization with LHCD has been extensively investigated with the new LH launcher. A broader profile of energetic electron current was obtained with higher \( N_b \) injection, as shown in Fig. 3(a). One the other hand, low \( N_b \) injection was effective for suppression of sawtooth activity. Figure 3(b) shows a clear correlation between the sawtooth stabilization period and the intensity of energetic electrons near the plasma centre. High power NB heated discharges with low \( N_b \) LHCD showed centrally peaked energetic electron current and ion temperature profiles during the sawtooth free periods of up to 1.8 s. Figure 4 shows a comparison of NB and NB + LHCD discharges. The appearance of the first sawtooth with relatively large inversion radius \( (r_s/a \sim 0.2) \) at the termination of the sawtooth free phase of NB + LHCD discharges and the existence of \( m = 1 \) oscillation in some marginally stabilized discharges suggest that \( q(0) \) may be less than or equal to unity during the sawtooth suppression period. From these observations, the existence of energetic electrons near the \( q = 1 \) surface or local current profile modification may be possible keys to the sawtooth suppression mechanism.
FIG. 3. (a) Radial profiles of hard X ray intensity for low \( N_t \) to high \( N_t \) LH injections. (b) Sawtooth suppression period, \( \Delta t_{ST} \), versus hard X ray intensity at the plasma centre, \( I_{HX}(0) \).
FIG. 4. Comparison of NB and NB + LHCD (sawtooth suppression) 1.5 MA discharges.
3.3. Limiter H mode with LHCD

The extension of the operation regime of LHCD has greatly benefited the confinement programme. The H mode was achieved in limiter discharges with LHCD for the first time. The threshold LH power was as low as the Ohmic heating power with hydrogen plasmas. A nearly steady state ELM free H mode with a duration up to 3.3 s was established without significant impurity accumulation [4].

3.4. Higher harmonic ICRF heating in combination with NB heating

ICRF heating accompanied by beam acceleration in the third harmonic regime was investigated. A typical operation regime is: $f = 131$ MHz, $P_{IC} < 2.3$ MW, $P_{NB} < 10$ MW, $B_T = 3$ T and $\bar{n}_e < 3 \times 10^{19} \text{ m}^{-3}$. Strong central heating accompanied by giant sawteeth is observed. The sawtooth period increases in proportion to $P_{NB}/\bar{n}_e$ with fixed ICRF power and reaches 410 ms. The energy confinement was enhanced by about 20% under the maximum operation condition [5].

3.5. Energy transport

Local heat transport has been analysed on the basis of radial profile measurements of electron density, and electron and ion temperatures. Most of the analysis was made in L mode discharges of $\bar{n}_e < 4 \times 10^{19} \text{ m}^{-3}$, where the electron–ion heat exchange is negligible in the power balance. Figure 5(a) shows the agreement of kinetic ($W_k$) and diamagnetic ($W_d$) stored energy, and the stored energy versus heating power of 1 MA discharges. The stored energy of ions ($W_i$) shows a very small increment compared with that of electrons ($W_e$) against the heating power. The transport analysis shows that the ion heat transport appears to be dominant in the heating power and $I_p$ scan. As is shown in Fig. 5(b), $\chi_i$ increases substantially with heating power and decreases with the plasma current, although $\chi_e$ showed little change. Lower X point divertor discharges with small minor radius ($a \sim 0.7$ m, aspect ratio $\sim 4.3$) show relatively small $\chi_i$ compared with the 1 MA limiter discharge with $a \sim 0.9$ m. In the hot ion enhanced confinement regime [6], where the peaked electron density profile of $n_e(0)/\langle n_e \rangle \sim 2.5$ is produced, $\chi_i$ is reduced more than an order of magnitude in the plasma centre region of $r/a < 0.3$.

4. STEADY STATE OPERATION RESEARCH

Major issues of steady state operation research are non-inductive current drive and particle and heat control. Fully non-inductive current drive capability is essential in a steady state reactor. High current drive efficiency is required for this objective. In JT-60, LHCD current of 1 MA was demonstrated. Further improvement of current drive efficiency was investigated with a new multijunction LH launcher. To minimize
the power required for the steady state current drive, bootstrap current was intensively investigated. Major problems in particle and heat control in the next step device are the exhaust of fusion produced \( \alpha \) particles and the reduction of heat flux to the divertor plate. These issues were investigated in high power heated lower X point discharges.

4.1. LHCD

LHCD experiments were carried out for a wide range of \( N_t (1.0-3.4) \). The current drive efficiency \( \eta_{\text{CD}} \) was proportional to the product of \( 1/N_t^2 \) and the accessible power fraction. The \( \eta_{\text{CD}} \) of \( 3.4 \times 10^{19} \text{m}^{-2} \cdot \text{A} \cdot \text{W}^{-1} \) has been achieved by the optimization of sharp \( N_t \) spectra. Plasma current of 1–1.8 MA was driven by 2–4.5 MW LH injection into lower X point discharges with \( \bar{n}_e \) up to \( 3.0 \times 10^{19} \text{m}^{-3} \). An order of magnitude improvement has been made over past experiments in the current drive product, which reached \( 12.5 \times 10^{19} \text{m}^{-2} \cdot \text{MA} \). Figure 6 shows the progress of current drive product and \( \eta_{\text{CD}} \). The achieved \( \eta_{\text{CD}} \) is approaching \( 5 \times 10^{19} \text{m}^{-2} \cdot \text{A} \cdot \text{W}^{-1} \), which is the efficiency necessary for ITER steady state operation [5].

4.2. Bootstrap current

The neoclassical bootstrap current was confirmed in a wide range of \( \beta_p \). The ratio of bootstrap current to the total current increased in proportion to \( \beta_p \). A significant bootstrap current fraction reaching 80% has been observed at \( \beta_p \sim 3.2 \) [6].

![FIG. 6. Progress of the JT-60 LHCD experiment in current drive product and current drive efficiency.](image-url)
Figure 7(a) shows the time dependence of the surface voltage. The calculated voltage with and without bootstrap current is shown. By using the measured time evolution of electron and ion temperatures and electron density profiles, the time dependent resistive diffusion equation and MHD equilibrium equation were solved incorporating the full neoclassical expression for the bootstrap current density as well as the beam driven contribution (only 3 kA because of perpendicular injection) [7]. The case including the bootstrap current shows good agreement with the experiment and shows a bootstrap current fraction of 70–80%. The existence of the bootstrap current was also demonstrated by keeping the primary OH current constant and measuring the decay time of the total current, as shown in Fig. 7(b). The enhancement of the L/R time in the high $\beta_p$ discharge (E10701) over that of the low $\beta_p$ discharge (E10704) is consistent with what is expected from the neoclassical theory with bootstrap current. The result is encouraging for the design of a high $Q$ (> 30) steady state tokamak reactor [8].

4.3. Helium exhaust

To simulate the $\alpha$ particle production in D–T plasmas, a 30 keV helium NB was injected into NB heated lower X point discharges, which produced a centrally peaked birth profile of $\alpha$ particles. Both the $H_2$ and the helium pressure in the divertor region increased in proportion to $n_e^3$. The helium pressure reached 0.03 Pa (10% of the $H_2$ pressure) at $n_e = 6 \times 10^{19}$ m$^{-3}$, where $\alpha$ particle density is 10% of the main proton density. A simple extension of the present result is promising for the helium exhaust in future devices. A pump speed of several tens of cubic metres per second will be sufficient for a 1000 MW fusion power reactor operating at $n_e = 1 \times 10^{20}$ m$^{-3}$ [9].

4.4. Formation of dense and cold divertor plasma

The capability of remote radiative cooling has been extensively investigated with high power NBI of $\sim$ 20 MW in lower X point discharges. Over a threshold main plasma density and safety factor, strong remote radiative cooling of up to 50% of the total input power was observed in JT-60. In this situation, discharges with 23 MW NBI showed no carbon burst during the whole injection period of 4 s. Measurement of $H_\alpha$ and carbon lines and heat flux density revealed that the divertor plasmas are dense and cold ($n_{e,div} \sim 2.4 \times 10^{20}$ m$^{-3}$, $T_{e,div} \sim 26$ eV), and the remote radiative cooling is a mixture of carbon and hydrogen radiation [9].

5. FUTURE PROSPECTS OF JT-60 UPGRADE RESEARCH

The primary objective of JT-60 Upgrade (Fig. 8) in the early phase of its experiments is impurity control under favourable energy confinement conditions.
FIG. 7. (a) Loop voltage, bootstrap current, safety factor and stored energy versus time. Calculated voltage is shown for cases with and without bootstrap current. (b) Comparison of decay time of the plasma current for low $B_T$ (E10704) and high $B_T$ (E10701) discharges with the primary OH current kept constant.
The divertor plates are designed to be toroidally continuous and to use high heat conduction C–C composite graphite with a height misalignment of adjoining tiles of less than 0.5 mm. A period of about 3 s without carbon burst is estimated under ~40 MW of injection power. The development of the divertor in particle and heat control will be carried out in collaboration with JFT-2M and DIII-D. After 1994 a second primary objective will be non-inductive current drive by ECH assisted NBCD with a 500 keV negative ion source NB in combination with bootstrap current and LHCD.
5.1. Confinement study near the thermal breakeven condition

The device capability of high power heated deuterium divertor discharges with high $I_p$ of 6 MA, $B_T$ of 4.2 T and aspect ratio of 4.0 is expected to extend substantially the present world confinement database in its relevance to a reactor grade plasma. Figure 9 shows the predicted fusion product based on the recently re-established ITER-89 power law confinement scaling. Since the scaling predicts a fusion gain proportional to $(R/a)^{1.4}$, the relatively high aspect ratio of JT-60 Upgrade ($R/a = 4.0$ compared with 2.5 in JET with the same plasma current of 6 MA) will substantially contribute to the confinement database for the next step device.

5.2. Non-inductive current drive

The regime of negative NBCD will be $n_e I_p \sim 5 \times 10^{19} \text{ m}^{-3} \cdot \text{MA}$ at $T_e(0) \sim 10$ keV. With the addition of bootstrap current, full to $\sim 70\%$ non-inductive current drive at $I_p = 2-3 \text{ MA}$, $n_e \sim 5 \times 10^{19} \text{ m}^{-3}$ and $\beta_p = 1.6-1$ will be the primary experimental objective. A major research issue in this operation regime is to realize long, stable discharges with a bootstrap current fraction of more than 50%. The primary research objective of LHCD will be the enhancement of $\eta_{CD}$ in the
high $T_e$ regime. In the case of LHCD + NBCD, absorption of LH waves by energetic ions might be a problem. Fully non-inductive current drive of 4 MA at $n_e \sim 5 \times 10^{19}$ m$^{-3}$ can be evaluated if the degradation of LHCD is in the range of $\sim 40\%$.

5.3. Energetic particle physics relating to burning plasmas

The ability to generate $\alpha$ particles due to the $D + \text{^3He} \rightarrow \text{^4He}$ (3.6 MeV) + p (14.7 MeV) reaction with a 500 keV D beam at which the fusion cross-section peaks will provide a good opportunity to study $\alpha$ particle behaviour prior to D–T operation. A real fusion reaction in the 1–1.5 MW range is expected. This will provide more reliable confinement properties of $\alpha$ particles compared with the present experiments producing fusion reactions of $\sim 100$ kW using ICRF. Energetic single particle confinement concerns the effect of non-axisymmetry of the magnetic field of reactor grade plasmas. The stochastic ripple diffusion effect may cause rapid energetic ion loss. The edge of the JT-60 Upgrade plasma will have $\sim 2\%$ ripple. Detailed orbit following Monte Carlo calculations of this process show $\sim 30\%$ of 120 keV fast ions lost for perpendicular injection in a high density discharge, compared with less than 10% loss for tangential injection. The experiment will provide an important database to evaluate the stochastic ripple diffusion process. The other important issue of burning plasma physics concerns the collective stability of the $\alpha$ particle population and its effect on confinement. The negative NBI experiment may be able to simulate the relevant physics before a burning plasma experiment is attempted. The expected parameter range will be $V \sim 1.4 V_{\text{Alfvén}}$ and $\beta_n \sim 1\%$ at $B_T = 3$ T for a 500 keV, 10 MW H beam, which seems to be well within the toroidal Alfvén eigenmode unstable regime [10].

6. SUMMARY

After five years of JT-60 operation, the 1989 experimental results have in particular shown progress in confinement and steady state operation research. Enhancement of confinement and fusion product accompanied by centrally peaked pressure profiles was demonstrated by profile control with pellet injection, LHCD and ICRF. Non-inductive current drive research provided high efficiency LHCD and a high fraction of bootstrap current, providing an important database for the design of a steady state reactor. Particle and heat control was successfully demonstrated in L mode lower X point discharges. After about one year of device modification work, JT-60 Upgrade will initiate a new regime of experiments: a confinement study in a new parameter regime with high $I_p$, $B_T$ and aspect ratio; non-inductive current drive by 500 keV negative NBI, bootstrap current and LHCD; and energetic particle physics relating to burning plasmas.
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REFERENCES


DISCUSSION

S.A. COHEN: When the cold, dense divertor plasmas were achieved, did you measure $\nabla V T$ or detect a hot electron tail at the divertor plate? The presence of such a tail would alter the divertor sheath potential and possibly the effectiveness of the thermal force in maintaining plasma purity.

M. NAGAMI: The cold, dense plasma was evaluated from measurements of the divertor heat flux and the absolute intensity of $H_a$ radiation.

R. PARKER: The high beta poloidal results, which indicate the presence of a large bootstrap current, are transient — that is, the current decays with time. Why? Could you combine these results with lower hybrid current drive to produce steady state conditions?

M. NAGAMI: The plasma current decays because the bootstrap current contribution to the total plasma current was 2/3 in this experiment. We have not yet tried the combined bootstrap and LHCD discharge.
DIII-D RESEARCH PROGRAM PROGRESS

DIII-D TEAM

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Abstract

DIII-D RESEARCH PROGRAM PROGRESS.

A summary of highlights of the research on the DIII-D tokamak in the last two years is given. At low $q$, toroidal beta ($\beta_T$) has reached 11%. At high $q$, $\epsilon \beta_T$ has reached 1.8. DIII-D data extending from one regime to the other show the beta limit is at least $\beta_T \geq 3.5 I/aB (MA, m, T)$. Prospects for using H-mode in future devices have been enhanced. The discovery of negative edge electric fields and associated turbulence suppression have become part of an emerging theory of H-mode. Long pulse (10 second) H-mode with impurity control has been demonstrated. Two methods, radial sweeping of the divertor strike points and gas puffing under the X-point, were each shown able to lower peak divertor plate heat fluxes by a factor of 2–3. $T_e = 17$ keV has been reached in a hot ion H-mode. Electron cyclotron current drive has produced up to 70 kA of driven current. Program elements now beginning are fast wave current drive and an advanced divertor program.

1. INTRODUCTION

The DIII-D research program has as its long range goal the integrated demonstration of high $\beta$ operation with good confinement and non-inductive current drive to provide a basis for tokamak progress toward steady-state reactors. This goal is pursued through several lines of research which are also important for their contributions to basic tokamak physics. We describe these areas and recent research accomplishments in them.

Plasma cross-section shaping and profile control are employed to probe $\beta$ limits and increase $\beta$. Recently in the low $q$ regime $q_{95} \sim 2.6$, $\beta_T = 11\%$ in a highly elongated ($\kappa = 2.35$) double-null divertor plasma was achieved [1], demonstrating the tokamak can sustain $\beta$ considerably in excess of that planned in future devices. The $\beta_T$ part of the integrated demonstration referred to above has been met with $\beta_T = 5\%$ reached at

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1 See the Appendix.
full $B_T = 2.1$ T [2]. The scaling of $\beta$, $\beta_T$ (\%) $\geq 3.5 \ I/\alpha B$ (MA, m, T), and detailed analyses continue to be well in accord with theory, laying a solid foundation for future device design [2,3]. At higher $q$, the normalized beta, $\beta_N \equiv \beta_T/(I/\alpha B)$, has reached 5 [2,3] and at still higher $q$, $\beta_p = 5.1$ and $\epsilon \beta_p = 1.8$ have been reached [4]. These results, achieved in the regime where bootstrap current can be large, are favorable for current driven reactor designs [5].

For enhanced confinement, DIII-D contributes to providing a basis for the use of H-mode in future devices [6]. A scaling for ELM-free H-mode confinement was worked out by a joint team from JET and DIII-D [7]. Considerable progress was made on elucidating the physics basis of H-mode. A sudden increase in edge poloidal rotation, corresponding to an increased negative edge electric field was discovered on DIII-D [8]. Subsequently detailed investigations have not conclusively shown edge rotational effects cause the L-H transition [9], but substantial evidence points to increased shear in edge rotation as the mechanism for edge turbulence suppression [10]. H-mode with flat density profiles has been shown able to sustain a hot ion regime; $T_i = 17$ keV has been reached [11]. Theory suggests two fundamental prescriptions for scaling confinement results to future devices, the Bohm or gyro-Bohm prescriptions depending on whether the turbulence is long scale ($a$) or short scale ($\rho_S$) respectively. An experiment comparing dimensionally similar discharges was done to distinguish between these two possibilities [12]; the gyro-Bohm model is favored.

For long pulse impurity control and heat flux handling, two important results were obtained. A 10 sec long nearly stationary H-mode [13], was produced by using ELMs to control impurity levels [14]. Two methods, radial sweeping of the divertor strike points and gas puffing under the X-point to encourage a radiative divertor plasma, were each shown able to lower time average peak heat fluxes a factor of 3 and 2 respectively [15]. These methods could be applied in combination.

For non-inductive current drive, two rf methods are being developed: electron cyclotron current drive (ECCD) and fast wave current drive (FWCD). A first result, up to 70 kA of driven current, has been produced using 1 MW of 60 GHz (fundamental) waves launched in X-mode from the inboard side [16]. The rf program also seeks to develop ECH as a heating tool; important transport results, ELM suppression studies, and H-mode studies have been carried out [17]. Near term plans call for ECH to provide some current drive and also the high $T_e$ regime in which FWCD can be efficient.
2. RECENT MODIFICATIONS

In the last two years, the following facility modifications, which are important in the results to be described, were implemented.

1. **Neutron Shielding** was added around the entire DIII-D machine pit. The resulting shielding factor of 300 allowed routine \( D^0 \rightarrow D^+ \) operation. The improved confinement from the isotope dependence \( \tau_E \sim \sqrt{M_i} \) was important in reaching high beta and high performance plasmas.

2. **Deuterium Neutral Beam Operation** enabled the output power of the four beamline system to be increased from 14 MW \( H^0 \) to 20 MW \( D^0 \), also an important factor in attaining higher beta and stored energy.

3. **Carbonization** of the vacuum vessel wall (half graphite and half Inconel) was done by means of a glow discharge in a mixture of helium and deuterated methane. Subsequent glow in pure helium was used to desorb deuterium from the carbon film. The result was a reduction in metallic impurity levels by a factor of 30. The carbonization enabled an increase in the plasma current at which we could reliably operate divertor discharges from 2 MA to 3 MA with H-mode, an increase in the peak achieved stored energy from 2.6 MJ to 3.6 MJ, and an increase in the ion temperature from 11 to 14 keV in hot ion H-mode in corresponding conditions. Peaked density profiles \( n(0)/\bar{n}_e = 2 \) were produced in H-mode and the ELM frequency was doubled [18].

4. **Ion Bernstein Wave Heating** (IBW) was initiated to get a local ion heating system for manipulation of H-mode plasmas that potentially did not produce ion tails and avoided the loading problems with fast waves and H-mode [19]. A 30–60 MHz, 2 MW generator feeding a pair of cavity type end fed loop antennas oriented along the toroidal field was used. Unfortunately, central ion heating was never seen and the loading (\( \sim 5 \) Ohms) was one to two orders of magnitude greater than theoretical prediction and never displayed any of the expected resonances at ion cyclotron harmonics. Edge electron heating was observed. Edge ion heating was correlated with increased amplitudes of parametric decay spectra. Impurity influx was customary, except with carbonization. This program has been discontinued.

5. A four strap fast wave current drive antenna has been installed in place of the IBW antenna. It is designed to handle 4 MW; the 2 MW generator will be used in initial tests of heating and current drive.
6. A baffle and ring electrode structure has recently been installed (Fig. 1) to enable particle and energy confinement experiments with divertor biasing and pumping.

3. HIGH TOROIDAL BETA $\beta_T$ AND NORMALIZED BETA $\beta_N$

The situation in regard to maximal $\beta_T$ limits is summarized in Fig. 2. Certainly the envelope of the data supports the conclusion that the beta limit is at least as high as $\beta_N = 3.5$; however, this conclusion is couched as an “at least” statement because the achieved $\beta_T$ in most cases on this
plot was limited by the product of available heating power and confinement time \((W = P \times \tau_E)\) rather than disruptive instabilities. Discharges identified as large open circles have been identified as clear candidates for external kink disruptions as evidenced by the \(m/n = 2/1\) locked mode precursor with fast growth time. Wall stabilization should be ineffective for locked modes (merely limiting the growth time to the vessel wall time constant, \(\sim 3\) msec in DIII–D) and so it is not surprising that our low-\(n\) mode stability calculations agree with those of Troyon on a limiting \(\beta_N = 2.8\) in these cases of free boundary kinks with no wall stabilization [20]. For a conducting wall at 1.5\(a\) from the plasma surface and nominal profiles, we find the low-\(n\) ideal limit at about \(\beta_N = 4\) [2,21].

Recently we have found that if in addition to the nearby conducting wall, profiles are optimized, the low-\(n\) limit can be raised to \(\beta_N \sim 5.5\) [22]. This possibility was first made clear in the calculations shown in Fig. 3.

**FIG. 2.** Equilibrium fit \(\beta\) versus \(I/a_B\) for DIII–D. The lines at \(\beta_T = 3.5\) (5.0) \(I/a_B\) are to guide the eye. The open circles locate ideal time-scale \(m/n = 2/1\) disruptions. Various \(\beta\) limit calculations are summarized in the curves (---, ---, ---, ...) with differing assumptions on the location of a conducting wall \((r_w/a)\).
FIG. 3. Calculated ideal $n = 1$ mode stability boundaries (i) versus internal inductance $\ell_i$ for the pressure (ii) and pressure gradient profiles (iii) shown versus normalized volume radius $\rho$. Mode displacement vectors peak near the edge for "external" modes and near the center for "internal" modes.

For broader pressure profiles, but more peaked current profiles ($\ell_i \sim 0.93$), access to $\beta_N \sim 5.5$ was predicted for $q_{95} = 3.2$ and $q_0 = 1.05$. Quite independently, an experimental campaign succeeded in producing stable discharges at $\beta_N = 5.0$ at about $I_N = I/a_B = 1-1.5$ ($q_{95} \sim 4-5$) as shown in Fig. 2. Stability calculations at that $q$ value (4-5) [2] also predict an optimized profile path to low-$n$ ideal stability and so we have drawn the short curve segment in the figure representing our calculated low-$n$ stability limit with optimized profiles. Further calculations are needed to see the range of $I_N$ for which $\beta_N \sim 5.5$ can be stable. For high-$n$ modes, our ideal MHD calculations in the ballooning representation have predicted an optimized profile limit of $\beta_N \sim 4.5$ [21], in good agreement with Sykes' result [23]. It remains to be seen whether these high values of $\beta_N$ can be reached over a broad range of $I/a_B$ and so redefine the expected
beta limit up to $\beta_N \sim 5$. Such an outcome could be important since it may allow future devices to be designed for operation at $\beta_N \sim 3.5$ [5] instead of the now customary cautious choice of $\beta_N \sim 2$.

With a beta limit scaling of the form $\beta_N \leq \text{constant}$, the task of maximizing the absolute value of $\beta_T$ becomes that of maximizing the normalized plasma current $I_N$ subject to the constraint $q_{95} > 2$ and providing sufficient power and confinement time. Using single-null divertor plasmas a maximum $I/a_B = 2.3$ was reached and a maximum $\beta_T = 7.4\%$. Double-null discharges allowed about the largest divertor plasma cross-section that would fit in the vacuum chamber and increased triangularity which allowed a larger $I/a_B = 3.3$ and a maximum $\beta_T$ of 9.3\%. This 9.3\% value was reached in a 60 msec long ELM-free phase following the L-H transition and the subsequent ELMs clamped $\beta_T$ to 8\% for 0.9 seconds [2].

4. HIGH ELONGATION AND HIGHEST $\beta_T$

Since the largest cross-section plasma in DIII-D has $\kappa \sim 2$, in order to further increase $I/a_B$, it was necessary to decrease $a$ in order to increase $\kappa$. Studies of the axisymmetric stability limits in DIII-D found that the ideal axisymmetric limits coincided with the predictions of GATO to within a few percent [24]. Between $\kappa = 2$ and 2.5, the simple rigid body shift model ($m = 1$) becomes inadequate and both the experiment and the calculations show that the vertical stability limit at $\kappa \sim 2.5$ has a large $m/n = 3/0$ component in the plasma motion. It was found that coils near the inboard midplane are optimal for vertical feedback control, since they have minimal interaction with the passive image current flow patterns in the vacuum vessel which dominate the vertical stability [24]. The principal importance of these studies was close (within a few percent) agreement between experimentally observed vertical stability limits and theoretically predicted limits (GATO).

With this basis of understanding and control system improvements, plasmas with $a = 0.56$ m and $\kappa = 2.35$ were operated at low $B_T = 0.75$ T to assess their $\beta$ capability. With $P_{\text{NBI}} = 19$ MW, $\beta_T = 11\%$ was reached [1]. Plasma current was 1.29 MA; $I/a_B = 3.1$ and $q_{95} = 2.56$. The equilibrium for this discharge is shown in Fig. 1. Various measures of $\beta_T$ were equilibrium fitting (10.7\%), diamagnetism (10.8\%), and kinetic profile integration with the fast ion contribution computed by standard methods (11.3\%). Besides the higher NBI power, the proximate cause why this discharge reached such high beta seems to have been the absence of sawteeth on the beta rise. Kinetic profile data show the thermal beta on axis
to be 20%; standard beam deposition codes estimate the fast ion beta on axis to be an additional 20%. These results raise the issues of FLR effects and fast ion stabilization of interior modes, subjects of current investigation. However, this result does establish that the tokamak can contain an absolute value of $\beta_T$ (11%) much in excess of that planned in CIT and ITER ($\beta_T \sim 4\%-6\%$).

5. ATTACHED EQUILIBRIA

Vertical instabilities in highly elongated plasmas have produced very large forces on the vacuum vessel in DIII-D and JET. The vertically unstable plasma drifts down at a speed limited by vessel image currents, limiting on the bottom of the vessel and shrinking in cross-section until it disrupts. We have found that while drifting the equilibria are “attached” to the vessel in the sense shown in Fig. 4. In a 1 MA discharge up to 0.3 MA flows on the open field lines, outside the last closed flux surface, completing the circuit through the bottom of the vacuum vessel.
These attached currents have been detected both by equilibrium fitting of magnetic data and directly by the resistive shunts under the divertor tiles and divertor tile Rogowski coils with excellent agreement between the two methods. The force arising from the poloidal component of this current crossed with $B_T$ accounts for the observed up/down motion of the vacuum vessel [14]. Such large wall currents must be carefully considered in future machine designs.

6. HIGH BETA POLOIDAL

Probing for the possible existence of a second stable regime, at very high $q$ ($q_{95} = 18$, $I = 0.54$ MA, at $B_T = 2.1$ T), 14 MW NBI was used to produce the very high value of $\beta_p = 5.1$ and $\epsilon \beta_p = 1.8$ [4]. These plasmas are approaching the equilibrium limit, with a magnetic axis shift of 0.22 m out of $a = 0.61$ m equal to the inverse aspect ratio 0.36 and the poloidal field on the inboard side weakened to 1/3 of the value on the outboard side. The high $\beta_p$ and confinement equal to $2.3 \times$ ITERP-89 are realized during the ELM-free phase. During ELMs, confinement is reduced and $\epsilon \beta_p$ was maintained near 1.0. High-$n$ stability calculations show, as was the case for previous results along this line [25], that with the motional Stark effect measured $q_0 = 1.3–2.0$, most of the plasma cross-section, except the very center, is in the transition region from first to second stability.

7. CONFINEMENT SCALING

In the last two years, JET and DIII-D carried out a joint program of H–mode confinement scaling studies aimed at discovering the dependence on machine size [7]. Both devices operated carefully designed single-null divertor plasmas with the same shape ($\kappa = 1.8$) and toroidal field (2–2.5 T). Confinement was studied in the ELM-free phase to avoid difficult to quantify effects of ELMs on confinement. Because both machines have the same aspect ratio 2.7, separate dependences on $a$ and $R$ were not determinable; the result is quoted in terms of a choice of linear dimension. Dimensionally correct scalings are discussed in Ref. 7. In engineering variables, the result is

$$\tau_B(\sec) = C L_p^{1.03\pm0.07} P_L^{-0.46\pm0.06} L^{1.48\pm0.09} \text{ (MA, MW, m) ,}$$

$$C = 0.106 \pm 0.011 \text{ for } L = R$$
$$= 0.441 \pm 0.044 \text{ for } L = a .$$
FIG. 5. Solid lines are the single fluid diffusivities calculated from experimental data versus normalized volume radius \( \rho \) for \( B = 1 \) T and 2 T dimensionally similar plasmas. On the basis of the 1 T result, the predicted 2 T results for Bohm (open circles) and gyro-Bohm (full circles) scaling are shown.

In another approach to deducing fundamental aspects of transport scaling, it has been shown that almost all magnetized plasma diffusion mechanisms can be cast in the form of a gyro-Bohm (gB) or Bohm (B) scaling, depending on whether the normative length scale in the problem is the short length \( \rho_S \) or the long length \( a \) respectively [12]. Local diffusivities can be written as \( \chi_{gB} = (c_S/\alpha)\rho_S^2 F_{gB} \) or \( \chi_B = c_S \rho_S F_B \), the form factors \( F \) depending only on dimensionless parameters. Any two discharges for which the dimensionless parameters are the same, and therefore \( F \), are said to be dimensionally similar and in that case one should observe either the local scaling \( \chi_{gB} \propto B^{-1} a^{-1/2} \) or \( \chi_B \propto B^{-1/3} a^{1/3} \). Two dimensionally similar L-mode discharges were constructed in DIII-D with \( a = 0.65 \) m, \( A = 2.7, \alpha = 1.70, q_{95} = 3.8 \), the same \( \beta = 1.96\% \pm 0.03\% \), and the same \( \nu_{\text{MIN}}^* = 0.128 \pm 0.001 \), but varying \( B = 1.05 \) T (1 MA) to \( B = 2.1 \) T (2 MA). The experimental results of local transport analyses for a single fluid \( \chi \) are given as the smooth curves in Fig. 5. The curves given by symbols are the predictions of the 2 T result based on scaling the 1 T result according to either the gyro-Bohm or Bohm prescriptions. The gyro-Bohm scaling is in better agreement with the experimental result at 2 T than the Bohm scaling.
FIG. 6. Ion and electron temperature and electron density profiles in the highest $T_i$, hot ion H-mode discharge.

8. HOT ION H-MODE

At modest currents (~1 MA) and with helium glow wall conditioning, it is possible to keep the Ohmic plasma density low, $\sim 1 \times 10^{19}$ m$^{-3}$. The application of strong beam heating results in a prompt L-H transition. In the subsequent long ELM-free period, although the density is rising, it remains low enough that electron-ion coupling is not large and
the ion temperature readily pulls away from the electron temperature. In this hot ion H-mode regime, \( T_i(0) \) as high as 17 keV has been reached; profiles for that discharge are in Fig. 6. The hot ion H-mode has produced the highest DD neutron rate of \( 2.7 \times 10^{15} / \text{sec} \) corresponding to \( Q_{DD} = 0.00026 \).

In going from similar limiter L-mode discharges to hot-ion H-mode discharges, transport analysis showed [11]:

1. \( \chi_e \) improves the most, about a factor of three;
2. \( \chi_\phi \) and \( \chi_e \) improve equally and are about equal for \( \rho > 0.3 \);
3. \( \chi_i \) improves only inside \( \rho < 0.5 \);
4. \( \chi_i \) is significantly less than \( \chi_e \) inside \( \rho = 0.5 \) and is within error bars of neoclassical.

9. L-H TRANSITION PHYSICS

Besides the usual H\( _a \) drop, two additional clear signatures of H-mode were found, an abrupt increase in edge plasma poloidal rotation speed and an abrupt drop in edge plasma turbulence levels (Fig. 7) [26,27]. The drop in turbulence levels occurs within 100 microseconds of the transition and is seen in both magnetic fluctuations and density fluctuations detected by a microwave reflectometer system [28]. The change in poloidal rotation speed (\( v_p \times B_T \) term) is the dominant factor in an increasingly negative radial electric field via the force balance equation [8]:

\[
E_r = \frac{1}{n_i Z_i e} \nabla P_i - (\vec{V}_i \times \vec{B})_r \quad (\text{Fig. 7}).
\]

The theory of Shaing [9] suggests the sudden increase in rotation speed arises from a bifurcation in the poloidal flow equation in which edge ion orbit loss is the mechanism to charge the plasma negatively; in which case the enhanced rotation should penetrate a few poloidal ion Larmor radii into the plasma. These observations also prompted the theory suggestion that shear in rotation could suppress turbulence by shearing apart turbulent eddies [10]; in which case one expects overlap between the region of enhanced shear and turbulence suppression.

We upgraded our edge Charge Exchange Recombination system to provide 1.5 cm edge spatial resolution and added a continuously variable frequency reflectometer system to make detailed edge measurements. Results to date show good agreement between the location of the rotational shear layer and the zone of improved confinement [29]. On the question of causality, we have some cases where the rotation change begins a few msec before the L-H transition, but no cases in which the rotation change follows the edge turbulence drop.
FIG. 7. Dominant $\vec{v} \times \vec{B}$ term in the radial electric field becomes more negative at L–H transition ($t = 1500$ ms) and less negative at each ELM (arrows). Magnetic and density fluctuations drop abruptly in H-mode.

10. EDGE LOCALIZED MODES (ELMs)

Three kinds of ELMs have been identified on DIII–D, presented in reverse order of their identification as a sub-species but in increasing order of the severity of their effects on the discharge. The existence of different types of ELMs has complicated our studies of whether ELMs are primarily pressure or current driven instabilities.

Type III. These are small amplitude ELMs which occur when the NBI power is just above threshold. As the power is increased, their frequency decreases and they disappear. The edge pressure gradient for Type III ELMs is well below the $n = \infty$ ideal ballooning limit.

Type II. These are high frequency, low amplitude ELMs which appear when the edge plasma region is put into the connection region between the first and second stable ballooning regimes by increased elongation and/or triangularity in low to moderate current plasmas ($I_p/B_T < 0.5$) [30].

Type I. Popularly called “giant ELMs,” these play the major role in H–mode discharges. The edge pressure gradient prior to these ELMs
is consistently at the ballooning limit [31]. Their frequency \( f \) increases with power and decreases as current is raised [32]. We have also found that ECH applied at the plasma edge can significantly affect the ELM frequency [17]. The role of ELMs in confinement is still a subject of controversy. In many discharges, a brisk rise in stored energy in the ELM-free phase appears to be completely arrested by the first ELM, implying the ELMs play a dominant role in H-mode confinement [33]. However, diamagnetic loop measurements show that the energy loss per ELM \( \delta E \) decreases with power and increases with current such that the ELM power \( = f \times \delta E \) is roughly constant at a low level 0.5–1.0 MW compared to input power >4 MW. These results were confirmed in a few cases by infrared TV measurements at the divertor plate [15].

11. TEN SECOND H-MODE

The efficacy of giant ELMs in conjunction with flat to hollow density profiles in preventing impurity accumulation has been well documented [34]. The ELM frequency can be continuously adjusted through plasma shape, current, and input power. An outstanding example is the 10 second long H-mode in Fig. 8; impurity levels are decreasing slowly in time [14]. The confinement enhancement over ITERP-89 is 1.4 at late times in the discharge when the ELM frequency is high, compared to an average enhancement factor of 1.7 for the DIII-D data in the ITER H-mode data base. The density is rising at a rate about 1/10 of the NBI particle input rate, indicating substantial particle absorption into the graphite.

12. HIGH RECYCLING DIVERTOR

Previous studies [35,36] have shown that the DIII-D divertor plasma operates in the high recycling regime (ion flux at the divertor plate/ion flux out of the core plasma > 10, \( n_{e,\text{div}} > n_{e,\text{sep}} \), and \( T_{e,\text{div}} < T_{e,\text{sep}} \)). Recent B2 code modelling agrees with previous ONETWO/DEGAS modelling that the density evolution in H-mode requires an inward pinch term and a ratio of \( D/V \sim 0.05 \text{ m} \). The strong density rise in H-mode is accounted for by a factor of 2 improvement in \( \tau_p \) but also by about a factor of 2 increase in the core plasma fueling rate. DEGAS calculations show this surprising result to arise from a factor of 2 reduction in divertor recycling flux being overcome by about a factor of 3-4 increase in the penetration of neutrals through the scrape-off layer because that layer is much thinner in H-mode compared to L-mode [36].
FIG. 8. Ten second H-mode impurity control via ELMs.

FIG. 9. (a) Heat fluxes across the single-null divertor target showing the reduction caused by deuterium gas injection between the separatrix strike points. (b) IR camera image of the outer separatrix strike point swept ±6 cm in major radius.

13. RADIATIVE DIVERTOR AND X-POINT SWEEPING

The high local heat fluxes at the separatrix strike points are difficult design issues in CIT and ITER. Previously, we had shown that the X-point could be swept sufficiently to lower the peak heat flux by a factor of 3 without affecting other aspects of the plasma operation. The ±6 cm
sweeping heat flux pattern on the divertor plates as viewed from above by an IR camera is shown in Fig. 9(b). Recently, we have employed gas puffing in the region between the separatrix strike points to make the divertor plasma more radiative. Deuterium gas injected at a rate of 120 torr/1/sec into a discharge with 14.5 MW NBI lowered the peak heat flux from 4 MW/m$^2$ to 2 MW/m$^2$ [Fig. 9(a)] with very little other effect on the plasma. Nitrogen injection was also fairly successful. With 20 MW NBI, the nitrogen injection reduced the target heat flux from 3 MW/m$^2$ to less than 1.5 MW/m$^2$ with only a 5%-15% drop in $\tau_E$ and $Z_{eff}$ increasing from 1.8 to 2.4 [15]. The combination of the radiative divertor and X-point sweeping is no doubt possible.

14. ELECTRON CYCLOTRON CURRENT DRIVE (ECCD)

The first step in the rf current drive program was taken with the observation of up to 70 kA of ECCD [16]. About 1.0 MW of 60 GHz rf power was launched in the fundamental X-mode from the high field side of DIII-D into discharges with Ohmic plasma current between 200 and 500 kA. These experiments were the first to be done in plasmas with strong single pass absorption (99%) and slowing down times for the current carrying superthermal electrons which are much shorter than the energy confinement time. Experiments were done with waves launched both with and against the plasma current. The loop voltage drops at constant current were compared to plasma resistivity changes using ONETWO. The bootstrap contribution was less than 10% of the total current. The loop voltage changes in excess of the resistive part give the black dot symbols in Fig. 10 for the experimentally determined driven current. The experimental $I_{rf}$ exceeds that predicted by Fokker-Planck calculations with no remaining electric field (open circles). Fokker-Planck calculations estimating tail enhancement by the remaining electric field show this effect may account for the increased ECCD efficiency.

15. ADVANCED DIVERTOR PROGRAM (ADP)

A baffle and ring electrode structure has just been installed in DIII-D (Fig. 1). The outer leg of the separatrix can be rested on this ring and current driven into or drawn from the scrape-off layer from a power supply (600 V, 20 kA). This current can either aid or oppose the main plasma current. An increase in the edge plasma current density is expected to produce local second stability and so possibly to stabilize ELMs [31,37]. Outward cross-field transport in the scrape-off layer may be either increased
or decreased depending on the sign of the applied $E_p \times B_T$ drift, perhaps affecting particle confinement or at least the scrape-off layer width [38]. Helicity injection current drive can be tested. The gas baffle between the ring and the wall enables divertor pumping experiments. Initially, studies of neutral pressure buildup will be made positioning the outer separatrix strike point either near or on the ring and puffing gas under the baffle. A cryo-pump to supply 50,000 l/sec under the baffle is being designed in collaboration with JET.

16. CONCLUSIONS

In high beta research the achieved beta values in divertor discharges have been pushed up to $\beta_T = 11\%$ in the low $q$ regime and $\beta_p = 5.1$ ($\epsilon \beta_p = 1.8$) at high $q$. At intermediate $q$, $\beta_N$ as high as 5 has been found stable and is theoretically expected to be stable for optimized profiles. The near term high leverage issue is to seek to establish whether the beta limit for optimized profiles rather generally lies at $\beta_N > 5$. Such a possibility would enhance prospects for current driven reactors operating at moderate to high $q$. 

FIG. 10. Current driven by ECH, measured (•) and Fokker-Planck predictions with (×) and without (○) electric field effects on the electron distribution function. $B_T = 2.1 \ T$, $I \times 10^{19} \ m^{-3} \leq n_e \leq 2.5 \times 10^{19} \ m^{-3}$, $200 \ kA \leq I_p \leq 500 \ kA$. 

\[ R_m = R_{\text{res}} = 1.60 \ m \]
\[ P_{\text{ECH}} = 0.9 \ MW \]
H-mode confinement in the high recycling open divertor geometry of DIII-D is about 1.7 times L-mode as represented by ITERP-89 scaling. Local transport studies favor a gyro-Bohm rather than the Bohm fundamental scaling prescription. In the hot ion H-mode regime, $T_i = 17$ keV has been reached and transport analysis reveals the principal improvement in H-mode from L-mode is a reduction of $\chi_e$. DIII-D results favor the idea that the H-mode originates from a sudden increase in poloidal plasma flow at the edge, consistent with an increased negative E-field, and resulting in sufficient shear in the flow to stabilize turbulence. Sufficient understanding of the physics and phenomenology of ELMs exists to utilize ELMs for impurity control, which has resulted in 10 second long H-mode and continuously declining impurity levels. Divertor X-point sweeping and radiation enhancement by both fuel and impurity gas puffed between the X-points have been shown able to lower peak divertor heat fluxes by factors of 3 and 2 respectively. Carbon blooms have not been observed in DIII-D despite local peak heat fluxes on the divertor tiles of 4 MW/m$^2$ and 50 MJ total heating energy injected. These results provide support for the effective use of H-mode in future devices and demonstrate that techniques are available to reduce divertor plate peak heat loads.

A clear observation of electron cyclotron current drive up to 70 kA has been made and is in accord with theory calculations. Methods of holding low density H-mode are being developed. The Advanced Divertor Program has installed a biasable ring and gas baffle structure for density control experiments. An in-vessel cryo-pump is being designed in collaboration with JET. Ion Bernstein wave heating did not produce central heating; only edge heating, parametric decay, and impurity influx were observed. The IBW antenna has been replaced with an ICRF fast wave current drive antenna as a first step toward full current drive by fast waves and EC waves with ECH (2 MW at 60 GHz and 2 MW at 110 GHz) used to produce the high temperatures needed for efficient current drive. An additional 5 MW, 110 GHz ECH system is being planned in cooperation with JAERI. An 8 MW ICRF fast wave system is also planned.
REFERENCES


DISCUSSION

S.A. COHEN: How did the quantity of gas injected (to effect power load reduction) compare with the particle flow onto the divertor plate? Also, did $n_e$ rise during the gas injection, and did $n_eT_e$ stay constant (at the divertor)?

R.D. STAMBAUGH: The deuterium puffed in was 120 torr·L/s, compared to typically 1000 torr·L/s of recycling flux. The line average density rose only slightly (~ 10%). The decreased heat flux to the divertor plate was reflected mainly in a drop in $T_e$ at the divertor plate with little change in $n_e$ at the divertor plate.

G. GRIEGER: I refer to the case with large currents through the target plates when one starts with a strongly elongated plasma. The currents then have to be carried by particles flowing onto the plate, ions on one side and electrons on the other. Did you find the potential distributions in front of the target plates which would be needed for this to happen? Otherwise you might need strong potential fluctuations.

R.D. STAMBAUGH: The halo currents in disruptions are driven mainly by high transient one turn voltages. In quiescent plasmas, we find that the naturally
occurring currents in the scrape-off layer are consistent with the thermoelectric effect (see IAEA-CN-53/G-I-3, Vol. 3).

A.K. SEN: You mentioned suppression of edge fluctuations by rotational shear. Do you have any measurements of the rotational shear length, and how do these compare with the fluctuation wavelengths?

R.D. STAMBAUGH: Yes, such measurements are covered in detail in IAEA-CN-53/A-VI-4 (Vol. 1).

A. GIBSON: You showed an ELMy H mode with an impressive duration of \(\sim 10\) s. How did the energy confinement time in this discharge compare with the L mode scaling and DIII-JET H mode scaling times?

R.D. STAMBAUGH: Early in the H phase, when the ELMs were infrequent, the confinement enhancement factor over ITERP-81 was 2.0. Late in the H phase, when the ELMs were more frequent, the enhancement factor came down to 1.4, indicating that we had paid some confinement penalty for the impurity control.

R.J. TAYLOR: DIII has a performance comparable to JET per unit cost and has similar operating regimes in terms of the normalized parameter \(n\tau_P\): \(10^{20}\) for DIII and \(10^{21}\) for JET. If you draw a line through these machines, what would you do to JET to expand it to a reactor — in terms of size, field, and so on?

R.D. STAMBAUGH: Our results suggest that the gyro-Bohm prescription is the correct way to draw a scaling line through DIII-D and JET towards CIT and ITER. Such considerations are covered in IAEA-CN-53/D-IV-7 (Vol. 2).

P. SMEULDERS: You mentioned briefly that you can control the ELMs with ECRH. Can you describe the effect of this, and the result on the ELMs?

R.D. STAMBAUGH: ECH applied just inside the separatrix increases the ELM frequency. ECH just outside the separatrix decreases the ELM frequency. The difference was a factor of three in ELM frequency (see IAEA-CN-53/E-I-2, Vol. 1).

R.W. CONN: You showed that high B is strongly favoured with respect to decreasing \(\chi\). Also, DIII-D, TFTR and JT-60 have shown high bootstrap currents, which favour a high aspect ratio. Have you done studies of the trade-off between plasma current and aspect ratio with respect to the effect on energy confinement time?

R.D. STAMBAUGH: We are trying to run a double null divertor with \(A \sim 5\) to get some information on high aspect ratio in H mode. The configuration has been made with Ohmic heating. However, the outer strike points are on relatively unprotected vessel surfaces, so we are reluctant to apply NBI power. We have more work to do on this difficult shape control problem before we can hope to get a meaningful result.
APPENDIX
THE DIII-D TEAM
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PERMANENT ADDRESS

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(l) University of California at Los Angeles, USA
(m) University of California at Irvine, USA
(n) University of Maryland, USA
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(p) Japan Atomic Energy Research Institute, Japan
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INTERPRETATION AND MODELLING OF ENERGY AND PARTICLE TRANSPORT IN JET


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Abstract

INTERPRETATION AND MODELLING OF ENERGY AND PARTICLE TRANSPORT IN JET.

The study of energy and particle confinement in JET plasmas has been performed by means of various interpretive and predictive techniques. The paper deals with the most recent and relevant results obtained, concentrating on local rather than on global analysis. Electron and ion energy transport and particle transport are studied and their relationship is examined. Comparisons with the predictions of theoretical models are presented, with special attention to transport coefficients derived from the theory of ion temperature gradient driven turbulence and to the critical electron temperature transport model of Rebut and co-workers.

1. Introduction

In this paper we report on the results of studies carried out to assess the local energy and particle transport properties on the basis of JET experimental data. The important issue of global confinement and its relationship to local transport models is discussed extensively in another contribution to this conference [2] where JET data are analysed together with data from other tokamaks collected in the so-called ITER database. The global confinement of JET plasmas can also be found in other contributions to this conference [3-5].

A proper understanding of plasma transport requires the derivation of a complete transport model applicable in all relevant regimes. It must be capable of predicting the evolution of plasma profiles in the present and next step generation of tokamak devices.

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We have not reached this goal, but we have made much progress in:

- identifying the most important phenomena to be explained, pointing out correlations between the transport of electron and ion energy and the transport of particles;
- assessing various models proposed to explain and predict the plasma performance in tokamak devices.

In Section 2 we address the problem of electron and ion energy transport. We quantify the absolute and relative magnitudes of the electron and ion heat transport coefficients. A comparison with the predictions of models based on the theory of the ion temperature gradient driven turbulence is presented. We summarise the results obtained and assess the validity of the critical electron temperature gradient model of Rebut et al. (R-L-W in the following) [1].

The evolution of electron density profiles is analysed in Section 3, and it refers mainly to full transport code simulations based on an extension of the R-L-W model to include particle transport.

Section 4 presents the results of an integrated analysis of the propagation of the electron temperature and density perturbation following a sawtooth crash and their implications for transport models. An impurity transport analysis which implies a departure from simple conventional modelling assumptions and the relevance of neoclassical transport theory, is discussed in Section 5. Some concluding remarks are given in Section 6.

2. Electron and Ion Heat Transport

There are regimes in JET where the electron and ion contributions to heat transport can be separated with the aid of interpretive codes. Among these regimes are the hot ion L and H-modes, sawtooth free discharges with strong RF heating (a "diagnostic" low power NBI is used to determine $T_i$ by charge exchange recombination spectroscopy) and pellet fuelled discharges with combined heating [6,7]. Figure 1 illustrates results for two typical cases. The first is a hot ion H-mode discharge with plasma current $I_p = 4$MA, toroidal field $B_T = 2.8$T, neutral beam injected power $P_{NBI} = 17.9$MW, peak electron density $n_e(o) = 4.8 \times 10^{19}$m$^{-3}$ and peak electron and ion temperatures $T_e(o) = 8.8$keV and $T_i(o) = 22.3$keV. The second is a monster sawtooth discharge with $I_p = 3$MA, $B_T = 3$T, $P_{NBI} = 2.6$MW, ICRH power 8.9MW, $n_e(o) = 4.5 \times 10^{19}$m$^{-3}$, $T_e(o) = 8.8$keV and $T_i(o) = 6.1$keV. The figure illustrates the spatial dependence of the electron and ion heat diffusivities $\chi_e$ and $\chi_i$, evaluated using the interpretive code TRANSP under the assumption of a diagonal transport matrix with no heat pinch term.
The ion transport is clearly anomalous, except perhaps in the central region of the hot ion H-mode. (See also [6-9] for pellet fuelled cases.) We also find as a general trend that $\chi_i \geq \chi_e$ in the outer region of the plasma while $\chi_e$ becomes comparable or more important in the central region. Plots of the temperature increase in the central region as a function of the power input per particle to electrons and ions (Fig. 2) illustrate in a crude, but rather general way, that as $T_{e0}$ increases $\chi_e$ must become larger in the central plasma region; this is not so for $T_{i0}$ and $\chi_i$. 

FIG. 1. Radial dependence of $\chi_e$ and $\chi_i$ for a hot ion H-mode (case a, pulse 20981) and an RF heated non-sawtoothing discharge (case b, pulse 19739).
FIG. 2. Increase of the ion (case a) and electron temperature (case b) as a function of $P/n_i$: $\Delta T$, $P$ and $n$ refer to $\rho < 1/3$. Sawtooth free L and H-mode plasmas with $2 \leq P_{\text{ion}} \leq 25$ MW, $I \leq n_e < 5 \times 10^{19}$ m$^{-3}$ and various combinations of ICRF and NBI are considered.
FIG. 3. JET data (referring to sawtooth free, L-mode, 3 MA pulses) compared with the predictions of one of the published versions of \( \eta_l \)-mode theory [13]: (a) ratio of measured \( \frac{\eta}{\eta_{cr}} \) to theoretical threshold value as a function of \( \frac{T_{\text{io}}}{T_{\text{eo}}} \); (b) ratio of predicted ion heat flux to total measured heat flux at \( \rho = 0.4 \), as a function of peak ion temperature.
Several attempts to relate the observed heat transport to predictions of models derived from the theory of electrostatic micro-instabilities have failed to reproduce JET results [7-10]. The most recent study [8] has considered models of the anomalous ion energy transport derived from the theory of VT_j-driven turbulence [11-13]. The analysis shows that, while there is some qualitative agreement (e.g. in pellet fuelled auxiliary heated discharges [14]) between theoretical predictions and experimental findings, there is a serious quantitative disagreement, namely:

- all models predict too low ion energy transport in the region $\rho > 0.7$ ($\rho \leq 1$ is a normalised radius) even when fully developed turbulence is taken into account;
- all models predict a large $\chi_i$ in the central and intermediate plasma region. Hence one expects $T_i$ to be determined by $\eta_i = n_i VT_i/T_iVn_i$ close to the instability threshold $\eta_i^{cr}$. Thus to reconcile theory with JET results a substantial increase of $\eta_i^{cr}$ is required (Fig. 3).

An extensive campaign of simulations using predictive transport codes has allowed the assessment of heat transport models [15]. Ohmic and L-mode discharges have been considered, covering the following range of parameters: $I = 3-5\text{MA}$, average electron density $\bar{n}_e = 1.5-7.10^{19}\text{m}^{-3}$, $B_T = 2.2-3.4\text{T}$, auxiliary power (ICRH, NBI or combined) up to $25\text{MW}$. These computations show that among the theory-based models adapted empirically to simulate experimental results, the R-L-W model emerges as rather good and complete; it covers electron and ion heat and particle transport (see also Section 3). These results confirm previous findings, carried out on a more restricted set of discharges which also included the H-mode regime [9], [16].

We recall that the R-L-W model predicts the existence of a critical electron temperature gradient $VT_{e}^{cr}$ such that the electron heat flow is given by:

$$q_e = -n_e \chi_e^{RLW} VT_e \left(1 - \frac{VT_{e}^{CR}}{VT_e} \right)$$  \hspace{1cm} (1)

when $|VT_e| > |VT_{e}^{cr}|$. The expression for $VT_{e}^{cr}$ and $\chi_e^{RLW}$ can be found in Ref. [1]. Transport is assumed to be neoclassical when $|VT_e| < |VT_{e}^{cr}|$ or $\nabla q < 0$, $q$ being the safety factor. Moreover:

$$\chi_i = \frac{Z_i}{\sqrt{1 + Z_{\text{eff}}}} \frac{VT_i^{b}}{VT_i} \chi_i^{b}, \quad D_e = \chi_e^{RL} \left(1 - \frac{VT_{e}^{CR}}{VT_e} \right)$$  \hspace{1cm} (2)
De is the electron diffusion coefficient, Z_i and Z_{eff} the ion charge and the plasma effective charge.

If, in an interpretive analysis, the $\nabla T_e^{cr}$ term in Eq. (1) is not taken into account explicitly, the resulting $\chi_e$ must be compared to $\chi_{ecr}^b$. On the other hand $\chi_{eRLW}$ applies to the analysis of heat pulse propagation (see Section 4). It should be noted that no critical ion temperature gradient is predicted for ion energy transport, consistent with the results shown in Figs. 1 and 2. Similarly no inward particle pinch is predicted for a pure plasma.

The main deficiencies with the R-L-W model have been found in the outer region of the plasma, especially at low density, where $|\nabla T_e^{cr}|$ tends to exceed the observed $|\nabla T_e|$. It is expected that the model, based on a single phenomenon [1], has to be modified here. Atomic physics processes may affect the model [17] and phenomena related to MHD instabilities might be important. An empirical solution to this problem, adopted in the predictive 1^1/2-D code JETTO, is to reduce the anomalous transport gradually when $|\nabla T_e|$ approaches and becomes smaller than $|\nabla T_e^{cr}|$. We also remark that a quantitative validation of the R-L-W model in the central region of the plasma is subject to large uncertainties owing to the dependence of $\chi_{eRLW}$ on the local shear and the practical difficulty of measuring $Vq$ in this region. A similar remark applies to other models and in particular to $\eta_i$ related transport coefficients which depend sensitively on the local shear length.

3. Simulation of the Plasma Density Evolution

JETTO code simulations have been used to study simultaneously the evolution of density and temperature profiles in the ohmic and L-mode plasmas described in Section 2. The following expression (based on the R-L-W model) have been used for the flux of the hydrogenic species:

$$\Gamma_i = -D\nabla n_i + n_i (v_{an} + v_{w})$$

$$D = \alpha \chi_e^b, \quad v_{an} = -\alpha_{in} \frac{2Dr}{a^2}$$

$v_w$ and $v_{an}$ are the neoclassical (Ware) [18] and anomalous inward pinch velocities; $\alpha$ and $\alpha_{in}$ are constants to be determined; $a$ is the minor radius.
FIG. 4. Comparison of computed and experimental $T_e$ and $n_e$ profiles for a sawtoothing RF discharge (case a, pulse 19617, $B_T = 3.1$ T, $I = 3$ MA, $P_{RF} = 8$ MW) and a sawtoothing NBI discharge (case b, pulse 20334, $B_T = 3.1$ T, $I = 3$ MA, $P_{NBI} = 8$ MW). In case b, the peaked density profile was computed with the same $\nu_m$ as in case a; the flat one was computed with $\nu_{in} = 0$. 
The impurity ion density profiles needed to compute the electron density $n_e$ are not modelled but evaluated from $Z_{\text{eff}}$ (from visible bremsstrahlung or charge exchange recombination spectroscopy). It is assumed that the impurity ions are fully ionised and this is valid for most of the plasma, when dominated by relatively low Z impurities ($Z \lesssim 8$).

We find that the time scale of the evolution of $n_e$ requires $\alpha = 0.3-0.5$. There are cases, in particular with RF heating alone and $Z_{\text{eff}} \gtrsim 2.5$, where an anomalous inward term seems to be required in the region $\rho \gtrsim 0.3-0.4$, and $v_{\text{an}}/v_{\text{ware}} \gtrsim 10$. Figure 4a illustrates such a case. Various regimes however, do not need an anomalous inward velocity, as first pointed out in [19] for discharges with pellet injection. Other cases are found in ohmic low $Z_{\text{eff}}$ discharges and with NBI heated discharges where a high density is reached, developing a flat or even hollow density profile. An example corresponding to a relatively clean plasma ($Z_{\text{eff}} \lesssim 1.5$) is illustrated in Fig. 4b, where for comparison we also show the peaked profile of $n_e$ that would be obtained with the same inward pinch as in the previous case.

4. Analysis of Heat and Density Pulse Propagation Measurements

The analysis of fast transients provides a method for determining a linearised matrix of transport coefficients from measured data and complements interpretive and predictive studies. Recent work at JET has been based on an analysis of the heat and density pulses following sawtooth crashes which takes into account the coupling between heat and particle transport [20]. The main result of the analysis is the determination of the 2x2 diffusion matrix in a system of linearised particle and thermal diffusion coefficients $D_P$ and $\chi_{eP}$ with $D_P/\chi_{eP} \lesssim 0.1$ are derived in the outer plasma region $\rho \gtrsim 0.65$ from this kind of analysis. These values are consistent with the values of $D$ and $\chi_{eRLW}$ found in the simulation of the same pulse (Fig. 6). In fact a linearisation of such a model shows that $D$ and $\chi_{eRLW}$ are dominant diagonal terms and can be compared to $D_P$ and $\chi_{eP}$.

The pulse propagation analysis shows that there is a linear coupling between particle inward pinch and negative temperature gradient. The sign and magnitude of the coupling are correlated to the initial density decrease (see Fig. 5b at $r/a = 0.69$) coinciding with the location of the maximum temperature perturbation. Such a
FIG. 5. Measurements and simulations of the propagating perturbation of $T_e$ (case a) and $n_e$ (case b); $\Delta T_e$ is normalised to the central temperature before the sawtooth crash; $\Delta n_e$ is given by the phase changes of the reflectometer. Pulse 19617 is considered, as in Fig. 4a.
term is not found linearising the R-L-W model for a pure plasma. It is possible that it is due to an anomalous inward flux of impurities proportional to $\nabla T_e$, related to a thermoelectric force in the direction parallel to the magnetic field. This has not been considered in the model so far.

5. Transport of Impurities

The study of impurity transport has largely been based on the calculated evolution of impurity density profiles in prescribed experimental plasma profiles. Computed and experimental emissivities from individual lines [21,22] and soft X-ray profiles [23] and radiation profiles are compared. The conventional approach assumes the impurity flux $\Gamma_I$ to be given by:

$$\Gamma_I = -D_I \nabla n_I + n_I v_I$$

$$v_I = -\alpha_I D_I r / a^2$$

(4)
where $n_i$ is the impurity density. The diffusion coefficient $D_i$ and the convective velocity $v_i$ are to be determined from the simulations and are assumed to be the same for all ionisation states.

In many cases the simple approach with $D_i$ and $\alpha_i$ radially constant fails. For example impurity transport appears to be reduced to a level close to neoclassical in the central region of "good confinement" in pellet fuelled discharges [21] and also in the central region of ohmic and RF heated discharges [23].

The simple approach (4) with constant coefficients also fails for the analysis of H-mode discharges: a rather slow diffusion coefficient $D_\parallel \approx 0.1 m^2/s$, and a spatial variation of $v_i$ which depends on the plasma temperature and density profiles [22] is needed. Figure 7 shows the empirically determined spatial profile of $v_i$ during the H-phase of a high density pulse ($n_e \geq 7 \times 10^{19} m^{-3}$, developing a hollow profile, $Z_{eff} \approx 2$, decreasing with time, $I_p = 3.1 MA$, $B_T = 2.2 T$, $P_{NBI} = 8 MW$). Before density steady state, in a phase lasting 3.5 seconds, the velocity $v_i$ has to be outward in the region $\rho \lesssim 0.8$ in order to explain the nickel emissivity lines and radiated power; the empirical convective velocity $v_i$ is consistent, within the large uncertainties of the analysis, with neoclassical theory [24]. This situation is
reminiscent of the absence of an anomalous inward pinch term in the modelling of the electron density evolution in L-mode plasmas as shown in Fig. 4b.

6. Conclusions

The analysis carried out at JET shows that no pure theoretical model is sufficiently correct or complete enough to explain all features exhibited by the wide variety of plasma regimes found in JET. We underline the problems associated with a quantitative assessment of such models using the theory of instabilities driven by $\nabla T_e$ as an example. Time dependent simulations of a variety of phenomena evolving on different time scales indicate a correlation between energy and particle transport. In particular, we find that $D/\chi_e$ as normally defined in interpretive codes is larger than $D^P/\chi_e^P$ from pulse propagation analysis. We also find that neoclassical theory may be relevant for impurity transport under various circumstances. These results are generally consistent with the picture of the energy and particle transport assumed by the critical electron temperature model of Rebut et al. [1]. This model, which has been successfully tested in a variety of JET plasma conditions, is a good candidate for predicting tokamak plasma performance.

Acknowledgements

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DISCUSSION

F. ROMANELLI: You have pointed out two of the difficulties of using present $\eta_i$ mode theories to explain set data, namely the high value of the quasi-linear flux and the low value of the $\eta_i$ threshold. However, the $\eta_i$ threshold will be higher if parallel ion dynamics is taken into account. In addition, numerical simulations of $\eta_i$ mode turbulence show a reduction of the flux with respect to the quasi-linear estimate.

A. TARONI: I take your point. However, what I wish to illustrate is that the predictive capability of theoretical models based on the theory of $\eta_i$ instabilities and turbulence is rather poor. This seems likely to remain true, even though I cannot exclude the possibility that the theory might be improved.
R.J. GOLDSTON: Could you tell us whether the factor \((T_e/T_i)^{1/2}\) in the ratio \(\chi_i/\chi_e\) of the RLW model is part of the theoretical model, or part of the empirical adjustments which you referred to?

A. TARONI: To my knowledge the expression for \(\chi_i\) in the RLW model has the same theoretical basis as that of \(\chi_e\). A derivation of both \(\chi_e\) and \(\chi_i\) was given in a paper presented by Rebut and co-workers at the IAEA Conference in Nice (see Plasma Physics and Controlled Nuclear Fusion Research 1988, Proc. 12th Int. Conf. Nice, 1988, Vol. 2, IAEA, Vienna (1989) 191). For further details, you should address the authors of that paper.

K. ITOH: I would like to ask you about the off-diagonal term in your 2 \(\times\) 2 transport matrix. How large is it compared to the diagonal term? If it is finite, does that not contradict the statement that the anomalous (inward) pinch is not necessary in many JET discharges?

A. TARONI: The off-diagonal term in the 2 \(\times\) 2 transport matrix — derived from the analysis of heat and density pulse propagation — shows that there is indeed a non-negligible inward pinch for the case I have shown. However, I have also pointed out that for the same case (and similar cases with rather high \(Z_{\text{eff}}\)) an inward pinch is also found from our analysis of the evolution of the electron density.
ADVANCES IN TRANSPORT UNDERSTANDING USING PERTURBATIVE TECHNIQUES IN TFTR


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Abstract

ADVANCES IN TRANSPORT UNDERSTANDING USING PERTURBATIVE TECHNIQUES IN TFTR.

The parametric dependence of transport in TFTR plasmas is investigated using perturbative techniques. By perturbing the current in L-mode plasmas, the thermal transport outside $r \sim a/2$ is found to be insensitive to the plasma current or its shape. As a consequence, it is found that the global confinement can be strongly enhanced or reduced relative to L-mode scaling at fixed $I_p$ by modifying the current profile. The density profile has been perturbed in supershot plasmas to force the plasma far above the stability threshold for ITG-modes. The ion transport in supershots is found to be not governed by marginal stability to the ITG-mode. Density pulse propagation studies of L-mode plasmas indicate that the electron particle transport varies as $T_e^{1.5-2.5}$ and that the diffusivity of trace amounts of helium is similar to the thermal diffusivities. Both measurements are consistent with electrostatic transport models.

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Local transport has been analyzed in ohmic and auxiliary-heated TFTR plasmas subjected to a variety of perturbations. The perturbations separately vary parameters (e.g. $L_n$ and $L_s$) that are correlated in steady-state situations. Careful analysis during these variations allows the dependence of the transport on individual parameters to be determined and compared with theoretical predictions.

This paper presents results from experiments that dynamically vary (1) the current profile to identify the source of $I_p$ scaling in plasmas having L-mode confinement, (2) the density and $T_i$ profiles to test Ion-Temperature-Gradient mode marginal stability in peaked density-profile supershots [1], and (3) the density profile to measure the temperature dependence of electron particle transport and to measure helium transport coefficients.

1. Current-Profile Perturbations

The energy confinement time in neutral beam heated plasmas has been found to increase roughly linearly with the plasma current $I_p$ (so-called $I_p$ scaling) in both the L-mode and H-mode regimes in a large number of experiments. This is not theoretically predicted by most transport mechanisms applicable to the collisionless interior region of tokamak plasma. In particular, most electrostatic fluctuation-induced transport models typically predict a strong toroidal field $B_t$ dependence (which is not observed), little or no direct $I_p$ dependence, and various dependencies on the magnetic shear. In consequence, “multi-regime” models[2,3] have proposed that $I_p$ scaling is primarily due to resistive-ballooning turbulence[4] in the outer regions of the plasma, where it is predicted to have a strong poloidal field $B_p$ dependence.

In TFTR, neutral beam heated plasmas obtained with a high limiter recycling coefficient show energy confinement times in good agreement with Goldston L-mode scaling[5]. The high recycling coefficient can be achieved either by saturating the carbon limiter with deuterium gas, or by introducing helium (which is not significantly pumped by carbon) into the plasma. These plasmas have broad density profiles, in contrast to the peaked density profile enhanced confinement supershot regime, obtained with low limiter recycling coefficients, which shows no dependence on $I_p$[6]. In steady-state TFTR plasmas having L-mode confinement, transport analysis indicates[7] that both the electron and ion thermal diffusivities ($\chi_e$ and $\chi_i$) vary with $I_p$ across the entire plasma cross-section.

In order to determine the local dependence of the thermal transport on the plasma current and its profile, and thus determine the origin of current scaling, a series of experiments contrasted plasmas in approximate steady-state with cases where $I_p$ was rapidly changed from 2 MA to 1 MA (or vice-versa)
FIG. 1. Time evolution of (a) \( I_p \), (b) \( W_{\text{tot}} \), measured by the diamagnetic loop, (c) \( \lambda_i \), and (d) \( V_{\text{surf}} \) for deuterium plasmas with \( I_p = 1 \) MA, \( I_p = 2 \) MA, and \( I_p \) decreased from 2 to 1 MA. For these plasmas \( B_i = 4.8 \) T, \( R = 2.45 \) m, \( a = 0.8 \) m, and \( \bar{n}_e = 3.4 \times 10^{19} \) m\(^{-3}\).

FIG. 2. Ratio of the measured \( \tau_E \) (from the diamagnetic loop) to the L-mode scaling prediction versus \( \ell_i \) for \( I_p = 1 \) MA plasmas with \( \bar{n}_e = (3.4-8.1) \times 10^{19} \) m\(^{-3}\) and \( P_B = 6-13 \) MW, (x) in steady state, and for \( I_p \)-decrease cases 1.5 sec after starting the reduction of \( I_p \) from 2 to 1 MA at a rate of (□) 1 MA/sec, (○) 2 MA/sec, (△) 3.3 MA/sec.
during 2 seconds of balanced co- and counter-tangential neutral-beam heating, as shown in Fig. 1. The rapid change in $I_p$ produced large transient variations of the measured surface voltage $V_{\text{sur}}$ and internal inductance $\ell_i$ indicating that the $j_{||}$ and $q$ profiles were perturbed. The approximate value of $\ell_i$ is obtained by subtracting the diamagnetic $\beta_{p\perp}$ from $\Lambda \equiv \beta_p^{\text{eq}} + \ell_i/2$. Kinetic calculations indicate that $\beta_p^{\text{eq}}/\beta_{p\perp} \sim 1.05$ for these plasmas. The rate of $I_p$-increase was limited to 1 MA/sec, due to the onset of edge MHD activity\cite{8}. The rate of $I_p$-decrease was varied from 1 to 3.3 MA/sec without any detected MHD activity. During the steady state and $I_p$-decrease conditions, the line average density $\bar{n}_e$ was maintained at a constant value by gas puffing. However, during the $I_p$-increase conditions, the density rose without external gas puffing. The energy confinement times $\tau_E$ of the steady-state plasmas are in good agreement with L-mode scaling. Following the perturbation, the energy confinement time was observed to not follow the L-mode relationship ($\tau_E \propto I_p$), but relaxed towards it on a long time-scale ($\gg \tau_E$) similar to the resistive equilibration time (which is several seconds for the case shown). Since $B_p$ and $q$ near the edge of the plasma depend on $I_p$ and not (to lowest order) on the current profile, this indicates that $I_p$ scaling is not due to a direct dependence on the edge value of these quantities.

The $I_p$-decrease experiments have been performed at a variety of densities ($\bar{n}_e = 3.4 - 8.1 \times 10^{19} \text{ m}^{-3}$), beam heating powers (6 - 13 MW), and with deuterium and helium background plasmas, to vary the achieved $T_e$ and thus the resistive equilibration time. In all cases, $\tau_E$ responds to the change in $I_p$ on a long time-scale, which is shorter for the colder (high density or low power) plasmas. The enhancement of $\tau_E$ over L-mode scaling (for $I_p = 1$ MA) at the end of the beam pulse (1.5 seconds after the beginning of the decrease in $I_p$) is directly correlated with the $\ell_i$ at that time, as shown in Fig. 2. For the lowest densities and higher beam powers, after completion of the $I_p$-decrease (independent of decrease rate), the plasma showed evidence of transition into the H-mode (decreased $H_\alpha$ emission, small ELMs) and a 5% increase in stored energy, as observed on JIPP T-IIU\cite{9}. In Fig. 1 this occurs at $t = 4.5$ sec, $\sim 10 \tau_E$ after the start of the $I_p$ decrease.

The evolution of the current profile has been calculated by the time-dependent transport analysis code TRANSP\cite{10,11} by numerically evolving the poloidal field diffusion equation using the measured time evolution of $Z_{\text{eff}}$ and $I_p$, and the $T_e$, $n_e$, and $T_i$ profiles to calculate the neoclassical resistivity, bootstrap current and beam-driven currents. The calculated evolution of $V_{\text{sur}}$, $\ell_i$, and the radius of $q = 1$ are in good agreement with the measured values, as previously found for ohmic\cite{12} and beam-heated plasmas\cite{13}. Measurements of the $q$ profile using lithium pellet injection\cite{14} are consistent with the calculation, but have large enough statistical uncertainties that they do
FIG. 3. Calculated radial profiles of $q$ and $j_i$ for the plasmas of Fig. 1 with $I_p = 1 \text{ MA}$, $I_p = 2 \text{ MA}$, and at several times during the evolution of the plasma where $I_p$ decreased from 2 to 1 MA.

not constrain the calculation. The calculations indicate, see Fig. 3, that a skin current forms that resistively penetrates to $\sim a/2$ (for the plasma shown) during the available beam-heating pulse. This is confirmed by the observed evolution of the sawtooth inversion radius, indicative of the $q = 1$ radius, which is just beginning to change at the end of the beam-heating pulse. The changes in the $q$ profile are calculated to transiently increase (decrease) the shear $\dot{s} = r \partial_r \log q$ by a factor of $\sim 2$ during the $I_p$-decrease (-increase) sequence. During the change in $I_p$ and the resistive penetration of the current perturbation, the measured $T_e$, $T_i$, and $n_e$ profiles are not affected, see Fig. 4, until the current perturbation reaches the plasma core. As a consequence, the experimentally inferred $\chi_i$ and $\chi_e$ do not change, implying that the transport in the outer region ($r > a/2$) cannot depend strongly on $j_{||}$, $q$, or $\dot{s}$. In
FIG. 4. Measured radial profiles of $T_i$ and $T_e$ at sawtooth peaks for the plasmas of Fig. 1 with $I_p = 1 \text{ MA}$, $I_p = 2 \text{ MA}$, and at several times during the evolution of the plasma where $I_p$ decreased from 2 to 1 MA.

In particular, these observations imply that transport due to resistive ballooning modes [4] (which varies as $\sim q^{5/3}/s$) or the Rebut-Lallia model[15] (which varies as $(\nabla T_e - C)q/s$ with $C$ varying as $J_{\|}^{1/2}/q$), as presently understood, cannot be significant in these plasmas in this region.

These observations indicate that L-mode $I_p$ scaling is due to the current in the core of the plasma, perhaps near the $q = 1$ surface. Thus, the observed steady-state dependence of $\chi_e$ and $\chi_i$ on $I_p$ across the entire profile must be due either to an additional parameter, correlated with $I_p$ in steady state, or a non-local transport dependence on the current profile, perhaps due to a global instability or radial turbulence propagation. In addition, they indicate that $\tau_E$ can be substantially enhanced (or suppressed) relative to L-mode scaling for a fixed $I_p$ value by modifying the current profile.
2. Test of ITG-mode Marginal Stability

Previous studies of steady state transport in supershots[16] have shown that $\chi_i$ and $\chi_e$ are roughly equal, are much larger than neoclassical predictions, and are reduced for plasmas with peaked density profiles. In addition, it was found[17] that the ion temperature and density profile shapes were correlated such that the plasma was close to the theoretical [18,19] marginal stability for ion temperature gradient driven turbulence (ITGDT). Due to the large values of $T_i/T_e$ (2 – 3 for supershots) the flat-density corrections to the critical $L_T \equiv -(\partial_e \log T_i)^{-1}$ threshold for ITG-mode stability were found to be important. Since most theoretical estimates predicted more ITGDT transport[20] than is observed, it was expected that ITGDT should be able to enforce near marginal stability.

To test the ITGDT marginal stability of these plasmas, the peaked density profiles of supershots have been transiently broadened, using either a deuterium pellet or a helium gas puff. The gas pulse perturbed the outer half of the plasma, while the pellet size could be chosen to produce perturbations at various radii. If marginal stability to ITGDT is controlling the transport in these plasmas, the ion thermal transport should adjust during the perturbation to attempt to keep the plasma near the stability boundary.

The supershots studied had $I_p = 1$ MA, $B_T = 4.8$ T, $R = 2.45$ m, and 14 MW of balanced co- and counter-tangential $\sim 100$ keV deuterium neutral beam injection. This produced a plasma with $n_e(0) \approx 5 \times 10^{19} \text{ m}^{-3}$, $n_e(0)/\langle n_e \rangle \approx 2.3$, $T_e(0) \approx 8$ keV, $T_i(0) \approx 25$ keV, and $\tau_E \approx 2.7 \tau_E$.

The consequences of the two types of perturbations are the same: the density profile is locally broadened ($L_n \equiv -(\partial_e \log n_e)^{-1}$ increased by a factor $\sim 8$), but the $T_i$ profile becomes slightly steeper ($L_{T_i}$ dropped), and the plasma is driven far from the theoretical ITG stability boundary, as shown in Fig. 5. The theoretically predicted[18,19] critical-$L_T$ for ITG-mode linear instability increases during the perturbation, making the plasma more unstable, due to a decrease of $T_i/T_e$ and the increase in $L_n$ during the perturbation. This variation of the theoretical critical-$L_T$ has been confirmed by numerical calculations[21] for the measured experimental parameters.

The time-evolving thermal energy transport in these plasmas has been analyzed by TRANSP using the measured temperature and density profiles. The analysis indicates that $\chi_i \equiv Q_i/n_i \nabla T_i$ is not changed by the perturbation (and if anything, drops slightly). Analytic predictions of the $\chi_i$ for fully turbulent ITGDT[20] are a factor of 3 to $> 30$ times larger than the observed values (varying with minor radius), and are large enough that ITG marginal stability should have been enforced. Comprehensive kinetic quasi-linear numerical calculations in toroidal geometry[22] of drift modes driven
by both ITG and trapped-electron dynamics have been carried out using the experimentally measured plasma parameters. Initial results indicate that the growth rate and wavelength of the most unstable mode are approximately unchanged by the perturbation. However, the calculated real frequency drops by a factor $\sim 3.5$ and the propagation direction changes from the electron-diamagnetic direction (in the target supershot) to the ion-diamagnetic direction (during the perturbation), indicating the predicted domination by ion dynamics. Similar calculations, in toroidal geometry, by a code solving the local dispersion relation with ion dynamics[23] and by a gyrokinetic ballooning code with ion dynamics[21] also indicate that the growth rate and wavelength for the most unstable mode are almost unchanged by the perturbation. These linear calculations indicate that the total amount of transport should not be changed by the perturbation, and suggest that nonlinear theories of ITG turbulent transport might be brought into agreement with the experiment by a better treatment of toroidal and kinetic effects.

We conclude that the anomalous ion thermal transport in supershots is not controlled by ITGDT marginal stability, as presently understood. In particular, the observed [16,17] parametric variations of $\chi_i$ and $\chi_e$ in the supershot regime (both decrease with increasing $T_i$, $T_e$, or $\beta_P$, contrary to almost all theoretical predictions) must not be artifacts of ITGDT marginal stability.

FIG. 5. Time evolution at $r = 0.3$ m of (a) measured $L_T$, during a supershot perturbed by a deuterium pellet and predicted marginally stable $L_T$ from (b) Ref. [19], (c) Ref. [18], and (d) numerical solution [21].
3. Electron and Helium Particle Transport

Previous density perturbation experiments in ohmic plasmas[24] have been extended to neutral-beam heated plasmas in order to determine the temperature scaling of the electron particle transport. In this experiment[25], a small amount of helium (0.75 torr-l) was puffed into steady-state plasmas heated by balanced co- and counter-tangentially injected deuterium neutral beams, with \( P_b = 0, 4.5, 9.0, \) or 14 MW. The ohmic heating power is 1.1 MW in the plasma with \( P_b = 0 \), and is 0.6 MW when \( P_b = 14 \) MW. For these plasmas, \( B_t = 4 \) T, \( I_p = 1.5 \) MA, \( R = 2.58 \) m, \( a = 0.93 \) m, and \( n_e = 3.1 \times 10^{19} \) m\(^{-3} \). The resulting density profiles are very similar, the central \( T_e \) increases from 2.2 to 4.4 keV, \( T_i/T_e \) increases from 0.8 to 1.25, and \( Z_{eff} \) increases from 2.8 to 3.6 with heating power. The target plasmas were doped with helium to ensure high limiter recycling and near L-mode confinement.

The perturbed electron particle flux \( \delta \Gamma \) was calculated from the continuity equation, the time evolution of the density profile, as measured by a 10 channel far infrared interferometer array, the neutral-beam fueling, as calculated by TRANSP, and the edge source, assuming a fixed global particle confinement time of 0.1 s. The gas puff is found to only significantly perturb \( n_e \) and \( V_{\text{rad}} \), so the perturbed flux is expanded as

\[
\delta \Gamma = \left( \frac{\partial \Gamma}{\partial V_{\text{ne}}} \right) \delta V_{\text{ne}} + \left( \frac{\partial \Gamma}{\partial n_e} \right) \delta n_e,
\]

and the coefficients \( \partial \Gamma/\partial V_{\text{ne}} \) and \( \partial \Gamma/\partial n_e \) are fit to the data in the manner of Ref. 24. These coefficients are often interpreted as the particle diffusivity \( D \) and convective velocity \( V \), respectively. However, in previous studies[24,26] they have, themselves, been found to be functions of \( n_e \), indicating that the particle transport is nonlinear, as predicted by many theoretical transport models[27]. In this case, each fit coefficient must be expected to include contributions from any diffusive and convective processes[28]:

\[
\begin{align*}
\frac{\partial \Gamma}{\partial V_{\text{ne}}} &= -(D) - \frac{\partial D}{\partial n_e} \langle V \rangle - \frac{\partial V}{\partial n_e} \langle n_e \rangle, \\
\frac{\partial \Gamma}{\partial n_e} &= \langle V \rangle + \frac{\partial V}{\partial n_e} \langle n_e \rangle - \frac{\partial D}{\partial n_e} \langle V \rangle,
\end{align*}
\]

where \( \langle \rangle \) denotes equilibrium values.

The radial profiles of these coefficients have been determined for each of the conditions in the temperature scan, as shown in Fig. 6. The central core and outer edge of the plasmas have not been analyzed due to uncertainties from sawtooth induced particle transport and the wall recycling source. At each radius, the coefficients increase with increasing \( T_e \), which increases with heating power. This variation with \( T_e \) has been fit as an exponential
FIG. 6. Radial profile of the fit coefficients (a) $\frac{\partial T}{\partial \varphi} n_e$ and (b) $\frac{\partial \Gamma}{\partial n_e}$ for each of the conditions in the temperature scan.

FIG. 7. Radial profile of exponent in fit to the form $T_n^e$ for (○) $\frac{\partial T}{\partial n_e}$, (●) $\frac{\partial \Gamma}{\partial \varphi} n_e$, (×) $\chi_{avg}$, (■) toroidal quasi-linear calculation [22], and (---) slab quasi-linear calculation [29].
dependence, $T_e^p$, and the radial profile of the exponent for each coefficient is shown in Fig. 7. It is also found that the single fluid thermal diffusivity, $\chi_{\text{avg}} \equiv -(Q_e + Q_i)/(n_i \nabla T_i + n_e \nabla T_e)$, calculated from the steady-state power balance, increases with $T_e$. A similarly defined temperature exponent for it is also shown. The temperature exponents observed for the particle-flux fit coefficients and for the thermal diffusivity are between 1.5 and 2.5. These values are in the range predicted (1.5 to 3.5) for trapped-electron microinstabilities[29] for the experimental plasma parameters. In addition, two quasilinear numerical codes were used to calculate the temperature exponent from microinstability theory for the experimental plasma conditions. The first is the comprehensive kinetic microinstability calculation in toroidal geometry of Ref. 22. For $r < 0.65$ m, $\gamma_i > 2$ for all of the plasmas, and this calculation indicates that ITG-mode behavior should dominate. The temperature exponents of the quasi-linear transport coefficients were calculated and the average of the temperature exponents for the four experimental conditions at $r = 0.5$ m is shown in Fig. 7. The average temperature exponent for the four discharges was also calculated over the radius $r = 0.35$ – 0.6 m with a slab quasilinear model[29] of trapped-particle transport assuming that the ITG-mode is dominant. For both models, the numerically calculated temperature exponents are close to the experimentally measured values.

In addition, the radial variation of $\Gamma$ is consistent with the electrostatic-turbulence forms $(eT_e^{1.5-2.5}/n_e L_{n_e}^2)\nabla n_e$ and $(eT_e^{1.5-2.5}/n_e L_{n_e} L_{T_e})\nabla n_e$, where $\epsilon = r/R$. The flux is consistent with either form alone, or a linear combination, together with just the neoclassical Ware pinch.

Similar experiments have been performed to measure the transport of trace amounts of helium[30] in order to empirically predict the behavior of helium ash in future devices. Small amounts of helium ($n_{He}/n_e \sim 0.02$) were puffed into deuterium plasmas with $R = 2.45$ m, $B_t = 4$ T, $I_p = 1.4$ MA, $\bar{n}_e = 2.5 \times 10^{19}$ m$^{-3}$ heated with 7 MW of co-tangential neutral beams. The central plasma parameters were $T_e \sim 5$ keV, $T_i \sim 10$ keV, $Z_{\text{eff}} \sim 4$, toroidal velocity $v_\phi \sim 2.5 \times 10^5$ m/sec, and $\tau_B = 1.3 \tau_i^{1/2}$. The time evolution of the $n_{He}$ profile was measured directly using charge exchange recombination spectroscopy, and analyzed using the impurity transport code MIST. The data are consistent with a linear transport process, including a large anomalous “pinch” ($\sim 2$ m/sec inwards at $r \sim a/2$) that increases exponentially radially (in magnitude). This contrasts with the nonlinear electron transport considered above, but may be consistent with theoretical expectations for a trace impurity in turbulence determined by the background plasma parameters. The inferred $D_{He} \sim \chi_i \sim 2 \chi_e$ across the radial profile, roughly as expected for electrostatic mechanisms and similar to previous measurements of $D_e$ [24], where $\chi_i$ and $\chi_e$ were determined from the steady-state power balance.
4. Summary

Transport has been analyzed in auxiliary heated TFTR plasmas subject to perturbations which (1) varied the current profile in L-mode plasmas to investigate the source of $I_p$-scaling, (2) varied the $n$ and $T_i$ profiles to test whether supershot transport is controlled by ITG-mode marginal stability, and (3) perturbed the $n$ profile to measure the temperature variation of the particle transport and to measure the transport of trace amounts of helium.

The local thermal transport and global confinement in the L-mode are found to be insensitive to the current and current profile outside $r \sim a/2$, but do depend on the current in the core of the plasma. The data show that, for fixed $I_p$, $\tau_E$ can be substantially increased (or decreased) relative to L-mode scaling by peaking (or flattening) the current profile.

The anomalous transport in supershots is found to not be controlled by ITG-mode marginal stability. The observed ion transport does not increase when the plasma is far beyond the theoretical ITG threshold. This may be consistent with initial results from linear numerical calculations in toroidal geometry including kinetic effects.

The particle and thermal transport in L-mode plasmas has been found to increase with increasing $T_e$, roughly as $T_e^2$, in agreement with numerical calculations of transport due to drift-type electrostatic microinstabilities. The particle diffusivity for helium transport in a deuterium background plasma has been measured to be similar in magnitude and shape to the thermal transport coefficients, as expected for electrostatic processes.

Acknowledgments

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REFERENCES


DISCUSSION

S.I. ITOH: In the current ramp-down experiment, did you observe any changes in toroidal and/or poloidal rotation? The radial momentum balance equation tells us that if you change $I_p$ and, accordingly, $B_p$, either $U_t$ or $U_p$ or both will change, provided that $E_{\tau}$, $\nabla P_{\tau}/n$ remain unchanged. If $U_t$ and $U_p$ remain unchanged, this implies a change in $E_{\tau}$.

M.C. ZARNSTORFF: We have not investigated any possible changes in the poloidal rotation. The measurements indicate that any changes in toroidal rotation are relatively small.
R.J. TAYLOR: Your poloidal rotation must be measured even if it is at the diamagnetic drift level. Do you believe that its value can affect your analysis?

M.C. ZARNSTORFF: As I said, we have not investigated any possible changes in the poloidal rotation. It would have to change throughout the region outside \( r \sim a/2 \) to affect our conclusions. In addition, it would be somewhat surprising if it changed in such a way as to precisely counteract the changes in \( q, \delta, \) or \( j_b \).

A.K. SEN: This question relates to your \( \eta_i \) mode study. The reason for the disagreement between theory and experiment may be that the formulas used for \( \eta_{\text{crit}} \) are not general enough. For example, these formulas do not include the different roles of \( \eta_{iA} \) and \( \eta_{iA} \), as well as the effect of the non-thermal ion population.

M.C. ZARNSTORFF: In general, your comment is correct. However, the initial results of numerical calculations including the non-thermal ions indicate that they do not change the ITG mode stability significantly for these plasmas.

B. COPPI: Could you elaborate on your statements about the importance of the current in the central region of the plasma column? Do your results imply that the central current density values are important?

M.C. ZARNSTORFF: The data indicate that the current in the central region is controlling \( I_p \) scaling in the L mode. We cannot say yet whether it is the current density or current profile in this region that is important. The confinement is observed to change roughly as the \( q = 1 \) radius changes.
TRANSPORT IN AUXILIARY-HEATED, HOT-ION H-MODE AND L-MODE DISCHARGES IN THE DIII-D TOKAMAK*

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Abstract

TRANSPORT IN AUXILIARY-HEATED, HOT-ION H-MODE AND L-MODE DISCHARGES IN THE DIII-D TOKAMAK.

Local thermal and angular momentum transport in the bulk of neutral-beam-heated, hot-ion L-mode and H-mode plasmas in DIII-D has been studied. This work has been done at reactor-relevant central ion temperatures up to 14.5 keV and normalized beta values up to $\beta_n = 3$, which approach the beta limit seen in DIII-D. Significant improvement in local transport has been seen after the L to H transition. This demonstrates that the confinement improvement in H-mode is not simply due to the creation of a transport barrier at the plasma edge. The dominant improvement is in the electron thermal and angular momentum transport. Ion transport is at the neoclassical level in the center of the H-mode discharges. The local transport in the outer half of the plasma improves markedly within 10 ms after the L to H transition. The electron density profile flattens in this region on the same time scale. This connection leads to the speculation that the confinement improvement is caused by this flattening. Comparison of experimental results with the predictions of electrostatic, drift-wave-induced transport shows significant disagreement. Magnetic-fluctuation-based models agree better with experiment, but still exhibit significant discrepancies. Complete, quantitative agreement between theory and experiment will require further improvement in the theory.

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1. INTRODUCTION

Improving energy confinement and better understanding of energy transport are major goals of the tokamak program. We have been working towards these goals by investigating local transport in hot-ion H-mode and L-mode discharges in the DIII-D tokamak. Improved understanding of the H-mode, first discovered on ASDEX [1], is particularly important because the H-mode is one of the most robust and ubiquitous of the improved confinement regimes seen over the past few years. The H-mode is particularly reactor compatible, since it is an improved confinement regime that can persist for long periods of time without the need for extreme wall pumping. H-modes lasting as long as 10 seconds have been demonstrated in DIII-D [2].

The most obvious confinement improvement at the L to H transition occurs at the plasma edge, where a transport barrier forms just inside the separatrix after the transition [3,4]. Creation of this barrier is connected with decreased edge density fluctuations, magnetic fluctuations and changes in the edge poloidal rotation and electric field [5-10]. Based on this, one might be tempted to conclude that all the H-mode confinement improvement is caused by the edge transport barrier. However, there is also a significant improvement in energy confinement throughout the bulk of the plasma after the transition [5,11-15]. For the present work, the bulk of the plasma is the region inside magnetic flux surface label \( \rho = 0.85 \); the edge transport barrier exists in the region outside of this [10].

The goal of the present work is to investigate the bulk confinement improvement in DIII-D at the L to H transition in order to demonstrate its existence more clearly and to try to determine its cause. Our experiments have shown that, at constant external conditions, the bulk confinement improvement the H-mode is due to a simultaneous improvement in electron thermal diffusivity \( \chi_e \) and ion thermal diffusivity \( \chi_i \), although the improvement in \( \chi_e \) is the larger of the two. Angular momentum diffusivity \( \chi_\phi \) also improves by about the same amount as \( \chi_e \). In addition, by doing a detailed time history across the L to H transition, we find that the bulk confinement improvement takes place in about 10 milliseconds after the transition, a time significantly shorter than the global energy confinement time for the discharges in question but similar to the time that it takes the electron density gradient to evolve in the outer half of the plasma. The apparent correlation between confinement improvement and density profile evolution has motivated us to consider theories which predict decreased transport for flatter density profiles [16-20]. One of the most elaborate theories to date of electrostatic, drift-wave driven transport [17] is, however, neither qualitatively nor quantitatively consis-
tent with the experimental results. The magnitude of the diffusivities as well as the improvement across the L to H transition agrees better with the electromagnetic-fluctuation-based Rebut-Lallia model [18,19] than it does with the electrostatic drift wave theory [17], although there are still discrepancies here, too. Quantitative agreement between theory and experiment will require further improvements in the theory.

Our transport experiments have been done in the hot-ion mode of operation in order to better separate the electron and ion transport channels. Utilizing the improved confinement produced by deuterium beam injection into deuterium H—mode discharges [7], we have produced hot-ion discharges with ion temperatures $T_i$ in the center of the plasma up to 17 keV at moderate line-averaged densities of $\bar{n}_e = 3 \times 10^{19} \text{m}^{-3}$. Most of our transport work has been done for central $T_i$ around 12 to 14.5 keV in H—mode and 7.5 to 11 keV in L—mode. The H—mode shots in this sequence also reach normalized beta values $\beta_N = 3$, which approaches the beta limit seen in DIII—D [2]. Accordingly, we are studying transport at reactor-relevant temperature and beta values.

2. COMPARISON OF H—MODE AND L—MODE TRANSPORT WITH CONSTANT EXTERNAL CONDITIONS

Most previous comparisons of bulk transport in H—mode and L—mode were done by analyzing the transport in the L—phase and the H—phase of one shot or a series of similar shots [11—13]. Because the plasma density typically rises after the L to H transition, this results in transport comparisons that are made at systematically different densities. If the local transport depends on density, this systematic difference could have affected the confinement comparison. We have made detailed comparisons of energy and angular momentum transport between deuterium L— and H—mode discharges with the same external conditions: line averaged density $\bar{n}_e = 3.5 \times 10^{19} \text{m}^{-3}$, plasma current $I_p = 1.0 \text{MA}$, toroidal field $B_T = 2.1 \text{T}$, and neutral beam input power $P_B = 8.7 \text{MW}$.

One of the major sources of uncertainty in the $\chi_e$ and $\chi_i$ inferred from transport analysis is the electron-ion transfer term. This term is of the form [20]

$$3 \frac{m_e}{m_i} \frac{n_e}{T_e} (T_e - T_i)$$
Fig. 1. Comparison of electron and ion temperature profiles for hot-ion H-mode and L-mode discharges. The curves are plotted as a function of flux surface label $\rho$, which is proportional to the square root of the toroidal flux enclosed within a given flux surface. The electron temperature comes from Thomson scattering measurements and electron cyclotron emission measurements; the ion temperature is obtained by charge exchange recombination spectroscopy.

Operation in the hot-ion mode minimizes uncertainties caused by this term by operating at low density and high temperature, which minimizes magnitude of the electron-ion transfer by minimizing the coefficient $n_e/T_e$ (proportional to $n_e^2/T_e^{3/2}$). In addition, as is shown in Fig. 1, utilizing 75 keV deuterium beam injection also couples significant amounts of power directly into the ion channel, resulting in $T_i \gg T_e$ [15]. In this case, the relative error in determining $T_e - T_i$ is almost the same as the relative error in $T_i$.

Accessing the hot-ion mode requires starting with low Ohmic target densities, typically $1 \times 10^{19}$ m$^{-3}$, since the density rises after the start of neutral beam injection in both L-mode and H-mode. Such densities are easiest to obtain at modest plasma currents. The results presented here were obtained at 1 MA. In other experiments, we have obtained hot-ion H-mode plasmas at currents up to 1.5 MA.
In order to insure that plasma shape did not affect our transport comparisons, we chose to use an elongated limiter L-mode plasma and a double-null divertor H-mode plasma with basically the same interior flux surface shapes. Both discharges had elongations of about 1.8 [21]. It was necessary to use a limiter discharge for the L-mode plasma because a divertor discharge with this much input power would have gone into the H-mode so quickly that a hot-ion mode with an $n_e$ comparable to the H-mode plasma's could not have been obtained. Lower power operation would have increased the time to the transition, but would also have reduced the difference between $T_i$ and $T_e$.

In the H-mode shot, $n_e$ rises linearly in time until the onset of the Edge Localized Mode (ELM) clamps it. The density in the L-mode shot soon reaches a steady state. Comparing transport at the same $n_e$ requires time-dependent transport analysis, which demands a complete time history of profile measurements. These were obtained by combining Thomson scattering measurements from two series of repeat shots into complete time histories of the 300-millisecond-long interval bracketing the time when $n_e = 3.5 \times 10^{19} \text{ m}^{-3}$ in both the H-mode and L-mode shots [21].

Although both the L-mode and H-mode shots had sawtooth oscillations in the soft X-ray flux in the Ohmic phase of the discharge, the sawteeth disappeared during the neutral beam injection phase where the transport was analyzed. Accordingly, sawteeth do not affect the transport. In addition, the ELM-free portion of the H-mode was analyzed [15]; transport is, consequently, unaffected by ELMs.

The ONETWO code [22] has been used to analyze the transport in these discharges. Because of the large toroidal rotation speeds found in these shots [21] (about 500 km/sec in the center of the H-mode plasmas), the fast-ion slowing down package in the code had to be upgraded to account for the birth of the fast ions in a rotating plasma. At these rotation speeds, this significantly affects the partition of input power between the electron and ion channel. In addition to utilizing the code to calculate the basic diffusivities, we have made error estimates for the diffusivities as described by Groebner, et al. [23].

Transport analysis shows that the dominant energy loss in the ion channel is due to conduction and electron-ion energy exchange. The dominant loss in the electron channel is electron conduction. Radiation loss in the electron channel is only 30% of the conduction loss at $\rho = 0.85$ and is significantly smaller throughout most of the plasma. During the density rise phase where transport is analyzed in these discharges, the net particle
Fig. 2. (a) Comparison of experimentally inferred $\chi_e$ and $\chi_\phi$ as a function of magnetic flux surface label $\rho$ in limiter L-mode and divertor H-mode plasmas. (b) Comparison of experimentally inferred $\chi_i$ and neoclassical $\chi_i$ in the same plasmas as in (a).

flux is quite small; essentially all the particle source from the beams and the thermal neutrals is consumed in fueling the plasma density rise.

The inferred diffusivities $\chi_e$, $\chi_i$, and $\chi_\phi$ are compared for the L-mode and H-mode discharges in Fig. 2. The most dramatic improvement from L-mode to H-mode is in $\chi_e$, which improves by about a factor of three throughout most of the plasma, with an even larger improvement near the plasma edge. The improvement in $\chi_\phi$ is similar to that in $\chi_e$. Indeed, outside of $\rho = 0.3$, $\chi_e$ and $\chi_\phi$ are basically equal within the error bars in both L-mode and H-mode. The $\chi_i$ improves only inside of $\rho = 0.5$; there is no change in $\chi_i$ within the error bars in the outer half of the plasma. An additional important feature of these inferred diffusivities is that $\chi_i$ inside of $\rho = 0.3$ in the hot-ion H-mode agrees with the predictions of Chang-Hinton neoclassical theory [24] within the error bars. Finally, it is notable that $\chi_i$ is significantly smaller than $\chi_e$ inside of $\rho = 0.4$ in both L-mode and H-mode discharges. These experimentally inferred diffusivities will be compared with the predictions of various theories in Section 4.

One puzzling feature of the experimental results is the similarity of $\chi_e$ and $\chi_\phi$. Since all the angular momentum in the plasma is contained
in the ions, one might think $\chi_i$ and $\chi_\phi$ should be more similar than $\chi_\phi$ and $\chi_e$. However, this is true only in the very center of the plasma.

The results in Fig. 2 clearly indicate that the bulk transport of energy and angular momentum in the divertor H-mode is significantly reduced relative to that in the limiter L-mode plasma. Accordingly, the confinement improvement in H-mode is not simply due to the creation of a transport barrier at the plasma edge.

3. TIME SCALE FOR TRANSPORT IMPROVEMENT AT THE L TO H TRANSITION

Having confirmed a bulk confinement improvement in H-mode, we were also interested in determining how fast this confinement improvement occurred. To investigate this, we performed a transport analysis of a detailed time history taken across the L to H transition in a divertor plasma identical to the one used in the previous study. The beam power was increased to 10.5 MW for this investigation in order to produce a sharper, more well defined L to H transition.

As is shown in Fig. 3, the $n_e$ and $T_i$ profiles into $\rho = 0.5$ show measurable changes in 10 to 20 milliseconds after the transition. Consequently, we are interested in doing transport analysis in an interval of about 100 milliseconds bracketing the transition. Because the Thomson density profiles were assembled shot by shot, considerable care was taken to maintain reproducible shots. In addition, the plasma profiles used in the transport analysis were smoothed by first fitting the experimental profile data to cubic splines, then using linear interpolation in time to smooth the knot values for the cubic splines. A separate linear function was fit to the knot values prior to and after the L to H transition. This procedure allowed us to detect systematic changes in the profile values and their first time derivative across the transition, but it smoothed out minor variations which could have caused spurious oscillations in the inferred diffusivities. This procedure detects trends that continue for 10 to 20 milliseconds or longer while rejecting higher frequency variations.

As is shown in Fig. 4, the $\chi_e$ and $\chi_\phi$ in the region outside $\rho = 0.5$ show significant drops in the 10 milliseconds after the transition. As the estimated error bar in Fig. 4(a) shows, any changes inside of $\rho = 0.5$ during this time are within the experimental error. As might be expected from Fig. 2, $\chi_i$ shows a smaller drop; in addition, the time scale is perhaps a bit longer. It is interesting to note that, as is shown in Fig. 3, the $n_e$ profile flattens in the same region where the bulk transport improvement
Fig. 3. (a) Electron density profile and (b) ion temperature profile as a function of magnetic flux surface label $\rho$ at various times spanning the L to H transition. The density profile is determined from Thomson scattering supplemented by and normalized to the results from four chords of two-color, CO$_2$ interferometry. The curves are labeled with the measurement time in milliseconds relative to the L to H transition. Notice the flattening of the electron density profile soon after the transition in the region around $\rho = 0.6$. Data points are shown on the curves with the highest density and highest central $T_i$.

occurs and that this flattening occurs on the same time scale as the change in the bulk transport. Although the confinement improvement evident during the early H–mode is most obvious outside of $\rho = 0.5$, the results in the previous section indicate that that improvement ultimately penetrates to the center.

The data in Fig. 4 demonstrate that confinement improves over a significant portion of the plasma within 10 ms after the L to H transition. This is further evidence that the confinement improvement is not localized just to the edge transport barrier.

4. COMPARISON WITH THEORY

The results in the two previous sections suggest two important questions. First, what is the cause of the bulk confinement improvement after the L to H transition? Second, what governs the underlying transport after the transition?
The simultaneously change of the $n_e$ profile and the bulk transport motivates consideration of electrostatic drift wave driven transport, since most of the drift wave models predict that $\chi_e$ decreases [16,17] as the electron density gradient scale length $L_n$ increases. Indeed, the $\chi_e$ from many of the drift wave models [16] vanishes in the flat density gradient limit that, as Fig. 3 shows, is reached by the H-mode plasmas about 150 to 200 milliseconds after the L to H transition. In addition, electrostatic modes are of interest because they should produce $\chi_e$, $\chi_i$, and $\chi_\phi$ values that are quite similar; this is what is seen in both L-mode and H-mode plasmas in the region around $\rho = 0.6$. Finally, the ion temperature gradient (ITG) driven mode can persist even in the flat density gradient limit, thus providing a candidate for anomalous transport in that limit.

In order to confront the question of whether drift-wave-based models could quantitatively match our results, we have compared one of the most elaborate of the recent drift wave models [17] with our experimental results from Section 2. This model includes the effect of ITG-driven turbulence as well as the effect of collisional and collisionless trapped electron modes. As can be seen in Fig. 5, the most obvious feature of

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Fig. 4. Experimentally inferred (a) $\chi_e$ and (b) $\chi_\phi$ as a function of magnetic flux surface label $\rho$ at various times across the L to H transition. Both diffusivities decrease significantly outside of $\rho = 0.5$ within 10 ms after the transition. Representative error bars are shown on a portion of one $\chi_e$ graph. Error bars at other locations on that curve and for other curves would be similar.
the graphs is that the theoretically predicted $\chi_e$ and $\chi_i$ are larger than the experimentally inferred values by factors of 10 to 400 in the center of the plasma. In addition, in both H-mode and L-mode, the theoretical $\chi_e$ and $\chi_i$ have the wrong spatial dependence; the theoretical values decrease strongly with increasing $\rho$ while the experimental values increase with increasing $\rho$. Such discrepancies in magnitude and spatial dependence are not unique to DIII-D, but have also been seen, for example, in L-mode and H-mode plasmas in JET [25,26] and in L-mode and supershot plasmas in TFTR [27,28]. Finally, when all the functional dependences are included, the $\chi_e$ actually increases from L-mode to H-mode over most of the plasma. This electrostatic drift-wave theory does not even match our qualitative expectation of lower $\chi_e$ in H-mode, because the electron temperature dependence overcomes the improvement caused by a flatter density profile.

The transport estimates in Ref. [17] are based on quasilinear theory. More accurate simulations [29,30] indicate that the quasilinear estimates are usually too large. Accordingly, improvements in the theory based on, for example, the inclusion of the effect of coherent structures [30] may result in better agreement with experiment than the present, quasilinear estimates [17].

Since the magnitude of the predicted drift-wave-driven transport is much too large to match the experimental value, we need to consider the possibility that the transport is strong enough that it has reduced the plasma gradients so that the turbulence is only marginally excited. We can examine this possibility by considering the threshold for ITG-driven turbulence. In both our H-mode and L-mode plasmas, the density gradients are flat enough and the ratio $T_i/T_e$ large enough that the flat density gradient limit of the ITG theory is the appropriate one [31,32]. This limit is expressed in terms of a critical ion temperature gradient scale length $L_{T_i}$.

In Fig. 6, we compare the experimental $L_{T_i}$ for the H-mode and L-mode plasmas in Section 2 with the theoretical predictions of Romanelli [31,32] and an improved version of Dominguez and Rosenbluth [33]. As can be seen there, the measured gradient scale length is a factor of two to three steeper than that required to give instability in Romanelli's theory. In other words, the plasma ion temperature gradient is not clamped at the marginal stability threshold given in Refs. [31,32]. The improved Dominguez and Rosenbluth prediction is much closer to the experimental result. We have evaluated the results in Ref. [33] for our case by creating an interpolation function which simply multiplies together the dependence on safety factor $q$ and $Z_{eff}$ given in Ref. [33]. In addition, we have
Fig. 5. Comparison of theoretically predicted $\chi_e$ [(a) and (c)] and $\chi_i$ [(b) and (d)] with the experimentally inferred values for the hot-ion divertor H-mode [(c) and (d)] and hot-ion limiter L-mode [(a) and (b)] discussed in Section 2. The theories considered are those of Romanelli and Briguglio [17] and Rebut et al. [18,19]. The uncertainty in the experimental measurement is shown by the shaded band around the experimental result.
Fig. 6. Comparison of theoretically predicted marginal stability threshold for ITG turbulence with experimental values for (a) H-mode plasma and (b) L-mode plasma from Section 2. Romanelli's prediction is based on Refs. [31,32]; Dominguez and Rosenbluth's is an improved version of the results in Ref. [33]. The Dominguez-Rosenbluth prediction does not extend into $p = 0$ because the calculation is not valid for $q < 2$. The plasma is stable to the mode above the theoretical threshold line and unstable below it. The uncertainty in the experimental result is shown by the shaded band around the experimental result.

improved the result in Ref. [33] by multiplying that critical $L_T$ by the factor $2/(1 + T_i/T_e)$ to account for the dependence on $T_i/T_e$ seen in other ITG theories [31,32,34].

Leaving aside the precise threshold value, an additional puzzle for ITG modes is the region of the hot-ion H-mode plasma inside of $p = 0.3$. As is shown in Fig. 2, this is the region where $\chi_i$ is equal to the value predicted by Chang-Hinton neoclassical theory [24]. However, as is shown in Fig. 6, in the region around $p = 0.2$ to 0.3, $L_T$ has already reached its minimum value. If the minimum $L_T$ were set by the onset of ITG-driven turbulence, how is it possible that the $\chi_i$ in this region is neoclassical?

Another broad class of anomalous transport models involves electromagnetic turbulence. In attempting to explain the H-mode confinement
improvement with a theory based on electromagnetic turbulence, the difficulty is in finding a model which has the proper functional dependence to give decreased transport in the H-mode. The resistive ballooning mode models [35,36], for example, would predict little decrease in transport in the bulk of the H-mode plasmas because the increase of the theoretical $\chi_e$ with temperature cancels the decrease caused by flattening the density gradient. One electromagnetic-based model which does predict lower transport in the flat density gradient limit is the Rebut-Lallia model [18,19]. Even though the theoretical foundation of this model has not been completely clarified [31], the functional form for $\chi_e$ and $\chi_i$ given in this model predicts a decrease in transport as the density gradient flattens.

As can be seen in Fig. 5, the predictions of the Rebut-Lallia model match the experimentally inferred diffusivities significantly better than those of the drift wave model. The radial dependence of the predicted $\chi_i$ and $\chi_e$ is similar to the experimental quantities. In the H-mode, there is reasonable quantitative agreement with $\chi_e$ and $\chi_i$ inside of $\rho = 0.8$. Outside of $\rho = 0.8$, the predicted values rise sharply because of the steep H-mode edge density gradient; the experimental results show no similar increase. In the L-mode, only the $\chi_i$ prediction agrees with the experimental value; the theoretical $\chi_e$ is too small. This indicates that the model is unable to explain all of the confinement improvement from L-mode to H-mode, since the predicted improvement in $\chi_e$ is less than that observed.

The anomalous transport in the Rebut-Lallia model [18,19] occurs only in regions of the plasma where the electron temperature gradient is steeper than a critical gradient prescribed by the theory. If the measured gradients are close to the critical gradients, the precise form of the threshold function used in the theory [18,19] might impact the comparison with experiment. However, we have found that, in both our hot-ion H-mode and L-mode plasmas, the measured electron temperature gradient exceeds the critical gradient by at least a factor of 3 over the region of the plasma where the theoretical predictions have been given in Fig. 5. In H-mode, the measured gradient is a factor of 10 to 40 above the theoretical critical gradient over most of the plasma. Accordingly, threshold effects are not a problem in evaluating the theoretical prediction.

Our comparison with theory indicates that, although we can find models with the correct qualitative dependence to explain the H-mode confinement improvement, the theories will have to be significantly improved before we can obtain quantitative agreement.
5. CONCLUSIONS

We have clearly demonstrated that the bulk confinement in the plasma improves after the L to H transition. Accordingly, the H-mode confinement improvement is due both to a transport barrier at the plasma edge and to a confinement improvement in the bulk of the plasma. All three diffusivities $\chi_e$, $\chi_i$ and $\chi_\phi$ improve after the transition; $\chi_e$ and $\chi_\phi$ improve more than $\chi_i$. In the center of the hot-ion H-mode plasma, $\chi_i$ is at the neoclassical level. Significant reductions in $\chi_e$, $\chi_i$, and $\chi_\phi$ occur within 10 milliseconds in the region of the plasma outside of $\rho = 0.5$. During this same time, the $n_e$ profile flattens in the same region. This connection leads us to speculate that the confinement improvement is due to the flattening of the density profile. Comparison of our results with a theory based on electrostatic, drift-wave turbulence [17] show significant quantitative and qualitative disagreement. The electromagnetic turbulence-driven model of Rebut, et al. [18,19] agrees better with the experiment, but there are still areas of significant disagreement. Quantitative agreement between experiment and theory will require significant improvement in the theory.

REFERENCES


[22] ONETWO Code, to be published.


**DISCUSSION**

R.J. TAYLOR: Can you not explain the improvements at the edge by the increase in edge density alone, and perhaps by improved edge stabilization?

K.H. BURRELL: The density profiles for the L- and H-mode plasmas compared in Section 2 are such that the plasmas have the same local density near $\rho = 0.5$. In spite of this, $x_e$ differs by a factor of three. Because of this, we
conclude that the bulk confinement improvement between L- and H-mode is not due to a change in local density since it exists even at a point where the densities are the same. The local density gradient is not the same in these plasmas; this is why we have investigated theories where transport improves in the flat density gradient limit.

F. ROMANELLI: Can you tell us anything about the safety factor profile in these discharges?

K.H. BURRELL: We have information on the q-profile in these shots only from external magnetic measurements coupled with MHD equilibrium calculations. We have used this information, for example, in comparing the predictions of Refs [18, 19, 33] with our experimental results. As is mentioned in the caption to Fig. 6 in the paper, the work in Ref. [33] is not valid for q < 2. The point where the lines in Fig. 6 stop is the q = 2 surface.

J.G. CORDEY: After the transition to the H-mode in JET, as well as the change in the density profile, there is also an accompanying change in the current profile due to the bootstrap current, i.e. a strong current broadening. Do you see a similar broadening of the current in DIII-D, and could this be the cause of the improved confinement? Also, is the change in the current profile taken into account in comparisons with theoretical models?

K.H. BURRELL: At the L-H transition in DIII-D, the plasma internal inductance drops, indicating a broadening of the current profile. This broadening of the profile was taken into account in the comparison that we made with theoretical models. The paper by Zarnstorff et al. (IAEA-CN-53/A-II-2, this volume) calls into question all transport theories which depend on the safety factor profile or magnetic shear profile in the outer half of the plasma. Accordingly, this provides some evidence against the idea that the bulk transport improvement in H-mode is caused by a change in the current profile. More definitive results could be obtained by repeating Zarnstorff’s experiment in the H-mode.

C.S. CHANG: You have shown that there is an abrupt change of poloidal rotation speed at the edge, at the onset of the L-H transition, and you also showed that $\chi_\phi$ drops significantly near the transition time. Is there any abrupt change of toroidal rotation speed at the edge (similar to that of poloidal rotation) at the onset of the L-H transition?

K.H. BURRELL: A more complete discussion of changes in the plasma edge at the L-H transition can be found in R.J. Groebner’s paper (IAEA-CN-53/A-VI-4, this volume). We see an abrupt change in both edge toroidal and poloidal rotations at the L-H transition. If we define the direction of the plasma current as the positive direction for toroidal rotation, then the edge toroidal rotation becomes negative at the L-H transition. This is consistent with a more negative edge electric field after the transition. The work of R.D. Stambaugh, cited as Ref. [2] in my paper, contains further discussion of this point. The negative toroidal rotation exists only at the very edge of the plasma and only for 10 to 20 ms after the transition.
EXPERIMENTAL DETERMINATION OF THE TRANSIENT TRANSPORT AND OF FLUCTUATIONS RELEVANT TO TRANSPORT IN ASDEX

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Abstract

EXPERIMENTAL DETERMINATION OF THE TRANSIENT TRANSPORT AND OF FLUCTUATIONS RELEVANT TO TRANSPORT IN ASDEX.

Particle transport was studied in ASDEX with modulated puffing of the discharge gas and of impurities. The energy transport is investigated by numerical simulation of the heat pulse after the sawtooth crash. Small scale density fluctuations are investigated in the confinement region with far infrared scattering and reflectometry and in the edge plasma with Langmuir probes and $H_\alpha$ diagnostic. In addition to a diffusive component of the particle transport, a strong inward drift is observed in all discharges. In ohmic discharges the transport coefficients decrease and saturate like $1/T$ with increasing density. They are smaller in deuterium than in hydrogen. In the improved ohmic confinement (IOC) regime mainly $D$ in the outer region is reduced. $D$ increases proportionally to the heating power in L-mode discharges. The improvement of particle confinement in the H-mode is explained by an increase of the inward drift at the edge rather than a decrease of $D$. The impurity diffusion coefficient is independent of the impurity mass and charge. In ohmic discharges, it varies with $n_e$ like the bulk diffusion coefficient, is independent of $B$ or increases weakly with $B$ and increases with $I_p$. In L-mode discharges, $D_{imp}$ increases linearly with the heating power. The electron thermal conductivity determined by heat pulse propagation exceeds the stationary value by a factor of 3–4, assuming merely diffusive heat transport. Convection does not significantly reduce this factor. However, non-diagonal terms

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in a general transport equation system may remove the discrepancy. Drift-type turbulence with a remarkable radial asymmetry is found. The Doppler shift due to plasma rotation complicates the interpretation of frequency spectra. No separate ion feature could be identified. In the L-mode an increase of the fluctuation level together with a broadening of the frequency spectra and a shift of the k-spectra to small values are observed. During L-H transitions the fluctuation level drops immediately. However, during ELM free H-phases the fluctuation level can in some cases begin to grow again. The particle transport at the edge can be explained by the fluctuating n·E × B flux derived from probe measurements. Like the core turbulence, these flute-like fluctuations show an inboard-outboard asymmetry. This suggests that bad curvature is a key element of the driving mechanism. Because of the high correlation of density fluctuations along the magnetic field lines, the interaction with the target plates may be important.

1 INTRODUCTION
The mechanism responsible for the anomalous transport of energy and particles generally observed in tokamaks is not yet well understood, even though this phenomenon is as exciting as it is important for the development of a fusion reactor. It is accepted by the majority of the fusion community that microscopic fluctuations are involved and that they lead not only to enhanced heat and particle diffusion coefficients but also to anomalous off-diagonal transport coefficients. This paper reports on work performed in order to obtain better insight both into spatial and parameter variations of the coefficients using transient particle and heat fluxes and into the nature of fluctuations suspected as driving the transport.

2 TRANSPORT IN THE BULK PLASMA
2.1 Method
Small density perturbations about equilibrium, induced by sinusoidal modulation of the gas valve, were analyzed for different radial channels of the ASDEX HCN-laser-interferometer. The measured amplitudes and phase shifts are compared to solutions of the particle conservation equation, \( \partial n/\partial t = - \nabla \cdot \Gamma + P \). A transport law with a diffusive and a convective component \( \Gamma = - D \nabla n - Vn \) is assumed. The coefficients \( D(r) \) and \( V(r) \) are determined with a crude radial resolution by a numerical fitting method [1].

2.2 General observations
Discharges with different heating methods, densities, plasma currents, toroidal fields and hydrogen isotopes have been investigated [2]. In all cases a strong inward particle drift \( V \) was observed. Transport coefficients are generally larger in the outer region of the plasma.
2.3 Ohmic discharges

A general decrease in transport coefficients both in the interior and the outer part of the plasma with increasing density is observed. However, the decrease saturates at higher density, an effect especially pronounced in the diffusion coefficients. Furthermore, there is a strong isotope effect, that is all coefficients but $V$ in the outer region are much smaller for deuterium compared to hydrogen. Coefficients in the interior are significantly lower than those in the exterior at nearly every density.

The dependence of $D$ and $V$ on other parameters is more complex and less strong. There is a clear tendency for $D$ in the center to increase with $I_p$, probably due to increased sawtooth activity. For the outer portion of the plasma, the only systematic effect is that $D$ increases with toroidal field.

The most significant difference between the Improved Ohmic Confinement regime (IOC) and the saturated regime is a reduction of the outer $D$ by a factor reaching 4 at the highest densities.

2.4 L-mode discharges

In co-injection heated discharges, $D$ in the outer region increases proportional to the neutral injection power. The central value of $D$ increases strongly at small power and has a general weak tendency to increase with further increasing power. The increase of $D$ in the outer region with heating power is very large at low plasma current, but less pronounced with higher plasma currents. Diffusion in the core is largely unaffected. The current dependence of particle transport in L-mode is however mainly an effect of $q$, as shown by a magnetic field variation at constant $I_p$.

2.5 H-mode discharges

A steady-state H-mode with grassy ELMs was investigated. The value of $D$ in the confinement zone during the H-mode remains at its elevated L-mode value. In contrast, the convective velocity in the periphery increases strongly, by a factor of 3-5. This explains the improved particle confinement and is consistent with the density pedestal observed.

3 TRANSPORT OF IMPURITIES

3.1 Method

A harmonic analysis method similar to the one discussed above for the background plasma was applied to impurity transport [3]. In this case trace impurities such as SiH$_4$, H$_2$S, and HBr are puffed into the plasma with sinusoidal modulation in time using a programmed valve. Harmonic analysis of the measured line radiation in several spectral ranges from visible to X-ray
allows determination of phase relations and Fourier amplitudes at the modulation frequency.

3.2 Theory

The solution of the transport equation for each impurity is calculated assuming a sinusoidal time dependence and a simple model for the transport coefficients, with diffusion $D=\text{const.}$ and drift $V\sim r$. By comparing the calculated phase shifts and Fourier amplitudes at a given frequency with the corresponding measured values, it is possible to determine the transport parameters within the scope of the simplified model. Since the Fourier amplitudes of the total impurity density could not be measured, the analysis is restricted to the determination of the diffusion coefficient. This is possible from the measured phase alone due to its weak dependence on the drift velocity.

As a matter of basic interest, we note that the method of harmonic analysis offers also the possibility of determining both transport quantities $D$ and $V$ as functions of the radius, provided that the amplitude $A$ and phase $\varphi$ of the total impurity density have been accurately measured. Integration of the source free transport equation leads to the integral transforms

$$D(r) = -\frac{\omega}{r\varphi'(r)} \int_0^r dr' \frac{A(r')}{A(r)} \cos(\varphi(r')-\varphi(r))$$

$$V(r) = -\frac{\omega}{r} \int_0^r dr' \frac{A(r')}{A(r)} \sin(\varphi(r')-\varphi(r)) + \frac{A'(r) D(r)}{A(r)}$$

3.3 Parameter range

Diffusion coefficients, derived by means of the analysis described above, were determined for ohmic and NI-heated L-mode discharges with varying plasma parameters [3], [4].

3.4 Ohmic discharges

We observed that the diffusion coefficient is independent of the charge (and mass) of the ions as expected from the predictions of neoclassical theory. On the other hand, we find that the diffusion coefficient decreases with increasing electron density. Furthermore, at fixed density the diffusion coefficient seems to be constant or even increase slightly with the toroidal magnetic field. It also increases with the plasma current. These results are in complete contradiction to the neoclassical theory of impurity transport. Thus the results obtained can only be understood in terms of an additional
anomalous diffusion dominating the transport processes. This behavior is in striking contrast to the non-stationary accumulation phases found in discharges with improved energy and particle confinement [5].

3.5 L-mode discharges

For L-mode discharges the dependence of the diffusion coefficient on the NI heating power was investigated. The results show an approximately linear increase with the injected power.

As a final remark, we note that the diffusion coefficients presented agree well with results from earlier measurements using laser ablation and impurity doped pellets.

4 HEAT PULSE TRANSPORT

4.1 Method

The local electron temperature is measured at four radial positions outside the mixing radius by the ECE diagnostic and the signal is sampled at 20 kHz to provide sufficient time resolution of the sawtooth crash and the resultant heat pulse. In Ohmic discharges the sawtooth period is 5 -10 ms. Boxcar averaging of the temperature perturbation permits the electron thermal conductivity, $\chi_e$, to be determined. The time dependent temperature perturbation at each radial position is fitted by the numerical solution of the heat pulse propagation equation as a forced boundary value problem [6]. The first channel outside the mixing radius is used as the time dependent boundary condition. In contrast to particle transport, there is no generally accepted ansatz for the heat transport equation.

4.2 Results

Assuming a transport law $Q = - n\chi \partial T/\partial r$ for electrons and ions, heat pulse values of $\chi_e$ are generally a factor of 3 higher than equilibrium values. In this approximation the density pulse associated with a sawtooth crash is neglected, because the relative amplitude of the density perturbation in these Ohmic discharges is considerably smaller than the relative amplitude of the temperature perturbation ($\delta n/n \approx 0.1 \delta T/T$). The perturbed particle flux is the quantity of interest as this also contributes to heat pulse transport. By considering a particle and heat flux of the form $\Gamma = - D \partial n/\partial r - V n$, $Q = 3 \delta \Gamma T - n\chi \partial T/\partial r$, where $D$ is the particle diffusion coefficient and $V$ is the inwards particle drift velocity, it is then possible to include the enhanced heat flux due to a finite density perturbation. The equilibrium values of $D$ and $V$ are determined from the zero order particle flux as described above.

On TEXT the diffusion coefficient from density pulse measurements was found to be a factor of three larger than the value calculated from
equilibrium [7]. Allowance for an increased value of D and a relative perturbation value of $\delta n/n = 0.2 \delta T/T$ show that the enhanced heat flux due to finite density perturbation leads to a reduction of 20% in the inferred value of $\chi_e$. Work is in progress to consider the case in which the temperature and density perturbations are coupled by the off-diagonal terms of the transport matrix. In this case particle and heat flux terms take the form $\Gamma = -D\partial n/\partial r - \alpha n/T\partial T/\partial r$, $Q = -\alpha T\partial n/\partial r - n\chi \partial T/\partial r + 3/2\Gamma T$. Significant enhancement of the heat flux is found when $\alpha > 0.5\chi_e$ is assumed.

5 TURBULENCE IN THE CORE

5.1 Diagnostics

Far infrared scattering and microwave reflectometry have been used to investigate the fluctuations in the plasma core which might cause the observed enhancement of particle and heat transport.

Wavenumber and frequency spectra of electron density fluctuations are measured with far infrared laser scattering along chords at different radii, including a horizontal chord through the plasma center and a vertical tangent at the outer edge. The system can be operated in homodyne and in heterodyne mode, where spectra for frequencies of both signs, corresponding to the different directions of propagation, are resolved. Reflectometry detects density fluctuations with spectral resolution at a selectable radius.

5.2 Results

Under ohmic conditions a number of parameter variations including electron density, electron temperature, toroidal magnetic field, and filling gas ($H^+$, $D^+$, $He^{++}$) were performed. Fig 1 shows a typical homodyne FIR $k$-$\omega$ spectrum from an ohmic hydrogen discharge at a chord radius of $25 \text{ cm} = 0.63 \alpha$.

The spectra observed fall into the wavenumber range predicted by numerical driftwave simulations. However, gyroradius scaling of $k^{\text{max}}$ in the sense $k^{\text{max}} \propto \rho_0^{-1}$ is not found. In the dominant frequency range the rms value of the scattered power increases with the mean electron density, consistent with a fluctuation level determined by the mixing length criterion. In ohmic shots an asymmetry of the frequency spectra in the dominant wavenumber range is seen which increases with increasing line density. Since no effective spatial resolution across the minor radius exists at $k^{\text{max}}$ this indicates an in-out asymmetry. At high k the fluctuations can be localized on the outer side, indicating propagation predominantly in the electron diamagnetic drift direction. In the SOC regime at high densities no separate feature emerges which would indicate an $\eta_1$-mode.
Fig. 1. Scattered signal power as function of wavenumber and frequency measured in an Ohmic discharge at a chord radius of 25 cm.

With neutral beam injection heating the k-ω spectra as well as the fluctuation levels change in a complex manner which depends sensitively on the heating scenario. In L-shots an increase of the fluctuation level together with a broadening of the frequency spectra and a shift which is possibly due to plasma rotation effects [8], [9] are measured. The k-spectra are shifted towards lower wavenumbers with no resolvable maximum.

The L-H transition manifests itself as a sudden change of the fluctuation spectra. The total scattered power decreases significantly to about ohmic level. The frequency spectra broaden and shift with respect to the L-phase. It must be pointed out, however, that in some cases an increase of the fluctuation level is detected by FIR and reflectometry during quiescent H-phases although the confinement remains unchanged. Fluctuations other than drift-type modes might be involved [10].

6 TURBULENCE IN THE EDGE REGION

6.1 Diagnostics

Langmuir probes were used to measure density and potential fluctuations in the edge region and in the divertor with good spatial and temporal resolution. Neglecting temperature fluctuations both parameters are derived from the ion saturation current and from the floating potential
respectively. Multiple probes and poloidal probe arrays with up to 17 pins yield simultaneously density and potential fluctuation measurements and details of their spatial evolution.

The observation of fluctuations in Hα emission from the edge with 16 photomultiplier channels allowed the temporal and spatial evolution of the edge density fluctuations to be studied in those cases where probes could not be used or an array with increased separation between channels was necessary.

6.2 Particle Transport

The radial particle flux due to the interaction of the fluctuating density with the fluctuating ExB drift was determined from:

$$\bar{\Gamma}_r = \frac{1}{2B_t} \int_0^\infty k_{\text{pol}}(v)|P_{n\Phi}(v)| \sin(\alpha_{n\Phi}(v)) \, dv$$

where $P_{n\Phi}$ is the cross-power spectrum of density and potential and $k_{\text{pol}}$ is the average poloidal wave number at frequency $v$. Within the limits of the accuracy of this method, the calculated particle flux is in agreement with the flux estimates from the confinement time [11]. We conclude that these fluctuations cause most if not all of the radial particle flux in the edge region.

Fig. 2 shows spectra of the different components contributing to the particle flux in comparison to the spectra of the density fluctuations at different radii. In the vicinity of the separatrix there is a strong low frequency component of fluctuations not contributing to the transport. This might indicate the transition to a region where a different type of fluctuation dominates the particle transport.

An evaluation based on the spatial Fourier transform indicates that fluctuations with wavenumbers around 2 cm$^{-1}$ contribute most strongly to the particle flux.

6.3 Propagation and correlation lengths

Outside the separatrix the fluctuations propagate in the ion diamagnetic drift direction with a velocity of 500 - 1500 m/s corresponding to the local plasma velocity. Inside the separatrix the direction of propagation changes due to a strong shear of the poloidal plasma velocity.

Correlation lengths perpendicular to the magnetic field are of the order of one cm. Between a probe in the divertor and one in the midplane, a correlation above 80% with zero time lag over a distance of 10 m was found. Within the limits of accuracy of measurement, both probe tips were located on the same field line. The edge turbulence is essentially 2-dimensional for frequencies up to at least 50 kHz.
Fig. 2  Spectra of the different factors contributing to the particle flux. Left: Transport spectra (solid), cross-power spectrum of density and potential fluctuations. Right: Average phase angle between density and potential fluctuations (solid), mean poloidal wavenumber.

6.4 Asymmetry between low and high field side

In double null discharges no fluctuations are observed on the inboard edge of the torus. However, fluctuations with characteristics similar to the low field side are observed in single null discharges, that is when there is a connection along field lines to the outer edge. The turbulence seems to be driven at the outer edge, probably by the unfavourable curvature in this region.

7 CONCLUSIONS

The diffusion coefficients of the bulk plasma and impurities appear to vary with plasma parameters in exactly the same way as each other and in a similar way as the energy transport. Electrostatic turbulence involving convection with a scale length of the order of a banana width, which is typically 1 cm, might be able to produce this effect. Drift type fluctuations with k-spectra of this scale length are found in all discharges. There are some indications that they are connected with the changes in the particle and energy transport, as shown by the reduction of the fluctuation level at the L-H transition. Coupling between diffusion and inward particle drift and the
properties of the transient heat transport are experimental observations which impose boundary conditions on theoretical models. Experimental evidence for the relation between fluctuations and transport exists only at the edge outside the separatrix, where the fluctuations seem to be of a different nature. Their strong correlation in the region along the field lines suggests that their interaction with the target plates should play an important role.

REFERENCES


DISCUSSION

R.W. CONN: Since you take your measurement at the midplane boundary — and since we know that poloidal asymmetry in the flow is strong — how can you conclude, as you did, that fluctuation driven transport accounts for essentially all the particle transport? (You said it explains your global $\tau_p$ measurement.) And a second question. Do you have any evidence of DC, steady convection at the edge in addition to the fluctuation driven flow? We see large flows of this kind in the continuous current tokamak (CCT). Results will be reported by Taylor later at the Conference.
H. NIEDERMAYER: We do not conclude that the fluctuations account for essentially all the particle transport, but for a major part of it — and possibly all. The relatively large error bars due to unknown asymmetries, the use of an approximate formula and the fact that temperature fluctuations are disregarded — all these things make a more precise statement unwise. We do not have any need for or any indication of another mechanism which would contribute significantly to the radial particle transport, but we cannot exclude the possibility that such a mechanism exists.

There is no evidence for stationary particle flows at the edge in addition to the fluctuation driven radial diffusion, poloidal and/or toroidal rotation of the plasma and the streaming-off into the divertor along field lines outside the separatrix. We believe that the symmetry of our magnetic field is adequate to keep magnetic islands inside the separatrix small enough so that radial particle flows along the field lines can be neglected. Magnetic probe measurements at the edge show that B-field fluctuations do not contribute significantly in good steady state discharge phases. Flows in bad discharges with high magnetic activity and during ELMs have not been investigated. As we do not yet know the findings from CCT, we cannot comment on them.

R.J. GOLDSTON: Burrell presented very interesting data which strongly suggest that the changes in outer region transport at the H–L transition are due to changes in the profiles in that region. This would imply that local edge fluctuations do not drive fluctuations and transport in the interior. Do you have information on the time scale for H–L changes in $\chi_e$, $\chi_i$ or $\chi_\phi$ at the edge of ASDEX?

H. NIEDERMAYER: Edge profiles and the energy flux across the separatrix change on a millisecond time-scale, indicating changes in edge transport coefficients on the same time-scale. The bulk transport reduces within 10 ms.
ANALYSIS OF SATURATED OHMIC CONFINEMENT IN THE FRASCATI TOKAMAK

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Abstract

ANALYSIS OF SATURATED OHMIC CONFINEMENT IN THE FRASCATI TOKAMAK.

The energy confinement time $\tau_E$ for ohmically heated discharges in the Frascati Tokamak (FT) has been studied in a wide range of plasma parameters. It is found to increase linearly with the average density $n_e$ below $1 \times 10^{14}$ cm$^{-3}$ and to saturate at higher densities. The analysis of the experimental data shows that at the same density at which saturation is reached $Z_{\text{eff}}$ becomes nearly equal to 1 and the electron and ion temperatures become nearly equal. The saturated $\tau_E$ shows a strong isotope effect with a reduction of a factor 1.5 in operation in hydrogen. The sawtooth period is closely correlated with $\tau_E$ and displays the same saturated behaviour and isotope effect. The ion thermal diffusivity $\chi_i$ is found to be compatible with neoclassical theory up to the maximum density but a neoclassical multiplier of 3 would only fit the highest $n_e$ data. $\chi_e$ first decreases with the inverse of $n_e$ but remains nearly constant in the saturated region. Therefore the Ohmic saturation in FT seems to be primarily caused by electron transport.

1. INTRODUCTION

Early studies of Ohmic transport in the Frascati Tokamak (FT) [1] showed a linear increase of the energy confinement time $\tau_E$ with the line averaged density $\bar{n}_e$ up to the maximum density, in contrast with experimental findings in Alcator-C [2], where a saturation of $\tau_E$ at high density was observed. Recently, however, results obtained in ASDEX [3] have shown that a regime of improved Ohmic confinement

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can be found in gas fuelled discharges at high density, in contrast with the normally occurring saturated behaviour.

We have therefore resumed our studies of Ohmic confinement in FT with the aim of documenting the behaviour of $\tau_E$ in a wide range of plasma parameters. For this purpose a large data set has been collected, which is described in detail in Section 2. The most noticeable result of our analysis is that $\tau_E$ saturates at high density in FT, too. The conditions for this saturation and the parametric dependence of $\tau_E$ in the linear and saturated region are analysed in Section 2. The analysis of the ion heat balance is presented in Section 3, where we show that our data are consistent with a neoclassical ion transport with an upper bound of three times neoclassical at high density. The picture which then emerges corresponds to a $\tau_E$ saturation caused primarily by electron transport. A final discussion is given in Section 4.

2. EXPERIMENTS

Data analysed in this work have been collected during the last two years of experiments in FT. They cover a large span of the operational space of the tokamak, namely

$$2.5 \leq B_T \leq 8 \, T$$

$$2 \times 10^{13} \, \text{cm}^{-3} \leq \bar{n}_e \leq 3 \times 10^{14} \, \text{cm}^{-3}$$

$$2.2 \leq q_a \leq 5$$

Deuterium was used as filling gas for most of the experiments. Only for $B_T = 6 \, T$ have experimental runs with hydrogen been analysed.

Kinetic data have been used for the evaluation of the plasma energy. Electron temperature profiles were measured by Thomson scattering and electron cyclotron emission (ECE). Electron density profiles were only available for a limited number of shots because of instabilities in the absolute calibration of the Thomson scattering system. These data show that $n_e$ profiles are well fitted by a generalized parabola with the external exponent $\alpha_n$ ranging from 0.7 to 1.2.

The total electron energy was obtained from profile integration with a parabolic density profile using the line averaged density as measured by a HCN interferometer and density at the limiter position extrapolated from Langmuir probe data.

Ion temperature profiles were obtained by solving the ion heat balance equation with given electron profiles and neoclassical heat conductivity. This procedure is justified by the fairly good agreement obtained with the experimental data from fast neutrals diagnostics and, for deuterium discharges, from neutron emission rate measurements (see Section 3).
The behaviour of $\tau_E$ with increasing density is shown in Fig. 1(a) for deuterium discharges with $B_T = 6$ T and $q_a$ ranging from 2.5 to 5. It reproduces the findings of other tokamaks with a linear increase at low density and a saturation at high density. A close inspection of the linear region reveals a weak dependence of $\tau_E$ upon $q_a$ for $q_a > 3$ but an important reduction for $2 < q_a \leq 3$.

Confinement saturation takes place for $\tilde{n}_e > 1 \times 10^{14}$ cm$^{-3}$ when $q_a > 3$. It is correlated with a distinct behaviour of other plasma parameters, namely the sawtooth repetition time $\tau_{ST}$, the effective ion charge $Z_{eff}$ (measured from plasma resistivity and using neoclassical corrections) and the ratio of central ion and electron temperatures $T_i(0)/T_e(0)$. As shown in Fig. 1, our data suggest that saturation of the Ohmic confinement time sets in when, with increasing density, the value of $Z_{eff}$ is reduced to a constant and when ions and electrons have nearly equal temperatures. Also noticeable is the close similarity of $\tau_E$ and $\tau_{ST}$, including the saturation at high density.

The analysis of electron temperature profiles shows that they remain self-consistent in both the linear and the saturated regime, with $(T_e)/T_e(0) \sim 1/q_a$ and the location of sawtooth inversion radius $r_s$, measured by a soft X ray camera, given by $r_s/a \sim 1/q_a$.

Increasing the toroidal field to 8 T gives a slight increase of the saturation density, again maintaining the correlation with $Z_{eff}$, $T_i(0)/T_e(0)$ and $\tau_{ST}$.

A density scan with $B_T = 2.5$ T shows only a linear increase of $\tau_E$ until the density limit is reached at $\tilde{n}_e \sim 8 \times 10^{13}$ cm$^{-3}$. This scan is characterized by a linear increase of $T_i(0)/T_e(0)$ with density and $Z_{eff}$ still decreasing up to the maximum $\tilde{n}_e$.

Changing the filling gas from deuterium to hydrogen leads to different confinement, as documented in Fig. 2. The saturated value for hydrogen is consistent with a reduction of a factor of $\sqrt{2}$, as observed in other devices.

In the linear region, however, similar confinement is obtained for the two isotopes and saturation for hydrogen is reached at lower density. This is related to the lower $Z_{eff}$ and higher $T_i(0)/T_e(0)$ observed in hydrogen in the linear region. The close relation between $\tau_E$ and $\tau_{ST}$ is maintained in hydrogen, too. Besides this strong isotope effect, the value of $\tau_E$ in the saturated region shows weak dependences upon other macroscopic parameters. In fact, once saturation is reached, $\tau_E$ remains constant up to $\tilde{n}_e \sim 3 \times 10^{14}$ cm$^{-3}$ within the experimental error bars ($\pm 15\%$). Varying $q_a$ from 2.5 to 5 also leads to no appreciable difference in the saturated value of $\tau_E$, although at low $q$ saturation is reached at higher density.

The influence of $B_T$ is more difficult to ascertain from our data. Saturation was in fact only reached with density scans at $B_T = 6$ T and $B_T = 8$ T and, although data at higher field are systematically higher, the differences between the two scans are only slightly larger than the experimental errors. On the other hand, data obtained with $B_T = 2.5$ T clearly show that any dependence of the saturated $\tau_E$ upon $B_T$ has to be lower than linear. This result can be compared with those reported for D-III [4], where a linear $B_T$ dependence is observed for divertor discharges but a weaker one for limiter operation.
(a) \( \bar{n}_e (10^{13} \text{ cm}^{-3}) \) vs. \( \varepsilon_e \) (mA)

- \( 2.5 < q_e < 3.0 \)
- \( 3.0 < q_e < 3.5 \)
- \( 3.5 < q_e < 4.0 \)
- \( 4.0 < q_e < 4.5 \)
- \( 4.5 < q_e < 5.0 \)

(b) \( \bar{n}_e (10^{13} \text{ cm}^{-3}) \) vs. \( \vartheta_{ST} \) (mA)
FIG. 1. Variation of (a) energy confinement time, (b) sawtooth period, (c) $Z_{eff}$ and (d) $T_{i}(0)/T_{e}(0)$ as a function of line averaged density for discharges with $B_T = 6$ T in deuterium.
3. POWER BALANCE ANALYSIS

In most of the discharges we are discussing and especially at high density in the saturated region, the difference between electron and ion temperatures is too small to be reliably measured and therefore a direct determination of electron and ion thermoconductivity from the experimental data turns out to be impossible.

We used a different approach whereby the ion heat balance equation is first solved using experimental electron profiles and a theoretical expression for the ion thermoconductivity $\chi_i$. The results of this calculation are then compared with the available experimental information. This consists of the total neutron emission rate $\Phi_n^{ex}$ for deuterium discharges and of energy spectra of neutral particles produced by charge exchange reactions in the plasma and measured by an energy analyser with mass selection.

The total neutron emission rate has been calculated from the profile obtained, assuming negligible dilution, as indicated by spectroscopic diagnostics which identify the dominant impurities at high $Z_{eff}$ in metals. For neutral particles, the experimental slope of the measured spectrum $T_f^{ex}$ has been compared with the simulated value $T_f^{i}$ obtained by running a neutral transport code with profiles calculated from the power balance equation. This procedure minimizes the impact of systematic errors on the results of our analysis.

The experimental uncertainty on $T_f^{ex}$ is mainly due to errors in the calibration of the analyser and is of the order of 10%. The error on $\Phi_n^{ex}$ is only due to the
calibration of the neutron counters. This has been done by inserting a radioactive source in the tokamak and, independently, by comparing with foil activation measurements during a plasma shot. The two methods give results differing by a factor of 2 [5] and therefore we have adopted as a calibration factor the mean value of the two results with an error bar of ±50%.

The results obtained with a neoclassical heat conductivity \( \chi_{NC} \) as given by Chang and Hinton [6] are compared with experimental data from fast neutrals and neutron diagnostics in Figs 3(a) and 3(b), respectively. In both cases the theory gives fairly accurate predictions over the whole density range. If we introduce a neoclassical multiplier \( \alpha \) and assume \( \chi_i = \alpha \chi_{NC} \), already for \( \alpha = 3 \) the calculated values are outside the error bars of experimental data for \( \bar{n}_e \leq 1.5 \times 10^{14} \text{ cm}^{-3} \).

At higher densities, however, owing to the strong electron–ion coupling, the ion temperature profile is constrained to remain close to the electron one and differences between the two cases are small. Anyway, for \( \bar{n}_e = 2.0 \times 10^{14} \text{ cm}^{-3} \), predictions with \( \alpha = 3 \) are at the upper limit of the error bar for neutron flux measurements.

These results are in contradiction to a picture of Ohmic saturation based only on the onset of ion transport and indicate the possibility that a significant departure of the electron transport from Neo-Alcator behaviour is taking place [7, 8]. They are therefore a good test for theories recently worked out for the threshold of the \( \eta_i \) instability and the associated electron and ion heat fluxes [9].

For a limited number of shots we have increased the neoclassical \( \chi_i \) with the contribution of \( \eta_i \) turbulence as given in Ref. [9]. For \( \bar{n}_e = 5 \times 10^{13} \text{ cm}^{-3} \) the data are consistent only with \( \eta_i \) modes below threshold (\( \eta_i^m \approx 1 \)), and this can be obtained with \( \alpha_n \geq 1.5 \). When the density is increased to \( 1 \times 10^{14} \text{ cm}^{-3} \), the mode turns out to be clearly above threshold for the experimental values of \( \alpha_n \), and the results are consistent with an ion heat flux reduced by a factor of 3 with respect to the quasi-linear estimate. Comparisons performed at \( \bar{n}_e \geq 2 \times 10^{14} \text{ cm}^{-3} \) are less conclusive than for the case of enhanced neoclassical \( \chi_i \). Transport code simulations show that in this density range the quasi-linear estimate for \( \eta_i \) transport is a factor of 2 to 3 larger than neoclassical transport.

For constant toroidal field and assuming neoclassical heat conduction, \( \chi_i \) evaluated at half-radius is nearly independent of \( \bar{n}_e \) and takes a value of about \( 10^3 \text{ cm}^2\cdot\text{s}^{-1} \) for \( B_T = 6 \text{ T} \). This is caused by the variation of \( Z_{\text{eff}} \) with density. In this case the power transferred collisionally from the electrons to the ions becomes of the order of 20% of the Ohmic power at the highest density. With a neoclassical multiplier of 3 the power dissipated by the ions would increase to 40% of the Ohmic input.

For purely neoclassical ions, the electron heat conductivity \( \chi_e \) at half-radius decreases with the inverse of \( \bar{n}_e \) in the linear region but remains nearly constant in the saturated region around a value of \( 5 \times 10^3 \text{ cm}^2\cdot\text{s}^{-1} \) for \( B_T = 6 \text{ T} \). With a neoclassical multiplier of 3, a 30% reduction of \( \chi_e \) would be obtained at high density while the product \( \bar{n}_e \chi_e \) remains an increasing function of \( \bar{n}_e \) in the saturated region.
4. DISCUSSION

In contrast with earlier results, in this study of OH confinement in FT we have observed a clear saturation of $\tau_E$ at high density. Our data indicate that the electron transport seems to remain dominant in the saturated regime and that the increase of $\tau_E$ in the linear region is correlated to the decrease of $Z_{\text{eff}}$ and to the increase of energy stored in the plasma ions. This can be explained with a temperature dependent $\chi_e$ and a nearly neoclassical $\chi_i$. 

FIG. 3. Ratio of experimental neutron emission rate $\Phi_n^e$ to the calculated values of $\Phi_n^c$ for $\chi_e = \chi_{NC}$ (open circles) and $\chi_i = 3\chi_{NC}$ (open squares) as a function of line averaged density (a); the same plot for $T_i^e$ and the corresponding calculated value $T_i^c$ (b).
We have observed a clear correlation of the sawtooth period and $\tau_E$. This has been reconstructed by simulations of the sawtooth instability with a numerical code based on reduced MHD equations and heat transport equation for the electrons [10]. By using a self-consistent $\chi_e$ this gives a scaling $\tau_{ST} \propto \tau_E^{0.6}/q_a^{0.5}$ [11].

We have also explored the possibility that the transport mechanism determining OH confinement is also responsible for the power degradation observed in additionally heated discharges. To this aim we have compared our data set with different scalings proposed for the case of additional heating and we have found a good agreement with a number of them, namely with plateau scaling, Rebut-Lallia scaling and collisionless trapped mode scaling [12].

Transport simulations with $\chi_e$ and $\chi_i$ given by collisionless trapped mode scaling can reproduce both the behaviour of $\tau_E$ and the electron temperature profile observed [13]. They also show that at high density $\chi_e$ can remain higher than $\chi_i$ in the saturated regime if the density profile is sufficiently peaked, thus reconciling our result with those of ASDEX and Alcator C, where an ion dominated saturation is observed together with flat density profiles [2, 7].

REFERENCES

B. COPPI: The Alcator experiments indicated quite clearly that the diffusion coefficients derived from the simple-minded quasi-linear theory of the $\eta_1$ mode were too high to explain the derivations. Have you compared your results with the results that we presented on this subject at the 1984 conference in London?

R. BARTIROMO: In our work we derive $T_i(r)$ for a given $\chi_i$ and then compare it with experimental data. It turns out that $\chi_i$ has to be reduced to a level such that the ion heat transport cannot justify the observed saturation of $\tau_E$. This is different from the case of Alcator C.
STUDY OF INSTABILITIES ACCOMPANYING THE DENSITY LIMIT ON T-10


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Abstract

STUDY OF INSTABILITIES ACCOMPANYING THE DENSITY LIMIT ON T-10.

Results of a study of instabilities arising with intense gas puffing near the $n_e$ limit in the ECH experiment on T-10 are presented. Some peculiarities of discharges with pellet injection are also described.

1. INTRODUCTION

The main goal of the T-10 experiments presented in this paper is to investigate processes leading to a density limit at fixed plasma current $I_p$ and longitudinal field $B_t$. In the ECH experiments, the $B_t$ value corresponding to on-axis heating is $B_t = 3$ T. The $I_p$ value was chosen to provide $q_L = 3.5$ to 4; this value allows one to investigate the processes at the plasma edge ($r > r_q = 2$) in detail. To increase the density, both intense gas puffing and H-pellet injection were used. 3.7 mm gyrotrons provided an HF power of 0.1 to 1.8 MW. Owing to ECH, which partially compensated the enhanced radiative losses at the plasma edge, the processes that arise because of plasma cooling — such as 'saturation' in $n_e$ rise and MHD activity burst — were delayed.
2. RESULTS OF ECH EXPERIMENTS WITH INTENSE GAS PUFFING

Traces of chord averaged $\bar{n}_e$ are displayed in Fig. 1. The ECH pulse was switched off at different times. We see that the rapid $\bar{n}_e$ rise accompanied by the MHD burst starts later with a longer ECH pulse.

Changes in $\bar{n}_e$ with intense gas puffing in the ECH discharges were described in Refs [1, 2]. This process can be broken down into four stages (Fig. 2).

The first stage is characterized by an increase in $\bar{n}_e$, degradation in confinement of the exterior plasma, and enhancement in the plasma-wall interaction. The intensity of the $D_\alpha$ line rises more rapidly than $\bar{n}_e$, and it radiates mainly near the inner wall where MARFEs exist. The electron temperature is decreased at the edge and the plasma cooling front moves inwards (Fig. 3) [1].

During the second stage, the $D_\alpha$ line intensity decreases together with the plasma density near the wall and at the limiter where the Langmuir probes were placed. (The electron temperature near the wall decreased slightly, if at all.) Thus, we may conclude that the plasma-wall interaction was reduced.

The $n_e$ profile peaks during the second stage (Fig. 4(a)). This may be explained either by an increase in the pinch velocity or, which seems to be more probable, by a decrease in the diffusion, which previously compensated for the pinch. Energy confinement in the plasma interior ($r < 15$ cm) is also improved (Fig. 5).

Many features of this stage are characteristic of the H-mode, but up to now we have no concrete evidence of their identity.

*FIG. 1. Traces of density averaged over the chamber diameter, $\bar{n}_e$, and of the $m = 2$ activity in ECH for different durations. $I_p = 220$ kA, $B_t = 3.05$ T, $P_{HP} = 0.9$ MW.*
FIG. 2. Traces of $\bar{n}_e$, intensity of deuterium line, $I_{D\alpha}$, power of radiative losses, $P_{rad}$, and $m = 2$ mode activity in discharge with intensive gas puffing.

FIG. 3. Time behaviour of edge electron temperature profile [1]. ECRH starts at $t = 520$ ms, gas puffing at 540 ms, the second stage at about 630 ms and the third stage at about 645 ms. $I_p = 200$ kA, $b_z = 3.0$ T, $P_{HF} = 2$ MW.
FIG. 4. (a) Variation of $n_e$ profile in the second stage (start at about 650 ms); $I_p = 220$ kA, $B_t = 3.0$ T, $P_{HP} = 0.3$ MW; (b) variation of $n_e$ profile: $I_p = 220$ kA, $B_t = 2.95$ T, $P_{HP} = 0.3$ MW.
FIG. 5. Variation of energy content in the plasma interior; same regime as in Fig. 1.

FIG. 6. $T_e$ profile at plasma edge, 5 ms before first MHD splash. The durations of the ECH pulses are different in different shots. The traces of $\bar{n}_e$ and MHD activity are shown in Fig. 7, for some of these shots. The absolute value of $T_e$ for $r = 25$ cm may be too low, because the plasma is not black.
FIG. 7. a) Amplitude of $m = 1$ sinusoidal oscillations as a function of radius for stationary stage of discharge with sawteeth (dashed line) and in the second stage (solid line, with $\sim 10$ ms shift of the maximum); b) phase shift of the oscillations for the latter case.

FIG. 8. Temporal traces of $T_e$ at different radii in the second and third stages. Zero lines are not shown. Sawteeth with $r_s = 6$ cm are changed by $q = 2$ internal disruption with $r_s = 18$ to 22 cm.
FIG. 9. Temporal traces of $T_e$ at different radii expanded for the interval between the first and the second (not shown) minor disruptions. At $t = 590$ ms, a small amplitude $q = 1$ internal disruption occurs with $r_e = 13$ cm.

The second stage ends by an MHD burst. Note that the burst is preceded by a certain value of grad $T_e$ at the periphery. As can be seen in Fig. 1, the second stage begins at different values of $\bar{n}_e$. The time of the start depends on the duration of the ECH. The central $T_e$ values and HF power profiles are also different. The reason for this difference is as follows: for longer ECH pulses and, hence, at higher $\bar{n}_e$, the HF power is deposited far from the plasma centre ($r \sim 12-15$ cm) because of refraction, while, for short ECH pulses, only Ohmic heating persists with a power deposition profile proportional to $j_z^2(r)$. Nevertheless, the $T_e$ profiles measured in both cases 5 ms before the first MHD burst prove to be the same at the edge ($r > 23$ cm, see Fig. 6).

The third stage is characterized by a well developed MHD activity (Fig. 2). Note that, as a rule, the MHD activity does not lead to a major disruption.
The specific features of the third stage are:

(a) The existence of $m = 1$ oscillations, which have a maximum at a larger radius ($r_{m=1}$) than the phase inversion radius $r_s$ of the preceding sawteeth: $r_s = 5$ cm, while $r_{m=1} = 10$ to 11 cm. The results of X ray measurements are shown in Fig. 7.

(b) The $m = 2$ and $m = 3$ modes appear virtually simultaneously with the $m = 1$ mode. X ray tomography gives evidence of the existence of $m = 1$ and $m = 2$ islands.

(c) Minor disruptions are seen (Fig. 8), manifesting themselves in $T_e$ drops in the plasma interior ($r < 19-20$ cm) with explosive $T_e$ growth at the periphery ($r > 20$ cm). The instability is analogous to the $q = 1$ sawtooth, but takes place.
near the $q = 2$ surface. Thus, it could be described as a $q = 2$ internal disruption. As a rule, it neither results in strong plasma-wall interaction nor causes a plasma disruption, although cooling occurs in the bulk of the plasma. The first $q = 2$ internal disruption also manifests itself by a splash of soft X rays near the $q = 2$ surface. A possible explanation would be the local reconnection of magnetic surfaces.

(d) Every internal $q = 2$ disruption enhances the sinusoidal $m = 2$ oscillations. The enhancement is followed by gradual damping.

(e) The ‘usual’ internal $q = 1$ disruptions are also sometimes observed together with $q = 2$ disruptions, but with a larger value of $r_s$. This also indicates an expansion of the region with $q \leq 1$. An example for the coexistence of islands is shown in Fig. 9. The internal $q = 1$ disruption occurs 4 ms after the first internal $q = 2$ disruption with $r_s \approx 21$ cm, whereas $r_s \approx 13$ cm. The sinusoidal oscillations also shown in the figure are caused by rotation of the $m = 1$ mode as indicated by X ray tomography.

(f) Cooling of the plasma interior during the first minor disruption proceeds as a wave process (Fig. 10). The wave velocity corresponds to a coefficient of electron thermal conductivity of $K_e \approx 5 \times 10^5$ cm$^2$·s$^{-1}$. This high value of $K_e$ is probably due to fast propagation of the wave across the $m = 1$ island.

Thus, in minor disruptions, the MHD structure undergoes some changes but is not destroyed to such an extent as in major disruptions.

The fourth stage is the major disruption, which results in an intense plasma-wall interaction without current interruption.

3. H-PELLET INJECTION IN ECH DISCHARGES WITH INTENSE GAS PUFFING

Pellet injection during the first stage results in:

(i) fast cooling of the plasma column as a whole;
(ii) flattening of the $T_e$ profile (Fig. 11);
(iii) stabilization of sawteeth.

Evaporation of the pellet takes place near $r = 12$ cm and is accompanied by an $n_e$ increase only within this surface. The characteristic time of particle accumulation is about 1 ms or less. The resulting $n_e$ profile is similar to the final profile in a gas puffing discharge (second stage, Fig. 4(b)). The injection also results in a marked improvement in the confinement of the plasma interior: $n_e$ and $T_e$ within the $r = 18$ cm surface virtually do not change for about 50 ms, which is the characteristic time for the experiments (Figs 12 and 4(b)). The absolute value of the central $n_e$ with injection can exceed the value achieved without injection under similar conditions. The third stage begins approximately at the same time whether there is injection or not (Fig. 13). However, the scenario of the third stage in discharges with
FIG. 11. $T_e$ profiles before and after pellet injection.

FIG. 12. Traces of $n_e$, $I_{De}$, and $m = 2$ activity in ECH discharge with intensive gas puffing and pellet injection.
FIG. 13. $T_e$ traces expanded at the moment of pellet injection.

FIG. 14. $T_e$ profiles during third stage of discharge with pellet injection.
injection differs if the H-pellet stabilizes the sawteeth. In this case, the $q = 2$ internal disruption cools the plasma layer given by $11 \text{ cm} < r < 18 \text{ cm}$, but does not change $T_e$ at the centre ($r < 11 \text{ cm}$), where smooth oscillations of small amplitude occur. In the initial two or three disruptions, an increase in grad $T_e$ ($r \lesssim 10 \text{ cm}$) during the crash (Fig. 14) does not lead to energy losses from the centre. Then, the major disruption occupies the interior in a manner similar to the situation in discharges with no pellets.

Langmuir probes placed near the wall record an $n_e$ rise at each minor disruption. About 5 to 7 ms before the first minor disruption, a drop in the ion saturation current on the electron side of the Mach probe together with a signal rise on the ion side are observed. Toroidal rotation towards the ion side can be a good explanation for the observation.

4. DISCUSSION

An analysis of the results presented in this paper leads to the following possible scenario (Fig. 2):

In the first stage, the plasma edge is cooled by an intense flux of particles. The flux increases with time because of the confinement degradation in the plasma exterior, together with high recycling. The plasma–wall interaction is asymmetrical along the minor azimuth, as a result of MARFEs. The increase in the flux results in a growth of the radiation intensity of the deuterium and impurity lines (the radiated power approaches 100% of the input power).

In the second stage, the plasma–wall interaction is reduced for two reasons:

(i) a significant part of the plasma energy is lost through radiation;

(ii) the plasma current is expelled from the cooling plasma edge to its 'hot' interior.

During the contraction of $j(r)$, grad $j$ appears to become high enough to provide a shear that is comparable in magnitude to that provided by the separatrix in the H-regime. High shear is favourable to a stabilization of the edge instabilities and to an improvement in confinement. (Another possibility is that grad $T_e$ or grad $(n_eT_e)$ — rather than grad $j$ — controls the edge instabilities.)

Suppression of the instability governing the transport results in an improvement of particle and energy confinement. The $n_e$ profile peaks because of an imbalance between pinch and diffusion outflux. (The latter decreases possibly together with an increase in the pinch velocity.) The profile of $T_e$ becomes slightly flatter. The sawtooth period increases. A skin layer of $j(r)$ moving towards the $q = 2$ surface produces a local maximum with $q \leq 1$. A rotating magnetic $m = 1$ island appears at $r = r_s$, which is larger than the phase inversion radius of the sawteeth (in the case under consideration, these values are 10 and 5 cm, respectively).

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$^1$ MARFE = multifaceted asymmetric radiation from the edge.
In the third stage, parameters are conducive to the growth of an \( m = 2 \) island and an internal \( q = 2 \) disruption with \( r_{\text{a2}} = 19 \) to \( 20 \) cm (as a result of flattening of \( j(r) \) in the plasma interior and cooling of the plasma edge beyond the \( q = 2 \) surface). This major disruption is accompanied by a marked loss of energy. The cooling propagates to the plasma centre. However, no strong interaction between the \( m = 1 \) and \( m = 2 \) MHD modes seems to occur, and there is no major disruption. The existence of enhanced \( \operatorname{grad} j \) near (but inside) the \( q = 2 \) surface is a possible explanation for this phenomenon. High \( \operatorname{grad} j \) is conducive to an increase in current density outside this surface during the internal disruption. Such an increase in \( j \) outside the island limits its growth.

In the fourth stage, the plasma cools after the \( m = 2 \) sawteeth, and a major disruption displaying the usual characteristic features occurs.

The EC heating itself does not govern the above mentioned processes (the density profile peaks both during ECH and afterwards), but ECH prolongates the first stage, apparently, owing to a partial compensation of the radiative cooling by an additional heat flux from the EC heated zone at the plasma edge. Simultaneously, ECH favours the above mentioned redistribution of the plasma current when the heating takes place near the \( q = 2 \) surface. This is the case if the heating zone shifts outwards with the rise in \( n_e \), because of refraction. Analogously, with off-axis ECH (\( r_{\text{ECR}} \sim 15-18 \) cm), the current redistribution and, hence, the second stage occurs earlier.

Pellet injection allows an increase in the \( n_e(0) \) limit without change in the edge conditions. This makes it clear that the above mentioned processes limit the mean density rather than the value of \( n_e(0) \).

Peaked profiles of \( n_e \), together with flatter \( T_e \) profiles, are produced by pellet injection. They are similar to profiles obtained in ASDEX [3], JET [4], and, possibly, in the supershot of TFTR. Such a configuration corresponds to better confinement of energy and particles. Improved confinement at the edge was shown to favour formation of these optimal profiles of \( T_e \) and \( n_e \) in the second stage of discharges with no pellets. However, the redistribution of the plasma current at the edge, which is a contrary process, results in the excitation of an \( m = 2 \) mode. Without pellets, the process of redistribution of the central parts of \( T_e \) and \( n_e \) stops before the stabilization of the sawteeth (due to the flat \( T_e \) profile). Stabilization of the sawteeth, i.e. an increase in \( q(0) \) above unity, would result in significantly better confinement at the centre due to either the suppression of the magnetic islands or stabilization of the ballooning modes.

5. CONCLUSIONS

(1) With intense gas puffing, the major disruption is preceded by a series of internal \( q = 2 \) disruptions which 'softly' expels the plasma energy to the wall, mainly through radiation.
(2) The internal $q = 2$ disruption seems to be governed by a restoration of $j(r)$ near the $q = 2$ surface. Restoration is the result of radiative cooling of the plasma edge.

(3) A high value of $\nabla j$ between the $q = 1$ and $q = 2$ surfaces reduces the interaction between the $m = 1$ and $m = 2$ magnetic islands. A weakening of the interaction prevents major disruptions.

(4) Redistribution of $T_e(r)$ (of $j(r)$?) at the plasma edge results in an improved confinement and in a reduction of the plasma-wall interaction. The duration of the improved confinement period is limited by a contrary process, which provokes an internal $q = 2$ disruption.

(5) Improved confinement in the bulk plasma is characterized by a peaked $n_e$ profile and a flat $T_e$ profile as compared to the self-consistent profiles in the L-regime.

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REFERENCES

PROFILE OPTIMIZATION
AND HIGH BETA DISCHARGES,
AND STABILITY OF HIGH ELONGATION
PLASMAS IN THE DIII-D TOKAMAK*

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Abstract
PROFILE OPTIMIZATION AND HIGH BETA DISCHARGES, AND STABILITY OF HIGH ELONGATION PLASMAS IN THE DIII-D TOKAMAK.

The maximum experimental volume averaged toroidal beta in DIII-D scales with the plasma current divided by the plasma minor radius and the toroidal field, \( \beta_T \propto I/aB \). \( \beta_T = 11\% \) has been achieved at high normalized current by operating at high plasma elongation, \( \kappa = 2.3 \), and high triangularity, \( \delta = 0.8 \), near the axisymmetric stability limit. In a limited range of \( I/aB \), \( \beta_T (\%) = 5I/aB \) (MA/m/T) has been obtained. Stability calculations indicate that the high beta stability limit depends on the profiles of the current density and the pressure.

1. INTRODUCTION

The maximum experimental beta values have been shown to scale with the normalized plasma current in many tokamak experiments, \( \beta_T < C_\beta \times I_N \) [1–12]. \( \beta_T \) is the volume averaged toroidal beta, \( \beta_T \equiv \frac{2\mu_0}{V_B^2} \int dV \rho_P \); \( I_N \) is the normalized current, \( I_N \equiv I/aB(MA/m/T) \); where \( I \) is the plasma current, \( a \) is the plasma minor radius, and \( B \) is the vacuum toroidal magnetic field at the geometric center of the discharge. Stability calculations for both ideal low \( n \) kink modes and high \( n \) ballooning modes have both predicted and confirmed this scaling with normalized current with the value of the scaling coefficient depending on the instability in question, the stabilizing influence of the wall or a conducting mantle, and the parameterizations of the current density profile and the pressure.

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profile used in the calculations [13-18]. The magnitude of the scaling coefficient \((C_P)\) determined in the different experimental devices has varied between 2.8 and 3.5 [1-12]. But, in each experimental device, the highest values of beta are achieved at the highest values of \(I_N\). This scaling leads naturally to the development of high beta experiments with low aspect ratio and strong shaping, since \(I_N \propto \frac{1}{\beta_0} \left(1 + \kappa^2(1 + \delta/3)\right) \left(1 + \frac{3}{2}(\kappa)^2\right)\) where \(\kappa\) is the plasma elongation, and \(\delta\) is the triangularity. Our purpose is to (1) maximize \(\beta_T\) through shaping control, (2) understand the physics of the \(\beta_T\) limits, and then to (3) define a plasma shape and plasma profiles which maximize the energy confinement, \(\tau_E\) at high \(\beta_T\).

As the extent of strong shaping has increased to obtain an increase in the normalized current and \(\beta_T\), axisymmetric \((n = 0)\) instabilities become the limit to stable discharge operation. Towards overcoming these limitations, we have developed a multiple time scale vertical stability control system based on a simple analytical model which includes the conducting vessel, the shaping coils, and rigid displacement of the plasma. This development allows routine operation of the discharge within a few percent of the ideal \(n = 0\) limit. The highest achieved beta in DIII-D \((\beta_T \sim 11\%)\), was obtained in a strongly shaped discharge operated near the axisymmetric stability limit.

The maximum beta depends on more than the magnitude of the normalized current. In DIII-D, \(\beta_T\) of at least 3.5 \(I/a_B\) has been obtained over nearly the entire operating range of normalized current and \(q\), \(0.4 < I_N < 3.1, 16\) \(q_{BS} > 2.5\). Furthermore, in a limited range of \(I_N\), beta significantly greater than 3.5 \(I/a_B\) has been obtained [19]. However, beta limiting instabilities, including fast disruptions, are observed at \(\beta_T\) as low as 3.0 \(I_N\). We explain these various results with ideal \(n = 1\) stability calculations completed for a single value of \(I_N\), which show that for broad pressure profiles, higher beta is accessible at higher values of the internal inductance, \(l_i\). \(l_i \equiv \frac{1}{2} \int dV \frac{B_p^2}{B_p^2}\), where \(B_p\) is the poloidal field and \(\langle B_p \rangle = \frac{\mu_0 l_i}{C_{pa}}\) is the average poloidal field at the plasma surface where \(C_{pa}\) is the poloidal circumference of the plasma. The dependence of the ballooning limit on the current profile has been shown elsewhere [16,20,21]. We show that in DIII-D the maximum achievable operational beta scales with the plasma internal inductance as well as the normalized current, and offer a simple scaling, \(\beta_T \leq 4 l_i I/a_B\).

In the next section, we will discuss the \(n = 0\) stability limits, including the development of the control scenario used to operate near the ideal axisymmetric limit. In Section III, we will present a brief summary of the recent experimental high beta results. We will present the results of \(n = 1\) ideal stability calculations which show the accessibility of high beta at high \(l_i\) in Section IV. In Section V, we will illustrate the scaling of the experimental beta values with \(l_i I/a_B\). We summarize in Section VI.

2. AXISYMMETRIC STABILITY

Significant progress has been made in the past two years on the understanding and control of axisymmetric modes in DIII-D. This work has enabled (1) routine operation of a double-null-divertor (DND) configuration at high elongation, \(\kappa \geq 2\), (2) achievement of \(\kappa = 2.5\) in limiter plasmas and \(\kappa \geq 2.4\) in DND
plasmas, and (3) operation to within a few percent of the ideal MHD stability limit for $n = 0$ modes.

The basis of this work is the understanding of the control problem developed using a filament model of the plasma [22]. This simple analytical model was pursued in order to learn how to implement a successful plasma control system rather than as a complete physical description of the problem. We have successfully addressed the destabilizing effect of the interaction of an active control coil with the stabilizing shell which was pointed out by Jardin [23]. The solution to this problem lies in the characterization of the poloidal distribution of the shell current induced by the plasma motion.

We considered a filament plasma and separate the poloidal control into symmetric and anti-symmetric parts (in $z$). In this model the symmetric part only contributes to the external field which has dipole and quadrupole components, their ratio being essentially the decay index, $n$, defined as $n \equiv -\frac{R_n}{B_x \delta R}$ where $B_x$ and its derivative are evaluated at the center of the filament. The plasma is surrounded by a shell (vacuum vessel) which provides a restraining force, reducing the acceleration of the plasma as its motion induces currents in the vessel. As shown in [22], the dispersion relation for this plasma-shell system is nearly separable into inertial and electromagnetic branches, with the electromagnetic branch characterized by the vessel $L/R$ time and a critical index, $n_c \equiv \frac{2 M_{np}^2 X_0}{\mu_0 L_{sv}}$, where $\Gamma \equiv \frac{k_{sv}}{\mu_0 X_0} + \frac{\beta}{2} + \beta_p + \frac{1}{2} M_{vp}^2$ is $\frac{d}{dx}$ (plasma-vessel mutual inductance), $L_v$ is the self inductance of the lowest eigenmode of the vessel, and $X_0$ is the radial location of the filament. The electromagnetic branch is characterized by a growth rate on the shell $L/R$ time scale, $\gamma = \frac{R}{L} \frac{n}{n + n_c}$, with a sharp transition at $n/n_c = -1$. In the control problem we need to address only the electromagnetic branch, $0 > n/n_c > -1$. In the limit of a rigid shift, $n/n_c = -1$ is the ideal MHD stability limit. The separability into two branches allows us to neglect inertia, reducing the order of the problem.

In our analysis the shell is characterized by eigenmodes of its poloidal current distribution with the orthogonality condition that the mutual inductance of any two modes vanish [22]. Each higher mode has a higher spatial frequency than the next lower one and thus a shorter $L/R$ time. The plasma-vessel interaction is dominated by the lowest order mode. Thus the higher order components are neglected. We add the circuit equation for an active coil with a proportional and differential controller and our system of plasma-shell-active coil is described by a second-order system of three coupled equations. These are Laplace-transformed and the characteristic polynomial is extracted for further study. Power requirements are deduced from numerical treatment of the analytic solutions with the addition of the non-linear response brought about by limitations in supply voltage and bandwidth.

The major results of this work are as follows: (1) as a consequence of plasma toroidicity the stabilizing shell currents flow primarily in the outboard region of the vessel (i.e., couple to the lowest vessel eigenmode), (2) because of the poloidal localization of the shell current, the poloidal location of the active coils is critical to operation near the ideal stability limit, (3) both the voltage and bandwidth requirements for the active coil scale exponentially with the quantity $n/n_c$, (4) the power supply bandwidth requirement is the growth rate of the instability with the active coil shorted, divided by $2 \pi$, (5) there is also a critical
index for the active coil beyond which derivative gain is required for a stable system, (6) no single active coil pair is suitable for both stabilizing the plasma at the ideal limit and providing the uniformity of radial field needed to adjust the equilibrium vertical position of the plasma. This last conclusion led us to the development of a hybrid control scheme, in which outboard (6A, 7A, 6B, 7B, Fig. 1) coils are used on a time scale of the vessel L/R time to provide equilibrium position control while inboard coils (2A, 3A, 2B, 3B) provide stability control on a time scale which is faster than the L/R time of the lowest vessel eigenmode.

Experimentally we have studied the system response [24] and compared the results to the model transfer function. In this work we measured the step response to a perturbation in z as a function of proportional and derivative controller gains and plasma decay index. From this we calculate the Laplace transform which is compared to the model calculations. The results reported here largely confirm the validity of the control model proposed in [22]. We find (1) the tokamak response is second order; (2) with an increase in the vessel toroidal resistance consistent with the high density of ports in the outboard region, the characteristic polynomial determined experimentally is in agreement with the model calculation; (3) the destabilizing effect of control with outboard coils (7A, 7B) predicted by the model is observed in the experiment; (4) the hybrid control system extends the range of decay index over which the plasma can be stabilized, and (5) with the hybrid control system \( \kappa = 2.5 \) is achieved in a limiter plasma.

A rigid-body stability calculation indicates that \( \kappa \geq 3 \) should be possible, whereas experimentally we have not exceeded \( \kappa = 2.5 \). This limitation to elongation was explored [25] and we find that the onset of non-rigid behavior in the vertical instability is responsible. As the current profile is broadened and the triangularity is increased experimental results begin to show a marked departure from the rigid-body limit. More complete calculations using GATO [26] show that we are operating at 97% of the ideal limit: Discharges calculated to be stable, are calculated to be unstable when the vessel radius is increased by 3%. Furthermore, the shape of the perturbation as computed in GATO is in excellent agreement with that observed experimentally, with a relatively small vertical shift of the magnetic axis and a much larger deformation of the outer region of the plasma toward the divertor coil [25]. A spectral decomposition of the perturbation of the instability from GATO shows that the plasma is destabilized by the coupling of an \( m/n = 3/0 \) mode to the \( m/n = 1/0 \) mode. For the experimentally realizable current profiles this behavior has served to limit the achievable elongation to \( \kappa = 2.5 \). It is important to note that, although the control system is based on a rigid-body model, we find it adequate to allow operation to within a few percent of the ideal MHD stability limit even when these non-rigid effects play a significant role in the evolution of the instability.

\( \kappa \) as high as 2.5 has been achieved with auxiliary power and ramping of the plasma current to broaden the current profile. A cross section of the flux surfaces of this discharge is shown in Fig. 1. In addition, we have sustained \( \kappa = 2.45 \) for many confinement times. In other conditions we observe vertical instability at much lower values of \( \kappa \). In Fig. 2, we show a set of discharges which were within 10% of the ideal stability limit when they became vertically unstable at
elongations ranging from 2.0 to 2.5. As shown in the figure, the ability to achieve high elongation depends critically on the current density profile (as parameterized by \(\ell_i\)). Part of the reason for the scatter lies in the details of shaping, but we have not as yet completely characterized the onset of the instability. Operation near the stability limit is difficult in that changes of 0.05 in \(\kappa\) (or 0.03 in \(\ell_i\)) are often the difference between sustaining the plasma and a vertical disruption. Without a precise guide as to the operating point relative to the ideal limit, tuning is difficult. On DIII–D this is done by trial and error and vertical disruptions are tolerated. On future larger devices a more refined approach will be required. For fixed plasma conditions, \(n/n_c\) seems a reliable guide to stability margin in DIII–D plasmas.

3. EXPERIMENTAL HIGH BETA RESULTS

As a result of the increased injection power into H-mode divertor discharges, improved wall conditioning, and the enhanced vertical control, higher values of toroidal beta, \(\beta_T\), and normalized beta, \(\beta_N \equiv \beta_T/I_N\), have been obtained in DIII–D. The addition of graphite tiles over the ceiling, the floor, and the inside wall, and the development of techniques to effectively condition these large graphite surfaces, has allowed full power injection into double-null divertor discharges [27]. Full power injection with deuterium neutral beams has increased the available power to 20 MW, as compared to \(\approx 14\) MW with hydrogen beams. In addition the confinement time with \(D^0 \rightarrow D^+\) is approximately 1.5 times larger.
FIG. 2. Maximum experimentally achieved elongation $k$ for plasmas near the ideal axisymmetric limit. The broken line is drawn to highlight the dependence of $k$ on $l_i$.

than confinement with $H^\circ \rightarrow D^+$. This increase in the product of power and confinement time, $P\tau_E$, along with the added capabilities to handle the power successfully in high current double null divertor discharges is largely responsible for the achievement of higher beta values, high values of normalized beta and high beta values at full field.

The highest beta values experimentally obtained at each level of the plasma current and toroidal field are shown in Fig. 3. Equilibrium $\beta_T = (\beta_\parallel + \beta_\perp)/2$ are shown. Values from the diamagnetic loop measurement, $\beta_T^{DIA} = \beta_\perp$, and from kinetic measurements, $\beta_T^{(KIN)}$, are generally slightly higher. Here we show $\beta_T > 9\%$ has been obtained on several discharges and a maximum $\beta_T \approx 11\%$ has been reached. As can be seen $\beta_T \approx 3.5 I_N$ has been reached over most of the operating range in $I_N$ ($0.4 < I_N < 3.1$). $\beta_T$ significantly exceeds $3.5 I_N$ over a limited operating range and a maximum value of $\beta_N = 5$ has been obtained. The access to higher $\beta_T$ in a double-null divertor discharge compared to a single-null divertor is a consequence of the larger achievable normalized current, $3.3 \text{ MA/m/T}$ in a double-null divertor compared to $2.5 \text{ MA/m/T}$ in single-null divertor, owing to the higher elongation and higher triangularity of the double null divertors.

The maximum value of $\beta_T$ obtained in DIII-D to date is $\beta_T = 11\%$, and was obtained with 19 MW of neutral injected power into a highly elongated, $\kappa = 2.3$, highly triangular, $\delta = 0.8$, double-null divertor H-mode discharge [28]. Parameters at the highest beta value, as obtained from the equilibrium analysis using EFITD are shown in Table 1. Input into the equilibrium analysis includes the pressure profile as determined from measured kinetic data and the calculated
fast ion pressure [29]. The total beta on axis is ≈40%; 20% arising from the measured thermal pressure and 20% arising from the calculated fast ion pressure. The ion pressure is calculated from the deposition of the beam ions and the classical slowing down time and does not include radial transport of the beam ions, and hence may be an over-estimate near the axis.

This $\beta_T = 11\%$ discharge, like the other examples in Fig. 4, remains at high beta for times much longer than the growth time for ideal MHD modes. In Fig. 4, we compare the temporal evolution of several high beta discharges. As seen in Fig. 4(a), beta remains above 10% for 50 msec, much longer than the Alfvén time, and the vessel time constant, $\tau_V < 5$ msec, and comparable to an energy confinement time. The high beta portion of this discharge occurs during a sawtooth free period as shown in Fig. 4(a), and the absence of sawteeth for this duration is believed to have been a major factor in the achievement of the high beta value. For example, another discharge, very similar to this 11% discharge but with sawteeth during the high power injection phase only reached $\beta_T = 9.4\%$. A continuous $m/n = 1/1$ mode is observed before the high power neutral injection and continues throughout the high beta phase of the discharge, and this mode may contribute to the absence of the sawteeth. The $\beta_T = 11\%$ discharge terminated in a disruption following the rapid growth $(1/\gamma = 20 \mu\text{sec})$ of a non-rotating $n = 1$ mode. An $m/n = 3/2$ mode is observed $\approx 20$ msec prior to the rapid growth of the external $n = 1$ and suggests the disruption may be caused by the overlapping of islands, similar to the scenario described by Chu [20].

FIG. 3. Experimental high beta results in DIII-D. Data shown represents the highest value obtained at each value of plasma current, $I$, and toroidal field. Highest $\beta_T$ value obtained for each condition does not necessarily represent a $\beta$ limit.
TABLE 1
DISCHARGE PARAMETERS AT $\beta_T = 11\%$

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$I$</td>
<td>1.28 MA</td>
</tr>
<tr>
<td>$B$</td>
<td>0.75 T</td>
</tr>
<tr>
<td>$R$</td>
<td>1.67 m</td>
</tr>
<tr>
<td>$a$</td>
<td>0.56 m</td>
</tr>
<tr>
<td>$\kappa$</td>
<td>2.34</td>
</tr>
<tr>
<td>$\delta$</td>
<td>0.85</td>
</tr>
<tr>
<td>$q_{95}$</td>
<td>2.6</td>
</tr>
<tr>
<td>$\beta_p$</td>
<td>0.96</td>
</tr>
<tr>
<td>$\beta_T$ (MHD)</td>
<td>10.7%</td>
</tr>
<tr>
<td>$\beta_T$ (DIAMAGNETIC)</td>
<td>10.8%</td>
</tr>
<tr>
<td>$\beta_T$ (KINETIC)</td>
<td>11.3%</td>
</tr>
<tr>
<td>$\beta_T$ (FAST ION)</td>
<td>0.33 $\beta_T$</td>
</tr>
<tr>
<td>$\beta(0)$</td>
<td>~40%</td>
</tr>
<tr>
<td>$\beta_T(0)$ (THERMAL)</td>
<td>~20%</td>
</tr>
<tr>
<td>$\beta_T(0)$ (FAST ION)</td>
<td>~20%</td>
</tr>
<tr>
<td>$W$</td>
<td>0.82 MJ</td>
</tr>
<tr>
<td>$P_{TOT}$</td>
<td>20.5 MW</td>
</tr>
<tr>
<td>$P_{NBI}$</td>
<td>19 MW</td>
</tr>
<tr>
<td>$\tau_E$</td>
<td>40 ms</td>
</tr>
<tr>
<td>$\bar{n}$</td>
<td>$5.9 \times 10^{19}$ m$^{-3}$</td>
</tr>
<tr>
<td>$T_i(0)$</td>
<td>2.5 keV</td>
</tr>
</tbody>
</table>

of the discharge development and ideal stability calculations are given by Lazarus et al. [28].

$\beta_N$ up to 5 has been achieved in DIII-D discharges following a negative ramp in the plasma current. The temporal evolution of the discharge that reached the highest value of $\beta_N$ is shown in Fig. 4(b). A negative current ramp begins at $t = 1.74$ sec. The plasma current is ramped down from 0.7 MA to 0.55 MA in 40 msec. Following the negative current ramp $\beta_N$ increases from 4 to 5. $\beta_N$ remains above 4.5 for more than 30 msec, much longer than relevant ideal time scales. The internal inductance, $\xi_i$, also increases during the current ramp from just above 1.2 to above 1.6 and slowly relaxes over about 250 msec to about the original value. Although some modest ELM activity occurs earlier in the shot, at $t = 1.81$ sec, large and frequent ELMs begin, confinement is degraded, and beta decreases. Continuous $n = 1$ and fishbone oscillations [30] occur during the high $\beta_N$ phase, as is observed on most high $\beta_N$ discharges. $\beta_T$ continues to decline in this shot due to the growth of $n = 2$ and $n = 1$ instabilities and a complete beta collapse follows [1,31]. Although the highest values of $\beta_N$ were obtained following
A negative current ramp, $\beta_N \approx 4.5$ has been obtained with a constant plasma current. In some of these high $\beta_N$ discharges, high frequency, $50 < f < 200$ kHz, high toroidal mode number, $5 < n < 10$ oscillations are observed. These modes could be marginally stable ballooning modes [19] or toroidicity induced Alfvén eigenmodes [32].

High $\beta_T$ values and high $\beta_N$ values are also maintained quiescently for many confinement times. Two discharges which represent such quiescent operation are

FIG. 4. Temporal evolution of several DIII-D high beta discharges. (a) Toroidal beta, $\beta_T$; soft x-ray emission from a central viewing chord, SXR; and neutral injection power, $P_{NB}$, for a $\beta_T = 11\%$ discharge. (b) Normalized beta, $\beta_N$; internal inductance, $l_i$; and $P_{NB}$ for a $\beta_N = 5$ discharge. (c) $\beta_T$, $l_i$, and $P_{NB}$ for a long pulse high $\beta_T$ discharge. (d) $\beta_N$, $l_i$, and $P_{NB}$ for a long pulse high $\beta_N$ discharge.
shown in Fig. 4(c) and 4(d). As shown in Fig. 4(c), $\beta_T$ is maintained near 8% for $\approx 1$ sec, and in 4(d) $\beta_N$ is maintained above 3.4 for longer than 1.5 sec. Macroscopic discharge parameters such as the plasma current, the plasma shape, the electron density and the internal inductance remain nearly constant during the high beta period of both discharges. $q_{95}$ for the discharge in Fig. 4(d) is 3.8, which illustrates reliable high $\beta_N$ operation at moderate values of $q$. As a comparison, the discharge shown in Fig. 4(c) is at $q_{95} = 2.4$ and $\beta_N = 2.8$. Discharges at $q_{95} = 3.1$ have been operated quiescently at $\beta_N > 3$ for longer than 1 sec.

With the increase in the product of $P_{TE}$ made available with the operation of deuterium neutral beams, high $\beta_T$ and high $\beta_N$ have been achieved at the maximum toroidal field in DIII-D, 2.1 T. At a plasma current of 2.8 MA, $\beta_T = 5\%$ was reached in a double-null divertor H-mode discharge. The stored energy in this discharge was 3.6 MJ. At lower plasma current, ($I_N = 0.8$) a discharge with $\beta_N = 3.4$ was achieved with a central ion temperature of 10 keV. A central ion temperature of 17 keV was obtained at $\beta_N = 2.8$ in a 1.2 MA discharge. These high beta values illustrate high beta stability at reactor relevant temperatures.

4. DEPENDENCE OF STABILITY BOUNDARIES ON PRESSURE AND CURRENT PROFILES

The achievement of $\beta_N = 5$ indicates there is more detail in the $\beta_T$ limit than just the scaling with $I/aB$. In the range of $1 < I_N < 1.4$, or equivalently $6 > q_{95} > 4$, $\beta_N$ significantly greater than 3.5 is experimentally obtained. In other discharges, beta limiting instabilities, including disruptions that result from the fast ($\approx 100$ $\mu$sec) growth of an $n = 1$ mode are observed at much lower values of $\beta_N$, as low as $\beta_N \approx 3$. In a series of discharges with nearly identical global parameters, plasma shape, total current, auxiliary heating power, density, and internal inductance, dramatically different MHD behaviour is found. Some discharges run stably at $\beta_N \approx 3.5$, some show beta saturation or beta collapse as a result of low $n$ pressure driven modes [1,33], and some show a fast disruption following a sawtooth crash [34]. Significant toroidal rotation is observed in these discharges with high power neutral beam injection, and so it is expected that the stabilizing influence of the wall remains nearly the same in all cases.

These varied results motivated us to evaluate the dependence of the stability boundaries on the profiles of the pressure and current density. Stability against the ideal $n = 1$ kink was calculated for a set of equilibria with the same discharge shape, $q_{95}$, and $q_0$. The pressure and current density profiles were systematically varied [35]. The discharge shape was that of a single-null divertor (from experimental equilibria) with $q_{95} = 3.2$. We chose three pressure profiles for the comparison by specifying the pressure gradient as a function of the normalized poloidal flux, $\psi$: (1) $p'(\psi) = 1 - \psi$, (2) $p'(\psi) = 1 - 0.625\psi$, and (3) $p'(\psi) = \psi^2 - \psi^3$. $\psi = 0$ on axis and $\psi = 1$ at the boundary. The three resulting pressure profiles are shown in Fig. 5. The diamagnetic current was specified as a 3rd order polynomial of $\psi$, $f'f = \gamma_0 + \gamma_1 \psi + \gamma_2 \psi^2 + \gamma_3 \psi^3$. The toroidal
current density given by $j_\phi = R f' + \mu_0 f f'/R$, and hence the $q$ profile, is varied by varying the coefficients in $f f'$ consistent with the constraints of fixed $q_0$, constant current, $I$, and fixed current density at the separatrix (X-point). For these three cases the current density at the separatrix was set to zero. At the same $\ell_i$ and same $\beta_T$ the $q$ profiles for the three different pressure profiles are nearly indistinguishable. A fourth case was studied with the current density at the separatrix set to 50% of the average current density, (4) $j_\phi(R_x, \psi = 1) = 0.5 I/S$, and $p'(\psi) = 1 - 0.625 \psi$, where $R_x$ is the major radius of the X-point and $S$ is the area of the plasma cross section. The $n = 1$ kink mode stability was evaluated with $q_0$ held at 1.05 to maintain Mercier stability and to avoid the internal kink at low beta. Stability was evaluated with the MHD stability code GATO [26], with a conducting wall at $r_w = 1.5a$, the approximate position of the wall for the DIII-D vessel.

The stability boundaries in ($\beta_N$, $\ell_i$) space are shown in Fig. 5. As shown in the figure, the maximum $\beta_N$ stable to $n = 1$ ideal kinks varies from $\sim 3$ to 5.8. With peaked pressure profiles, Case (1), the maximum $\beta_N$ stable to kinks is 3, and shows little variation with $\ell_i$ for $0.6 \leq \ell_i \leq 0.8$. For Case (2), the maximum in the stability boundary is at $\beta_N = 4.4$ and $\ell_i = 0.8$. For Case (3) a strong dependence of the kink beta limit on $\ell_i$ is seen; the stability limit increases
with $\ell_i$ and the highest $\beta_N$ is $\beta_N = 5.7$ at $\ell_i = 0.95$. The kink mode boundary for Case (3) exceeds that imposed by ideal ballooning modes at high $\ell_i$. Finite current density at the boundary, Case (4) eliminates the stable region at high $\ell_i$ and significantly reduces the $\beta_N$ stable to $n = 1$ kinks for all $\ell_i$.

The structure of the modes at the stability boundary varies with $\ell_i$. At low $\ell_i$ the modes exhibit primarily an external structure: the amplitude of the perpendicular displacement is largest near the boundary. At high $\ell_i$, the modes exhibit primarily an internal structure: the amplitude of the displacement is large at the center and decreases toward the boundary. At intermediate values of $\ell_i$, the modes are “global” in the sense that the displacement occurs over the entire plasma. This trend is seen for all the cases studied, although the transition occurs at different $\ell_i$ and the details, particularly at the stability boundary at high $\ell_i$ are dependent on the different pressure profiles. The high $\ell_i$ boundary is due to internal modes and is limited at low central pressure and central pressure gradient (approximately Case 3) by low shear $[\psi/q (dq/d\psi)]$ near the axis. As $\ell_i$ is increased with fixed $q_{95}$ and $q_0$, more current is forced near the center, and the $q$-profile becomes “flatter”. The maximum $\ell_i$ for all pressure profiles, even at low $\beta_N$ is set by the inability to obtain monotonic $q$-profiles within the constraints on the equilibria. The stability boundary for Case (4) at $\ell_i \geq 0.9$ is a result of the reduced central shear brought about by the edge current density. With total current held fixed, current must be removed from the interior to increase the current density near the boundary. Moving current density from the interior to the boundary results in a decrease in $\ell_i$. To maintain the same $\ell_i$, current must be added near the axis as well as the edge. As a result, increasing the edge current density with constant $q_0$, constant $q_{95}$, and constant $\ell_i$, reduces the shear near the axis as well as the shear at the boundary. One would expect then an increased likelihood of internal modes with an increasing edge current density, as well as an increased likelihood of external modes driven by the finite edge current density and current density gradient.

5. BETA-LIMIT SCALING

Motivated by the predicted dependence of the $n = 1$ kink limit on $\ell_i$ for broad pressure profiles, previous work on the dependence of the ballooning limit on the current profile [16,20,21] and the observation that the highest values of $\beta_N$ were obtained following a negative current ramp, we have evaluated the dependence of the operational beta limit in DIII-D on the internal inductance, $\ell_i$. We show in Fig. 6, the dependence of $\beta_N$ on $\ell_i$ for two selected $q$ ranges. We have chosen the lowest $q$ for which sufficient data is available and the range of $q$ for which high $\beta_N$ is obtained. The range in $q$ is limited to 0.4 to eliminate the expected dependence of $\ell_i$ on $q$. As seen quite clearly for these cases, $\beta_N$ is higher for higher $\ell_i$. This trend is true for other values of $q$ chosen, but in some cases not as clearly, particularly at $q \approx 3.2$.

The data shown in Fig. 3 is replotted as a function of $\ell_i I/aB$ in Fig. 7, and we find that the maximum experimentally obtained $\beta_T$ is well represented by $\beta_T(\%) < 4\ell_i I/aB$ (MA/m/T). The different symbols used in Fig. 7 represent
different values of $\beta_N$, and of particular note, is that the inclusion of $\xi_i$ in the scaling seems to much better represent the data in the region of $1 < I_N < 1.4$. This data represents a wide variety of operating conditions, including both single-null and double-null divertor discharges over a range of elongation and triangularity, and the scaling represents the data very well except possibly at low $I_N$. At low $I_N$ and thus low plasma current, $I$, the H-mode discharges are at much lower density and high $\beta_T$ is at high poloidal beta, $\beta_p$. These are conditions where significant fishbone activity is generally observed in DIII-D [30], and one might expect lower $\beta_T$. This scaling does not imply that all the beta limit physics is encompassed in $\xi_i$ and $I_N$, and the scaling with $\xi_i$ should be taken as a guideline, much in the same spirit as the scaling with $I_N$.

Although the scaling $\beta_T \propto I/aB$ has been proven to be a fairly accurate representation of the $\beta_T$ limit, other more detailed dependencies need to be investigated. The non-circular cross section of DIII-D has allowed us to build a database including measurements of the internal inductance and from them to propose a dependence of the $\beta$ limit on this measure of the current profile width. Dependencies of a similar order on pressure profile, indentation and triangularity, and the stabilization of fast particles might be expected.

6. DISCUSSION

Operation with high power deuterium injection into double-null divertor discharges with good vertical control has allowed an increase in the experimentally obtained values of $\beta_T$ and $\beta_N$. The highest values of $\beta_T$ are still reached at the high values of $I_N$, and increasing $I_N$ has required an understanding of the $n = 0$ axisymmetric stability and control problem. The model with rigid plasma
displacement has provided adequate understanding of the control problem to develop a control scenario which allows operation within a few percent of the ideal axisymmetric limit.

The achievement of $\beta_T \approx 11\%$ illustrates that discharges with high values of $\beta_T$ are ideally stable at low values of $q$ with high values of central pressure, $p(0)/(\langle p \rangle) > 3$. The high central pressure, although perhaps not maintainable in steady state, might be important for achieving ignition. The power density near the axis of 3 MW/m$^3$ is comparable to that expected in a reactor. This achievement of $\beta_T \approx 11\%$ extends the range of $I_N$ over which $\beta_T(\%) > 3.5 I/aB$ (MA/m/T) up to $I_N = 3.1$.

Our stability analysis shows the strong dependence of the calculated beta limit on the profiles. This analysis shows high values of $\beta_N$ are accessible at high $\ell_i$ with appropriately broad pressure profiles. One might speculate that if the current density profile is sufficiently peaked, transport owing to internal kink modes and/or ballooning modes will adjust the pressure profile to the appropriate broad profile. This points out the importance of current density profile control.

The achievement of $\beta_N \approx 5$ represents a significant advance in the experimental high beta program. This result, taken with the scaling with $\ell_i$ and the $n = 1$ stability calculations, strongly suggests that with profile control, high beta values at moderate $q$ should be obtainable and sustainable. The operation $\beta_T \approx 8\%$ and $\beta_N \approx 3.5$ in near steady state conditions demonstrates that high beta is not a transient phenomenon. These results support Kikuchi's view of
a modest toroidal field (≈7–8 T) test reactor operating in the first regime of stability at modest $q \approx 5$ [36].

On a less positive note, we find computationally that finite current density near the plasma boundary should significantly limit the operational beta range. In most experimental high beta discharges to date, the pressure profile and the current density profile can be developed somewhat independently and Ohm’s Law (locally) and Faraday’s Law are satisfied. However, in low collisionality high beta plasmas, the bootstrap current becomes a significant part of Ohm’s Law, and the pressure profile and the current density profile must be calculated self consistently for long pulse plasmas. The stability boundary represented by Case(4) in Fig. 5 represents the edge current density that is expected in a low collisionality H–mode discharge, consistent with the edge pressure gradient near the limit imposed by the first regime ballooning limit. Removing the ballooning limit (second regime operation), does not remove the kink limit that results from the broad profiles. Of particular importance is the reduction in the central shear that results from the addition of the edge bootstrap current. We note that the current density profiles in JET plasmas are shown to broaden following neutral injection and the transition to H-mode and that the bootstrap current is calculated to be rather large [37]. We speculate that the high bootstrap current calculated for the JET $\beta_N = 2.8$ discharge leads to a low shear near the axis and the “beta clipping” observed [10]. Much more work needs to be done on the stability analysis with current density profiles having finite edge current density and finite edge current density gradient.

We propose that the scaling $\beta_T(\%) \leq 4 \ell_i I/aB \ (MA/m/T)$ well represents the maximum achievable $\beta_T$ in DIII–D. This scaling emphasizes the importance of the current density profile in achieving high $\beta_T$, and represents significant progress over the more simple scaling with $I_N$. This scaling is in qualitative agreement with the kink calculations presented here and the ballooning calculations done previously. Low $\ell_i$ is needed for the development of strongly shaped plasmas for increasing $I_N$, and thus $\beta_T$, while this scaling indicates that $\ell_i$ should remain high for high $\beta_N$. Note that Fig. 2, for the $n = 0$ limit and Figs. 5 and 6 for the $n = 1$ limit show this conflicting trend which should lead to interesting beta optimization in strongly shaped plasmas. The achievement of $\beta = 11\%$ followed a fine line between $n = 0$ and $n = 1$ stability limits.

This scaling with $\ell_i$ points out the need to develop current drive schemes to control the current density profile as well as simply provide steady state non-Ohmic current drive. A promising approach for increasing beta lies in developing broad pressure profiles through heating deposition profile control while maintaining peaked current density profiles (high $\ell_i$), with current drive.
REFERENCES


DISCUSSION

S.A. COHEN: Do you observe 'hard' or 'soft' $\beta$ limits in your experiments, and how do they vary as a function of $I/a_B$?

T.S. TAYLOR: We observe both 'hard' and 'soft' $\beta$ limits in DIII-D. For $I/a_B \leq 1$ (or $q_{95} \geq 5$) soft $\beta$ limits ($\beta$ saturation and/or collapse) are observed. For $I/a_B \geq 1.7$ (or $q_{95} \leq 3$) hard $\beta$ limits are observed: a fast disruption is associated with rapid growth of a non-rotating $n = 1$ mode following a sawtooth crash. Both types of limit are observed in the range $1 \leq I/a_B \leq 1.7$ (or $5 \geq q_{95} \geq 3$).

R.J. HAWRYLUK: In your highest $\beta$ case, which was an H mode discharge, what is the ratio of $\tau_E/\tau_L$? More generally, in your high $\beta$ discharge do you observe a degradation in the H mode enhancement factor?

T.S. TAYLOR: In general, we observe no degradation of confinement with increasing $\beta_T$ or $\beta_N$ beyond the scaling with power, $\tau_E \propto P^{-0.5}$, for $\beta_N \leq 3$. In some discharges above $\beta_N \sim 3$, some $\tau_E$ degradation occurs as a result of increased ELM and internal MHD activity. The confinement for $\beta_T = 11\%$ is approximately 70% higher than that expected for L mode discharges at these values of $I/a_B$ (at low $q$).

P. SMEULDERS: What happens when you increase the power beyond that required to reach the beta limit, in particular with respect to the occurrence of disruptions?

T.S. TAYLOR: Disruptions and beta saturation or collapse occur with increasing frequency as $\beta_N$ is increased. If the power is increased beyond that necessary to reach the maximum $\beta_T$ obtained experimentally in given operating conditions, then saturation, collapse or disruption is observed.
HIGH POLOIDAL BETA EXPERIMENTS WITH A HOT ION ENHANCED CONFINEMENT REGIME IN THE JT-60 TOKAMAK

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Abstract

HIGH POLOIDAL BETA EXPERIMENTS WITH A HOT ION ENHANCED CONFINEMENT REGIME IN THE JT-60 TOKAMAK.

A high poloidal beta regime with hot ion enhanced confinement has been extended up to \( \varepsilon \beta_p \sim 1.2 \) with nearly balanced perpendicular neutral beam injection for long durations up to 5 s in JT-60. The enhanced toroidal rotation was observed to have high rotational shear associated with the peaked ion temperature profile, suggesting a large negative potential at the centre of the order of the central ion temperature, on the basis of neoclassical rotational equilibrium. The confinement improvement has been limited by high \( \beta_p \) collapses with a fast time-scale with a critical normalized beta \( \beta_N \sim 1.1 \) in the regime of \( \varepsilon \beta_p \sim 0.7 \). The MHD observations are consistent with low \( n \) ideal kink/ballooning analysis associated with the formation of a magnetic pitch minimum. The experimental resistive loop voltages are in good agreement with the neoclassical predictions with bootstrap currents including fast ion effects for a wide range of \( 0.1 \leq \beta_p^{\text{iso}} \leq 3.2 \). The bootstrap current fraction is shown from the 1.5-D time dependent transport analysis to reach \( \sim 80\% \) for a stable plasma with \( \beta_p^{\text{iso}} = 3.2 \).
1. INTRODUCTION

In tokamak research, improvement in energy confinement and its sustainment in steady state for auxiliary heated plasmas remain the critical issues for a future tokamak reactor. After the modification of the JT-60 divertor geometry to the lower single null configuration [1], a hot ion enhanced confinement plasma was obtained in 1989 by exploring an extremely low plasma current regime, where high $\beta_p$ experiments have been carried out utilizing long pulse beams. A variety of MHD activities have been observed to be closely related to the confinement properties, including beta limits. The substantial amount of neoclassical bootstrap current plays an important role in the MHD stability. While several tokamak experiments have recently addressed high $\beta_p$ regimes, for example on TFTR [2-4] and DIII-D [5], no experimental regime in which the bootstrap current fraction is dominant has been explored except for the present regime in JT-60. So the characteristics of a plasma with a large bootstrap current fraction in a collisionless regime will be highlighted in this paper.

2. EXPERIMENTAL CONDITIONS

In JT-60, $I_p$ scan experiments have been conducted in the range $0.3 \leq I_p \leq 1.2$ MA. The plasma was operated with a lower single null divertor configuration having a relatively high aspect ratio of 4.5, a major radius of $R = 2.9$ m, a minor radius of $a = 0.65$ m and an ellipticity of $\kappa = 1.4$ at toroidal fields of $B_t = 4.0$ T or 4.5 T (usually 4.5 T). Neutral hydrogen beams with almost fixed power up to ~21 MW and beam energy of 65 keV were injected for 3-5 s under nearly balanced perpendicular conditions; the balance ratios of co/counter beam power defined as $\Gamma = (P_{co} - P_{ci})/(P_{co} + P_{ci})$ are 0.1-0.25, and the injection angles are $\pm 15^\circ$. The first wall was repeatedly degassed by helium discharge cleaning, so that the target hydrogen plasma was obtained with low electron density ($\bar{n}_e \sim 0.5 \times 10^{19}$ m$^{-3}$). The density was increased by large beam fuelling typically up to $\bar{n}_e \sim 4 \times 10^{19}$ m$^{-3}$.

The $\beta_p^{\text{dia}}$ values are determined by magnetic measurements. The non-circular cross-section of the JT-60 configuration allows us to separate the equilibrium poloidal beta ($\beta_p^{\text{eq}}$) and the internal inductance ($l_i$) in equilibrium fits to external magnetic measurements, using an equilibrium code. The total poloidal beta value including pressure anisotropy is expressed as $\beta_p = (\beta_p^1 + 2\beta_p^\perp)/3 = (\beta_p^{\text{dia}} + 2\beta_p^{\text{eq}})/3$. The pressure anisotropy may be evaluated by a measure defined as $\gamma = (\beta_p^1 - \beta_p^\perp) / (\beta_p^1 + \beta_p^\perp) = (\beta_p^{\text{dia}} - \beta_p^{\text{eq}}) / \beta_p^{\text{eq}}$. Ion temperature ($T_i$) and toroidal rotation velocity ($V_{\phi}$) of the plasma are measured by charge exchange recombination spectrometers (CXRS). Electron temperature ($T_e$) is measured by Thomson scattering and Michel-son interferometer for electron cyclotron emission (ECE). Line electron density and local electron density are measured by a far infrared laser interferometer with two vertical chords and by Thomson scattering, respectively.
3. PLASMA OPTIMIZATION AND MHD EFFECTS

The peak values of ion and electron temperatures \((T_i(0), T_e(0))\) and the L mode enhancement factors \((\tau_E/\tau_E^L)\) to the Goldston L mode scaling [6] as a function of the plasma current are shown in Fig. 1; the \(\tau_E\) values are obtained from the total stored energy. The \(\tau_E/\tau_E^L\) values are saturated up to 1.6 for the regime of \(I_p \leq 0.5\) MA. The plasma pressure becomes anisotropic as \(I_p\) and \(n_e\) decrease; typically, \(\gamma \sim 0.1\) at \(I_p \sim 0.5\) MA. Thus, the best performance is found to be obtained in the range \(I_p \sim 0.5-0.7\) MA, where \(T_i \sim 12\) keV and \(T_e \sim 6\) keV have been achieved.

The MHD observations are characterized by continuous low \(m\) resistive modes, as shown in the bottom part of Fig. 1(a). The sawteeth were completely suppressed during NBI below \(I_p \sim 1.0\) MA. With \(I_p \leq 0.7\) MA, continuous \(m = 2\) modes were most often observed on soft X ray signals, to which Mirnov signals were sometimes observed to be correlated rotating in the ion diamagnetic drift direction with the same mode frequency (typically 5 kHz). The toroidal mode number is considered to be unity from the mode location. Following the change of the continuous modes, the soft X ray emission profiles are observed to be peaked within a mode rational surface corresponding to \(q = 1, 2\) or 3. This enlargement of the mode rational surface as \(I_p\) decreases seems to play an important role in the observed enhancement of the confinement, since the ion temperature profile also tends to become peaked within the mode rational surface.

Waveforms for a typical high \(\beta_p\) discharge with \(\tau_E/\tau_E^L = 1.6\) are shown in Fig. 1(b). The discharge was operated with nearly balanced injection \((\Gamma = 0.25)\) at an approximate cylindrical safety factor of \(q_{cyl} = 10\) having a low collisionality \(\nu_e^* \sim 0.2\) and \(Z_{eff} \sim 3-4\) (dominated by carbon impurities), while no clear impurity accumulation was observed during NBI; \(q_{cyl} = 5a^2B/(RI_p[MA]) (1 + \kappa^2)/2\). The \(\beta_p^{dia}\) value is measured to reach 3.2, for which the \(\beta_p\) value is calculated to be 2.9 (or \(\epsilon \beta_p = 0.69\)). While the degree of pressure anisotropy is \(\gamma = 0.1\), an orbit following Monte Carlo code calculation indicates a significant fraction of beam components around half of the \(\beta_p\) value. Strongly peaked profiles for the density and ion temperature are also observed with peaking parameters of \(T_i(0)/\langle T_i \rangle = 3.5\) and \(n_i(0)/\langle n_i \rangle = 2.6\), though the broader profile is apt to be maintained for the electron temperature insensitive to the safety factor. While no significant MHD activity was observed up to \(\sim 6.1\) s, the confinement properties were deteriorated with continuous \(m = 2\) oscillations from \(\sim 6.1\) s up to the beam turn-off. The tomographic reconstruction analysis for the soft X ray signals shows \(m = 2\) islands with a full width of \(\sim 10\) cm located at \(r \sim 25\) cm.

4. ENHANCED TOROIDAL ROTATION

A strong coupling between ion temperature and plasma flow velocity has been observed in conjunction with enhanced confinement. As shown in Fig. 2(a), the central toroidal rotation velocity \(V_\phi(0)\) is strongly enhanced in the co direction to the
FIG. 1. (a) Peak ion and electron temperatures and L mode enhancement factors ($r_p/r_e^L$) as a function of plasma current from $I_p$ scan experiments, together with observations of dominant MHD activity. (b) Waveforms of plasma current ($I_p$), beam power ($P_{NB}$), $\beta_p^{\text{ Shia }}$, ion temperature ($T_i$) and toroidal rotation velocity ($V_\phi$) for a typical high $\beta_p$ discharge.
plasma current with high rotational shear as the $T_i$ profile becomes peaked, in spite of nearly balanced injection in JT-60. This feature can be seen by comparing $V_\phi$ and $T_i$ profiles for hot ion enhanced confinement with balanced injection, L mode confinement with balanced injection and L mode confinement with pure co-injection.

In order to investigate the relationship between the $V_\phi$ and $T_i$ profiles, the $V_\phi$ values at $r/a ~ 0$ and $r/a ~ 0.4$ are plotted against the $T_i(0)$ values in Fig. 2(b). An offset linear correlation between the $V_\phi(0)$ and $T_i(0)$ values can be seen in this figure. It should be remarked that the $V_\phi (r/a ~ 0.4)$ values approximately maintain zero independent of $T_i(0)$, indicating that the rotational shear increases with the $T_i(0)$ value. The MHD effects are also shown to reduce the $V_\phi(0)$ and $T_i(0)$ values. Note that the ratios of the rotation velocity to the ion thermal velocity, $V_\phi(0)/V_{th_i}(0)$, are less than 20%.

From neoclassical theory [7], the toroidal rotation velocity can be related to the ion temperature in the expression

$$V_\phi = \frac{qR}{rB_i} \left[ E_r - \frac{1}{Z_i e} \frac{\partial T_i}{\partial r} \left\{ \frac{1}{\eta_i} + 1 - (\beta_1, g_{2i}) \right\} \right]$$

where $Z_i e$ is the charge, $E_r$ is the radial electric field, $\eta_i = d(\log T_i)/d(\log n_i)$ and $(\beta_1, g_{2i})$ are the neoclassical coefficients depending on the collisionality regimes. Since the experiments show $V_\phi (r/a ~ 0.4) ~ 0$ and $(\beta_1, g_{2i}) ~ 1$ in a banana regime, $E_r$ at $r/a ~ 0.4$ may be simply written as

$$E_r = - \frac{\partial \Phi}{\partial r} = - \frac{1}{Z_i e \eta_i} \frac{\partial T_i}{\partial r}$$

The result suggests that a large negative electric field exists in the rotational equilibrium, typically $-40$ kV/m where $\eta_i ~ 1.0$. This negative electric field causes a large electrostatic potential at the centre that is negative to the edge potential, which is calculated to be of the order of the central ion temperature, that is $[\Phi(0) - \Phi(a)] ~ -eT_i(0)$.

The energy balances in the core plasma tend to be decoupled for electrons and ions, where the convection losses are comparable to the conduction losses. The local transport analysis shows that the enhanced confinement results from the improved ion heat transport through a severalfold reduction of neoclassical heat diffusivity in the central region. The momentum transport analysis shows that, although the toroidal momentum diffusivity $\chi_\phi$ is close to the ion heat diffusivity $\chi_i$ around $r/a ~ 0.4$, the central $\chi_\phi$ value unrealistically becomes negative because of the large convection losses. So it is suggested that the observed enhanced rotation involves some effect other than a direct response of the plasma flow due to the beam momentum input.
FIG. 2. (a) Toroidal rotation velocity and ion temperature profiles for hot ion enhanced confinement with balanced injection \( (P_c = 9.9 \text{ MW}, P_{ce} = 8.3 \text{ MW}, q_{opt} = 10) \), L mode confinement with balanced injection \( (P_c = 11.5 \text{ MW}, P_{ce} = 7.8 \text{ MW}, q_{opt} = 4) \) and L mode confinement with co-injection \( (P_c = 11.5 \text{ MW}, q_{opt} = 4) \). (b) Toroidal rotation velocity at \( r/a = 0 \) and \( r/a = 0.4 \) as a function of the peak ion temperature for the discharges with \( I_p = 0.48 \) and 0.57 MA and \( B_t = 4.5 \text{ T} \), where MHD effects are indicated.
5. HIGH $\beta_p$ COLLAPSES

The high $\beta_p$ regime in JT-60 has been extended up to $\epsilon \beta_p \sim 1.2$ or $\beta_p \sim 5$ at $q_{cyl} \sim 17$ without major disruptions, as shown in Fig. 3, where $\beta_N$ indicates the normalized beta defined as $(\beta_0 \mu_0 I_p a B_t) / (100 \epsilon \beta_p / (2 \pi q_{cyl}))$. However, the discharges have encountered internal disruptions, called high $\beta_p$ collapses, in the best confinement during discharges in the regime of $\epsilon \beta_p \sim 0.7$ (or $\beta_p \sim 3$), so that confinement improvement in this regime was limited. As shown in this figure, a critical beta was observed to exist around $\beta_N \sim 1.1$, substantially below the Troyon limit ($\beta_N \sim 2.8$).

The soft X ray signals near the $\beta_p$ collapse and the fluctuation levels at the $\beta_p$ collapse are shown in Fig. 4 ($I_p = 0.5$ MA and $B_t = 4.5$ T). It is shown that the $\beta_p$ collapse is a partial MHD relaxation unlike a normal sawtooth collapse with a fast time-scale ($\sim 100 \mu$s) without beta saturation or any precursor oscillation just before the collapse. These MHD features suggest that the $\beta_p$ collapse is triggered by ideal low n kink/ballooning modes rather than resistive pressure driven modes or ideal high n ballooning modes.

Time dependent MHD mode analysis for the continuous low m modes provides information on the location of the collapse and the current profile during NBI. The results for the discharge of Fig. 4 are shown in Fig. 5, where the discharge encounters the $\beta_p$ collapse at $t = 7.53$ s during constant beam power. A possible mechanism inferred from these observations is as follows. (1) Current profile broadening is developed with a slow time-scale related to resistive field diffusion during NBI. (2) Improvement of the confinement is initiated with the change of the dominant mode from $m = 2$ to $m = 3$ at $t \sim 6.5$ s. (3) Increase of the bootstrap current fraction enhances the current profile broadening. The disappearance of the $m = 2$ mode
implies a minimum $q$ value above 2 after $t \sim 7.0$ s. (4) The increasing bootstrap current fraction causes a hollow current profile with a magnetic pitch minimum just below the $q = 3$ surface. (5) Finally, the discharge encounters the stability boundary for the ideal kink/ballooning modes.

Results from MHD stability analysis for low $n$ ideal kink/ballooning modes using the ERATO-J code are shown in Fig. 6(a); the $q$ profile used here is modelled so as to be changed from monotonic to hollow profiles. The stability boundary at $q_{\text{min}} \sim 3$ is found to be significantly reduced, down to $\beta_p \sim 3$, in association with the presence of a magnetic pitch minimum. An experimental trajectory inferred from the mode analysis is shown to be consistent with the ideal kink/ballooning analysis. The radial profiles of different Fourier components are shown in Fig. 6(b) with the $q$ profile, indicating that the internal $m = 3$, $n = 1$ mode is most dominant. We have also carried out ideal high $n$ ballooning stability analysis for the $\beta_p \sim 3$ regime using the BETA code, showing that the experimental pressure gradients are close to the stability boundary near the pitch minimum and substantially below that in the other regions.

6. Bootstrap Current

In JT-60, considerable efforts have been devoted to the confirmation of bootstrap current through comparison between experiment and theory because of its increasing attractiveness for steady state tokamak operation [8]. The neoclassical the-
ory of bootstrap current including fast ion effects is used for the analysis. Theoretically predicted and experimentally measured values of resistive loop voltages as a function of $\beta_p^{\text{dia}}$ are compared for the cases of neoclassical theory with and without bootstrap current and Spitzer resistivity. The experimental resistive loop voltages were estimated by $V_R = V_{\text{sur}} - (\mu_0/4)R_p I_p dl/dt$, where $V_{\text{sur}}$ is the surface voltage from the magnetic measurements. As shown in Fig. 7, a remarkable agreement between experiment and theory is obtained for neoclassical theory with the bootstrap current over a wide range of $0.1 \leq \beta_p^{\text{dia}} \leq 3.2$. For measurement, the bootstrap current fraction is defined as the measured resistive loop voltage normalized to the neoclassical resistive loop voltage without bootstrap current. For theory, the bootstrap current fraction is defined as the neoclassical resistive loop voltage with bootstrap current normalized to the neoclassical resistive loop voltage without bootstrap current.

FIG. 5. Temporal evolution of the $\beta_p^{\text{dia}}$ value, the location of $m = 2$, $m = 3$ modes, including the inversion radii (open circles) and the mixing regions (bars) for the $\beta_p$ collapse and the following collapse, and the fluctuation levels of soft X-ray signals for the $m = 2$ and $m = 3$ modes.
FIG. 6. (a) Stability diagram for ideal low n kink/ballooning modes represented by $\beta_p$ as a function of $q_{\text{min}}$ in which an experimental trajectory up to a high $\beta_p$ collapse is indicated, and the $q$ profiles used for the MHD stability analysis. (b) Radial dependence of the different Fourier components and the $q$ profile for $\beta_p = 3.5$.

FIG. 7. Bootstrap current fraction as a function of the $\beta_p^{\text{dia}}$ value, showing the measured values and the predicted values from neoclassical theory with bootstrap current.
FIG. 8. Temporal change of the surface loop voltage as measured and from neoclassical predictions with and without bootstrap current, the bootstrap current fraction and the internal inductance values from bootstrap calculations and equilibrium fits to magnetic measurements.

A large bootstrap current fraction up to ~80% of $I_p$ has been demonstrated under negligibly small beam currents less than 1% of $I_p$, as shown in the above analysis. Typically, this was obtained for the high $\beta_p$ discharge shown in Fig. 1(a). For further confirmation, taking account of the effects of electric field diffusion during NBI, we have carried out a 1.5-D time dependent transport analysis for the same high $\beta_p$ discharge. As shown in Fig. 8, the measured surface voltage trace is in good agreement with the prediction from neoclassical theory with bootstrap current, and the bootstrap current fraction is shown to reach ~80%. The previous analysis is justified by the fact that the radial profiles of the longitudinal electric field are calculated to become relatively uniform because of large bootstrap currents in the hot ion core and a long discharge duration normalized to the resistive skin time. As shown in Fig. 8, the calculated $I_l$ values agree well with the values from equilibrium fits to magnetic measurements. The development of the large bootstrap current is found to cause a hollow current profile with a magnetic pitch minimum.
Thus, the bootstrap current is concluded to be substantial, as calculated from the neoclassical theory, implying that electrons along the field lines behave classically in spite of their anomalous behaviour in the perpendicular direction.

7. CONCLUSIONS

The high $\beta_p$ regime up to $e\beta_p \sim 1.2$ with hot ion enhanced confinement has been established by means of nearly balanced perpendicular NBI in JT-60. The observed features of the enhanced toroidal rotation show a strong coupling with the peaked ion temperature profile and suggest a large negative potential at the centre of the order of the central ion temperature with enhanced confinement. The confinement improvement in the regime of $e\beta_p \sim 0.7$ was limited by high $\beta_p$ collapses with a critical beta $\beta_n \sim 1.1$ substantially below the Troyon limit. The MHD observations are consistent with low $n$ ideal kink/ballooning analysis associated with the magnetic pitch minimum. Conclusive confirmation of neoclassical bootstrap currents has been achieved for a wide range of $\beta_p^{de}$, $0.1 \leq \beta_p^{de} \leq 3.2$, in which the bootstrap current fraction reaches $\sim 80\%$.

In the high $\beta_p$ regime with high $q_{min}$ much greater than unity, a hollow current profile produced by bootstrap currents is found to play an essential role in MHD stability. The experimental evidence of the bootstrap current dominated discharges will have a significant impact on the feasibility of a steady state operation scenario for a future tokamak reactor. Strong central active current drive as a current profile control scheme with bootstrap current plus neutral beam current drive proposed in the SSTR [9, 10] could offer a way to suppress large MHD relaxation such as the high $\beta_p$ collapse. A further extension of this regime will be attempted in deuterium experiments for the JT-60 Upgrade programme [11].

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DISCUSSION

K.H. BURRELL: Your equation relating the ion temperature gradient to the radial electric field contains the parameter η, which is the ratio of the ion density gradient scale length to the temperature gradient scale length. The algebra for this formula shows that the electric field depends only on the ion density gradient and is independent of the ion temperature gradient. Since the formula has no dependence on the ion temperature gradient, how is it possible to use it to relate the electric field to the ion temperature?

S. ISHIDA: In the hot ion enhanced confinement regime, we have observed that the ion temperature profile peaks with the density profile. The dependence of the ion temperature gradient on the radial electric field therefore tends to be emphasized compared to the rather small change of the η value at r/a ~ 0.4.

R.J. HAWRYLUK: How was the fast ion component of the bootstrap current calculated? Was a Fokker-Planck or a Monte Carlo model employed? And what fraction of the bootstrap current was due to fast ions?

S. ISHIDA: We deduced the coefficients of the bootstrap current driven by the fast ion pressure gradient on the assumption of an isotropic velocity distribution with small banana size and pitch angle scattering dominant collisions, and used a Monte Carlo model for the fast ions.

For the typical high β discharge shown here, the fast ion driven fraction is calculated as 25% of the total bootstrap current.

S.C. LUCKHARDT: Recent theoretical work by J. Ramos indicates that the kink–ballooning beta limit has an inverse dependence on the central q. For example, \( \beta_{\text{crit}} = 3I/q(0)aB \). You have reported a case where q(0) = 3 and the beta limit is \( \beta_{\text{N}} = 1.1 \), so it appears that your data may be consistent with Ramos’s scaling. Have you studied other cases with q(0) > 1? And do they follow this scaling?

S. ISHIDA: In our analysis, it is the magnetic pitch minimum near the \( q = 3 \) surface that is essential for the kink–ballooning stability rather than the central \( q \) value q(0). It should also be noted that the q(0) value was not known experimentally. We have not studied other cases with q(0) > 1, so we cannot say whether Ramos’s scaling is consistent with the experiment.
STUDY OF HIGH POLOIDAL BETA PLASMAS IN TFTR* AND DIII-D**

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Abstract

STUDY OF HIGH POLOIDAL BETA PLASMAS IN TFTR AND DIII-D.
Tokamak equilibria with poloidal beta larger than $R/a$ have been produced in neutral beam heated discharges in TFTR and DIII-D. Values of $\epsilon_B p, dia$ as large as 1.5 were attained in TFTR with a natural inboard poloidal field null produced inside the vacuum chamber causing a transition to a diverted discharge. In DIII-D $\beta_p=5.1$ and $\epsilon_B p=1.8$ were reached with significant equilibrium modification. Both machines operate near the equilibrium conditions.

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limit at high $q^*$ and approach $\beta_N \sim 4$ at lower $q^*$. Energy confinement was observed to be enhanced to more than 3 times L-mode values. In TFTR large stored energy up to 2.5 MJ at 0.85 MA was produced with a measured $Q_{DD} = 1.1 \times 10^{-3}$. Experiments on DIII-D and TFTR have shown sustained periods of negative loop voltage with the calculation of substantial bootstrap current.

I. Introduction

There has been considerable interest in the operation of tokamaks at high values of poloidal beta, $\beta_p$, [1] because of the technical advantage for fusion reactor operation at a reduced plasma current, the theoretical prediction of large values of bootstrap current which may allow a significant reduction in current drive requirements, and the possibility for improved confinement due to the substantial modification of the equilibrium. In this paper we describe the results of a study of plasmas with high poloidal beta produced in neutral beam heated discharges in DIII-D and TFTR using relatively low values of the toroidal plasma current.

Tokamak equilibria with poloidal beta, $\beta_p \equiv 2\mu_0 <p>/<<B_p>>, \mathrm{larger} \ than \ R/a$ have been produced in neutral beam heated discharges in TFTR and DIII-D, where $<p>$ is the volume averaged plasma pressure and $<<B_p>>$ is the outer flux surface line average of the poloidal magnetic field. Values of $\Lambda \equiv \beta_p + \beta_t/2 \geq 6$, and $\varepsilon \beta_p$ as large as 1.5 were attained in TFTR. In DIII-D $\beta_p=5.1$ and $\varepsilon \beta_p$ up to 1.8 was reached. These results are summarized in Fig.1 as a function of $q^* \equiv (5a^2B_0/R_0\beta_p)[1 + \kappa^2(1 + 2\delta^2 - 1.2\delta^3)]/2$ where the plasma current, $I_p$, is in MA, and $a$ is the plasma half-width across the mid-plane. Both machines have operated near the equilibrium $\varepsilon \beta_p$ limit for large values of $q^*$. In the oblate ($\kappa \sim 0.7$, $R/a=3$) plasmas studied in TFTR, this limit appeared to be at $\varepsilon \beta_p \sim 1.6$, while in the elongated ($\kappa \sim 1.9$, $R/a=2.5$) plasmas studied in DIII-D the limit was near $\varepsilon \beta_p \sim 2$.

These plasmas exhibited significant modification of the equilibrium as they approached the equilibrium $\beta$ limit as shown in Fig.2. In the TFTR experiments, when $\varepsilon \beta_p \geq 1.2$, a natural inboard poloidal field null moved inside the vacuum chamber causing a transition to a diverted discharge. The magnetic axis was shifted strongly outward to about 50% of the 0.8 m minor radius. While tokamak equilibria of this type were previously produced transiently in the HBT tokamak[2] and values of $\varepsilon \beta_p \sim 1$ without an inboard null were produced in HBT, Versator II[3], and DIII-D[4], these experiments are the first to demonstrate a natural, inboard poloidal field null at high $\varepsilon \beta_p$ sustained for many energy confinement times in a large, high field, neutral beam heated device [5]. Separatrix formation was confirmed by the direct measurement of $B_p$ on the inside wall (shown in Fig.3a), an abrupt drop in the $H_{\alpha}$ intensity seen from the chord viewing the inside wall on the mid-plane (shown in Fig.3a), a tangential viewing video camera of visible light from the plasma, and an MHD model fit to the equilibrium. Diverted discharges were
Fig. 1. Summary of $\epsilon_\beta p$ achieved in DIII-D and $\epsilon_\beta p_{dia}$ in TFTR as a function of $q^*$.

(a)

Fig. 2. Equilibrium poloidal flux achieved in: (a) DIII-D where the dashed flux contours model the ohmic phase of the discharge when $\beta_p = 0.7$ and the solid flux contours model the peak $\beta_p = 5.2$, and (b) TFTR at $\epsilon_\beta p = 1.4 \pm 0.13$ for a 400 kA discharge.
Fig. 3. (a) Current ramp down TFTR discharge 46465 with 0.85 MA ohmic phase decreased to 0.4 MA immediately prior to NBI; (b) DIII-D discharge 67700: NBI start at 1.5 sec.
maintained for up to 0.5 sec, limited by the length of the neutral beam pulse. At the maximum values of $\varepsilon \beta_p$ achieved, the plasma assumed a highly oblate shape with $\kappa \sim 0.7$.

In DIII-D, a double-null divertor configuration was studied with $\kappa \sim 1.9$. At the largest values of $\beta_p \sim 5$, the inboard $B_p$ was reduced to less than 1/3 of the outboard value. The position of the divertor x-points was shifted towards the mid-plane, while the magnetic axis was shifted outward 36% of the 0.61 m minor radius as shown in Fig.2a.

II. Experimental Parameters

The basic machine and plasma parameters used in these studies are shown in Table I. In TFTR two types of plasma formation were employed: (i) constant current and (ii) current ramp down. In the ramped down case (an example is shown in Fig. 3a), an ohmic plasma was formed in the range 1 MA to 1.6 MA of toroidal plasma current and held for about 1 sec. The plasma current was then ramped down rapidly (2.5 MA/sec) to a value in the range 0.3 MA to 0.85 MA with a maximum ratio of initial to final current of about 2 to 1 prior to neutral beam injection. The plasma current was then held constant during neutral beam injection. The highest values of $\varepsilon \beta_p$ were obtained in the ramped down plasmas where the final value of $I_p$ was in the range of 0.4 MA to 0.6 MA with peak $T_i = 21$ keV and $T_e = 8$ keV. In the DIII-D experiments constant current discharges were studied in a double-null divertor configuration with $\varepsilon = 0.36$ and $\kappa = 1.9$. Up to 14 MW of either co- or counter- NBI was used. The target density was about $1 \times 10^{19}$ m$^{-3}$, rising after the H-mode transition until the first ELM.

In the TFTR experiments significant levels of stored energy and neutron production have been observed at the largest currents studied. In discharges with $I_p = 0.7$ MA ($q^* \sim 7.5$) about 2 MJ of plasma energy has been maintained for about 0.4 sec with a neutron production rate of $1.45 \times 10^{16}$ sec$^{-1}$ using 20 MW of beam heating. When $I_p$ was increased to 0.85 MA ($q^* \sim 6.5$),

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2.5±0.3 MJ of plasma energy was transiently produced with a peak neutron production rate of 2.1±0.3 \times 10^{16} \text{ sec}^{-1} using 22 MW of beam heating corresponding to a Q_{DD} of 1.1 \times 10^{-3}. This is comparable to the performance of the best 'supershots' at these power levels [6] but was achieved at a significantly lower plasma current.

III. Stability

In the TFTR experiments most discharges did not end in disruption and the MHD activity was relatively low amplitude. However, as \( q^* \) was reduced, an operational limit of \( \beta_N \sim 4 \) was found and the incidence of disruption near the \( \beta_N = 4 \) boundary increased sharply, where \( \beta_N \equiv 100<\beta_a B_o/I_p \) with \( I_p \) in MA. As shown in Fig.1 these discharges have achieved significantly higher levels of \( \epsilon \beta_p \) and stored energy than previously reported supershot parameters in the range of plasma current studied. These plasmas have exceeded the previously reported beta limit for supershots [7] of \( \epsilon \beta_p \sim 0.7 \) by about a factor of two. The ramp down of the plasma current during the ohmic phase of the discharge was found to be necessary to attain these high levels of \( \epsilon \beta_p \). From a transport analysis we observe the major effect of this current ramp down on the ohmic equilibrium was to remove (and in some cases to reverse) the toroidal current density in the outer third of the plasma. The core of the plasma was largely unaffected. This modification of the current profile prior to NBI has the effect of leaving the central \( q \)-profile unchanged while increasing the global shear in the outer third of the plasma.

In the DIII-D experiments, the L-H transition occurred within 20 msec of NBI initiation. During heating, the plasma was quiet until \( \epsilon \beta_p \sim 1 \) was reached. After passing this threshold, a series of bursts of MHD oscillations were seen \((\delta B/B_p \sim 0.2\%; f \sim 35-50 \text{ kHz}; \tau_{\text{burst}} \sim 2-10 \text{ msec}; m/n = 2/1)\) as indicated in Fig.3b. During the bursts, a small decrease in neutron production was seen, and \( I_t \) decreases, indicating a broadening of the fast-ion profile. As the stored energy, \( W_p \), saturates, the MHD bursts disappear. An extended stationary state is prevented by the onset of ELMs. At higher power, up to 20% of \( W_p \) is lost at the first ELM event; at lower power, \( W_p \) varies smoothly. At higher \( \beta_p \), the ELM frequency approaches 1 kHz.

One of the key parameters needed to assess the expected ballooning stability is the value of the central \( q \) [8]. A diagnostic using polarimetry of the emission from an injected lithium pellet [9] has been used to measure the internal profile of the poloidal field in TFTR. This \( B_p \) profile, together with the external currents and magnetic field, is used to reconstruct a least-squares 'best fit' free-boundary equilibrium for a constant 0.3 MA discharge, which indicates that \( q_0 = 1.6\pm0.55 \) and \( \epsilon \beta_p = 1.4\pm0.13 \). A detailed MHD stability analysis of this best fit model equilibrium indicates that the flux surfaces lie primarily in the first region of ballooning stability although the edge flux surfaces may lie in the unstable transition region. A stability analysis of a 0.4 MA discharge \((q^* \sim 10; \epsilon \beta_p \sim 1.3)\) modeled with \( q_0 = 1.138 \) yields similar results and is shown in Fig.4. Since the central \( q \) has not been increased above 2, these plasmas, while
having a very high level of $\epsilon \beta_p$, are largely in the first stability regime with the outer part of the plasma at or near the predicted boundary for high-$n$ ballooning modes as shown in Fig. 4a.

Including a single-point, motional Stark effect measurement of the poloidal field in the DIII-D equilibrium analysis gives a range $q_o \sim 1.3 - 2$ for
the high $\beta_p$ discharges. A discharge with $\beta_p = 4.7$ and $e\beta_p = 1.7$ has been tested for high-$n$ ballooning stability using the CAMINO code, and to $n = 0$ and $n = 1$ modes with the GATO code. The analyses were done for fitted equilibria with $q_0 = 1.5$ and 2. The discharge was stable to $n = 0$ and $n = 1$ modes with a conducting wall at the location of the vacuum vessel, but was unstable when the wall was moved to $1.5a$, indicating the discharge was close to the external kink limit and consistent with the observed value of $\beta_N = 3.7$. This plasma was stable to localized ballooning on all flux surfaces. Shown in Fig.4b is a rescaled $q'$ vs $p'$ stability diagram for this discharge assuming $q_0 = 1.5$. The central portion of this plasma lies in the first stability region while the edge is in the transition to the second region of stability. If $q_0$ were taken to be equal to 2, the full cross-section would be in the transition region between first and second stability regions.

IV. Confinement

Experiments in both machines observe energy confinement larger than L-mode predictions. In the TFTR current ramp down cases, $\tau_E$ greater than $3\tau_{\text{TER-P}}$ has been observed in 0.85 MA discharges using 10 MW to 22 MW of NBI. Prior to the onset of ELM activity, the DIII-D discharges show good H-mode behavior ($\tau_E = 2.3 \pm 0.5\tau_{\text{TER-P}}$). During the ELM phase, the energy confinement time is lower with $\tau_E \approx 1.5\tau_{\text{TER-P}}$. Shown in Fig.5 is a summary of the observed enhancement relative to $\tau_{\text{TER-P}}$ as a function of $e\beta_p$. In the DIII-D experiments, no strong trend indicating an improvement of confinement with increasing $e\beta_p$ was seen, but there clearly was no degradation as $\beta_p$ was increased. In the TFTR experiments, the largest confinement improvement relative to $\tau_{\text{TER-P}}$ was seen at the higher values of $\beta_p$.

Fig.5. Ratio of $\tau_E$ to $\tau_{\text{TER-P}}$ for DIII-D and TFTR as a function of $e\beta_p$. 
V. Bootstrap Current

In the DIII-D co-injection discharges, the current is predominantly non-inductive (beam-driven and bootstrap). Initially the current profile becomes peaked, but with the onset of MHD activity the profile broadens and $I_i$ decreases (Fig.3b). In the counter-injection discharges, $I_i$ reaches the same value, but the loop voltage is large and positive, indicating a significant level of reversed NBCD. One goal of these experiments has been to study the evolution towards fully non-inductive tokamak plasmas. We have identified a power threshold which must be exceeded to sustain the plasma solely with NBI [10]. This threshold is exceeded only during the $e\beta_p>1$ phase of these discharges. The power needed to reach zero loop voltage is within 30% of the predicted value.

For the 0.4 MA TFTR discharge using balanced neutral beam injection shown in Fig.3a, the plasma was collisionless with the central values of $v^* = v^*e \sim 0.1$. A transport analysis of this discharge calculates that 40% of the current is due to neoclassical bootstrap current at the end of the discharge. In the most collisionless experiments with longer beam pulses up to 1 sec and higher currents (0.6 MA to 0.85 MA), transport analyses have shown up to 70% of the current due to bootstrap current for several hundred msec at the end of the beam pulse. In these discharges the plasma surface voltage was observed to remain negative for the duration of the beam pulse but to relax over a period of 0.5 sec from greater than -1 Volt to a value near to -0.2 Volts, but it has not yet reached a steady state.

VI. Discussion

The results of these experiments are noteworthy in five respects: (i) the range of operation of large neutral beam heated tokamaks has been extended to the equilibrium limit at high $q^*$ with no degradation in confinement while achieving significantly enhanced confinement over L-mode predictions; (ii) at low $q^*$ an operational limit (consistent with the onset of low-n kink modes) at $\beta_N \sim 4$ is observed, although with careful profile control DIII-D has reached $\beta_N \sim 5$ [11]; (iii) large stored energy and significant QDD have been produced with these discharges in TFTR where a current ramp down method for $q$-profile control has resulted in roughly a factor of two improvement in the stored energy achieved for a given plasma current; (iv) up to 70% bootstrap current has been calculated for these plasmas in TFTR with sustained negative loop voltage; and (v) this expansion of the range of operation of large tokamaks to near the equilibrium limit and $\beta_N \sim 5$ coupled with confinement improvements of up to 3 times L-mode provides support for consideration of low current operating scenarios in the next generation of fusion devices and in reactor designs.

Acknowledgements

We wish to thank M. W. Phillips of the Grumman Corporation for making available the EQGRUM and STBAL equilibrium and stability codes used in the TFTR data analysis.
References


DISCUSSION

K. ITOH: I would like to ask about the temporal evolution of the separatrix in TFTR for the equilibrium $\beta$ limit. If the plasma is perfectly conductive, the separatrix cannot come to the plasma surface and $\beta_p$ can be high, in accordance with the concept of the flux conserving tokamak. This indicates the importance of the time evolution of the separatrix in high $\beta_p$ plasma. Do you have any information which identifies the physics process governing the evolution of the separatrix? In this sense, is the real equilibrium $\beta$ limit observed?

G.A. NAVRATIL: In these experiments, with a separatrix inside the vacuum vessel forming a diverted discharge, we see the position of the separatrix as roughly stationary when $\beta_p$ is maintained at a nearly constant value. We do not see the separatrix moving further into the plasma. However, the current relaxation time for these plasmas is longer than two seconds and these plasmas are still changing their current distribution.

A. GIBSON: To what extent is the $\beta$ limit in these discharges a hard disruptive limit and to what extent is it a soft limit?

G.A. NAVRATIL: At high $q^*$, when we approach the equilibrium limit, we occasionally see strong internal MHD activity which results in a soft beta collapse on the time-scale of the energy confinement time. At low $q^*$, when we approach the $\beta_N = 4$ limits, we do see hard disruptions at the limit. In the case of the highest current studied, 850 kA, we have not yet pushed to the disruptive limit.
HIGH DENSITY REGIMES AND BETA LIMITS IN JET

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United Kingdom

Abstract

HIGH DENSITY REGIMES AND BETA LIMITS IN JET.

Results are first presented on the density limit in JET discharges with graphite (C), Be gettered graphite and Be limiters. There is a clear improvement in the case of Be limiters. The Be gettered phase showed no increase in the gas fuelled density limit, except with ion cyclotron resonance heating (ICRH), but the limit changed character. During MARFE formation, any further increase in density was prevented, leading to a soft density limit, which was a function of input power and impurity content with a weak dependence on q. Helium and pellet fuelled discharges exceeded the gas-fuelled global density limits, but essentially had the same edge limit. In the second part, results are presented of high β operation in low-B Double-Null (DN) X-point configurations with Be gettered carbon target plates. The Troyon limit was reached during H-mode discharges and toroidal β-values of 5.5% were obtained. At high beta, the sawteeth were modified and characterised by very rapid heat waves and fishbone-like pre- and post-cursors with strongly ballooning character.

I. OPERATION NEAR THE DENSITY LIMIT

Operation near the density limit has been systematically studied in JET for limiter discharges, with C limiters, with evaporated Be layers and with Be limiters [1]. The operating density in tokamaks is usually presented in the form given in Fig. 1. Each point represents the maximum obtained normalised density during a discharge with either Ohmic (OH), Neutral Beam (NBI), ICRH or combined heating. The broken lines marked by OH (C) and NBI(C) are the highest limits obtained in the previous campaigns with C walls and limiters [2] (nRq/B = 12 for OH and 20x10^19m^-2T^-1 with NBI). There are clear improvements due to the Be limiters and Be coated walls. After conditioning of the Be limiter, which led to a strong reduction of Cl, the limit was substantially extended beyond that with Be gettering, so that nRq/B = 33 was reached with combined heating, ICRH and NBI heating. Furthermore the limiting density increases with the applied power as we shall see shortly. Pellet fuelled and He discharges have exceeded the deuterium gas fuelled limits. There is strong evidence that the edge density determines the limit. The limit at low q (q~2), however, is unchanged and is still set by major current disruptions.

* See Appendix to IAEA-CN-53/A-I-2, this volume. Including: J.P. Goedbloed and G. Huysmans, FOM-Instituut voor Plasma Fysica, Nieuwegein, Netherlands; T.C. Hender and O.J. Kwon, UKAEA/Euratom Fusion Association, Culham Laboratory, United Kingdom.
FIG. 1. Operating density range for JET shown as normalised current \( q'^{-1} = \pi R B / S A B \) [m, MA, \( m^2 \), T] versus normalised density \( n, R / B \) \( (10^{19} \text{ m}^{-3}, \text{ m}, \text{T}) \). Comparison is made between carbon disruptions (two broken lines at left) and MARFEs with Be evaporation (solid symbols) and Be limiter (open symbols).

FIG. 2. Normalised edge density versus input power showing that the MARFE density limit occurs at the boundary of the existence region close to the curve \( n_{\text{edge}} R / B = 2.37 P^{0.5} \) \( (10^{19} \text{ m}^{-3}, \text{ m}, \text{T}, \text{MW}) \). B varies in the range 1.4 to 2.6 T.
There has not yet been such a systematic study for X-point discharges, but the behaviour is similar, with somewhat lower density limits. The highest densities so far in X-point discharges are obtained in H-modes with \( nRq/B = 20 \times 10^{19} m^2 T^{-1} \). At higher densities, typically when \( nRq/P_1 = 60\% \), an H to L-mode transition occurs and the density falls without causing a disruption. In limiter discharges the nature of the limit is different for C and Be limiters. With C, the limit is marked by an asymmetrical edge radiation (MARFE) that leads to a symmetric radiative collapse and ends in a hard disruption. With Be limiters the limit is generally marked by the appearance of a MARFE which, in gas fuelled discharges, is accompanied by a fall in recycling and reduction in density. This typically leads to a soft density limit with a relaxation oscillation of density, radiation and MARFE near the limit. The internal inductance of the plasma during the MARFE generally did not increase, indicating that the plasma was not contracting significantly, as in the pure C limiter during the radiation collapse. This is consistent with the absence of strong MHD fluctuations and disruptions in the Be gettered and limiter plasmas.

The line-averaged density could be substantially increased with pellet fuelling, providing the pellets penetrated deeply. The limit for both pellet and gas fuelled discharges can be unified by considering the edge density and the input power (Fig. 2). The dependence of the density limit on the total power \( P_t \) and in particular on the radiation power balance in the edge region of the discharge has been suggested in a number of papers [2-5]. These models suggest that the limit should increase approximately as \( P_t^{1/2} \). The MARFE limit for pellet and gas fuelled discharges does lie at the boundary of the existence region.

2. CONCLUSIONS: DENSITY LIMIT

Operation near the density limit in JET can be summarised as follows:

- The density limit for additionally heated discharges in JET is now independent of the heating method and exceeds \( nRq/B = 33 \times 10^{19} m^2 T^{-1} \) for Be limiters and \( 20 \times 10^{19} m^2 T^{-1} \) for the H-mode;
- The density limit has a different MARFE behaviour for Be limiters compared with C limiters and leads to a density pump-out returning the discharge to a stable operation region;
- The density limit in gas and pellet fuelled discharges increases with input power approximately as \( P_t^{1/2} \) and is determined by edge parameters, particularly the edge density;
- The high limit obtained in JET means that acceptable densities should be reached in next step devices provided that sufficient degree of impurity exclusion can be obtained.

3. BETA LIMITS

High \( \beta \) operation has been achieved at low toroidal fields (\( B < 1.2 T \)) where the Troyon limit [6] is reached with additional heating at power levels ~10MW below that at which
Fig. 3. Maximum toroidal beta ($\beta_\varphi = 2\mu_0 \langle p \rangle / B^2$) as a function of $q_e^{-1}$ (proportional to normalised current $I_p/B a [MA, T, m]$), for all JET discharges with poloidal beta $\beta_p > 0.4$. The line is the Troyon limit $\beta_{\text{Troyon}} = 0.028 I_p/B a [MA, T, m]$. The highest $\beta$ is 5.5%.

Carbon self-sputtering becomes important [7]. It has not yet been possible to surpass the Troyon limit as has been done in DIII-D [8].

Fig. 3 shows the maximum toroidal $\beta_\varphi$ as a function of $q_e^{-1}$ obtained for all discharges between 1986 and 1989. A steady state $\beta_\varphi$ of 5.5% has been reached for DNH-modes in a hydrogen plasma. In these discharges with Be coated walls, $\beta$ saturation is generally observed without disruptions. The saturation is related to MHD-modes, ELM’s and $n=1$ activity. Sawtooth and fishbone events occur and sometimes continuous $n=1$, 2, or 3 modes appear, which can lead to a $\beta$ decline. A peaked and roughly triangular $p(r)$ profile develops from an initially broad profile. The internal inductance decreases from ~1 to 0.7, which indicates a broadening of $j(r)$ towards those profiles used in the $\beta$-optimisation by Troyon [6]. The decrease of the inductance is calculated to be due to the bootstrap current, which is approximately 25% of the total current.

4. BETA SATURATION

The evolution of $\beta$ for the discharge with the highest $\beta$ obtained so far is shown in Fig. 4. Also shown is the MHD activity, central ion temperature and volume-averaged density as a function of time. The main $\beta$-limiting mechanism in this discharge is the
high-β sawtooth. Increased MHD \((n=1\) and \(n=3\)) activity (around \(t=15\) s) leads to a diminished rate of rise in β after the crash and to a decline in the central ion temperature and so contributes to the β saturation.

The high-β sawteeth differ from those at low β in two ways:
1. The associated heat pulse is very rapid with \(\tau_{hp} \sim 100\mu s\) instead of ~10 ms.
2. Dominant \((1,1), (2,1)\) and higher \(m\) pre- and postcursors are seen, similar to high-β fishbones but of twice the amplitude. The modes have a ballooning character near the outer edge with a ratio in amplitude from the low to high B-side of ~10 as seen by the X-rays. Similar to a normal sawtooth, a high-β sawtooth causes a flattening of the pressure profile within the \(q=1\) radius.

5. HEAT LOSSES

Like other H-mode discharges in JET the high-β discharge has a confinement time twice that of the Goldston L-mode [9]. The observed plasma energy \(W_\text{DEA}\) lies close to the energy \(W_0\) calculated from the effective power input and \(\tau_\text{p} = 2 \times \tau_0\) [10]. The fraction of the losses due to high β sawteeth is 10 to 15% and that due to the intermittently appearing MHD-modes 20 to 30% of the total energy losses. This is sufficient to prevent further β increase since the heating power \(P\) is close to the critical power required to reach the Troyon limit. The fishbones and especially the sawtooth events strongly affect the fast particle distribution as measured by the neutron emission with consequences for future α-particle heating.

The central neutron emission drops by 70% (Fig.5) and its total rate by 30% during a sawtooth [11]. Fishbones are observed which individually cause up to 10% drop in the global neutron emission. However they occur about 10 times more frequently than sawteeth and may contribute appreciably to the central loss of fast particles and energy. Relatively large heat losses (150 kW) have also been measured by the neutral particle analyser with losses that are proportional to the MHD mode amplitude. Measurements with a multi-channel O-mode reflectometer indicate that high frequency density fluctuations grow exponentially with \(P_N\). The measurements are carried out between 3.9 and 4.1 m with frequencies ~130 kHz well above the \(n=1\) MHD modes present, perhaps indicating high-n ballooning-mode activity.

6. BETA COLLAPSE

In a few JET cases, high β collapses occur triggered or preceded by large \(n=2\) (or sometimes \(n=1\) or 3) MHD activity with \(δB\sim15\) G at the edge, and differ from β-saturation in various ways [7]:
- a dominant \((3,2)\) and other coupled \(n=2\) modes are responsible as seen from SXR analysis,
- there is a drop in the electron density in contrast to the saturation due to the high β-sawteeth,
- the central ion temperature and the fast ions are not affected at first.
FIG. 4. Evolution of $\beta_\phi$, MHD mode amplitude $B$ (top left), $H_\alpha$ and magnetic field $B_\phi$ (bottom left), ion temperature $T_i$, volume average density $\langle n \rangle$ (top right) and injected power $P_{\text{inj}}$, radiated power $P_{\text{rad}}$ for the 5.5% $\beta$ discharge, a 2MA double-null H-mode at 0.9T with 11 MW 80 kV D-injection into a H plasma. $T_e(0)$ and $T_i(0)$ of 3.5 and 6 keV were obtained in these low $q$ discharges ($q_{95} = 2.2$ or $q_e = 1.6$) and $\kappa$ of 1.8. $Z_{\text{eff}}$ slowly increases in time from 1.3 and levels off at $-2.5$. The confinement time $\tau_E = 0.35$ s.
7. PLASMA STABILITY

The stability of the high-β discharges has been examined with various stability codes: ERATO [12], HBT [13], BALLOON [14] and FAR code [15]. These stability studies are discussed more fully in [16]. It is found that before a high-β sawtooth the central plasma over more than half its radius is close to or even above the marginal ideal ballooning stability threshold. The ideal n=1 internal kink is also found to be strongly unstable for $\beta_N \approx 1$ when $q_0 \leq 1$. This instability may be linked to the observed (1,1) instabilities which seem to cause the β-saturation. We have calculated the fast particle effects on the internal kink. It is found that at the β values reached, the fast particles can no longer stabilise the internal kink. The operation is outside the Porcelli-Pegoraro stable region in the $(\gamma_{\text{mid}}, \alpha_0, \beta_p)$ space with experimental values of (1.0, 0.5, 1.5) [17]. In addition, severe fishbone activity is expected in this regime, resulting from the coupling with high energetic beam ions above 40 keV. It is further found that in the

FIG. 5. Cross-sections of plasma neutron emissivity before and after a beta crash. The central emission drops by 70%. The integration time is 100 ms.
cases where the β-collapse occurs internal modes of either \( n=2 \) or \( n=1 \) structure appear to be responsible for the enhanced plasma losses. These modes have been simulated by the FAR code where the q-profile has been tuned to match the measured X-ray fluctuations over the plasma cross-section [16]. In the case where \( n=1 \) modes are dominant, the q-profile had to be relatively flat in the centre with \( q(0) = 1.1 \), supported by Faraday-rotation measurements.

8. CONCLUSIONS: BETA LIMIT

In low q discharges at high β, saturation of the plasma energy is observed without disruptions. Global \( n=1 \) modes in the form of high-β sawteeth and fishbones are generally responsible for this saturation. Occasionally, β-collapses occur which seem to be related to large \( n=2 \) (sometimes \( n=1 \) or \( 3 \)) MHD modes. Triangular temperature profiles exist at the limit, which together with the rather flat density profiles lead to constant \( V_p \) across the plasma. Such peaked pressure profiles are favourable for a fusion reactor. Both the fishbones and sawteeth strongly affect the fast particle distribution. This has important consequences for future α-particle heating, burn control and wall loading. The role of the ballooning limit in the inner part of the plasma is not yet clear. Generally good agreement between theoretically predicted internal modes and observations at the beta limit has been obtained. The role of the fast particles on the beta limit needs further study both theoretically and experimentally. Further experiments in JET are required to see if the beta limit remains a soft limit even at much higher input powers.

REFERENCES

DISCUSSION

B. COPPI: The JET team’s assessment of the relative importance of electron thermal energy transport and of the anomalous ion energy transport appears to have changed since the Nice Conference in 1988.

P. SMEULDERS: In these high beta H-mode discharges the NBI heats the ions predominantly. We usually find in JET that the electron heat conductivity is higher than that of the ions. It all depends on the regime in which the discharge is operated.

S.A. COHEN: Does the density limit in the divertor configuration depend on the density at the separatrix (midplane) or the density at the divertor plate?

P. SMEULDERS: This is an interesting question, but no information is available as yet on this subject.

Ya.I. KOLESNICHENKO: What was the slowing-down time for fast ions and the time between the sawtooth crashes in experiments with neutral beam injection and ICRH?

P. SMEULDERS: The high beta experiments described in this paper were with \( \approx 10-15 \) MW, 80 kV NBI alone. The sawtooth period was around 0.5 s. Slowing-down times were typically between 80 and 120 ms in the plasma centre.

R.J. GOLDSTON: I think you will find that \( nRq/B = 20 \) is not quite adequate for ITER. Can you comment on whether there is any hope to be derived from a power scaling of the H-mode density limit, similar to the scaling of the L-mode limit?

P. SMEULDERS: We do have data about the power scaling of the density limit in H-modes, and in that connection I can refer you to the paper presented by Tanga (IAEA-CN-53/A-IV-1, this volume). The same scaling as in the limiter configuration applies.

We also think that at higher values of \( q_0^{-1} \) (higher shaping), \( nRq/B = 20 \) can yield adequate density values for ITER.
TURBULENT FLUCTUATIONS AND TRANSPORT IN TEXT*

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Abstract

TURBULENT FLUCTUATIONS AND TRANSPORT IN TEXT.

Two techniques, heavy ion beam probe (HIBP) measurements of $\bar{n}/n$ and modulated gas feed have been used to investigate transport in the TEXT tokamak. The HIBP measurements indicate that the frequency averaged wave vector ($k$) of the turbulent fluctuations in the interior of TEXT is approximately $0.3 \text{ cm}^{-1}$, which is much smaller than expected, and that the radial correlation length is greater than 2 cm. The modulated gas feed technique implies an inward convective velocity for perturbations even greater than that for the equilibrium.

1. INTRODUCTION

Transport in the TEXT tokamak has been investigated by carrying out multi-point measurements of $\bar{n}/n$ using a heavy ion beam probe (HIBP) [1] and by using a modulated gas feed technique [2]. The localized HIBP measurements were carried out at all radii and cover the frequency range 0-500 kHz. The modulated gas feed experiments were used to determine a diffusion coefficient and a convective velocity. The results of both experiments are inconsistent with drift wave theory.

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1 Magnetic Fusion Science Fellowship Program administered by Oak Ridge Associated Universities, Inc., Oak Ridge, Tennessee, USA.
2. HIBP MEASUREMENTS

The HIBP provides direct simultaneous measurements of \( f_i / n \) from two small sample volumes (~0.1 cm\(^3\)). Using digital time series analysis, we calculate the coherence, phase difference and frequency averaged wave vector, \( \langle k \rangle \), for the turbulent fluctuations. \( \langle k \rangle \) is in the sample volume separation direction and in general has both \( \langle k_r \rangle \) and \( \langle k_\theta \rangle \) components. For the HIBP measurements, TEXT was operated at \( B_T = 2.16 \) T, \( \bar{n} = 1.6 \times 10^{19} \) m\(^{-3}\) and \( I_p = 200 \) kA.

2.1. Edge measurements

An example of an edge \((r/a \geq 0.8)\) measurement is shown by the bold lines in Fig. 1. The \((r/a, \theta)\) locations of the two sample volumes were \((0.85, 29.4^\circ)\) and

**FIG. 1.** Autopower spectra, coherence and phase angle for an edge (heavy line) and an interior (light line) measurement for a two point \( f_i / n \) measurement.
(0.84, 35.4°). Note that spatial aliasing occurs in the phase measurement at \( \sim 120 \text{ kHz} \). The \( \langle k \rangle \) for these data is \( \sim 0.9 \text{ cm}^{-1} \), which is about a factor of 2 lower than Langmuir probe and FIR scattering measurements. It should also be noted that in general the radial correlation length in the edge region is very short (< 1 cm).

### 2.2. Interior measurements

In the interior of the plasma (\( r/a \leq 0.6 \)) the results are quite different, as shown by the thin lines in Fig. 1. The \( (r/a, \theta) \) location for these two sample volumes is (0.55, 20.8°) and (0.53, 29.4°). \( \langle k \rangle \) for these data is 0.3 cm\(^{-1} \). The interior wave vectors are approximately a factor of 3 smaller than the edge wave vectors, and the radial correlation length in the interior is much larger (>2 cm).

A simplified dispersion relation for electron drift waves [3] can be written as

\[
\omega = \frac{\omega_{e} - \omega_{\perp}}{1 + k_{\perp}^{2}\rho_{e}^{2}} + v_{E} \bar{k}_{\theta}
\]  

where \( k_{\perp}^{2} = k_{r}^{2} + k_{\theta}^{2} \) and \( v_{E} \) is the steady state \( E \times B \) rotational velocity. As noted earlier, the measured \( \langle k \rangle \) in general has both radial and poloidal components, but no reasonable relation between \( k_{r} \) and \( k_{\theta} \) will permit agreement between the experimental results and Eq. (1).

The small \( \langle k \rangle \) measurements reported here are in clear disagreement with drift wave theory and with FIR scattering measurements [4]. FIR scattering quotes a mean poloidal wavenumber of \( \langle k_{\theta} \rangle = 3 \text{ cm}^{-1} \), which agrees with drift wave theory. The longer radial correlation lengths observed with the HIBP are in agreement with results reported by Fonck [5] for the interior of PBX-M. The difference between the HIBP and the FIR measurements is not understood and is under investigation. The two techniques have different sensitivities and spatial resolutions for various \( k \) values, so they may be observing different parts of the spectra or different phenomena. Note that edge fluctuations can modulate the beam intensity, which affects the interpretation of the interior measurements. Such modulation would be interpreted as a smaller \( \langle k \rangle \) in the interior than is actually present. Beam modulation cannot explain all of the characteristics of the interior signals and there is experimental evidence that the beam modulation is very small for low density plasmas.

### 3. MODULATED GAS FEED EXPERIMENTS

Particle transport coefficients for hydrogen have been determined for a wide variety of discharges in TEXT by a perturbative technique, modulation of the gas feed [2]. Both diffusion (D) and convective velocity (V) have been determined with radial resolution at least sufficient to distinguish central from outer regions. The
presence of both diffusion and convection precludes simple comparisons with fluctuations, but the coefficients may be examined for features which would be implied by possible theories. Two sorts of comparison are made. The first is scaling with global discharge parameters and radius to ascertain which parameters seem most important. The second is comparison with the consequences of assuming transport coefficients depending on local parameters [6], especially gradients, e.g. \( D(n, \partial n/\partial r, T, \partial T/\partial r) \). The observations span a broad range of densities, plasma currents and toroidal fields (\( 1 \times 10^{19} \text{ m}^{-3} < n < 6 \times 10^{19} \text{ m}^{-3} \), \( 100 \text{ kA} \leq I_p \leq 300 \text{ kA} \), \( 1.5 \text{ T} \leq B_T \leq 2.8 \text{ T} \)).

For parametric dependences, both \( V \) (for \( r > 0.4a \)) and \( D \) (everywhere) show a clear and consistent inverse density dependence. However, there are no similar clear scalings with \( q \), toroidal field or current. (This differs from the coupled temperature–density perturbations of sawteeth, whose propagation shows a marked \( q \) and a weak density dependence [7].) Experiments with ECRH evince no strong \( T_e \) dependence. Like other diffusivities, \( D(r) \) generally rises toward the edge, but not as strongly as the turbulence level. The equilibrium profiles broaden for low \( q \) but show little density dependence [8].

The results do not depend upon either the amplitude (2–15\%) or frequency of density modulation, which argues against a marginal stability model. However, \( D \) and \( V \) are characteristic of a perturbation about the equilibrium and are not necessarily those of the equilibrium. If the form of \( D \) is known, e.g. \( D \propto n^\alpha (\partial n/\partial r)^\beta \times T^\gamma (\partial T/\partial r)^\delta \), relations between the equilibrium and perturbed quantities may be inferred and compared with the observations. The class of simple drift wave models \( \alpha \sim -1, \beta \geq 1 \) and \( \gamma > 1 \) imply

\[
D = (1 + \beta)D_{eq}; \quad V \sim V_{eq}
\]  

(2)

Some more sophisticated models which include convection [9, 10] imply a form for the particle flux \( \Gamma = -D_0[(\partial n/\partial r) - c(n/T)(\partial T/\partial r)] \), with \( D_0 = T^{3.5}/B^2L^2_\alpha n \) and \( c \)

---

**FIG. 2.** Comparison of transport coefficients for perturbations with those for the equilibrium.
a constant which must be somewhat less than unity to match observed density profiles. Such a model implies

$$D = D_{eq}; \quad V = V_{eq}$$

(3)

The ratio of $D_{eq}$ and $V_{eq}$ can be determined from the equilibrium profile and may be compared with the ratio for the perturbed coefficients as the dimensionless parameter $S = aV/D$. Figure 2 shows $S$ from the perturbation plotted as a function of the corresponding $S_{eq}$ for the whole range of discharges, both values being evaluated for $r \geq 0.4a$. The value of $S$ for the perturbation is consistently larger than that for the equilibrium by factors of 1 to 3. The result contrasts markedly with the predictions of Eqs (2, 3), both of which imply $S < S_{eq}$. Generally speaking, it is most difficult to reconcile these observations with any model in which the diffusion is driven primarily by the density gradient.

REFERENCES


DISCUSSION

P.C. EFTHIMION: You observe a density dependence in your particle diffusivity. Doesn’t this density dependence therefore require a non-linear analysis of the particle flux? Non-linear transport would explain why your linear perturbation analysis does not agree with the equilibrium analysis. In addition, the large convective term $V$ which you observe may be caused by the non-linear particle flux. Our work over the last year has focused on using a non-linear analysis in perturbation experiments to measure the particle transport.

K.W. GENTLE: The parametric density scaling could be interpreted as a local density dependence, particularly for the diffusion coefficient, but this does not alter the analysis. Dependence of the transport coefficients on local plasma parameters is
explicitly included; this accounts for differences between $D'$ and $D_0$, the perturbed and equilibrium values, for example, in a simple drift wave model. The analysis of the density perturbations is a formally complete linear analysis, the validity of which is experimentally confirmed by varying the amplitude of the induced density perturbation from 2% to 15% and finding that the transport coefficients inferred do not change.
LOCAL TRANSPORT MEASUREMENTS DURING AUXILIARY HEATING IN TFTR


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Abstract

LOCAL TRANSPORT MEASUREMENTS DURING AUXILIARY HEATING IN TFTR.

Local heat and momentum transport is analyzed in beam-heated steady-state TFTR plasmas using measured profiles for temperature and rotation speed over a wide range of plasma conditions. As heating power is increased in high-recycling 'L-mode' plasmas, the thermal diffusivities $\chi_t(r)$ and $\chi_e(r)$ are found to increase strongly with temperature throughout the plasma confinement zone. Both $\chi_t$ and $\chi_e$ decrease with increasing plasma current in steady state L-mode plasmas. Dedicated major-radius and aspect ratio scaling experiments show that the global confinement time $\tau_E$ has a strong dependence on major radius and a weak dependence on minor radius in L-mode plasmas, $\tau_E \propto R^{-1.6} a^{-0.3}$. Kadomtsev dimensionless scaling of the ohmic $\tau_E$ obtained in Alcator C and PLT to TFTR at constant $\tilde{n}_e/B_0^2$, $q_{95}$, and $a/R$ underpredicts $\tau_E$ by about a factor of two, possibly due to variations in $L_n/a$. Ion heat and

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momentum transport in low-recycling TFTR 'supershot' plasmas shows an inverse dependence on temperature. Several independent measurements of radial particle transport indicate diffusivities of order $\leq 0.1 \text{ m}^2/\text{s}$ for energetic ions ($\sim 30 \text{ keV} - \text{ MeV}$), an order of magnitude less than the observed thermal transport rates.

1. Introduction

An extensive campaign of transport measurements has been carried out in the Tokamak Fusion Test Reactor [TFTR] to characterize the functional dependence of local heat and momentum transport in L-mode plasmas, and in peaked-density supershot plasmas with improved energy confinement. One objective of this study is to establish, on an empirical basis, the 'origin' of global L-mode energy confinement scaling ($\tau_\text{E} \propto I_p R^{1.75} a^{-0.37} F_b^{-0.5}$ [1]) from the parametric dependence of local heat transport coefficients ($\chi_e$, $\chi_i$). In particular it is desired to identify whether ions or electrons are the dominant power loss channel, which one (or both) causes the observed confinement deterioration with input power and the favorable scaling with plasma current, and whether these changes reflect variations in transport across the entire plasma cross section or only at the plasma edge or center. The ultimate objective of these studies is to identify the dependence of local heat flows on local plasma parameters and gradients for comparison with theory. Beam power scans in L-mode plasmas indicate a strong increase of $\chi_i$ and $\chi_e$ with heating power, and correspondingly with local temperature, that is responsible for the unfavorable $F_b^{-0.5}$ scaling of L-mode $\tau_\text{E}$. Increasing the plasma current is observed to reduce the transport across the entire confinement region in L-mode plasmas, but it still retains an unfavorable scaling with temperature at high current. By contrast, ion heat and momentum transport in high-temperature TFTR 'supershots' ($T_i \geq 10 \text{ keV}$) appear to decrease with increasing heating power, and correspondingly with increasing temperature, density, and $\beta$ [2].

The mechanism responsible for the transition from unfavorable to favorable temperature scaling is not understood, but suggestive evidence is provided by measurements of radial transport of energetic ions. Several independent measurements of the fast-ion radial diffusivity (for both passing and banana trapped ions) indicate $D_{\text{fast}} \approx 0.1 \text{ m}^2/\text{s}$, or about an order of magnitude less than the characteristic transport rates for the thermal plasma. Thus there appears to be some energy threshold beyond which ions are insensitive to the turbulence that drives anomalous transport in the thermal plasma. Irrespective of whether this threshold is related to $T_i/T_e$,
to the ratio of the ion orbit size to the turbulence correlation length, or to some other parameter, it is plausible that a significant fraction of the thermal ion population in TFTR supershots (which typically have $T_i/T_e \geq 1.8$) exceeds this threshold and consequently suffers less transport.

In addition to the local transport studies, we describe two scaling experiments that refine the empirical $\tau_E$ scaling relations from which the ignition requirements for BPX and ITER have been estimated. The first is a major-radius and aspect ratio study in beam-heated L-mode plasmas which extends the aspect ratio to 8. The second is a test of Kadomtsev dimensionless scaling comparing ohmic plasmas from the Alcator C, PLT, and TFTR tokamaks that collectively span a range of 200 in density and 30 in energy confinement time.

2. Plasma Transport Scans

Beam power, plasma current, and particle recycling ratio scans have been performed, holding other quantities constant to the extent possible. The measurements were taken during the quasi steady-state phase of beam-heated discharges that was typically established $\leq 0.5$ seconds after the start of neutral injection. Local coefficients for heat and momentum transport were inferred from local profile measurements of temperature, density, and rotation speed using the 1-D steady-state transport analysis code SNAP. $T_e(R)$ was measured by first harmonic ECE radiometry and by Thomson scattering; $n_e(R)$ by a 10-channel infrared interferometer, and $T_i(R)$ and $\nu^\phi(R)$ by charge-exchange recombination spectroscopy. Further details of the diagnostic set and the SNAP models for beam power deposition and local heat transport are described in references 3 and 4. To avoid difficulties posed by TFTR supershots for which the convective flows of heat can dominate the total energy flow [2], we will consider ‘total’ diffusivities in this paper, defined by $Q_i \equiv -\chi_i(\sum_j n_j)\nabla T_i$, $Q_e \equiv -\chi^e n_e \nabla T_e$, and $\Gamma^\phi \equiv -\chi^\phi(\sum_j n_j m_j)\nabla \nu^\phi$ where $Q_i$ and $Q_e$ are the total ion and electron energy flux, $\sum_j$ represents a sum over all ion species, and $\Gamma^\phi$ is the radial flux of toroidal momentum. Classical ion-electron coupling $P_{ei}$ was calculated to be a significant power transfer term in L-mode plasmas owing to the low temperatures[3,5] especially in the region $r/a \geq 0.7$. In regimes where the estimated measurement error in $(T_i - T_e)$ creates uncertainties in $P_{ei}$ that are comparable to the total energy flow through a flux surface ($\int Q_idA$ or $\int Q_e dA$), radial heat transport cannot be separated accurately into ion and electron channels and it is more illuminating to evaluate a single-fluid thermal diffusivity, $\chi_{avg} \equiv (Q_i + Q_e)/(\sum_i n_i \nabla T_i + n_e \nabla T_e)$. 
Several general trends were observed in the scans summarized in Table I. First, the dominant energy loss channel in all scans was through the ions; the ion energy confinement time was typically less than half that of electrons at both $r/a = 0.3$ and 0.7. The ions received between 1/3 and 2/3 of the beam power inside of $r/a = 0.7$ (the fraction was higher at the plasma center due to higher $T_e$), but including ion-electron coupling, the ions carried 60-90% of the total power flow through the $r = 0.7a$ surface. Second, the beam deposition profile did not change significantly throughout any of the scans because the variations in density were relatively modest. The ‘heating effectiveness’ parameter[6], which represents the ratio of stored energy obtained with a given heating profile to that expected for a delta function at the magnetic axis, varied less than 11% in every scan. Thus, less than 11% of the variation in integrated quantities ($W_e$, $W_i$, $\tau_{Ei}^{th}$, etc.) can be directly attributed to changes in beam penetration.

<table>
<thead>
<tr>
<th>SCAN</th>
<th>Range</th>
<th>$P_b$</th>
<th>$P_a$</th>
<th>$B_T$</th>
<th>$a$</th>
<th>$\bar{n}_e$</th>
<th>$T_{Te}$</th>
<th>$T_{Tc}$</th>
<th>$F_{ne}$</th>
<th>$Z_{eff}$</th>
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<td>2.9</td>
<td>3.9</td>
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<td>1.7</td>
<td>1.8</td>
<td>2.5</td>
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<tr>
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<td>2.2 – 20</td>
<td>1.8</td>
<td>3.8</td>
<td>93</td>
<td>3.2</td>
<td>3.2</td>
<td>2.8</td>
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<td>1.9</td>
<td>1.7</td>
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<tr>
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<td>3.6</td>
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<td>80</td>
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Table 1: Summary of local transport scans. Double column entries represent the values of each quantity at the start and end of the scan, respectively. Variable definitions: $P_b$ (MW); $P_a$ (MA); $B_T$ (Tesla); $a$ (m); $\bar{n}_e$ ($10^{19}$ m$^{-3}$); $T_{Te}$, $T_{Tc}$ (keV); $F_{ne}$ = density peakedness = $n_e(0)/\langle n_e \rangle$; $W_e$, $W_i$, $W_b$ = total electron, ion, beam stored energy (MJ); $\tau_{Ei}$, $\tau_{Ei}^{th}$ = thermal ion, electron energy confinement time at $r/a = 0.7$ (ms); $\tau_{Ei}^{th}$ = global energy confinement time = $(W_i + W_e + W_b)/(P_b + P_{Hi})$ (ms); Re = limiter recycling (H = high, L = low); $(h(p)p)$ = mean minor radius of beam power deposition expressed as a fraction of the plasma minor radius; $\eta$ = ‘heating effectiveness parameter’ [6] of the beam deposition profile; $F_{pe}$ = fraction of total power delivered to ions or electrons inside $r/a = 0.7$ (excluding ion-electron energy exchange); $F_{ei}$ = total power exchanged between ions and electrons inside $r/a = 0.7$ normalized to the absorbed heating power at the same radius; and $L_i$ = fraction of radial power flow at $r/a = 0.7$ carried by ions. Some power fractions do not add exactly to 1.0 due to viscous heating and radiation terms.
Fig. 1. (a) Variation in central temperatures, line-averaged electron density, thermal energy (\(W_{\text{th}}\)), and total energy (\(W_{\text{tot}} = W_{\text{th}} + W_{b}\)) with beam power at \(I_p = 1.2\) MA (scan A of Table I). (b) Single-fluid thermal diffusivity as a function of beam power. The error bars indicate the uncertainty in \(\chi_{\text{avg}}\) due to estimated systematic errors in the measurements of \(T_e(R), T_i(R), n_e(R), Z_{\text{eff}},\) neutral density, and other parameters. The ion temperature profile for the ohmic discharge was calculated assuming \(\chi_i = 2\chi_e.\)
2.1. Beam power scans

Scans A and B in Table I represent beam power scans at \( I_p = 1.2 \) and 1.8 MA in high-recycling deuterium plasmas for which the ohmic target density was adjusted to yield approximately constant \( \bar{n}_e \) at the end of injection (the time of transport analysis). Heating in these L-mode plasmas was very poor, e.g. less than a factor 1.6 increase in \( T_i \) and \( T_e \) at \( I_p = 1.2 \) MA despite a factor seven increase in beam power (Fig. 1a). The global energy confinement followed the L-mode power scaling \( \tau_E \propto P_b^{-0.6} \) [1], but only because the stored beam energy – which represented approximately half the total stored energy at high power – increased roughly linearly with beam power. The stored thermal energy increased only as \( P_{\text{tot}}^{0.29} \) [5].

Figure 1b illustrates the variation of the single-fluid thermal diffusivity \( \chi_{\text{avg}}(r) \) with beam power. The radial profile of \( \chi_{\text{avg}}(r) \) became somewhat steeper with the addition of the first 2.4 MW and thereafter remained fixed, while its magnitude increased rapidly with beam power. The ion energy confinement time in these scans was typically more than a factor of two lower than that of electrons at both \( r/a = 0.3 \) and 0.7. Energy transport by both ions and electrons increased rapidly with power (Fig. 2), although the uncertainty in \( \chi_i \) and \( \chi_e \) was large due at low power to strong ion-electron power exchange, which was calculated to exceed 50% of the

![Graph](image-url)  

**Fig. 2.** Dependence of \( \chi_e \) and \( \chi_i \) on heating power at \( r/a = 0.5 \) in (L-mode plasmas at \( I_p = 1.2 \) MA (scan A in Table I).
total input power at \( r/a = 0.7 \). At high power, the ion-electron power exchange was relatively smaller, and the dominance of the ion loss channel became clearer. The power dependence of \( \chi_i \) and \( \chi_e \) in L-mode plasmas can also be represented as a strong temperature dependence (Fig. 3). The implied temperature scaling of heat transport (\( \chi_{\text{avg}} \propto T^m \); \( m = 2 - 4 \)) is reminiscent of the observed temperature scaling of electron particle diffusivity in similar L-mode plasmas (\( D_e \propto T_e^{1.5-2.5} \) \([7,8]\)). Note it is found that at fixed heating power, both \( \chi_i(r) \) and \( \chi_e(r) \) increase with minor radius despite the decreasing temperature \([3]\), indicating that other variables such as \( q \), gradient scale lengths, or the fraction of trapped particles must also strongly affect transport.

Fig. 3. Variation of \( \chi_i \) and \( \chi_e \) with local temperature at \( r/a = 0.5 \) in L-mode beam power scans at \( I_p = 1.2 \text{ MA} \) and 1.8 MA (scans A and B). The range of beam power is \( \sim 2-20 \text{ MW} \).
An additional pair of beam power scans (scans C and D in Table I) were performed at $I_p = 1.3$ MA to study differences in heat and momentum transport between broad density, high-recycling L-mode plasmas and low-recycling plasmas which attained more peaked density profiles and considerably higher temperatures. The latter were nominally in the 'supershoot' regime at high power. Sets of co- and counter-directed beam sources (e.g. 3-co + 0-ctr; 4-co + 1-ctr; ...; 6-co + 3-ctr) were selected to provide approximately constant beam torque as the beam power was increased from 7 to 21 MW. By measuring the toroidal rotation profile, it was possible to infer a temperature scaling of the ion momentum diffusivity under conditions where the driving term (beam torque) was fixed, for comparison with the ion and electron heat diffusivity. One consequence of unbalanced injection was a correlation between density peakedness and beam power in the low-recycling plasmas ($F_{ne} = n_e(0)/\langle n_e \rangle = 2.1 \rightarrow 2.4$). These values are somewhat less than the density peakedness obtained in supershots with well-balanced injection ($F_{ne} \geq 2.5$). Correspondingly, these low-recycling plasmas obtained global $\tau_B$ less than the best supershot results, but compared to the L-mode regime they attained significantly higher central temperatures, and higher thermal and total energy confinement time ($\tau_B/\tau_B^{Lmode} \leq 2.6$).
High-recycling conditions were created for the L-mode scan by puffing helium gas into the ohmic target plasma ($n_e \approx 2.5 \times 10^{19} \text{ m}^{-3}$) without changing the low-recycling properties of the limiter with respect to deuterium. The high-recycling plasmas thus obtained had the broad density profiles, low temperatures, and global energy confinement times characteristic of the L-mode regime; $\tau_E$ remained within 40% of the L-mode value throughout the scan.

Figure 4 illustrates the scaling of central temperatures (b,d) and momentum density (c) with heating power for both low-recycling and L-mode
plasmas. The observed change in $T_{i}(0)$ is considerably greater than the change in $T_{e}(0)$, which results both from differences in the scaling of ion and electron thermal diffusivities, and from the increasing fraction of beam power delivered to thermal ions as $T_{e}$ increases. A clear indication of the very different scaling of ion transport with temperature (or beam power) in the two regimes is that the central momentum density decreases with heating power in the L-mode plasmas but increases with heating power in the low-recycling plasmas, at constant torque input.

The scaling of $\chi_{i}$ and $\chi_{\phi}$ with heating power in the L-mode and low-recycling plasmas is illustrated in Fig. 5 at three radii ($r/a = 0.72, 0.48,$ and $0.22$). As observed earlier in L-mode deuterium plasmas scans, the L-mode $\chi_{i}$ shown in Fig. 5 increases systematically with power across the entire confinement region, somewhat more strongly at the plasma edge than at the center. Similarly, the L-mode $\chi_{\phi}$ increases with heating power (less strongly than $\chi_{i}$) over the entire cross-section. In the periphery of supershot plasmas ($r/a \approx 0.7$), $\chi_{i}$ was about a factor of two lower than in the L-mode, but it retained the unfavorable scaling with power. By contrast, $\chi_{i}$ was essentially independent of power at $a/2$ in low-recycling plasmas, while near the plasma center ($r/a = 0.22$) the temperature dependence reversed, and $\chi_{i}$ actually decreased with increasing power (or $\beta$, $n_{e}$, $T_{i}$, or $T_{e}$) up to $\sim 12$ MW, and remained constant thereafter. If the ion power balance is evaluated by subtracting a convective term $q_{i-\text{conv}} \equiv \frac{3}{2} \Gamma_{i} T_{i}$ from $Q_{i}$ as shown in Fig. 5d (where $\Gamma_{i}$ is the radial ion flux), the ion thermal diffusivity $\chi_{i}$ decreases monotonically and strongly (factor $\sim 4$) with heating power at the plasma center. Further out in radius ($r/a \geq 0.4$) the convective term becomes a smaller fraction of the total ion heat flux, so $\chi_{i}(r)$ and $\chi_{i}(r)$ scale similarly with power. In summary, the high ion temperatures attained in low-recycling plasmas appear to arise from a reduction, relative to the L-mode level, of ion heat transport which retains an unfavorable dependence on heating power (or variables correlated with it, such as $F_{ne}$ or temperature) in the plasma edge but a favorable dependence in the center.

The behavior of $\chi_{\phi}$ is of interest in these scans because $\chi_{\phi}$ has been found to correlate with $\chi_{i}$ in uni-directional low-recycling plasmas[4], and thus may corroborate the inference of reduced ion transport with increasing heating power. Ion-electron coupling is generally small terms in the momentum balance, thus $\chi_{\phi}$ can be determined more accurately than $\chi_{i}$ or $\chi_{e}$. Of particular interest is whether $\chi_{\phi}$ increases with heating power, as it does during uni-directional injection in low-recycling [4] and L-mode[3] plasmas, if the beam power is partially balanced. As shown in Fig. 5, the variation of $\chi_{\phi}$ with beam power is similar to that of $\chi_{i}$; it increases with heating power throughout the L-mode plasmas and in the edge of
the low-recycling plasmas, but clearly scales inversely with power at the center of the low-recycling plasmas. Transport of momentum by particle convection remains a small term in the momentum balance throughout the power scan, so the inferred $\chi_p(r)$ changes only slightly if a convective term $\Gamma_{p\text{-conv}} = \sum_i R_i m_i v_i \phi$ is subtracted from the total momentum flux $\Gamma_p$ (Fig. 5g,h). Thus, the transport of both momentum and heat by thermal ions appears to be reduced in the core of high-temperature, low-recycling plasmas.

The electron thermal diffusivity $\chi_e$ shows a substantially weaker variation with $P_b$ or temperature in the low-recycling power scan than is observed for $\chi_i$. Over a wide power range ($4 \rightarrow 22$ MW), $\chi_e$ increased by only $\sim 10 - 50\%$ in the region $r/a \geq 0.4$ and by a factor $\sim 2$ near the plasma center. If $\chi_e$ is corrected for convective energy transport using $q_{e\text{-conv}} = \frac{3}{2} \Gamma_e T_e$, we find that $\chi_e(r)$ is almost independent of heating power (within $\pm 25\%$) throughout the region $r/a \leq 0.65$. While both $\chi_i$ and $\chi_e$ are lower in the low-recycling plasmas than in high-recycling plasmas at comparable heating power, the improvement in energy confinement is stronger in the ion channel (cf. Table I). Stronger improvements in electron energy confinement have been observed previously in plasmas with more highly peaked density profiles obtained with balanced injection[2]).

### 2.2. Plasma current scan

To study the characteristics of local energy transport that yield roughly linear $I_p$ scaling of global $\tau_E$ in L-mode plasmas, several plasma current scans were performed at nominally constant density and beam power ($P_b = 6, 11, \text{or } 13$ MW) in high-recycling deuterium and helium plasmas. Scan E in Table I summarizes a comparison of deuterium $1.0$ MA and $2.0$ MA plasmas with balanced injection at $11.3$ MW. The central ion and electron temperatures, the thermal stored energy (obtained by integrating the measured density and temperature profiles), and the calculated beam stored energy increased significantly with current. Similar increases in $\tau_E^b$ were observed in all the $I_p$ scans; $55 - 60\%$ in deuterium plasmas and $60 - 70\%$ in helium plasmas for a factor-of-two change in $I_p$.

The ion and electron power flows were easily separable in these discharges over most of the confinement region. For example, in scan E less than $10\%$ of the power deposited inside the $r/a = 0.7$ flux surface was calculated to be exchanged between ions and electrons. Ion conduction remained the dominant power loss channel, typically responsible for $2/3$ of the total power flow at $r/a = 0.7$. Both $\chi_i$ and $\chi_e$ appeared to improve uniformly across the entire plasma cross section (Fig. 6) rather than
changing only in the plasma core (as might have been expected from reduction of the sawtooth region) or in the cooler edge plasma. Similarly, the momentum diffusivity $\chi_\phi$ was observed to decrease with $I_p$ uniformly across the plasma cross section (roughly as $\sim 1/I_p$) in helium L-mode plasmas with 6.5 MW of partially unbalanced neutral beam power, and in deuterium L-mode plasmas with 12 MW of co-injected beam power. The error bars shown in Fig. 6 represent the absolute uncertainty in $\chi_i$ and $\chi_e$ due to estimated systematic errors in the diagnostic measurements of $T_e(R)$, $T_i(R)$, $n_e(R)$, $Z_{\text{eff}}$, etc. Although the $1-\sigma$ error bars almost overlap (suggesting only marginal confidence in our conclusion that $\chi_i$ and $\chi_e$ decrease roughly equally with $I_p$), the uncertainties in the changes in $\chi_i$ and $\chi_e$ are smaller, because the systematic errors should be very similar for the two plasmas. The integrated confinement times $\tau_{\text{bi}}$ (32$\rightarrow$54 ms) and $\tau_{\text{be}}$ (103$\rightarrow$154 ms) at the plasma half-radius show comparable improvements for ions and electrons in deuterium plasmas.
Given the error bars shown in Fig. 6 it is also not immediately clear that transport is improved equally across the plasma cross section. The increase in $\tau_E^{th}$ with $I_p$ is roughly constant with radius, e.g. $\Delta \tau_E^{th} = 54\%$ at $r = a/2$ and $58\%$ at $r = a$, however $\tau_E^{th}(r)$ can be affected by changes in $\chi_i$ and $\chi_e$ both inside or outside the flux surface of radius $r$. To characterize the energy confinement within a particular plasma region, it is useful to divide its energy content into two components – a ‘pedestal’ term which represents the stored energy that would result even if $\chi_i$ and $\chi_e$ were infinite inside the region ($W_{ped} \equiv \frac{3}{2} V(n_e T_e + n_i T_i)|_r$), and a ‘core’ term that represents the additional stored energy due to nonzero temperature gradients within the region ($W_{core} \equiv \frac{3}{2} \int_0^r (n_e T_e + n_i T_i) dV - W_{ped}$). Note that in this representation, $W_{core}$ is entirely determined by transport coefficients within the plasma region, while $W_{ped}$ is determined entirely by transport coefficients outside the region (since they determine the energy density at the region’s surface). At the $r = a/2$ surface, $W_{core}$ and $W_{ped}$ increased by about the same amount (43% and 60%, respectively) as $I_p$ was increased from 1.0 to 2.0 MA in scan E, indicating that comparable improvements in energy confinement were realized in the regions $r < a/2$ and $r > a/2$.

Figure 3 plots the measured thermal diffusivities at $r/a = 0.5$ for power scans at 1.2 and 1.8 MA. Selecting temperature as the abscissa clarifies the very favorable scaling of local heat transport with $I_p$. Note that a straightforward comparison of $\chi_i$ and $\chi_e$ for different plasma currents at constant density and heating power would seriously underestimate their dependence on $I_p$, because higher temperatures would be realized at higher $I_p$, leading to increased transport which would partially compensate the reduction associated with increased current. Interestingly, recent L-mode current-ramp experiments [7] have shown that $I_p$ (or equivalently, the local $q$ or shear), appears not to be the parameter directly responsible for reduced transport in the region $r > a/2$, because large perturbations to the current density in that region had no effect on the local transport coefficients.

3. Scaling of $\tau_E$ with Plasma Size

Plasma size scans were performed in beam-heated, high-recycling L-mode discharges to determine the scaling of energy confinement with major radius and aspect ratio [10,11], an important design parameter for BPX and ITER and performance projections. The plasma size was varied by more than a factor of two ($a = 0.4 \rightarrow 0.9$ m, $R = 2.08 \rightarrow 3.20$ m) and the corresponding range in aspect ratio was $2.8 \rightarrow 8.0$. The plasma current was varied from 0.4 - 1.5 MA to maintain $q_{cy} \approx 3.1$ and the beam power was varied from 3 to 18 MW. An analysis of the energy confinement scaling
Fig. 7. Comparison of the $\tau_{E}^{\text{dia}}$ measured in the plasma size-scaling experiments with the ITER-P confinement regression (units: MA, $10^{19}$ m$^{-3}$, Tesla, m, MW, seconds). The data points represent the average of the kinetic and magnetic measurements, and the error bars represent their difference.

derived from magnetic measurements of $W_{\text{tot}}$ was presented in Ref. [10]. This section describes new analysis based on kinetic measurements of the stored energy, which is needed to identify the scaling of thermal energy and to corroborate the magnetic measurements. Relatively small but systematic differences between the kinetic and magnetic determinations of $W_{\text{tot}}$ lead to slightly stronger scaling of $\tau_{E}$ with major radius from the kinetic measurements.

The kinetic measurements were based on analysis of the measured density and temperature profiles using the SNAP code. In about half the discharges a full $T_i(R)$ profile was measured by charge-exchange recombination spectroscopy, and in the remainder only the central temperature was measured. For these, SNAP calculated $T_i(r)$ assuming $\chi_i = c\chi_e$ where $c$ was chosen to reproduce the central $T_i$ measurement. Power-law regressions were performed on a subset of the data constrained by the requirements $P_b \geq 4$ MW and $W_{\text{beam}}/W_{\text{tot}} \leq 0.3$. The results are:

\[
W_{\text{tot}}^{\text{dia}} (\text{magnetic}) = 0.045 I_p^{1.11} R^{1.62} a^{0.06} P_{\text{tot}}^{0.48}
\]

\[
W_{\text{tot}}^{\text{dia}} (\text{kinetic}) = 0.0436 I_p^{1.07} R^{1.78} a^{-0.17} P_{\text{tot}}^{0.37}
\]

\[
W_i + W_e (\text{kinetic}) = 0.0486 I_p^{1.11} R^{1.61} a^{-0.22} P_{\text{heat}}^{0.33}
\]
where $P_{\text{tot}} = P_b + P_{OH}$ and $P_{\text{heat}} = P_{be} + P_{bi} + P_{\text{therm}} + P_{OH}$. The typical statistically-determined uncertainty in the exponents for $I_p$, $R$, $a$, and $P$ are 0.09, 0.13, 0.12, and 0.06 respectively. The thermal stored energy was found to vary only as $P_{\text{heat}}^{0.33}$, in good agreement with the behavior over a wider power range in large ($a \geq 0.8$ m) plasmas [5]. Figure 7 compares the measured stored energy with the ITER-P scaling expression [9]. With a single adjustment of the multiplicative factor, which appears to be 15-20% too high, the ITER-P scaling regression provides a good fit to the data over a wide range of aspect ratio. Note, however, the parametric scaling of $\tau_E$ on individual parameters ($I_p$, $R$, $a$, $P_b$) deduced from these scans differs from the ITER-P scaling by more than the estimated error.

The observed dependence on plasma current is about the same for all of the regression expressions, and the dependence on minor radius is very weak. Due to the modest densities employed in these experiments, the average radius of beam deposition varied by less than 35%, so the energy confinement did not show a positive correlation with minor radius even when corrected for the “heating effectiveness” of the heating profile [6]. The weak scaling of $\tau_E$ with minor radius is contrary to the $\sim a^2$ scaling expected for a diffusive heat transport mechanism, and suggests that local heat transport is strongly influenced by toroidicity. This inference is supported by local transport analysis in these plasmas, which finds that the single-fluid thermal diffusivity $\chi_{\text{avg}}(r)$ is better characterized as a function of $(r/R)$ than $(r/a)$ [11].

4. Dimensionless Scaling

Another approach to the extrapolation of confinement from existing plasmas to future tokamaks of larger size or higher toroidal field involves the comparison of nondimensionally similar discharges ($a^{5/4}B/Z_{1}^{1/4} = \text{constant}$), as pointed out by Kadomtsev[12]. As observed recently[13], density scan data from Alcator C and PLT which are close to nondimensional similarity but which have quite different $\tau_E$ and $\bar{n}_e$ overlay nicely when plotted on Kadomtsev axes, $B\tau_E$ versus $\bar{n}_e/B^{1.6}$. Ohmic density scans have been performed at low toroidal field on TFTR with nondimensional parameters including $\varepsilon$ and $q_a$ chosen to match those in published PLT and Alcator C data. Plotting the dimensional results for $\tau_E$ vs. $\bar{n}_e$ gives orders-of-magnitude differences between the three machines (Fig. 8a). But they overlay fairly well when plotted on Kadomtsev axes (Fig. 8b-c) even though the tokamaks are of very different sizes (TFTR is twice the linear size of PLT and 4.5 times the size of Alcator-C) and magnetic fields (to maintain dimensionless similarity, TFTR operated at $B_T = 1.3$ and 1.5 T
Fig. 8. Comparison of energy confinement time versus density (a) in dimensional units for 3 tokamaks of very different sizes, (b) in Kadomtsev dimensionless variables for dimensionally similar TFTR and PLT plasmas, (c) in Kadomtsev dimensionless variables for dimensionally similar TFTR and Alcator-C plasmas.
compared to 3.2 T for PLT and at 10 T for Alcator-C). There is a residual discrepancy of a factor of two on the Kadomtsev plot, with TFTR performing twice as well as Alcator-C would predict at high densities. This apparent discrepancy in nondimensional scaling may be due to a poor match of some other nondimensional parameter (e.g. atomic physics such as neutral penetration affecting $a/L_n$) which influences confinement. For example, the density profiles in these ohmic TFTR plasmas were much more peaked than the Alcator-C plasmas.

In a separate experiment, density scans were performed in a pair of different-size plasmas at $q_a = 3.7$ and $\varepsilon = 5.2$ which were nondimensionally similar within TFTR ($R_1, R_2 = 2.05, 3.01$ m; $a_1, a_2 = 0.40, 0.59$ m; $B_1, B_2 = 5.8, 3.6$ T). Here, scale lengths and $T_i/T_e$ were reasonably well matched. The observed variations in $\tau_E$ were consistent with nondimensional scaling, although the range of dimensional quantities was much narrower in this TFTR-only scan. Additional scans in beam-heated TFTR discharges have been performed to assess the confinement scaling with gyroradius divided by minor radius [14] and with collisionality.

5. Fast Ion Confinement

Radial transport of fast ions (fusion products, RF-generated hot-ion tails, or beam ions) at the rate typically observed for thermal particles and heat would have serious consequences both for $\alpha$ heating losses in burning plasmas and for the interpretation of present confinement experiments. Previous neutral particle charge-exchange measurements during low-power tangential neutral beam injection on TFTR [15,16] established a rather low upper bound on the diffusivity of the passing fast-ion population (50-90 keV), $D_{\text{fast}} \lesssim 0.05$ m$^2$/sec, more than an order of magnitude less than that of the thermal plasma. This result has now been corroborated and extended to banana-trapped fast ions and to MeV ions in hotter, higher-$\beta$ plasmas using a variety of measurement techniques. These include measuring the 14 MeV neutron yield of the triton (fusion product) population, measuring the current and time dependence of fusion products lost from the plasma, and measuring the time decay of neutrons and charge-exchange flux following a short beam pulse.

Fast tritons are produced by the $d(d,p)t$ reaction at 1.0 MeV, with $\geq 80\%$ confined neoclassically for $I_p \geq 0.85$ MA. The tritons slow down primarily by electron drag through the peak in the cross-section for $t(d,n)\alpha$ reactions producing 14 MeV neutrons. The triton 'burnup' fraction (the
Fig. 9. Ratio of measured triton burnup (integrated yield of 14-MeV neutrons from copper foil activation divided by the integrated yield of 2.5-MeV neutrons from fission detectors) to the simulation of the classically expected value, plotted versus the time-averaged slowing down time for each of 47 discharges. The curve is the burnup expected if an anomalous loss time of 0.5 seconds were affecting the burnup.

The ratio of 14 MeV neutron production to the triton birth rate was determined from the integrated yield of 14-MeV neutrons measured by copper foil activation[17] was compared to the classical expectation from a fully time-dependent simulation[18] (Fig. 9). The ratio of measurement-to-classical calculation was less than one, and tended to decrease with longer triton slowing-down times. It was not observed to correlate with plasma current, $\beta$, $q_{\text{out}}$, or even fishbone MHD activity. The triton burnup in low-temperature L-mode plasmas with short thermalization times may not be inconsistent with negligible radial transport, due to uncertainties in the absolute calibration of the neutron measurements ($\pm 27\%$ on the burnup[19]) and the effects of hot ions and finite orbit radii[18], which might raise the burnup ratio by $\leq 25\%$. However, the burnup was significantly less than unity in high-temperature supershot plasmas with correspondingly longer slowing down times, consistent with diffusion of order $D_{\text{fast}} \sim 0.1 \text{ m}^2/\text{sec}$ (a result similar to that reported for JET[20]).

Radial diffusion of fast ions at this rate in the burning plasmas envisioned
Fig. 10. Escaping MeV ion flux (1 MeV tritons plus 3 MeV protons) versus plasma current, compared to model calculations. The models are normalized to the data at $I_p = 0.8$ MA. For $D_{\text{fast}} = 0$ the curve represents the calculated first-orbit loss, while for $D_{\text{fast}} = \infty$ the curve represents the additional loss of all interior counter-passing ions.

for BPX or ITER would not result in serious alpha heating losses, due to the shorter thermalization times.

The loss of MeV DD fusion products (tritons and protons) has also been measured using a set of scintillation detectors located near the wall $45^\circ - 90^\circ$ below the outer midplane[21]. Over a current range up to 1.6 MA and beam power up to 30 MW, the observed MeV ion loss was usually consistent with the expected neoclassical first-orbit loss. Clear exceptions to classical fast-ion confinement occurred during strong coherent MHD, when the time-averaged loss increased by up to $\times 3$ at 1.6 MA, and during sawtooth crashes, when large $\sim 100$ $\mu$sec bursts of MeV ions were occasionally observed. In absence of strong MHD, the loss fraction measured at 90° decreased by about a factor of 5 between 0.8 and 1.6 MA, approximately as expected from the reduced banana width at high current. Figure 10 shows the escaping MeV ion flux (1 MeV tritons plus 3 MeV protons) as
a function of plasma current (from Ref. [21]), compared to model calculations using various spatially-uniform diffusion coefficients $D_{\text{fast}}$. The data generally lie in the region $D_{\text{fast}} \leq 0.1 \text{ m}^2/\text{sec}$, as expected theoretically from the orbit-averaging effect [22-25]. However, the model calculations do not yet account for the energy dependence of the particle detectors and thus neglect the effect of the energy spectrum of the diffusing MeV ion population.

Finally, short pulses (~ 20 ms) of intense deuterium beams (90 keV; 100-250 kJ) were injected into ohmic deuterium plasmas to study radial transport of both passing and trapped beam ions. The decay rate of the measured global neutron emission compares well with classical slowing-down predictions which ignore radial transport of the fast ions (Fig. 11), similar to previous experiments on DIII-D [26,27]. The theoretical predictions of the neutron and charge-exchange measurements were made by a Fokker-Planck package [28] in the TRANSP transport analysis code, which includes collisional drag, energy diffusion, and pitch angle scattering, and an approximate model for the neutral density profile which assumes poloidal symmetry. A fast-ion radial diffusion coefficient $D_{\text{fast}} > 0.1 \text{ m}^2/\text{sec}$ can be ruled out because it would have carried the fast ions from the hot plasma center out to colder regions yielding a much more rapid decay rate of total DD neutron emission than was measured. This result
Fig. 12. Time history of the charge-exchange flux at $E = 30$ keV following a short heating-beam pulse, as measured by vertically-viewing neutral particle diagnostics at $R = 2.44$ m compared to the flux calculated for various assumptions of fast-ion radial diffusivity.

applies principally to passing ions because the beam ions were born on passing orbits over most of the plasma cross-section, and the pitch-angle scattering times were long ($\sim 450$ ms at the plasma center). The small diffusion rate was confirmed by measurements from an array of collimated neutron detectors [29] which provided time evolution of the radial profile of DD neutron emission. Charge-exchange measurements of the trapped beam ion population along two vertically-viewing sightlines implied a similarly small upper bound on $D^f$. The time decay of the charge-exchange flux near the plasma center ($R = 2.44$ m) was in reasonable agreement with the $D^f = 0$ simulation (Fig. 12), and substantially slower than the $D = 0.1$ m$^2$/s simulation. The measured flux at $R = 2.97$ m (tangency minor radius $= 0.6a$) persisted longer than the $D^f = 0$ prediction and required a $D^f \sim 0.1$ m$^{-2}$/s to move a fraction of the central beam ions out to this radius.

6. Summary

A variety of 'single-parameter' scans have been carried out to measure equilibrium heat and momentum transport over a wide range of parameters in ohmic and beam-heated TFTR plasmas. A strong temperature (or heating power) dependence of $\chi_i$ and $\chi_e$ has been identified in high-recycling plasmas which affects the entire confinement region, thus giving rise to the unfavorable power scaling of the L-mode. Most of the radial heat flow is carried by ions (conduction and convection) rather than by electrons. The
effect of increased plasma current in L-mode plasmas is to globally reduce both $\chi_i$ and $\chi_e$. In view of the strong scaling of $\chi$ with temperature, the effect of current appears to be quite strong. Beam power scans in low- and high-recycling plasmas with unbalanced beam torque have clarified the differences in local transport obtained in L-mode and supershot plasmas; transport in the edge region exhibits an unfavorable temperature (or heating power) dependence in both regimes, but in low-recycling plasmas the ion thermal and momentum transport reverse this dependence near the plasma center and become smaller with increasing temperature.

Detailed aspect ratio and major radius scans in L-mode have demonstrated a strong dependence of $\tau_E$ on major radius and a weak dependence on minor radius ($\tau_E \propto R^{-1.6}a^{-0}$), suggesting a strong role for toroidicity in the mechanism underlying anomalous heat transport. Initial studies of Kadomtsev dimensionless scaling suggest that projections from Alcator C and PLT would underpredict the TFTR ohmic $\tau_E$ by factors of 1.5 - 2.5. This result may be considered evidence that plasma physics controls transport in tokamaks – since the tokamaks collectively span a range of 200 in density and 30 in $\tau_E$ – but the residual discrepancy suggests that more accurate projections will require greater care to control other dimensionless parameters such as the ratio of gradient scale length to system size.

Several experiments have demonstrated surprisingly low rates of radial diffusion ($D_{\text{fast}} < \chi_{\text{avg}}/6$) for both passing and trapped energetic ions, suggesting that the turbulence responsible for driving thermal transport has a far smaller effect on energetic ions. This validates a key assumption in $\alpha$-particle confinement projections and in transport analysis of thermal energy transport (it does not address transport driven by collective $\alpha$-driven instabilities nor the ripple transport) and may rule out some classes of long-wavelength turbulence models[16]. Finally, it is natural to consider at what energy the ions become subject to anomalous radial transport. The observation that $\chi_i$ reverses its adverse temperature dependence obtained in L-mode plasmas when the central ion temperature increases beyond $\sim 10 \text{ keV}$ in supershot plasmas (correspondingly, $T_i/T_e \geq 1.8$) [30] invites speculation that the improved supershot energy confinement arises from an increasingly large part of the thermal ion population becoming exempt from turbulence-driven transport. This behavior would be consistent with expected orbit-averaging effects[22–24] on ion transport if the turbulence scale length were set by $T_e$ rather than $T_i$.

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References


DISCUSSION

R.J. TAYLOR: At high neutral beam power the ‘blip’ experiments may not yield the same neoclassical diffusion and slow-down. How can you extrapolate these results?

S.D. SCOTT: Yes, these blip experiments have established only that fast ion radial diffusion is small compared to thermal transport in a near Ohmic plasma. They do not necessarily imply similar behaviour in strongly heated, higher β plasmas. In principle, similar beam blips could be added to a plasma heated with substantial amounts of quasi-stationary beam power, but it would be difficult to distinguish the time decay of neutron emission unless one employed $D^0$ blip beams and $H^0$ quasi-stationary beams. The triton burnup and lost triton measurements performed over a range of beam power, including high power, also indicate small values of the fast ion radial diffusivity, $D_{fast} \sim 0.1 \text{ m}^2/\text{s}$. 
HIGH PERFORMANCE H MODES IN JET

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Abstract

HIGH PERFORMANCE H MODES IN JET.

In JET the scientific properties and technical basis of good confinement regimes have been evaluated in the light of the potential extrapolation of such regimes to reactor requirements. The main experimental H mode results are discussed, highlighting global confinement scaling, low q regimes, the role of the target plate material, the density limit, sawtooth suppression and the hot ion mode.

I. INTRODUCTION

The H mode in JET has been demonstrated with single and double null X point configurations, which in general are marginally limiting at the X point target tiles. H modes have been achieved with NB heating, ICRF heating and combined NB and ICRF heating. The power threshold for the H mode and the global energy confinement time do not depend on the heating method. In the ELM free H mode there is an improvement of about a factor of 2 in the global confinement time compared with JET limiter L modes, up to a total additional power of 25 MW.

The development of the JET H mode, towards steady state conditions, depends on wall conditioning and on the material of the divertor target plates, which determine the amount and type of impurities released and affect the time evolution of plasma density. The substantial reduction of $Z_{eff} (\approx 2)$ and the improvement of $n_D/n_e$ to 0.8–0.9, produced by routine beryllium gettering of the graphite tiles, were probably mainly due to the nearly complete removal of oxygen and oxygen generated carbon sputtering. The reduction of nickel was mainly due to the beryllium gettering of the ICRF antenna screens. As a consequence of lower radiation losses and improved density control, longer ELM free H mode phases have been achieved ($\approx 5.4$ s).

The transport of impurities in the JET H mode is characterized by a balance between neoclassical effects and anomalous transport leading to a buildup of impurities in the plasma ($\tau_{imp}/\tau_E > 1$) [1]. Further local analysis of energy transport in high power H modes confirms the reduction in thermal conductivity (energy flux/grad T) across the whole plasma cross-section reported earlier [2].

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1 See Appendix of IAEA-CN-53/A-I-2, this volume.
In the JET H mode the density limit corresponds to the radiative collapse which ends the H phase.

With additional power well above the threshold, the H mode can occur simultaneously with other plasma regimes such as the monster sawtooth or the hot ion mode.

In the hot ion H mode with NB(D) injection at moderate plasma densities \((n) = (1-4) \times 10^{19} \text{ m}^{-3}\) \(T_i\) is 2 to 3 times larger than \(T_e (T_i \gtrsim 22 \text{ keV}, Q_{DD} = 2.5 \times 10^{-3})\). The values of \(Q_{DT}\) for D-T simulated versions of the same discharges are above 0.5 for times of the order of one energy confinement time.

2. OPERATIONAL REGIMES

The H mode in JET has been demonstrated in single null \((I_p < 5 \text{ MA})\) and double null \((I_p < 4.5 \text{ MA})\) configurations, generally with the plasma limiting on the X point target tiles, but also with the plasma in contact with the inner wall. In single null configurations the H mode has been achieved with NB at 80 keV and 140 keV in both deuterium and hydrogen target plasmas. In double null configurations, where it has been possible to obtain good RF coupling, the H mode has been achieved with ICRF alone in dipole with hydrogen minority, and in combination with NB(D) in deuterium plasmas. The JET H mode is characterized by a transition to an ELM free period lasting several seconds. The duration of the H mode is considerably shorter for power in excess of 10 MW owing to a strong carbon influx from overheated graphite dump tiles [3].

The threshold power for the H mode was similar for double and single null configurations, with ICRF and NB heating. The H mode power threshold was lower with a well conditioned vessel. The power threshold, while scaling approximately linearly with the applied toroidal field as reported earlier [4], does not show a clear dependence on plasma current or plasma density. Scans of the plasma radial position show a minimum power threshold if the gap between the plasma and the limiter or between the plasma and the inner wall is greater than 5-8 cm. For shorter distances the threshold power increases. H mode was achieved even with the plasma in contact with the inner wall protection tiles, which required a threshold power of 10 MW at \(B_T = 2.2 \text{ T}\). Plasma recycling and power load were distributed on the inner wall. We think that this regime is similar to the inner wall H mode achieved in DIII-D [5].

Figure 1 shows the plot of additional heating power versus the gap between the plasma and the inner wall. L marks the discharges which had not undergone an H transition, while H marks the H modes. The H mode power threshold is also a function of the position of the X point as determined from magnetic diagnostics [6]. The threshold for H mode has only a weak dependence on the location of the X point within 10 cm outside the surface of the dump plate, with the plasma in a marginal limiter configuration, or 10 cm inside, with the plasma in an X point configuration.
3. GLOBAL CONFINEMENT

The main properties of the global confinement time of the JET H mode have been extended to higher additional power. The general trend of the global energy confinement time as a function of the total loss power is shown in Fig. 2. The two main trends of increase with plasma current and degradation with heating power continue, at high powers, the trend reported earlier [4]. The scaling of the global confinement time is similar in single null and double null configurations. From the analysis of the data it appears that the global confinement time has only a weak dependence on the toroidal field (as $B_t^{0.3 \pm 0.07}$). In operation at low values of $q_{90,95}$ it appeared that there was not an abrupt deterioration of confinement in the approach to $q_{90,95}$ of 2, as shown in Fig. 3 [7]. The scaling of the global confinement time observed in the JET H mode is similar to that observed in other tokamaks. A size scaling has been derived by combining data from JET and DIII-D [8]. A large database with data from JET, DIII-D, ASDEX, JFT-2M, PDX and PBX has been created; the scaling obtained by the analysis of the data of these tokamaks is $\tau_E = 0.07A^{0.51}P^{-0.48}B_t^{1.72}k^{0.48}(R/a)^{-0.13}$, when the toroidal field and density dependence are suppressed [9]. As an example a plot of the JET data versus the combined H mode scaling is shown in Fig. 4. A detailed analysis of the local transport of JET H mode is also presented in Ref. [10].
FIG. 2. Global energy confinement time versus total loss power for the 3 MA and 4 MA discharges with \((dW/dt)/P < 0.3\). Solid symbols refer to pellet fuelled discharges.

FIG. 3. Global energy confinement time as a function of the safety factor at 95% of the flux surfaces. The data refer to a set of 3 MA discharges in deuterium with NB heating in a narrow power range around 9 MW. The range of the toroidal field is between 1.2 and 2.4 T.
4. APPROACH TO STEADY STATE CONDITIONS

The ELM free H mode is a transient effect which is terminated either by radiative collapse, for additional input power below 10 MW [11], or by carbon influx, for power in excess of 10 MW.

4.1. Carbon influx

A strong carbon influx, which terminates the H mode, enters the plasma when the surface temperature of the dump plates exceeds 2500°C. The onset of the carbon influx can be delayed by radial sweeping of the X point position and/or by a strong gas puff. With a strong puff, long ELM free H mode phases have been obtained exceeding 5 s, as shown in Fig. 5(a, b). The main effect of a strong gas puff is a reduction of the surface temperature of the X point tiles (Fig. 5(c, d)) in the heated zone where the power is deposited by fast ions in drift orbits [3, 6]. The reduction in temperature is probably caused by reduced ion temperature and/or increased ion
FIG. 5. Comparison of long pulse H mode (a) with added gas and (b) the same plasma and similar NB power with no added gas. The early carbon influx in the shot without gas can be clearly seen. (c) and (d) show the measured maximum surface temperatures of the dump plate carbon tiles for the discharges in (a) and (b), respectively.
collisionality. The reduction in temperature leads to a significant drop in carbon sublimation and radiation enhanced sputtering. The gas puff causes increased divertor radiation losses with a general reduction in conduction losses.

4.2. Radiative collapse

The impurity confinement time, as determined by laser blow-off measurements and impurity transport computer simulation [12], is much longer than the energy confinement time, with typical values up to 4 s [1, 12]. As an example the time evolution of the concentration of Ni(XXV) and Ni(XXVI) obtained by laser blow-off of a nickel coated target is shown in Fig. 6. The results of transport analysis show that there is an outward flow of impurities [1, 12]. The plasma behaves like a 'leaky' integrator. For intrinsic impurities, in the H mode, the outward flow is an order of magnitude smaller than the inward flow, which is generated by the plasma interaction with the dump plates. Consequently, during the ELM free H mode the concentration of the intrinsic impurities increases continuously as a function of time, until the total radiated power becomes excessive [11]. However, if the inflow of intrinsic impurities could be reduced by at least a factor of 5, for example by the use of a pump divertor, the simulations show that a balance could be reached after a few impurity confinement times at levels of radiation compatible with the input power.

FIG. 6. Time evolution of the ratio of the nickel intensities during the H mode. The nickel is injected by laser blow-off at $t = 11$ s. From top to bottom: volume average electron temperature, traces of density normalized Ni(XXVI) and Ni(XXV), volume average plasma density.
5. EFFECTS OF DUMP PLATE MATERIAL ON THE H MODE

The comparative properties of graphite and beryllium gettered graphite have been studied.

With graphite tiles oxygen and carbon were the dominant impurities; radiative losses by oxygen were dominant at high densities. Among other impurities nickel, mainly generated by the ICRF antenna screens, accounted for a 10% fraction of the radiated power. With beryllium gettered graphite tiles the concentration of oxygen was reduced to a negligible amount. The reduction of carbon concentration can be explained as being due to the absence of oxygen sputtering generated carbon impurities, as in the case of a beryllium gettered graphite limiter [13].

Figure 7 shows the trend of $Z_{\text{eff}}$, as measured by visible bremsstrahlung, versus the total loss power (total input power minus the time derivative of the plasma energy). An average reduction of 1–2 units at all levels of power is seen. Consequently, charge exchange spectroscopy measurements show that with beryllium gettered graphite dump plates the ratio of electron to deuteron density was 0.8–0.9, while without beryllium gettering it was 0.5–0.6.

6. DENSITY LIMITS

During the H mode generally the plasma density increases continuously until a radiative collapse of the H mode occurs [11], precipitating an H to L transition and sometimes a full plasma disruption. For the H mode the density limit coincides with the high density prior to the H to L transition. The limit on the plasma density is caused by the fact that the power which is radiated by the bulk of the plasma approaches the input power. It is therefore natural to expect the density limit to scale with the square root of the total input power with a scaling similar to that observed in recent JET limiter discharges [14, 15]. The values of the maximum volume
FIG. 8. Volume average electron density at the end of 3 MA H mode pulses versus total input power. The crosses refer to pulses with $Z_{eff}$ less than 3.0, while the squares refer to pulses with $Z_{eff}$ larger than 3.0. The line is $n_e \text{(10}^{19} \text{ m}^{-3}) = 2.12 \times P^{0.5} \text{(MW)}$.

FIG. 9. Hugill plot for Ohmic and additionally heated X point discharges. Diamonds represent ohmically heated plasmas, crosses neutral beam heated plasmas and asterisks combined ICRF and NB heated plasmas.
average electron density prior to the H to L mode transition have been plotted versus the total input power in Fig. 8 for a series of 3 MA H mode discharges. The experimental points with $Z_{\text{eff}}$ smaller than 3 (crosses) have a higher density than those with $Z_{\text{eff}}$ larger than 3. For reference a Hugill plot of the non-disrupting H mode pulses is shown in Fig. 9 for discharges with beryllium gettering. The line $nRq/B_o = 20 \times 10^{19} \text{m}^{-2} \text{T}^{-1}$ encompasses the values for the discharges with moderate additional heating (<10 MW) and fuelled by gas puffing. The lower limit $nRq/B_o = 12 \times 10^{19} \text{m}^{-2} \text{T}^{-1}$ refers to Ohmic X point discharges. Preliminary results have been achieved with central pellet fuelling producing more peaked density profiles. The values of the pellet fuelled discharges are not shown in Fig. 9.

7. SAWTOOTH STABILIZATION

Sawtooth suppression has been observed in H mode discharges with NB and ICRF heating. The time evolution of a series of discharges showing sawtooth suppression during H mode is presented in Fig. 10.

With NB heating in excess of 8 MW, at least twice the power threshold for H mode transition, injected into a relatively low density Ohmic deuterium target, the H mode is accompanied by a period of sawtooth stabilization with a duration of 0.6–0.8 s. In this phase a modest enhancement (10–15%) of central ion and electron temperatures is observed.

![Fig. 10. Time evolution of central electron temperature in H modes with suppressed sawtooth: pulse No. 14834, NB heating; No. 19796, NB and ICRF heating; Nos 19995 and 20231, ICRF heating.](image-url)

FIG. 12. Time evolution of ion and electron temperatures for a sawtooth suppressed H mode with combined ICRF and NB heating.
With ICRF heating during the H mode [16], sawtooth suppression occurs routinely with ICRF input powers in excess of 7 MW. The maximum duration of the monster sawtooth has been 2.5 s. The start and end of sawtooth suppression were not correlated with the H mode phase, but sometimes the monster crash caused an H to L transition. The temperature peaking factor obtained in sawtooth suppressed H modes is enhanced by approximately 50%. A series of electron temperature profile shapes with sawtooth suppressed H modes is shown in Fig. 11. Here the peaking factor ranges between 3 and 4 (with electron pressure peaking factors between 4 and 5) with values of cylindrical \( q = 3.2 \) and average densities in the range \( (2-4) \times 10^{19} \text{ m}^{-3} \). It should be noted that in this case the shapes of the electron temperature profiles are similar to those obtained in the case of limiter monsters [17] and that the value of the edge plasma temperature is not very high.

With combined ICRF and NB heating, sawtooth suppression in the H mode has also been achieved, as shown by one of the traces in Fig. 12. In this pulse, \( P_{RF} = 2 \text{ MW}, P_{NB} = 6 \text{ MW} \) and \( (n) = 2.5 \times 10^{19} \text{ m}^{-3} \). Polarimetric analysis of the safety factor radial profile indicates that the central value of \( q \) is driven below unity in a way similar to other monster sawteeth, while estimates of the content of fast particles confirm the agreement with the theoretical expectations of sawtooth stabilization.
FIG. 14. Time evolution of high fusion yield pulse No. 20981. From top to bottom: central ion temperature $T_i$ and total neutron yield $Y_n$; plasma diamagnetic energy $W_{dia}$ and volume average electron density $\langle n_e \rangle$; $D_{\alpha}$ intensity near the X point; neutral beam power $P_{NB}$ and radiated loss power $P_{rad}$. The carbon influx at 11.5 s is followed by loss of the H mode.

FIG. 15. TRANSP code simulation of a D-T version of pulse No. 20981 obtained by using the same set of experimental measurements which are a good representation of the D-D case. The only changes are the species mix and the injection energy. This figure shows total fusion power and thermal and beam thermal contributions, assuming 15 MW of D (at 140 keV) injection on a target tritium plasma. Only D-T reactions have been considered.
8. HOT ION H MODE

In X point configurations with NBI in a low density Ohmic target it has been possible to produce simultaneously a hot ion plasma and an H mode transition. This regime is characterized by very high ion temperatures in excess of 20 keV, while the electron temperature was 8–10 keV. For large values of NB power and low densities the ion temperature profile is very peaked. The ratio of central to volume average ion temperature reaches values of 4–5.

The D–T performance of these pulses is examined by considering the standard \( n_D^{3/2} T_i \) diagram, shown in Fig. 13. Here the Q curves are for parabolic profiles of density and temperature raised to the power of 1/2 and 3/2, respectively. The Q curves are rather insensitive to the form of the profiles. The time evolution of plasma parameters for pulse No. 20981 is shown in Fig. 14.

The time behaviour of the pulse was simulated by the 1\( \frac{1}{2} \)-D TRANSP code for the actual conditions of NBI and background plasma. A good check on the consistency of these data is that of the predicted and measured neutron yield. The D–T simulations are completed by rerunning the code with the same measured profiles and replacing the background deuterium plasma with a tritium plasma or with a 50–50 D–T mixture. The results are shown in Fig. 15. Here the time evolution of the fusion power is shown in its components for the case of deuterium injection in a tritium plasma. The NB power was 17 MW, the time derivative of plasma stored energy was 8 MW. With Q defined as \( Q = \frac{P_{\text{th-th}}}{P - \frac{dW}{dt}} + \frac{P_{b-th} + P_{b-b}}{P} \), the peak value corresponds to \( Q = 0.77 \) [18].

9. CONCLUSIONS

(a) The experiments at JET have shown that the ELM free H mode is a transient plasma regime. The analysis of the results of impurity injection experiments and of the spectroscopical data suggests that if one could control the plasma density and drastically reduce the influx of impurities the H mode could reach a steady state transport equilibrium.

(b) The H mode is achieved in a configuration with a magnetic separatrix, which can be marginally limiting on the X point dump plates or on the inner wall protection plates. A parameter related to the shear at the plasma edge for this marginally limiting configuration could be the ratio between \( q_\phi \) and \( q_{\text{cyl}} \). H mode has been achieved with \( q_\phi \) larger than or equal to 7–8 while \( q_{\text{cyl}} \) was 2–3.

(c) In the H mode the global energy confinement time shows an improvement of approximately a factor of 2 over the limiter L mode independently of the heating method and in the range of additional power up to 25 MW.

(d) The radiation collapse, which terminates the H mode, determines the values of the density limit, which scales with the square root of the input power.
(e) An improvement in plasma purity and a reduction of fuel dilution have been achieved with beryllium gettering of graphite tiles.

(f) In a low density Ohmic target and largely with ICRF heating it has been possible to achieve transient stabilization of sawteeth during the H mode.

(g) With NBI in low density deuterium plasmas a hot ion plasma has been created during the H mode. In these discharges the highest plasma thermonuclear reactivities have been obtained.

REFERENCES


DISCUSSION

Y. KAMADA: You said that the central $Z_{eff}$ in your pellet experiment was close to unity. How long does the clean condition last?
A. TANGA: The clean condition lasts for a time of the order of one second.

G. FUSSMANN: What is the reason for the radiation collapse if the only significant impurity influx is due to carbon?

A. TANGA: Yes, most of the bulk radiation is produced in the outer 20% of the minor radius where the plasma density is also high.
RECENT RESULTS OF H-MODE STUDIES ON ASDEX

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Abstract

RECENT RESULTS OF H-MODE STUDIES ON ASDEX.

In a comparative study of various confinement regimes the H-mode demonstrated the best performance. Confinement enhancement factors (above ITER 89-P L-mode scaling) in the range of $1.6 < f_H < 2.8$ have been achieved with values depending on the divertor configuration, the wall condition, ELM behaviour and the plasma ion species. Long-pulse H-phases, with ELMs, of up to 3.5 s with constant confinement time, recycling and impurity characteristics are achieved. H*-mode operation is possible without a loss of current scaling at $q_a$ values as low as 2.2. The $\beta$-limit is the same with and without ELMs. Murakami parameters are similar in H- and L-modes ($M < 10 \times 10^{17} \text{m}^{-2} \cdot \text{T}^{-1}$). A high perpendicular edge rotation velocity $v_\perp$ is observed in the H-phase. $v_\perp$ develops gradually during the H-phase and increases with NJ power. Though edge turbulence drops initially after the H-transition, edge fluctuations appear again in quiescent phases which reduce the edge electron temperature. An ELM precursor has been identified in the frequency range of $100 < f < 200 \text{kHz}$. ELM induced losses set in during an turbulent phase following the precursor. ELM characteristics are the same at low or high edge pressure, for singular or 'grassy' ELMs.

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1 INTRODUCTION

On ASDEX we have developed a sequence of confinement regimes in order to explore their inherent transport physics and also, to assess their performance:

(1) Degraded confinement regimes with medium peaked density profiles \{ Saturated Ohmic Confinement [SOC], L-mode (co-NI) \},
(2) Improved confinement regimes with peaked density profiles \{ Improved Ohmic Confinement [IOC], pellet refuelled discharges [PD], L-mode counter neutral injection [ctr-NI] \} [1],
(3) H-mode with flat density profiles [2].

SOC and IOC are obtained under ohmic conditions (OH), L-mode co and ctr, and H-mode with NI, and PD both under OH and NI conditions. All improved regimes are distinguished by a substantial rise in the confinement time (≥ 50%). Transport analysis has benefitted from the possibility to compare regimes: SOC and IOC, L-mode (co-NI) and either H-mode or ctr-NI, gas-puff and pellet refuelling. The peaking of the density profile preferentially reduces the ion transport [1]; the edge transport barrier of the H-mode, first described in [3], leads to a substantial reduction of the electron transport. Momentum balance analysis of plasma rotation played a decisive role in the identification of the loss channel which causes the difference between an improved and a degraded regime [1].

The peaked density regimes have the following disadvantages: The IOC regime is limited to low power fluxes across the separatrix [5]. Higher power fluxes lead to the SOC or the L-regime. Ctr-NI plasmas with improved confinement never reached steady state and always disrupted. PD show a strong degradation of confinement with auxiliary heating power toward the usual L-mode level [1]. Both with ctr-NI [6] and in the quiescent phases of PD impurities accumulate with approximately neo-classical rates.

This comparative study showed that the best performance, the largest operational range, and the best prospects for satisfying the various reactor requirements can be expected from the H-mode. In the following, we will discuss the progress we made in H-mode studies.

2 CONSTRAINTS ON THE DIVERTOR CONFIGURATION FOR GOOD H-MODES

The major requirements for access to the H-mode at low power - divertor configuration and ion-drift to the X-point [7] - are well documented. The importance of recycling [8] and impurity control will be addressed here.
Tab. 1. Exponents of the regression analysis in the form \( \tau_E (ms) = \text{const} \times I_p^{a} \times B_t^{b} \ldots \), H-mode power threshold, number of analyzed shots for the various conditions. ELMy cases: \( W = 0 \), H* case: \( 0.25 \leq W \leq 0.33 \). * in \( A_i \) column: only \( H^0\rightarrow D^+ \) was studied.

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<th>( \text{const} )</th>
<th>( I_p ) (kA)</th>
<th>( B_t ) (T)</th>
<th>( \bar{n}_e ) ((10^{3} \text{ cm}^{-3}) )</th>
<th>( P_{\text{tot}} ) (MW)</th>
<th>( A_i )</th>
<th># data</th>
<th>( P_{\text{thr}} ) (MW)</th>
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<tr>
<td>DV-I st.st.</td>
<td>0.82 ±0.06</td>
<td>0.71 ±0.06</td>
<td>0.18 ±0.05</td>
<td>-0.29 ±0.05</td>
<td>*</td>
<td>150</td>
<td>1.1</td>
<td></td>
</tr>
<tr>
<td>DV-I open H</td>
<td>2.04 ±0.1</td>
<td>0.45 ±0.11</td>
<td>0.2 ±0.09</td>
<td>-0.6 ±0.14</td>
<td>0.46 ±0.21</td>
<td>75</td>
<td>1.8</td>
<td></td>
</tr>
<tr>
<td>DV-II closed H bor</td>
<td>0.67 ±0.05</td>
<td>0.68 ±0.05</td>
<td>0.36 ±0.07</td>
<td>-0.49 ±0.09</td>
<td>0.27 ±0.08</td>
<td>217</td>
<td>1.3</td>
<td></td>
</tr>
<tr>
<td>DV-II closed H* bor</td>
<td>0.16 ±0.05</td>
<td>1.08 ±0.14</td>
<td>-0.68 ±0.14</td>
<td>0.59 ±0.07</td>
<td>60 ±0.07</td>
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Recycling aspects were studied for two divertor chamber configurations (divertor chambers DVI / DVII with large / small volumes [9]) and for various stages of "leakiness" of DVII to the main chamber (bypass conductance larger or smaller than the divertor neck conductance). The wall material was either the original stainless-steel wall or carbon or boron layers deposited on it [10]. Table 1 gives the results of \( \tau_E \) regression analysis (from \( \beta_{\text{dia}} \)) and the H-mode power threshold for the various conditions. The divertor configuration affects the parameter scaling of \( \tau_E \). Other studies indicated that also the quiescent H-phase has lower confinement in the "leaky" divertor and that carbonization leads to lower confinement quality.

The distribution of the confinement enhancement factor \( f_H \) of \( \tau_E \) with respect to the ITER89-P L-mode scaling is plotted in Fig. 1 for the various conditions considered. The results demonstrate the necessity of having a closed divertor and the advantage of low-Z metallic wall coating. The best confinement is achieved under quiescent conditions with \( D^0 \rightarrow D^+ \) injection, boronization and DVII - closed. \( f_H \) is = 2.8.
Fig. 1. Distribution of \( f_H \), the ratio of the measured \( \tau_E \) with that from the ITER 89-P L-mode scaling. 1) DV II " leaky ", 2) DV I stainless-steel walls, closed, 3) DV II \( H^* \), boronized walls, \( D^0 \rightarrow D^+ \), 4) DV II ELMy, boronized walls, analyzed at \( \beta_{\text{MAX}} \). 5) DV II ELMy, boronized walls, long pulses ( \( \Delta t > 20 \tau_E \)).

3 LONG QUASI-STEADY-STATE H-MODES

Attempts to obtain the quiescent H-mode ( \( H^* \) ) under steady state conditions on ASDEX were unsuccessful. In addition to the different divertor configurations and wall conditions we tried to develop it under stationary conditions with pumping or non-pumping target plate material ( Ti or Cu ) and with Ti gettering within the divertor chambers. The longest quiescent H-phases were obtained with boronized walls ( duration \( \leq 220 \text{ ms} \) ).

Quasi-steady state H-discharges with ELMs require an optimized and constant ELM frequency as is obtained after boronization [11]. The optimum ELM frequency ( as a trade-off between impurity level and confinement degradation ) is obtained by careful radial and vertical positioning of the plasma column. A movement of 0.5 cm in the vertical direction and 1 cm in the horizontal direction decides between H and \( H^* \). A long quasi-steady state H-phase of 3.5 s ( \( \approx 80 \) confinement times ) is depicted in Figure 2. \( \tau_E = 44 \text{ ms} \); \( \tau_E/I_P = 0.16 \text{ s} / \text{ MA} \); \( f_H = 1.9 \); \( \beta / \beta_{\text{Troyon}} = 0.7 \); ELM period = 5.5 ms. Experimental traces are selected which demonstrate steady state conditions for confinement, recycling, and impurities. The longest H-phase ( with beam faults, however ) was 4.3 s. Figure 1 also shows \( f_H \) obtained under ELMy conditions with boronized walls. Two cases are plotted: (1) The usual case with analysis at \( \beta_{\text{MAX}} \); (2) a subset of long quasi-steady state H - phases ( \( \geq 20 \times \tau_E \) ) were selected. In order to assess the confinement quality during the pulse, data at 3 moments
are taken: at the beginning, in the middle and at the end of the H-phase: \( f_H = 1.9 \). To achieve steady state conditions, a degradation of \( f_H \) has to be accepted. A specific study showed that a degradation of 20 % can provide quasi-steady state.

Fig. 2. Example of a quasi-steady state H-discharge. Symbols are standard; \( Z_{\text{eff}} \) is from Bremsstrahlung; \( P_{\text{tot}} = 1.65 \, \text{MW} \).
4 EXPLORATION OF THE H-MODE IN CORNERS OF THE OPERATIONAL WINDOW

4.1. Low-q operation

With boronized walls ELM-free and ELMy H-modes have been studied over a large operational domain: $220 \text{kA} \leq I_p \leq 460 \text{kA}$, $1.25 \text{T} \leq B_t \leq 2.8 \text{T}$, $2.2 \leq q_a \leq 4.5 \ [12]$. The best fully developed ELM-free H-phases are obtained at values of $q_a \approx 3$. For $q_a$ higher than 3.3 the ELMs are frequent. For $q_a < 2.4$ the ELM-free phase is interrupted by a hard disruption and the plasma energy cannot reach the values obtained for $q_a > 2.5$. The power threshold to obtain ELM-free H-mode is somewhat higher than for ELMy H-discharges.

The confinement analysis of the ELM-free discharges is performed when $P_{IO}$ is still rising, when the $dW/dt$ correction is 25% - 33% of the total heating power and the radiated power is small. The low $q_a$ discharges - terminated by disruptions - are included. For $P_{tot} > 2\text{MW}$ the quality factor $\tau_E / I_p$ without ELMs and $D^0 \rightarrow D^+ \approx 0.24 \ s/\text{MA}$, for $H^0 \rightarrow D^+$ it is $\approx 0.20$; with ELMs it is 0.14 for hydrogen injection. $\tau_E / I_p$ is almost independent of $I_p$ and $B_t$. In particular, operation at low $q_a$ and high current does not show any pronounced degradation.

4.2. $\beta$-limit studies with and without ELMs

A Troyon factor of 2.8 (based on $\beta_{dia}$) is obtained both for quiescent H-phases and the phases with ELMs \ [13 \]. $\beta$-limit operation with ELMs requires, however, a higher heating power.

4.3. Density limit in the H-mode

No detailed density limit program has been carried out. However, at low $q_a$ where the density limit for OH and L-mode plasmas is highest the H-mode density limit has been determined for the H-mode with ELMs. It is comparable to the corresponding L-mode limit. $\bar{n}_e = 1.05 \times 10^{19} \text{cm}^{-3}$ has been obtained corresponding to a Murakami parameter of $10 \times 10^{19} \text{m}^{-2} \text{T}^{-1}$. $f_{H}$ at the density limit is 1.85. With pellet refuelling of H-mode discharges values of $M$ up to 13 could be achieved. Owing to the superior particle confinement successful density built up with pellets is possible in the H-mode.

5 POLOIDAL ROTATION

The poloidal rotation of He, C, and B impurities has been measured at the plasma periphery \ [14 \]. The measurement on helium-like boron were the
most successful. This impurity developed intrinsically after boronization of the vessel wall. The plasma edge was explored from 36 to 40 cm (a = 40 cm). The measurements were carried out with a scannable mirror system which allows the determination of poloidal (v_p), perpendicular (v_\perp), and toroidal (v_t) velocities. The poloidal and perpendicular velocities, given below, are averages from data obtained from opposing viewing directions.

Velocity vectors, as measured in the three directions, are plotted in Fig. 3 for L and two quiescent H-phases with different toroidal field.
directions. (For technical reasons, $I_p$ had to be changed with $B_t$; the negative $B_t$ case is therefore studied in a quiescent H-phase with ctr-NI.) $v_t$ is basically identical for the three cases. (At smaller radii, $v_t$ is larger in the H- than in the L-phase.) In the L-mode $v_p$ and $v_\perp$ are small and point into the direction of $B \times \nabla p$; in the two H-phases $v_p$ and $v_\perp$ are large and point in the $-B \times \nabla p$ (electron drift) direction. $v_p$ and $v_\perp$ change sign when the direction of $B_t$ is changed.

The high poloidal velocity cannot be explained in terms of the pressure gradient of the impurity ion under study but indicates the existence of a strong inward directed radial E-field (confirming the finding of DIII-D [15]) of $|E_r| \leq 20$ kV/m. The gradient in $v_p$ is small; the diagnostic accuracy did not allow for the determination of a gradient in $E_r$.

$v_p$ and $v_\perp$ increase gradually after the transition and show a temporal variation similar to that of edge gradients. As an example, the edge ion temperature (from Doppler measurements) and $p_e'$ are plotted in Fig. 3. When ELMs occur or in ELMy H-phases $v_p$ and $v_\perp$ decrease and remain at

![Fig. 4](image_url)

**Fig. 4.** (a) Development of edge density (reflectometry) from the L-phase, 0.5 ms after the transition and then in a quiescent phase which developed after an ELMy phase; (b) radial variation of edge fluctuations from the L-mode during the $H^*$ phase for the same time points. $I_p = 460$ kA; $B_t = 2.8$ T; $P_{tot} = 2.0$ MW.
small values. Nevertheless, in this case also, \( v_p \) and \( v_\perp \) point in \(-B \times \nabla p\) direction. The highest values of \( v_p \) and \( v_\perp \) (reached toward the end of the H* phase) rise with heating power. The edge drift velocities have also been studied at two different \( B_t \) values (\( I_p \) was also changed to keep \( q_a \) constant). No distinct difference has been observed. These various observations indicate that the origin of \( E_r \) may be the pressure or temperature gradients at the edge. \( E_r \) corresponding to \( E_r = \frac{v_p}{n_i e} \) (\( v_p \) being the deuterium pressure gradient; estimated with the assumption that \( p_i' = p_e^- \)) is, however, not more than 50% of the measured one. The neo-classical flow velocity \( v_p = -\text{const.} \times \nabla T_i / e \times B \) is approximately 30% of the measured one assuming const. = 1.

The collisionalities at the edge prior to the transition are typically \( v_e^* > 5 \) and \( v_i^* > 10 \) (as deduced from edge temperature, density and \( Z_{eff} \) measurements). Apart from reflectometry (see Fig. 4), however, none of the edge diagnostics on ASDEX can resolve the actual gradients.

6 TRANSITION AND FLUCTUATION STUDIES

ECE radiometer measurements revealed that an H-transition which is triggered by a sawtooth occurs at the arrival of the heat pulse at the edge. Then the edge transport barrier [3] develops. The \( H_\alpha \) radiation decreases at various positions (main plasma chamber, upper and lower, inner and outer divertor chamber) simultaneously within \( \approx 80 \mu s \). In Fig. 4a the development of the edge density profile (as measured by broadband reflectometry [16]) is shown. The first H-mode profile is obtained \( \approx 1.5 \) ms after the H-transition. The transport barrier expands radially during the quiescent phase at constant density gradient. The edge fluctuation level decreases. The most striking observation is the rapid reduction of the edge turbulence measured within the transport barrier by FIR laser scattering. Radially propagating modes are investigated. The fluctuation level drops within \(< 1 \) ms to a low level comparable with that during the OH phase (despite of strongly increased edge density). Figure 4b shows the fluctuation level as measured with reflectometry and indicates the radial regions of low fluctuation level during the quiescent H-phase.

Although, originally reduced, the edge fluctuations can strongly increase in the course of the quiescent phase to levels even above those of the L-mode. They are detected by edge FIR scattering, reflectometry, edge soft-X ray emission, and edge electron temperature (radiometer). Unlike the broadband L-mode turbulence, the fluctuation frequency is around 45 kHz.
indicating the coherent nature of these fluctuations. They develop after a delay time in the quiescent phase and do not appear between ELMs. The edge $T_e$ is instantly reduced; the high $v_L$ or the density built up are not impaired. The reduction of $T_e$ may ultimately cause the transition back into the L-phase. (Other studies indicate that the back-transition is not caused by excessive impurity radiation.)

7 MHD STUDIES ON ELMs

ELMs are important to keep the ASDEX plasma clean and to produce steady state conditions. Furthermore, as described in paper [17] they contribute to the deposition of heat onto the target plates. Their MHD nature has still not been identified. The $T_e$-inversion point of ELMs is about 4 - 5 cm inside the separatrix: for $r > 36$ cm $T_e$ transiently increases; $r < 35$ cm it decreases. The inversion radius does not change with $q_a$ (in the range $3.1 \leq q_a \leq 4.2$).

The ELM frequency can vary strongly during the H-phase. Figure 5a shows a plasma with a single ELM shortly after the transition followed by a group of 4 ELMs, then 2 singular ELMs and, finally, a sequence of grassy ELMs. Further MHD effects are a sawtooth in the L-phase (1), one which triggers the H-transition (2), a sawtooth in the quiescent H-phase (4) and one amidst the group of 4 ELMs (5). This sequence of MHD effects was studied with an hf-Mirnov loop (digitizing time 300 ns). Figure 5 (1) - (8) shows the development of the frequency spectra as contour plots. The characteristic low-frequency $m \geq 2$ modes accompanying the $m = 1$ sawtooth precursor, ELM-precursors in the frequency range 100 - 200 kHz followed by a strongly turbulent phase with a frequency spectrum reaching up to 300 kHz at the occurrence of ELMs, are shown. The duration of the turbulent phase, which causes the particle and energy losses, is typically 200 $\mu$s. Hf-ELM precursor modes may also appear as transient oscillations without causing an ELM in a manner already familiar from $m > 1$ fishbones. The growth rate of an ELM precursor is typically 50 $\mu$s. The precursor to the first ELM ($f = 180$ kHz) in Fig. 5 (3) has the lowest growth rate. Figure 5 c shows the integrated Mirnov signal for the upper ELM in Fig. 5 (5). The successive turbulent phase is averaged and displays the gross plasma movement. Figure 5 b shows the $H_\alpha$ radiation measured with a fast detector in the main chamber. $H_\alpha$ rises simultaneously within 20 $\mu$s at all toroidal positions in the main chamber, in the upper and lower divertor and in the inner and outer divertor SOL.

The ELM features do not differ between ELMs occurring at low pressure shortly after the transition, singular ELMs or in grassy ELM phases.
Fig. 5. (a) $H_\alpha$ in the divertor chamber; the numbers indicate the time point of various MHD activity. (b, c) $H_\alpha$ signal in main chamber and Mirnov-loop signal ($B_r$) of the 3. ELM in (5). (1) - (8) Time development of frequency spectra for the various MHD effects: (1), (2), (4) sawteeth, (3) - (8) ELMs, (5) 3 ELMs and a sawtooth after the 1st ELM.
The edge electron pressure gradient is stabilized at a lower level by repetitive ELMs. In appropriately positioned discharges (as described above) ELMs can be suppressed and much larger $p_e$ (scaling with $I_p$) can be sustained at the edge.

Although we cannot present results on the topology of ELM precursors, the observations are in qualitative agreement with an edge stability analysis presented in Ref. [2].

8 CONCLUSIONS

A comparative study of various confinement regimes revealed the superiority of the H-mode. Although quiescent H-phases do not seem to be suitable for steady-state operation, they might nevertheless be the best scenario for ignition. For continuous burn the transition to the regular H-phase with ELMs might be necessary. Divertor design and wall material seem to be crucial for a large $f_H$ value and for the development of quasi-steady state conditions. For minimizing the ELM-induced losses, the ELM-frequency has to be optimized by proper positioning of the plasma with respect to the vessel. Operation with ELMs does not impair the density - or the $\beta$ -limit. H-mode operation well below $q_a = 3$ is possible with still good confinement. ELMs are characterized by a precursor in the frequency range of 100 - 200 kHz which initiates a short highly turbulent phase which causes enhanced losses. These features are observed in singular ELMs or in those which appear in a burst of rapid ELMs at low or high edge pressure.

Edge drift velocity measurements indicate the development of a radial E-field close to the separatrix. Whether it corresponds to the bulk ion pressure gradient cannot be resolved. There is, however, indirect evidence: the drift velocity increases gradually during the quiescent H-phase on a time scale similar to that of the edge gradients; both the edge gradients and the drift velocity are strongly reduced in ELMy phases. Edge fluctuations develop in the quiescent H-mode which reduce the edge electron temperature.

REFERENCES

DISCUSSION

H. BIGLARI: Since the sampling times of your spectroscopic measurements of poloidal rotation are significantly longer than the time-scale for the L–H transition (< 1 ms), do you think that anything conclusive can really be said on the basis of your data regarding the incidence of (sheared) poloidal flow and the suppression of turbulent fluctuations?

F. WAGNER: About 20% of the overall increase of $v_\perp$ observed during the quiescent H-phase could occur instantly without being detected.

R. R. WEYNANTS: I would like to point out that according to Shaing’s theory the electric field has to be about 3–4 times higher after the transition than the pressure gradient term before the transition. During the H-mode the electric field can relax to rather close to the pressure gradient term and the plasma will not go out of the H-mode.

K. ITOH: Have you studied the correlation between the frequency and the stationarity of the discharge, for example with regard to impurity content or density? Is there any threshold frequency of ELMs required for a stationary H-mode to be achieved?

F. WAGNER: There is an inverse relation between $Z_{\text{eff}}$ and ELM frequency. Our actual task in the experiment was to optimize the conditions in such a way that steady state conditions would prevail at sufficiently low $Z_{\text{eff}}$ and sufficiently high confinement time. The example shown had $Z_{\text{eff}} < 2$, $f_H \sim 1.9$ and $\beta/\beta_{\text{crit}} \sim 0.7$. 
R.J. TAYLOR: We are getting closer to agreement with you. What is your spatial resolution on the rotation measurements?
F. WAGNER: About 1 cm.
R.J. TAYLOR: Then I think we are in 100% agreement, and that you have a bifurcated rotation just like those on DIII-D and CCT.
PELLET INJECTION WITH IMPROVED CONFINEMENT IN JT-60

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Abstract

PELLET INJECTION WITH IMPROVED CONFINEMENT IN JT-60.

With pellet injection, the energy confinement for auxiliary heated (<25MW) discharges was enhanced up to 40% relative to gas fuelling. For the highest current (3.1MA), \( n_e(0)/n_e < 3 \) with \( n_e(0) < 2.7 \times 10^{20} \text{ m}^{-3} \) was sustained for 0.5-1 s after a series of pellet injections. Central plasma pressure reached 2 atm and achieved values of \( n_e(0)T_i(0)\tau_E \) (<1.2 \times 10^{20} \text{ m}^{-3} \text{skeV}) were doubled compared to gas puffing. The \( q=1 \) surface \( r_s \) plays a particular role in the improvement. The density and pressure profiles peak strongly inside \( r_s \) and the particle confinement at \( r<r_s \) appears to be better than that in the outer \((q>1)\) region. With decreasing sawtooth frequency, the density peaking factor and \( \tau_E \) increase. The electron pressure gradient inside \( r_s \) is locally marginal for the ideal infinite-\( n \) ballooning mode, while...

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the total pressure inside $r_s$ is consistent with the poloidal beta limit for the internal $n=1$ kink mode. At the sawtooth crash after the pellet injection, only a small fraction of the central energy is released.

1. INTRODUCTION

One of the main goals of fusion research is to produce high central pressure with improved confinement in a pure plasma. To study the feasibility of this condition, pellet injection experiments have been carried out on tokamaks and demonstrated improved confinement with peaked density profile [1]. This paper studies the effects of the density and pressure profiles on the confinement and the behavior of the instability limiting the central pressure and demonstrates that the $q=1$ surface $r_s$ plays a particular role in improved confinement.

2. GENERAL OBSERVATIONS

In JT-60, hydrogen pellets were injected into hydrogen or helium target discharges [2-4]. The pellet injector produces 4 pellets (3mm $\Phi$×2 and 4mm $\Phi$×2) with velocities approaching 2.3km/s [3] which allows us the direct fuelling in the core region even for high $I_p$ and high power heated discharges. The scenarios of pellet injection, plasma equilibrium and diagnostics were presented in Ref. [2].

Figure 1(a) gives the relationship between the stored energy measured with a diamagnetic loop $W_{DIA}$ and the absorbed power $P_{abs}$ for NB-heated limiter discharges. With pellet injection (closed symbols), $W_{DIA}$ was enhanced up to 30-40% relative to the gas-fuelled scaling at $P_{abs} \approx 10-15$MW. In this figure, only the pellet case having the highest improvement is plotted for each combination of $I_p$ and $P_{abs}$. Figure 1(b) shows electron pressure profiles $p_e(r) (= n_e(r) \times T_e(r);$ Thomson scattering, ECE and soft X ray (SX) data ) for pellet fuelling and gas fuelling. About 50-70% of the enhanced stored energy ($\Delta W$) was the contribution of the pressure peaked inside $r_s$. For each $I_p$, $\Delta W$ is roughly constant and independent of $P_{abs}$, which is explained by the observation that the pressure at $r<r_s$ is limited by $\beta$-limits. For $I_p=3.1$MA, central plasma pressure reached ~2 atm and achieved values of $n_e(0)T_i(0)\tau_E (<1.2 \times 10^{20}$m$^{-3}$skeV; $n_e(0)\approx 2.0-2.7 \times 10^{20}$m$^{-3}$, $T_i(0)\approx 2.5$-3.5keV, $\tau_E\approx 0.15-0.2$s ) were doubled compared to gas puffing. The stored energy became maximum at 0.4-1 s after the pellet injections; at that time the profile of $Z_{eff}$ was almost flat with $Z_{eff}(0)\approx 2$ and the total radiation loss $P_{rad}$ was ~20% of $P_{abs}$. The Murakami factor reached $10-13 \times 10^{19}$m$^{-2}$T$^{-1}$ which was higher than those for gas fuelling by factors of 1.5-2.0. The density limit can be explained with the power balance in the outer region of the plasma [5]. Figure 1(c) shows typical $n_e(r)$ for limiter discharges. For $I_p=3.1$MA ($q(a)=2.3$), the density peaking factor of $F_{np}<3$ ( $F_{np}=n_e(0)/<n_e>$; central /
volume averaged) with \(n_e(0) < 2.7 \times 10^{20} \text{m}^{-3}\) was sustained for 0.5-1 s after a series of pellet injections. In Fig. 1(c), change in \(n_e(r)\) from 0.5 s to 0.8 s after the injection for the 3.1MA discharge shows the central particle confinement time \(\tau_p^*\) is \(~3\)s which is more than ten times longer than \(\tau_E\). For \(I_p = 2.1\text{MA}\) \((q(a)=3.3)\), \(F_{np}\) reached 4.5 [2]. Figure 1(d) gives \(\tau_E\) as a function of \(I_p\) for the pellet and gas fuelled lower x point discharges. For gas fuelling, the increase in \(\tau_E\) with \(I_p\) is degraded at \(q_{\text{eff}}<3\), while, for the pellet fuelling, \(\tau_E\) increases linearly with \(I_p\) even at \(q_{\text{eff}}<3\) because of the sawtooth suppression [2].
When pellets penetrate inside \( r_s \), the sawtooth activity can be suppressed completely during 0.4-1s or the sawtooth frequency is reduced by up to one order of magnitude, then highly peaked \( p_e \)-profiles are obtained. Figure 2(a) gives the profile evolution of Abel inverted SX-emission (~\( p_e^2 \) in the pellet injected JT-60 plasmas) for a NB-heated 2.1MA limiter discharge (\( P_{\text{abs}}=13\,\text{MW} \)) from \( t=6.2 \) to 6.5s. Three pellets were injected at \( t=6.0\,\text{s} \). The SX-profile starts to peak from \( t=6.2\,\text{s} \) and evolves only inside a radius \( r_p \) (~28cm). We scanned \( I_p \) (0.5-3.1MA) and \( B_t \) (3.5-4.8T) to survey the dependence of \( r_p \) on \( q(a) \). Figure 2(b) shows SX-emission profiles normalized by the central values for \( q(a)=2.3, 3.3 \) and 4.1. Figure 2(c) gives the relationship between \( r_p \) and \( q(a) \). The solid line given by \( r_p=a/q(a) \) represents the \( q=1 \) radius \( r_S \) estimated by the sawtooth inversion for gas fuelling. In Fig.2(c), \( r_p \) increases linearly with \( 1/q(a) \) and the values agree systematically with \( r_S \); \( r_p=r_S \). This fact is confirmed by the spatial structure of the m/n=1/1 oscillation emerging in the sawtooth-free phase after the injection. The sawteeth resuming after pellets have the inversion radii given by the same formula \( r_S=a/q(a) \). Therefore it is concluded that the \( p_e \)-profile peaks strongly just inside \( r_S \).
3. β-LIMIT

After the pellet injection, the $p_e$-profile grows inside $r_S$; however, the pressure gradient $dp/dr$ is saturated at a certain value before the resumption of a sawtooth and the saturated value is independent of the NBI heating power [2]. At first, the saturation occurs locally at the outer portion inside $r_S$, then the saturation spreads inward (see Fig. 2(a)). The saturation is consistent with the ideal ballooning limit. Figure 3(a) shows the measured $dp/dr$ (Thomson and SX data) and the marginal values for the ideal infinite-$n$ ballooning mode calculated for assumed $q(0)$ of 0.98, 0.95 and 0.90 [6]. Within the assumed region of $q(0)$, the maximum $dp/dr$ reaches the marginal value inside $r_S$. The Mercier stability limit is slightly higher than that for the ballooning mode where
the experimentally obtained \( \frac{dp}{dr} \) is maximum. Figure 3(b) shows the results of B\(_t\)-scans for the limiter plasmas with almost the same \( q(a) \). The hatched region indicates the marginal values against the ballooning mode calculated for \( q(0)=0.90-0.98 \). The observed \( \frac{dp}{dr} \) is roughly proportional to \( B_t^2 \). For the simple assumption that the normalized pressure gradient \( \alpha=-2\mu_0 \frac{R q^2}{B_t^2} \times \frac{(dp/dr)}{B_t} \) is fixed at the marginal values against the ballooning mode, \( \frac{dp}{dr} \) is proportional to \( B_t^2 \) because \( q \approx 1.0 \) and \( R \) is fixed. Figure 3(c) gives the experimentally obtained \( \beta_p^1 \) (open circles) and calculated \( \beta_p^1 \)-limits for the ideal ballooning mode and the \( n=1 \) ideal kink mode, where \( \beta_p^1 \) is the poloidal beta determined inside \( r_s \) given by

\[
\beta_p^1 = (2\mu_0 / B_p^2(r_s)) \int_0^{r_s} (r / r_s)^2 ( -dp/dx ) dr
\]

The \( n=1 \) kink limit was calculated with the ERATO-J code for \( q(0)=0.95 \) using experimentally obtained pressure profiles. For each discharge, the maximum \( \frac{dp}{dr} \) is locally marginal for the ballooning mode (above the broken line), while the total pressure inside \( r_s \) has not reached the ideal ballooning limit fully optimized in the whole region inside \( r_s \). In many cases, the experimentally obtained \( \beta_p^1 \) is limited by the sawtooth crash, which is close to the calculated limit for the \( n=1 \) internal kink mode. The results of the ballooning mode analyses in this paper suggest that the ballooning limit is a soft limit which only saturates the pressure gradient.

4. SAWTOOTH AND CONFINEMENT

Figure 4(a) compares the SX-signals and \( n_e(r) \) for four discharges (\( q_a \approx 2.3; I_p/B_t=2.8MA/4.5T \) and 3.1MA/4.8T). Pellet penetration becomes deeper from top (1) to bottom (4). The ratio of the central SX-intensity just before to that just after the injection is a good measure of the pellet deposition in the core region. The values are (1) 73\%, (2) 40\%, (3) 23\% and (4) 21\%. Values of \( \beta_p^1 \) are (1) 0.07, (2) 0.22, (3) 0.24 and (4) 0.35 (±0.05). With deepening penetration, \( F_{np}, \beta_p^1 \) and the duration of sawtooth suppression increase. The sawtooth crash time after pellets is typically 300–600\( \mu s \), which is longer than that (100–200\( \mu s \)) for gas fuelling. For higher \( q \) and lower \( \beta_p^1 \) discharges, the sawtooth crash often has precursor \( m=1 \) oscillation and very small or no successor oscillation and flattens the central \( p_e(r) \) profile completely. In contrast, for lower \( q \) and higher \( \beta_p^1 \), the crash tends to follow incomplete reconnection and the successor oscillation becomes larger. In some cases the core region behaves as a rigid body during the crash (~300\( \mu s \)). With increasing \( \beta_p^1 \) or the density gradient in the core region, the sawtooth period increases, the release of the central kinetic energy at the crash decreases and the
FIG. 4. (a) Time histories of the SX-signals and $n_e(r)$ (Thomson). Pellet penetration becomes deeper from top to bottom. (b, c) Relationships between $\Delta W$, $F_{np}$ and $\tau_n$ for $I_p/B_n=2.8\,\text{MA}/4.5\,\text{T}$ and $3.1\,\text{MA}/4.8\,\text{T}$ ($q(a) \approx 2.3$) limiter discharges with $P_{abs}=16-19\,\text{MW}$. (d) Dependence of $\Delta W$ on $\tau_n/\tau_E$. (e) $q(a)$ versus $\tau_n/\tau_E$ at the saturation point.
crash changes its characteristics from a reconnecting to a non-reconnecting type [7].

Figures 4(b)–(d) give the relationships among $\Delta W$, $F_{np}$ and the sawtooth period $\tau_{st}$ averaged from the injection time to the time when $W_{\text{DIA}}$ becomes maximum for $I_p/B_t=2.8\text{MA}/4.5\text{T}$ and 3.1MA/4.8T ($q(a)\sim2.3$) limiter discharges with $P_{\text{abs}}=16$–19MW. For gas fuelling with the same operation parameters, $F_{np}=1.2$–1.3 and $\tau_{st}\sim0.07$–0.08s. Figure 4(b) shows that $\Delta W$ increases linearly with $F_{np}$. Figure 4(c) indicates that $F_{np}$ (and therefore $\Delta W$) increases with $x_{st}$ for $x_{st}<0.4\text{s}$ and saturates for $x_{st}>0.4\text{s}$. In Fig.4(d) $\tau_{st}$ is normalized by $\tau_E$. The saturation occurs at $\tau_{st}/\tau_E\sim1.4$. Figure 4(e) gives the relationship between $q(a)$ and $\tau_{st}/\tau_E$ at the saturation point, $(\tau_{st}/\tau_E)_{\text{strt}}$, obtained in the four sets of discharges with different $I_p$ [8]. The contribution of the sawtooth to the global confinement becomes larger for lower $q(a)$. For $q(a)=2.3$ and 2.7, the plasma does not reach the balloon/kink limit at the saturation point and the saturation seems to be determined by transport properties independent of the MHD modes, while for $q(a)=3.3$ and 4.1, the saturation occurs due to the $\beta$-limit.

5. TRANSPORT

To simulate the time evolution of the pellet injected plasmas, a 1-1/2 dimension transport code was developed. The diffusion coefficients were assumed to reproduce the experimentally measured profiles of the kinetic data and the dependence of $W_{\text{DIA}}$ on $I_p$ and $P_{\text{add}}$ (the auxiliary heating power ) for L-mode discharges in the high density regime (Fig.5(a)):

$$
\chi_e(r) = (\chi_e^{\text{OH}}(I_p) + 0.045P_{\text{add}}) \times q(r) \\
\chi_i(r) = \chi_i^{\text{NC}} + \chi_e \\
D_e(r) = (0.2 + 0.015P_{\text{add}}) \times q(r) \quad ; \quad (r>r_s) \\
= 0.1 \text{ m}^2/\text{s} \quad ; \quad (r<r_s)
$$

With this model, the time evolution of $n_e(r)$, $T_e(r)$ and $W_{\text{DIA}}$ for the pellet injection can be well reproduced using the same diffusion coefficients with the gas fuelling (Fig.5(b) and (c)) when the inward pinch velocity of $V_a=0.4$–0.5m/s ($V(r)=V_a r/a$ ) is assumed for the pellet fuelling [9]. Using this model, the improvement of the thermal diffusivity is not necessary to explain the enhanced $\tau_E$. The good particle confinement within $r_s$ yields the improved $\tau_E$ when the strong particle source is applied inside $r_s$ under the sawtooth-free condition.

For $I_p=3.1$MA with highly peaked $n_e(r)$, data of spectroscopy show the central accumulation of titanium. The behavior of the metal impurity can be explained well with the reduced particle diffusivity inside $r_s$. The content of the metal impurity is not significant ( <0.01% of $n_e$) to deteriorate the power
3.1 MA, 4.8T, limiter

\( P_{NB} = 6 \text{MW} \) (t<6.05s)
\( 14 \text{MW} \) (t>6.05s)

**FIG. 5.** (a) Profiles of \( D_c, \chi_t, \) and \( \chi_e \) used in the transport simulation \((I_p=3.1 \text{MA}, P_{abs}=14 \text{MW})\). (b, c) Comparison of measurement and calculation for \( n_e(r) \) and \( T_e(r) \) (Thomson), \( W^{DMA} \), line integrated density at \( r=0.5a \) and \( T_e(0) \) (ECE) \((3.1 \text{MA}, P_{abs}=17.4 \text{MW})\).

balance; \( P_{rad}(\text{total}) \sim 20\% \) of \( P_{abs}(\text{total}) \) and \( P_{rad}(0) \sim 50\% \) of \( P_{abs}(0) \). The profile of \( Z_{eff} \) is almost flat throughout the improved stage of confinement and there is no clear evidence of the light impurity accumulation.

### 6. CONCLUSION

With pellet injection, \( \tau_E \) for auxiliary heated discharges was enhanced up to 40\% relative to gas fuelling. For \( I_p=3.1 \text{MA}, n_e(0)/<n_e> <3 \) with \( n_e(0)<2.7 \times 10^{20} \text{m}^{-3} \) was sustained for 0.5-1s after the pellet injection. Central plasma pressure reached 2 atm with \( Z_{eff}(0) \sim 2 \) and values of \( n_e(0)T_i(0)\tau_E \) \((<1.2 \times 10^{20} \text{m}^{-3} \text{keV}) \) were doubled compared to gas puffing. For the improvement, \( r_s \) plays a particular role. The density and pressure profiles peak
inside \( r_s \), where the particle confinement appears to be better than that at \( r>r_s \). With decreasing sawtooth frequency, the density peaking factor increases and \( \tau_E \) is enhanced. For the divertor configuration, the degradation of \( \tau_E \) at \( q_{\text{eff}}<3 \) recovers to the level for the \( q_{\text{eff}}>3 \) region because of the sawtooth suppression. Simulation analyses of transport show that the enhanced stored energy and the effects of sawteeth on \( \tau_E \) can be explained with reduced particle diffusivity inside \( r_s \). The electron pressure gradient inside \( r_s \) locally reached the marginal values for the ideal infinite-\( n \) ballooning mode, while the total pressure inside \( r_s \) is consistent with the beta limit for the internal \( n=1 \) kink mode. At the sawtooth crash after the pellet injection, only a small fraction of the central energy is released and the crash does not follow the fully reconnecting style.

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REFERENCES


DISCUSSION

R. GIANNELLA: Could you describe in greater detail what you observe on the soft X ray emissivity profile at the sawtooth crash in your pellet fuelled discharges?

Y. KAMADA: For lower \( \beta_1 \) discharges, the sawtooth crash is of the fully reconnecting type, and the central pressure profile is flattened completely by the crash. In contrast, for higher \( \beta_1 \) discharges the crash tends to follow incomplete reconnection, and the release of central kinetic energy at the crash is very small.
STUDY OF LIMITER H- AND IOC-MODES
BY CONTROL OF EDGE MAGNETIC SHEAR
AND GAS PUFFING IN THE JIPP T-IIU TOKAMAK

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Abstract

STUDY OF LIMITER H- AND IOC-MODES BY CONTROL OF EDGE MAGNETIC SHEAR AND
GAS PUFFING IN THE JIPP T-IIU TOKAMAK.

Two types of improved confinement regime, H-mode and IOC (improved Ohmic confinement)
mode, are studied in circular limiter plasmas on JIPP T-IIU. When rapid ramp-down of plasma current
(CRD) is used during auxiliary heating, the threshold heating power required for the L-H transition is
reduced by 30–50%, compared with that in the case without CRD. This is interpreted as being due to
enhancement of the global magnetic shear near the plasma edge by CRD. This model, based on edge
magnetic shear is also applicable to the limiter H-mode obtained without CRD by high heating power.
In the limiter configuration, the IOC-mode is obtained in ohmically heated, high density plasmas by gas
puff control. This improvement may be attributed to a reduction of anomalous transport since radiation
loss has only a minor effect.

1. LIMITER H-MODE

The trigger mechanism of the L-H transition is still unclear, while
many tokamaks have achieved the H-mode in divertor and limiter
configurations since its discovery in ASDEX [1]. Recently, transition
models based on a change in the edge radial electric field [2,3] have
been developed. They seem to partly explain the data on the edge
radial electric field obtained experimentally [4,5] and the dramatic
suppression of edge fluctuations in the high frequency, short
wavelength region at the transition [6]. The models, however, cannot
explain the experimental observation, in ASDEX [7], of coherent
magnetic fluctuations with low mode numbers ( m = 3 or 4, n = 1),

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which are localized near the plasma edge, excited during the L-phase and suppressed just before the transition. In Ref. [7], a new transition model is proposed in which the transition is governed by a change in the toroidal current density ($J_t$) near the edge. We test the new model in the limiter tokamak JIPP T-IIU (R = 91 cm, a = 23 cm, $B_t$ = 2.7-3.0 T), modifying the edge current density by rapid ramp-down of the plasma current (CRD) [8]. The limiter configuration is advantageous for studying the effect of magnetic shear on the transition because the magnetic shear can be simply changed by modifying the current density profile with the same external magnetic configuration.

![Graph](image_url)

**FIG. 1.** Effect of rapid current ramp-down on L-H transition in ICRF heated deuterium plasmas with about 10% hydrogen minority, where $P_{RF} = 1.2$ MW and $B_t = 2.8$ T. $I_a$ denotes $H_a/D_a$ emission.
Figure 1 shows the effect of CRD on the L-H transition, where the heating power is set below the threshold in the case without CRD [9]. The transition does not always correlate with high edge electron temperature \( T_{eb} \) just before the transition. The dependence of the heating power (RF or RF+NBI) on plasma current \( I_p \) at the transition is investigated for H-modes with and without CRD. The threshold power in the case with CRD is by 30-50% lower than that in the case without CRD for the same \( I_p \) [8]. It is concluded that the transition is not determined by the decreased \( I_p \), but by the modification of the edge \( j \) profile.

Figure 2(a) shows the H-mode with CRD, where the transition is initiated without any sawtooth event and without any rise in \( T_{eb} \) before the transition. The magnetic shear near the edge \( s_{mb} \) (at \( r/a = 0.8 \) and 0.9, where the transport barrier determined experimentally is located at \( r/a = 0.8-0.9 \) [8]) calculated from a magnetic diffusion equation is substantially increased by CRD. When it increases up to a certain value \( s_{mb} = 1.8-2.0 \), the transition readily occurs.

![FIG. 2. (a) H-mode triggered by CRD without sawtooth event and \( T_{eb} \) rise before transition, where \( P_{RF} = 1.3 \) MW and \( P_{RF} = 0.6 \) MW. \( s_{mb} \) is the magnetic shear at \( r/a = 0.8 \) (solid curve) and 0.9 (dashed curve), calculated for the \( Z_{eff} \) value from visible bremsstrahlung; (b) model profiles of \( j_\phi \), \( q \) (safety factor) and \( s_m \) (global magnetic shear) for H-mode initiated by CRD.](image-url)
FIG. 3. (a) Power spectrum calculated from Mirnov probe signal $B_\theta^\text{HI}$ from $t = 225$ to 230 ms in the H-mode without CRD, where $P_{RF} = 1.1$ MW and $P_{NI} = 0.75$ MW. The spectral peak $C$ corresponds to $m/n = 3/1$ unstable mode whose rational surface is located near the edge. Time evolution of the $m/n = 3/1$ mode filtered from signals of Mirnov probes on the inner and outer midplane ($15 \leq f \leq 35$ kHz), edge and central electron temperatures, and $H_+/D_+$ emission measured by a detector viewing outer edge; (b) model profiles for H-mode obtained without CRD.

increase in $\tilde{S}_{mb}$ up to this value means that the current channel is almost completely detached from the limiter, as shown in Fig.2(b). On the other hand, rapid ramp-up of $I_p$ easily quenches the H-mode obtained at very high heating power ($P_{RF} + P_{NI} = 2.6$ MW). This is interpreted by the assumption that $\tilde{S}_{mb}$ is appreciably reduced by the ramp-up.

We shall now study whether the H-mode without CRD, obtained with high heating power, is governed by the same mechanism, i.e.,
edge magnetic shear or edge $j_\phi$ profile. When strong electron heating occurs or a substantial amount of bootstrap current is generated by high power heating in the plasma volume, the plasma current $I_p$ tends to increase on a time-scale shorter than the resistive diffusion time of the plasma. Then, a reversed toroidal electric field near the edge may be induced to conserve the poloidal magnetic flux. This field may detach the current channel from the limiter, keeping $I_p$ constant, as shown in Fig.3(b). This change in $j_\phi$ profile is inferred from the time evolution of the low mode number magnetic fluctuations driven by $\nabla j_\phi$ as observed in ASDEX [7]. If the rational surface of the mode is located within the detached region ($i_\phi = 0$), it may easily be stabilized, because of reduced $\nabla j_\phi$ and enhanced $s_{mb}(\sim 2)$ there. Figure 3(a) shows the time evolution of the $m/n = 3/1$ mode observed in H-mode without CRD, where the rational surface is located near the edge around $r/a = 0.8-0.9$. Suppression of the mode can be explained by the above discussion. Note that, as shown in Fig.3(a), the $m/n = 3/1$ mode can be most unstable in the course of the transition (L-to-H, or H-to-L, or ELM) between the L-type and the H-type $j_\phi$ profile, accompanying ELM(-like) spike in the $H_\alpha/D_\alpha$ emission.

The above mentioned experimental evidence of JIPP T-IIU suggests that the transition between L- and H-phase is governed by the edge current density or the related magnetic shear rather than by the edge radial electric field or its shear.

2. Limiter IOC-Mode

The improved confinement mode in ohmically heated high density plasmas (IOC-mode) was discovered in the divertor configuration of ASDEX by gas puff control [10]. In JIPP T-IIU, the IOC-mode has been obtained in the limiter configuration by the same technique [11]. The transition is triggered by a sudden reduction of the gas puff rate in ohmically heated, high density plasmas at high $q(a) > 6$ (Fig.4(a)). After the transition, the electron density profile starts to peak, and the electron and ion temperatures increase. Then, the plasma confinement is improved considerably. The improvement may be attributed to a reduction of anomalous transport since the radiation power loss has only a minor effect on this ohmically heated plasma. However, density fluctuations related to electron and ion modes, measured by FIR laser scattering [12], exhibit only a small change, as shown in Fig.4(b). Further studies on the fluctuations of the density and plasma potential as well as other plasma parameters in this limiter IOC-mode are required.
FIG. 4. (a) Temporal evolution of limiter IOC discharge, where the vertical dotted line indicates the IOC transition; (b) time variation of electron density fluctuations obtained by FIR laser scattering with heterodyne detection in the other IOC shot.

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References

DISCUSSION

A. TANGA: What is the increase in energy and particle confinement time in the limiter H-mode?

K. TOI: There is a clear improvement in energy and particle confinement, but it is not significant. The improvement of energy confinement in the H-phase is ~15% compared with the L-phase.

H. BIGLARI: If I may offer an alternative scenario, it is conceivable that, as you ramp down the current, the q-value at the edge goes through a low (m,n) resonance, thus instigating a region of non-ambipolarity. A non-ambipolar radial current might then be induced which would torque up the plasma in the poloidal direction. The sheared poloidal rotation thus generated will then quench ambient turbulence and allow transition to an enhanced confinement state. This scenario is entirely consistent with our theoretical predictions, and I strongly encourage you to test this hypothesis by measuring the poloidal rotation and fluctuations at the edge.

K. TOI: We compared the difference between the threshold power for the transition with the current ramp case and the low current case without ramp-down. The threshold power with CRD is obviously about 50% lower than without CRD in the low Ip case, but your theory predicts the opposite. Also, we have identified the coherent mode (m = 3, n = 1) whose resonance surface is located near the edge. If the mode is in a non-linear saturated regime, the mode frequency may indicate the $E_r$ magnitude near the edge. However, in our H-modes this frequency hardly changes until the transition. This also suggests that our results are not in line with your theory.
LIMITER H-MODE EXPERIMENTS ON TFTR


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Abstract

Limiter H-mode experiments on TFTR.

Limiter H-modes with centrally peaked density profiles have been obtained on TFTR using a highly conditioned graphite limiter. The transition to these centrally peaked H-modes takes place from the supershot to the H-mode rather than the usual L- to H-mode transition observed on other tokamaks. Bi-directional beam heating is required to induce the transition. Density peaking factors, $n_e(0)/\langle n_e \rangle$, greater than 2.3 are obtained and at the same time the H-mode characteristics are similar to those of limiter H-modes on other tokamaks, while the global confinement, $\tau_E$, can be $>2.5$ times L-mode scaling. The transport analysis of the data shows that transport in these H-modes is similar to that of supershots within the inner 0.6 m core of the plasma, but the stored electron energy (calculated using measured values of $T_e$ and $n_e$) is higher for the H-mode at the plasma edge. Microwave scattering data for the edge plasma shows broad spectra at $k = 5.5$ cm$^{-1}$ which begin at the drop in $D_a$ radiation and are strongly shifted in the electron diamagnetic drift direction. At the same time, beam emission spectroscopy (BES) shows a coherent mode near the boundary which propagates in the ion direction with $m = 15-20$ at 20-30 kHz. During an ELM event these apparent rotations cease and Mirnov fluctuations in the frequency range of 50-500 kHz increase in intensity.

Introduction

The circular limiter H-mode plasmas on TFTR$^{1-3}$ are of interest for a number of reasons. First, under low recycling conditions the transition is from the supershot$^4$ to the H-regime rather than the usual L- to H- transition.$^5$ A new and important consequence is that the resulting H-mode density profile is highly peaked at the center. Peaked density profiles are of interest to advanced tokamaks such as CIT and ITER. Reactor design studies indicate that the degree of enhanced confinement required for ignition is lowered as the temperature and density profiles become more centrally peaked. Second, the good global confinement of the supershot is largely retained in these H-modes, often with a small increase in $\tau_E$ during the transition, but usually decreasing at the start of ELMs. Third, even with the peaked electron density profile the usual H-mode characteristics are observed. Thus, the experimental findings from the TFTR H-modes will be useful for other tokamaks. Fourth, on TFTR
aspects of transport and confinement of limiter H-mode, L-mode, supershot, and high beta poloidal plasmas may be compared.

**Obtaining TFTR Limiter H-modes: Characteristics**

Limiter H-modes, both with the usual L- to H- transitions and with supershot to H- transitions have been realized on TFTR. These regimes are obtained under different limiter conditions, high recycling and low recycling, respectively. However, since the resulting confinement is significantly enhanced over L-mode and the density profile is highly peaked for the supershot to H-mode case, it has been studied in more detail and will be the subject of this paper.

The operational parameter range for obtaining limiter H-modes on TFTR is plasma current $I_p = 0.8 - 1.7$ MA, toroidal field $B_t = 3.0 - 5.2$ T, major radius $R_p = 2.45 - 2.60$ m, minor radius $a_p = 0.79 - 0.95$ m, and neutral beam injection (NBI) power, $P_b \sim 11 - 28$ MW. Beam heating pulses of 0.5 - 2.0 sec have been used. Information about operational and plasma parameter effects on the threshold power for the transition has been obtained. (1) H-modes have only been obtained with neutral beam injection heating. (2) Bi-directional NBI is required, with modestly counter dominated injection being favorable. No H-modes have been obtained with all co- or all counter-NBI. (3) The threshold power for the transition is linearly dependent on $I_p$ and (4) the threshold power appears to be lowered when a) the current is ramped down, reaching a minimum just before beam turn on, or b) carbon pellets or short duration, medium to high flow, deuterium or helium gas are injected during the plateau of the heating pulse.

The need for nearly balanced-NBI can be seen qualitatively from Fig. 1. Plotted are the $D_\alpha$ signals for a discharge which had 2 secs of balanced-NBI, and for a discharge which was heated by 1 sec of counter-NBI followed by 1 sec of balanced-NBI. The H-mode transition (indicated by drop in $D_\alpha$ signals) in each case begins shortly after the start of balanced-NBI, and the
ELMs begin shortly after the transition. The $I_p$ dependence can be seen from Fig. 2 which is a plot of $P_b$ vs. $I_p$. The bottom of the envelope of the data points is an indication of the threshold beam power. The data of Fig. 2 include a large part of the H-mode database and thus the range of $I_p$ and $B_\phi$ values indicated earlier. No strong $B_\phi$ dependence of the threshold power for the transition has been observed. The line averaged density, $n_e$, is usually $\geq 2 \times 10^{19} \text{m}^{-3}$ for H-modes obtained under low recycling conditions. The effect of co-counter-NBI mix (i.e. net beam momentum) is thought to be due to ion orbit loss effects which lead to the proper radial electric field, $E_r$, needed to trigger the transition. This will be discussed in more detail in a later section.

Characteristic features of the TFTR limiter H-modes agree in general with those of other tokamaks; however, there are also differences. Figure 3 shows the time evolution of some of the plasma parameters typical of the TFTR limiter H-mode: $D_\alpha$, $n_e$ (edge), $T_e$ (edge), $n_e(0)/<n_e>$, $\Phi_\theta$, and $\tau E/\tau E_{L-mode}$.
FIG. 2. Plot of $P_b$ versus $I_p$ for the TFTR limiter H-mode. Data for a few supershots are also included. An apparent dependence on $I_p$ of the threshold power, $P_{th}$, required for a transition to the H-mode is indicated.

The features seen here, and usually observed in other tokamaks, are the drop in $D_\alpha$ light (the drop in CII light is usually larger than that for the $D_\alpha$) at the transition, with corresponding increases in edge $n_e$, $T_e$, electron stored energy ($W_e$), and $T_i$. The global energy confinement is enhanced over L-mode and can be $> 2.5$ times the value for Goldston low-mode scaling.

Peaked density limiter H-modes were first observed on TFTR, and are in contrast to the flat $n_e$ profiles characteristic of divertor H-mode discharges. During the transition, there is usually an increase in $n_e$ across the entire profile; however, the largest fractional increase is at the plasma edge (outer 0.2 m). A time series of $n_e$ profiles showing the evolution of a TFTR limiter H-mode is shown in Fig. 4. The maximum peaking factors, $n_e(0)/<n_e>$, were $\sim 2.5$ and 2.2 for the supershot and H-mode phases of this plasma respectively. For comparison, a similar set of profiles for a supershot are also shown. The transition usually occurs within 150 ms to $> 600$ ms after the start of neutral beam injection. This delay correlates inversely with neutral beam power. The $D_\alpha$ drop in TFTR (see Fig. 3) is not as large (and often it is undetectable) as for
divertor H-modes, but the characteristic increases in edge $n_e$ and $T_e$ still take place. Similar behavior has been observed for limiter H-modes on other tokamaks. Also, the $D_\alpha$ drop can proceed relatively slowly. In Fig. 3 the $D_\alpha$ signal decreases over a 50 ms period, compared to $\leq 2$ ms for H-modes in divertor tokamaks. H-modes of $\sim 1.5$ sec duration have been obtained and the duration appears to be limited only by the beam pulse length. On the other hand, the H-mode can be degraded by MHD activity, similar to what is observed during some supershots.
The Transition, ELMs, and Fluctuations

Plasma fluctuations characteristic of the H-mode and ELM activity have been studied. At the H-transition, magnetic fluctuations (as measured by a system of Mirnov coils) in the range 15-25 kHz increase, while the high frequency magnetic fluctuations in the range 150-200 kHz decrease (see Figs. 3e and 3f). The drop in $D_\alpha$ indicates a change in particle transport at the plasma edge and edge probes show a corresponding drop in floating potential during the transition. Also, at the transition density fluctuations near the plasma edge measured by the microwave scattering system\textsuperscript{7} show broadband activity at $k_\theta = 5.5 \text{ cm}^{-1}$ propagating in the electron diamagnetic direction. A contour plot of scattered power as a function of frequency and time is shown in Fig. 5a for a limiter H-mode with nominally balanced beam injection. This spectrum shows density fluctuations up to 1.8 MHz located near the plasma edge ($z = -0.7 \text{ m} \pm 0.2 \text{ m}, R = 2.85 \text{ m}$) and this is consistent with wave activity propagating poloidally at a velocity $v_\theta \sim 10^4 \text{ m/s}$. This feature begins to grow at the time of the transition, reaching a maximum frequency just before the ELM phase and persists for the duration of the H-mode. No similar activity is observed in the
FIG. 5. (a) Microwave scattering spectrum for a limiter H-mode. The transition is at \( \approx 4.04 \) sec (H-mode from \( \approx 4.04 \) to \( 4.42 \) sec) and ELMs begin at \( \approx 4.14 \) sec. Here \( k_\theta = 5.5 \) cm\(^{-1}\), \( z = -0.7 \) m, \( R = 2.85 \) m, \( I_p = 0.85 \) MA, \( a_p = 0.8 \) m, \( B_\theta = 4 \) T, and \( P_b = 12.7 \) MW (from 3.5 to 4.5 sec). (b) \( D_\alpha \) signal. (c) Spectrum digitized at fast rate. (d) Expanded \( D_\alpha \) signal.
central region of the H-mode discharges. Fig. 5c shows a similar contour plot over a narrow time range (4.2 - 4.25 sec) for the discharge of Fig. 5a,b. Narrowband fluctuations are enhanced dramatically at the time of the ELM while the shifted spectral feature, more apparent in the time averaged data, is associated with the period between the ELMs. This data is consistent with a slowing down of the poloidal rotation immediately after an ELM burst and a speeding up in rotation in the time window between the ELMs. The strongly shifted spectra are also characteristic of the occasional discharge with no ELMs. The BES system shows density fluctuations which are more coherent and which propagate in the ion diamagnetic direction. These disturbances start about 3 cm inside the scrapeoff with a frequency of 20-30 kHz and have a highly coherent poloidal structure with m = 15 to 20. At present, they are known to extend radially over at least 4 cm near the plasma edge. This coherent activity is associated with the period between the ELMs. The relation between the BES coherent activity and the scattered spectra is not understood at present.

ELMs can lead to a significant loss of stored energy resulting in a cessation of the rise of stored energy and density after the transition as is clear from Fig. 3. The studies on TFTR are important since in understanding the nature and physical causes of ELMs, it may be possible to suppress or control them. ELMs can play a useful role in impurity control and in maintaining a steady-state H-mode if the associated energy loss is not too great. In divertor tokamaks, such as DIII-D, single giant ELMs have been observed to cause deterioration in confinement to the extent that the plasma returns to the L-mode and \( v_0 \) decreases significantly.\(^8\) This extreme has not been observed on TFTR. The reason for this may be that since \( n_e(r) \) and the pressure profile, \( p(r) \), are peaked on axis, energy loss from the edge plasma due to ELMs is a smaller fraction of the total stored energy than for the flat density profiles of divertor discharges.

The ELMs are clearly observed on the Mirnov coils, soft x-ray signals, and the edge probe signal during the H-mode phase. The Mirnov coil system was used to characterize the mode responsible for the ELM. The data was digitized at 2 MHz and showed an increase in the intensity of high frequency
(50 to 500 kHz) magnetic fluctuations during an ELM. The oscillations were in phase toroidally for all of the coils in the upper half of the torus and 180° out of phase with the coils mounted in the lower half of the torus. The mode is not ballooning in structure as there is a dominant m/n = 1/0 MHD mode, or an up-down motion of the plasma, with amplitude of 0.1 to 1 gauss. The spectrum of a magnetic loop signal digitized at 2 MHz during an ELM is shown in Fig. 6a, and is compared to that for an ELM free period. The identification of the mode number as m/n = 1/0 (shown in Fig. 6b), is in contrast to the ASDEX results where m = 3 or 4 and n = 1. The TFTR ELM precursor is not outward ballooning in structure, unlike precursors on
ASDEX. High frequency precursor oscillations have also been observed on PBX;\textsuperscript{10} these oscillations also were not outward ballooning in character, but did not reflect the up-down motion of the plasma. Further studies are being done to determine if, in general, the same MHD mode and spatial origin are common for all occurrences of ELMs in TFTR.

Data from the grating polychromator (GPC), with a 10 \( \mu \)s time resolution, are shown in Fig. 7. Intense ECE spikes, with a 20 to 30 \( \mu \)s duration, sometimes precede the rise in \( D_\alpha \) emission and occur simultaneously with or slightly after the start of the high frequency magnetic oscillations monitored by the Mirnov coils. The spikes are consistent with the dumping of electrons (with energy 10 to 20 keV), from a radius 0.15 to 0.2 m inside the plasma edge. The disturbance associated with the ELM moves radially outward at a velocity of \( \sim 2 \times 10^3 \) m/s as is indicated by the data shown in Fig. 7b. The \( D_\alpha \) signal begins to increase at about the time the electron burst reaches the wall.

Chord averaged soft x-ray emission observed at a tangency radius of 0.7 m (\( a_p = 0.8 \) m) starts to drop approximately at the beginning of the increase in magnetic oscillations and the ECE spikes. Interestingly, the x-ray signal for a chord at \( r = 0.45 \) m drops 0.6 ms later, indicating that a precursor induced disturbance propagates slowly inwards from a region near the plasma edge.

**Transport Calculations and Comparisons with Theory**

Some qualitative comparisons of the TFTR limiter H-mode results with various theoretical models can be made. If the feature in the microwave scattering spectra at high frequency is interpreted as a poloidal drift in the electron diamagnetic direction, then, an inwardly pointing radial electric field can be inferred. This is in agreement with the predictions of Itoh & Itoh,\textsuperscript{11} and with \( v_\theta \) measurements on other tokamaks.\textsuperscript{12,13} The fact that there is no x-point within the TFTR plasma, along with the observation that slight counter
dominance in NBI is favorable, is in qualitative agreement with the neoclassical model (e.g. preferable loss of counter injected ions) for the transition. Qualitative agreement with Biglari, et al. is provided by the observation that the magnetic fluctuations at ~ 200 kHz decrease during the transition, possibly indicating stabilization of turbulence. Increase of fluctuations and modulation of the inferred poloidal rotation during the ELMs is also in qualitative agreement with the model of Biglari, et al.
As indicated earlier, there is evidence for a change in transport in the plasma edge at the transition as provided by several results in addition to the increase in $T_e$ and $n_e$, including changes in magnetic fluctuations and the apparent onset of edge poloidal rotation. However, the onset of ELMs results in a modest degradation of confinement and increase in transport at the plasma edge. Transport calculations (using the TRANSP code) have been carried out for the discharge of Fig. 3. Characteristics of the transport for the three different phases of the discharge can be compared. These are the supershot, which includes the time up to the transition at ~4.39 sec, the time from 4.4 to 4.52 sec, which includes the transition and a short ELM free period, and finally the ELM phase which lasts for the remainder of the H-mode and the beam pulse. Microwave scattering measurements show that except for the outer 0.2 m, density fluctuations, $\delta n/n$, are very similar during the H-mode (including the ELM phase) to those for supershots across most of the plasma cross section. TRANSP calculations show little difference in $\chi_i$ and $\chi_e$ for the supershot and H-mode phases of the discharge of Fig. 3, and results of the calculation are plotted in Fig. 8a and 8b along with values for an L-mode plasma with comparable values of $I_p$, $B_\phi$, and $n_e$. The values for the supershot and H-mode are significantly lower than those for the L-mode plasma. During the transition and short ELM free period the thermal energy confinement time rises slightly, whereas confinement of thermal deuterons is not affected significantly during and after the transition. Stored energies, $W_e$ and $W_i$ in the outer quarter of the plasma minor cross section increase up to the onset of ELMs. $Z_{\text{eff}}$ begins to rise from a value of ~3.4 at the transition, and plateaus at a value of ~3.9 (15% increase) at 4.65 sec or about 150 ms into the ELM phase.

Conclusions

The circular limiter H-modes on TFTR have density profiles with peaking factors up to ~2.3. Bi-directional NBI power is required in order to obtain the transition from the supershot to the H-mode. Global energy confinement can be enhanced over Goldston L-mode scaling by a factor of
$ \geq 2.5$. The threshold power for the transition is directly dependent on $I_p$. Features in the spectra from microwave scattering and BES measurements show the onset of poloidal rotations at the transition, which persist for the duration of the H-mode. During the transition from supershot to H-mode, low frequency magnetic fluctuations increase and high frequency fluctuations decrease. The bifurcation layer and the ELMs are found to originate within the layer of plasma between 0.03 to 0.2 m from the plasma surface. There are no strong low-m modes before the ELMs, however, high frequency (50 - 500 kHz) precursor magnetic oscillations have been observed for some ELMs and

FIG. 8. Comparison of transport between a supershot, an L-mode plasma, and a limiter H-mode. (a) $\chi_e$, (b) $\chi_i$. 
these were found to be \( m = 1, n = 0 \) standing waves which were not outward ballooning in structure. This is very similar to the quasi-coherent mode observed on PDX using infrared scattering.\textsuperscript{16} TRANSP code calculations show the transport to be similar for supershots and limiter H-modes, but different for L-mode plasmas. The main difference between limiter H-modes and supershots on TFTR is the reduced electron transport at the plasma edge.

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DISCUSSION

K. ITOH: Thank you for referring to our previous theoretical H mode model. As you will find in IAEA-CN-53/D-IV-12 (Vol. 2), the refined model predicts that a large negative \( E_z \) is associated with the H mode (\( E_z \) is negative). We would like to hear in future about \( E_z \) measurement in TFTR.
STUDIES OF IMPROVED CONFINEMENT ON JFT-2M

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Abstract

STUDIES OF IMPROVED CONFINEMENT ON JFT-2M.

The improved discharges achieved in JFT-2M are H-mode, improved L-mode, counter neutral beam injection and pellet injected H-mode. These discharges are characterized by two improvements: one is the H-mode which has sharp density and temperature gradients at the edge, and the other mode has peaked density, temperature and toroidal rotation profiles around the center. The improvement of the pellet injected H-mode achieved at optimum conditions of the pellet injection experiments may consist in a coupling of H-mode and peaked profile improvements. The study of L/H transition by poloidal rotation measurements shows that a radial electric field of about -150 V/cm is always formed at the edge during the H-mode of single null divertor and limiter configurations. The neutral energy spectrum measured by the time of flight method shows a fast change of the energy distribution at L/H and H/L transitions. This change in the energy distribution occurs faster than the Ho/Do drop at the L/H transition. A steady state H-mode with Ho/Do bursts is realized by applying an ergodic magnetic field. This demonstrates active controllability of the H-mode.

1. INTRODUCTION AND EXPERIMENTAL SET-UP

Many kinds of improved discharge have been discovered in numerous tokamak experiments. Some of them have been realized or found in JFT-2M limiter or divertor discharges. These are H-mode[1,2], improved L-mode (IL-mode)[3], counter neutral beam injection (CTR-NB) and pellet injected H-mode[4]. The differences of these improved confinement modes are characterized by the density profile. In particular, to understand the H-mode, a

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FIG. 1. Density profiles of improved shots: (a) CTR-NB and CO-NB at limiter configuration; (b) improved L-mode and H-mode at He<sup>6</sup> beam injected single null configuration; (c) D<sub>2</sub> pellet injected H-mode compared with the cases of shallow penetrated pellet and without pellet.
multichannel charge exchange recombination spectroscopy (CXRS) system[5] and a time of flight (TOF) low energy neutral analyser[6] are newly installed. The CXRS system has two sets of toroidal arrays (34 channels) for ion temperature and toroidal rotation velocity measurements, and two sets of poloidal arrays (23 channels) for poloidal rotation velocity measurements with a spatial resolution of about 0.4 cm. The TOF analyser can measure neutral energies of less than 1000 eV/amu with a minimum time resolution of about 200 μs.

The H-mode found in JFT-2M is normally burst free; then, there is no steady state. To control and to study its physics, an ergodic magnetic limiter (EML) coil[7] and a divertor bias system are installed.

2. CHARACTERIZATION OF IMPROVED CONFINEMENT

As is well known, the so-called electron temperature ‘pedestal’ is quickly formed at the L/H transition. The electron density and ion temperature profiles also show a pedestal at the edge. The increase in the total stored energy in the H-mode is largely due to pedestal formation[8]. On the other hand, IL-mode or CTR-NB does not show electron density and temperature pedestals at the edge. The improvement of CTR-NB is due to the increase of density and ion temperature and the confinement time is about 30 % higher than that in CO-NB heating. The density (Fig.1(a)) and ion temperature profiles show highly peaked profiles. The toroidal rotation velocity has a high shear and the velocity of the center reaches about 100km/s[9]. In this case, the sawtooth period increases up to 40 msec or more and then disappears. After that, the plasma sometimes undergoes a disruption due to a large impurity accumulation at the center. The IL-mode appears after the H/L transition. This means that this mode appears after a fast decay of the edge density. Then this mode shows a peaked density profile (Fig.1(b)). The ion temperature and the toroidal rotation velocity also show peaked profiles, even in the case of CO-NB heating. In this mode, the sawtooth period increases and then disappears. If a large sawtooth occurs after H/L transition, then the discharge shows H/L/H/L.. transitions. The improved mode found in JFT-2M is characterized by two kinds of improvement. One is the H-mode which has a sharp density and temperature gradient at the edge, and the other mode has peaked density[10], temperature and toroidal rotation profiles around the center. The improvement of the pellet injected H-mode achieved at optimum conditions of the pellet injection experiments may consist in a coupling of H-mode and peaked profile improvements. The critical condition is deep penetration of the pellet injection. In this mode, the density profile shows a
peaked profile with a pedestal (Fig.1(c)) and the stored energy is about 30% higher than in a gas fueled H-mode. However, this coupled improvement does not persist for a long time. As a result of the large impurity accumulation at the center, the stored energy rolls over and decays faster than in a gas fueled H-mode.

3. L/H TRANSITION STUDY BY POLOIDAL ROTATION AND TOF MEASUREMENTS

Recently, a radial electric field ($E_r$) near the plasma periphery has been found, both experimentally and theoretically, to play an important role in the L/H transition[11-13]. The radial electric field profiles inferred from poloidal/toroidal rotation and ion pressure profiles using the ion momentum balance equation are shown in Fig.2(a)[14]. The peak value of -150 V/cm at the H-mode phase is mainly a contribution from a change in the poloidal rotation velocity of about 10 km/s. A jump in the poloidal rotation velocity at the L/H and H/L transitions is observed at about 0.4 cm inside the separatrix, while a significant change in the ion temperature is observed at about 1.3 cm, more inside of the separatrix. The poloidal rotation increases in the electron diamagnetic direction in the H-mode of limiter or divertor configuration (NB only and NB+ ECH heating), regardless of the direction of plasma current and neutral beam injection. The shear of the radial electric field and poloidal rotation velocity in the H-mode is localized within the order of an ion poloidal gyroradius near the separatrix, in the region of ion collisionality $v_{*i} = 20-40$. It is noted that the particle and thermal transport barriers are produced in different regions of the plasma. These measurements (Fig.2 (b) and (c)) seem to indicate that the improvement in thermal transport correlates to negative $\partial E_r/\partial r$ or negative $E_r$.

FIG. 2. (a) Radial profile of electric field as a function of distance from separatrix for L-mode (open symbols) and H-mode (closed symbols). Negative $dE_r/\partial r$ lies inside separatrix; (b) gradients of electron temperature measured by ECE (circles) and electric probe (squares) and of ion temperature (triangles) at L-mode (open symbols) and H-mode (closed symbols); (c) gradients of electron density measured by Thomson scattering (circles) and electric probe (squares) and of brightness of C VI emissions (triangles) at L-mode (open symbols) and H-mode (closed symbols); (d) neutral energy spectrum for L-phase (open symbols) and H-phase (closed symbols); (e) neutral energy spectrum obtained by applying diagnostic gas puffing just before TOF measurement for L-phase (open symbols) and H-phase (closed symbols); (f) average energy $\langle E \rangle$ from TOF measurement versus $H_d/D_\alpha$ intensity around L/H transition. Time from symbol to symbol is about 200 μs.
The TOF neutral energy spectrum measured from the top of the machine shows a fast change in its average energy \( \langle (dE/\Gamma dE) dE/((dE/\Gamma dE) dE) \rangle \) and integrated flux \( \langle (dE/\Gamma dE) dE \rangle \) at the L/H and H/L transitions of limiter or divertor configuration (NB only and NB+ECH heating). Figure 2(d) shows the energy spectra during the L-phase and just after the H-transition. Comparing these two spectra, we find that the number of neutral particles whose energy is lower than 400 eV is decreased by the H-transition but the number of those higher than 500 eV does not change. As a result, a flat part of the energy spectrum can be seen during H-mode phase in the energy range from 300 eV to 500 eV. The slopes of the low energy (E < 200 eV) and high energy (600 eV < E < 1000 eV) parts do not change much, but the decrease in low energy flux and the constancy of the high energy flux lead to the observed fast change in the average energy. The decrease in low energy part is mainly due to the decrease of neutral density after the H-transition. To understand the change of ion energy distribution, a diagnostic gas puffing is applied to keep neutral density constant at the TOF measured position. There is a little change of integrated flux at the L/H transition by the diagnostic gas puffing and the spectrum shows an increase in high energy flux from an energy of about 200 eV (Fig.2(e)). The particle energy of about 200 eV corresponds to a collisionality of one for the edge plasma parameter. The change of the energy spectrum is not understood clearly but, perhaps, may be explained by negative electric field effects on collisionless banana particles. In the case of the diagnostic gas puffing, neutral flux increases, then the average energy is evaluated at the minimum time resolution (200 \( \mu \)s) of the TOF measurement. A plot of the average energy versus the \( H_\alpha/D_\alpha \) intensity (Fig.2(f)) shows that the change of the energy spectrum is faster than the decrease of the \( H_\alpha/D_\alpha \) intensity at the L/H transition.

4. CONTROL OF H-MODE BY EML AND EFFECTS OF DIVERTOR BIAS TO H-MODE

We try to control the H-mode by applying an ergodic magnetic filed produced by two local coil sets installed outside the vacuum vessel. The magnetic field structure calculated by Fourier analysis and field line tracing shows a broad poloidal mode spectrum due to the locality of the EML coil. It has the peaks of the poloidal mode numbers m = 5 and m = 11 near the plasma surface for the low-m and high-m connections respectively. By applying a high-m ergodic magnetic field, the suppression in density, radiation increase and a steady state H-mode can be realized with \( H_\alpha/D_\alpha \) burst. The total stored energy of the steady state H-mode shows some loss compared
FIG. 3. Power threshold and controlled region of the H-mode by high-m ergodic field. Symbols show the discharge without $H_\alpha/D_\alpha$ burst (open circles), the intermediate discharge with and without $H_\alpha/D_\alpha$ burst (open triangles), the steady state H-mode with $H_\alpha/D_\alpha$ burst (closed circles) and the L-mode discharge (crosses).

with that of the burst free H-mode. But it is only less than 10% due to the suppression of density increase (the total stored energy of the steady state H-mode is almost the same as that of the burst free H-mode at the same density). The profiles of the core plasma density and temperature are almost the same with and without the ergodic field. Only the difference can be seen as the decrease of the gradient of edge temperature pedestal and the decrease of toroidal rotation velocity. Figure 3 shows the dependence of the power threshold and the steady region on the high-m mode of EML coil current. The power threshold increases with increasing EML coil current (maximum 5.5 kA = 44 kAT), and the controlled region of the steady state H-mode can be seen as a belt above its power threshold. The controlled region increases with increasing the ergodic field. This demonstrates an active controllability of the H-mode.

The diver bias experiment starts May in this year (1990). An interesting preliminary result of an unipolar divertor biasing is the decrease of the
power threshold for the H-mode. There is no bifurcation during Ohmic heating but a drop in the Hα/Dα intensity and a change in the rate of increasing of the density are observed at negative unipolar biasing which is applied between divertor plate and vacuum vessel. This negative unipolar biasing contributes to the L/H transition power almost linearly, when its voltage is applied. The power threshold is reduced by about 200 kW (from 340 kW to 140 kW in deuterium plasma) at the voltage of about -80V. The measurement of poloidal rotation velocity inside the separatrix does not change in the divertor biased L-phase. It is understood that, although the applied negative electric field is only to outside separatrix, it furthers the H-mode transition.

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DISCUSSION

G. FUSSMANN: In your radial electric field measurements it turns out that the minimum of $E_r$ coincides with the minimum of the derivative of the C VI emissivity. Is this pure coincidence? Also, were other impurity ions involved in your investigations?

Y. MIURA: That coincidence may be purely accidental. We measure the poloidal rotation velocity from the C VI line only; no other impurity ion lines were used.

A. HASSAM: You said that the scale size of $E_r$ variation is of the order of the poloidal ion gyroradius. In your paper, you pointed out that $v^*$ is large, 30–40. One would therefore not expect $\rho_p$ to be a natural scale.

Are you suggesting that $\rho_p$ is a natural scale to explain $|dE_r/dr|^{-1}$, and, if so, can you comment on why this would be so?

Y. MIURA: When we change the plasma current, the scale length of the poloidal rotation profile at the edge changes without any substantial change in temperature. This is why we think the scale length is close to the ion poloidal gyroradius.

R.J. GOLDSTON: How quickly does the poloidal rotation change at the L-H transition?

Y. MIURA: The time resolution of our two-dimensional detector is 16.67 ms. Thus we can say that the change of poloidal rotation takes place within 16.67 ms.
DIVERTOR BAFFLING AND BIASING EXPERIMENTS ON DIII-D*


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Abstract

DIVERTOR BAFFLING AND BIASING EXPERIMENTS ON DIII-D.

First results of divertor baffling and biasing experiments in Ohmic and low power beam-heated plasmas using the DIII-D Advanced Divertor are reported. The Advanced Divertor commissioned recently is designed to study baffling, plasma biasing, and dc helicity injection current drive. Using this unique baffle configuration in Ohmic and low powered beam-heated plasmas, the authors have obtained pressures under the baffle in excess of 5 mTorr and driven a poloidal scrape-off layer current in excess of 4 kA. Pressure under the baffle is strongly influenced by biasing. With the \( E \times B \) drift in the direction of the baffle throat, the pressure under the baffle can be a factor of five higher than with the opposite polarity. Biasing increases the level of divertor region magnetic fluctuations.

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INTRODUCTION

The DIII-D collaborative Advanced Divertor Program (ADP) was launched two years ago with the principal objective of experimental examination of several proposed concepts for density control, enhancement of plasma performance, and current drive. The knowledge gained from these experiments can be applied in the design of ITER, CIT, and other future devices. The activity during the past year was focused on the installation of a combined baffle and electrode structure, shown in Fig. 1, and the supporting diagnostics. The toroidally continuous biasing electrode is insulated from the vessel and the rest of the baffle structure. This bias ring is designed for bipolar operation with a power supply up to 1 kV supplying a maximum current of 20 kA. Under the baffle, space is allotted for the future installation of a cryopump with a pumping speed of 50 000 l/s, which matches the conductance of the baffle throat. More detail on the design of the ADP and the planned experiments can be found in Refs. [1-9].

In this paper we report a brief summary of the results from the first week of DIII-D operation after the installation of the ADP. The results reported here are very new and observational in character, but nevertheless of great interest owing to the unique nature of these experiments.

![Fig. 1. A view of the DIII-D advanced divertor showing the bias ring and the baffle plate. Also shown is a cross section view of a cryopump, planned for future installation.](image-url)
FIG. 2. Diagnostic traces from an H-mode beam-heated shot with baffling: (a) loop voltage; (b) plasma current and neutral beam power; (c) line average electron density \(10^{19} \text{ m}^{-3}\); (d) divertor \(H_\alpha/D_n\) intensity; (e) neutral gas pressure under the baffle; (f) radial position of the X-point; (g) and (h) intensities of oxygen lines, viewed horizontally through the magnetic axis of the core plasma. When the divertor strike point is in the baffle throat the neutral gas pressure under the baffle exceeds 5 mTorr.

For all the plasmas described here, shortly after the plasma initiation, the plasma was formed into a single-null divertor configuration with the outer divertor strike point \(~10\) cm inboard of the baffle throat, and maintained in that position until 3 seconds into the discharge, at which time the X-point position was moved radially outward until the outer divertor strike point was at the baffle throat or riding up onto the bias ring. The radial position of the X-point is displayed in Figs. 2 and 3, in the boxes labeled R\_X\_POINT.
BAFFLE EXPERIMENTS

Experiments were carried out in Ohmic and beam-heated L- and H-mode plasmas with $I_p = 1$ MA and 2.5 MW of neutral beam-injected power. In all three cases, the pressure under the baffle rose sharply as the strike point approached the divertor throat. Pressures in excess of 5 mTorr have been observed. Neutral pressure under the baffle is measured by an ASDEX-type fast pressure gauge [10], calibrated against a capacitance manometer. Because of the saturation of the gauge amplifier above ~5 mTorr, pressures presented here represent only a lower bound on the pressure under the baffle. The data from an H-mode discharge are displayed in Fig. 2. In this discharge, during baffling the neutral pressure rises to ~5 mTorr, where it appears that the pressure is modulated by ELMs. This pressure rise is sufficiently large to allow effective density control with pumping at a speed of 50 000 l/s.

A surprising observation is that baffling modifies the H-mode characteristics. Before the strike point is moved to the baffle throat, as the Hα trace shows, the H-mode phases are separated by L-phases. In contrast, in the baffling position the L-phases disappear. The reason for this change in the H-mode plasma behavior is not understood at this time. It may be simply due to the change in plasma geometry. One may speculate that the effect is owing to a reduction of edge impurities as indicated by the O-V line intensity. A third possibility is a reduction of charge exchange angular momentum damping caused by recycling neutrals following ELMs [11]. The baffle volume is sufficiently large to provide ballast volume for time scales $\leq 10$ ms.

BIASING EXPERIMENTS

A principal goal of the first biasing experiments was to demonstrate that current indeed flows along the field lines. The bias ring current is measured at the ring feed points. The poloidal distribution of the return current is measured by an array of current monitors placed under the divertor and inner wall graphite protection tiles [5]. Figure 3 shows biasing results from an Ohmic plasma where a potential of ~300 V with respect to the vacuum vessel is applied to the outer divertor leg. The bias ring current of approximately ~4000 A is nearly fully accounted for by the current flowing from the graphite tiles into the scrape-off layer plasma at the inner strike point. The sign convention is positive current flows into the plasma from the tiles. Although this assures that the current flows along the field lines, the spatial resolution of the detectors (~15 cm) is too coarse to determine what fraction of the current flowed the long way around the plasma.
Diagnostic traces from a negatively biased Ohmic plasma: (h) applied voltage to the bias ring; (i) power supply current; (j) poloidal current, measured at the inboard divertor strike point; (k) Mirnov oscillations, as measured by a magnetic probe near the outboard divertor stripe point. The drop in the line average density at t = 300 ms can be accounted for by the neutral particles trapped under the baffle. During biasing the level of magnetic fluctuations increased by a factor of 2–3. Negative biasing enhances pressure buildup under the baffle.

The I-V characteristic curve of a negatively biased plasma (outer leg is negative) is displayed in Fig. 4. The scrape-off layer (SOL) plasma current shows some sign of saturation around –300 V, but the slope increases at –400 V, indicating a change in the plasma conditions. It should be noted that the biasing power at this point (0.8 MW) is comparable to the Ohmic heating power (0.63 MW), and thus can significantly modify the SOL plasma parameters. For a given bias voltage, the SOL current of L-mode beam-heated plasmas is slightly greater than that of Ohmic plasmas and nearly a factor of two larger than that of H-mode plasmas.
Pressure under the baffle is strongly influenced by biasing. With the $E_r \times B_T$ drift at the baffle throat directed into the baffle throat the pressure under the baffle can be a factor of five higher than with this drift directed out of the baffle throat. Enhancement of the baffle pressure due to the $E_r \times B$ drift was one of the goals of the ADP.

A result of biasing in Ohmic, L-, and H-mode plasmas is a factor of 2–3 increase in the amplitude of the high frequency magnetic oscillations in the divertor. The bottom trace in Fig. 3 shows the signal from a magnetic probe near the outer divertor strike point. The level of fluctuation does not increase when the strike point is moved onto the insulated bias ring until the application of bias voltage. Therefore, the absence of line-tying can be ruled out as the cause of enhanced fluctuations.

**CONCLUSIONS**

Observation of a high neutral pressure under the baffle in excess of 5 mTorr in H-mode plasmas is of great importance to the DIII-D program as well as ITER. This pressure is sufficiently high to allow density
control with the anticipated cryopumping speed and thereby a decoupling of density and current for transport studies in DIII-D H-mode plasmas. Helium ash removal is a critical issue for ITER and other future long-pulse ignition devices. Our observations are very favorable in regard to being able to exhaust helium with a moderate pumping speed.

As a first step in the divertor bias experiments we have demonstrated that a significant poloidal current can be driven in the SOL of the plasma along the field lines. The observation of increased edge magnetic fluctuations is indicative of turbulence caused by the bias current. This turbulence could enhance helicity transport across the field lines, a necessary feature of dc helicity injection current drive. We have demonstrated that the pressure under the baffle can be significantly enhanced by negative biasing.

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DISCUSSION

R.J. TAYLOR: When the pressure increases in the baffled region, does it decrease anywhere else, or are you ‘sourcing’ particles from the wall?

M.A. MAHDAVI: The rise in the baffle pressure is not due to a local wall source, since at $p^0 = 5 \times 10^{-3}$ torr, and a baffle throat conductance of 50 000 L/s, the plasma density would be doubled within 0.1 s. In our experiment, density in fact decreases slightly.

R.J. TAYLOR: Can you now estimate how much helicity drive you get at 1 MW of input?

M.A. MAHDAVI: No.

D. POST: These results are very encouraging for particle exhaust in ITER. Do you plan to repeat these experiments and expand on them using trace helium ($\sim 5\%$) to look at the helium concentration in the central, edge and divertor plasma and in the pumping duct with high power H-mode plasmas? The results will be very important in connection with helium exhaust requirements for ITER, NET, FER and other long pulse burning plasmas.

M.A. MAHDAVI: Yes, ORNL participants in the DIII-D Advanced Divertor Programme have proposed installing the diagnostics required for this experiment.

K.H. DIPPEL: Do you have an estimate of the particle flux which enters the baffled divertor region, where the pressure is measured? If so, how does this flux compare with the total plasma outflow?

M.A. MAHDAVI: The conductance of the baffle throat is 50 000 L/s, and therefore, at a pressure of $5 \times 10^{-3}$ torr, the particle flux entering the baffled volume is 250 torr·L/s. The core plasma outflow is typically of the order of 50–100 torr·L/s.

P.M. BELLAN: You are now driving 20 kA of current with 1 MW, a power which is equivalent to the Ohmic heating power. If you drive all the 1 MA current by helicity injection, then this power will increase. It should go up by the ratio of increase in current, i.e. by 50, or perhaps by the current squared, a factor of 2500. Could you comment on this large increase in required power?

M.A. MAHDAVI: To drive a current of 1 MA in the core plasma, we estimate that a power supply current of $\sim 20$ kA would be required. An increase in the power supply current by a factor of five can be achieved by increasing the neutral beam injection power, which increases the ion saturation current at the bias ring. The Ohmic heating loss in the SOL in this case should be of the order of 5 MW.
Anomalous losses, needed for inward transport of helicity, need not exceed ~ 3 MW. It should be emphasized that the ~20 kA of toroidal current shown here is only scrape-off layer current; we have not yet studied current drive in the core plasma.

G. FUSSMANN: What experimental evidence do you have that the current is actually flowing the way you expect, and not the short way round in the private region?

M.A. MAHDAVI: We have no direct evidence that current flows the long way round the plasma. Langmuir probe data should answer this question. However, in the data presented here, the outboard separatrix strike point was on an insulating surface; therefore, current must flow across the field lines first, before flowing along the field lines the short way. Our data show no evidence of significant cross-field current flow.
MODELLING IMPURITY CONTROL BY PLASMA FLOWS IN THE JET PUMPED DIVERTOR

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Abstract

MODELLING IMPURITY CONTROL BY PLASMA FLOWS IN THE JET PUMPED DIVERTOR.

A pumped divertor is planned for the extension phase of JET (1993–1996) in order to control impurities during long pulse, high power operation. Target produced impurities tend to migrate along the field lines, out of the divertor, because of temperature gradient forces. This is opposed by the friction forces arising from the streaming of hydrogen into the divertor. This problem is studied by using both analytic and numerical methods. For the highest power anticipated, impurities will be trapped by ‘natural divertor recycling’ only at high scrape-off layer (SOL) densities approaching $10^{20} \text{ m}^{-3}$. At lower densities, a forced flow or ‘external recirculation’ must be induced, e.g. by shallow pellet injection or gas puffing, to ensure impurity retention. Steady state then requires removal of an equal amount of neutral flow from the divertor, which imposes fairly severe pumping requirements. Less flow is needed, for a given SOL density, as the ion power flow is reduced.

I. Introduction

The aim of the planned extension of JET (1993 - 1996) is to demonstrate effective methods of impurity control in operating conditions relevant to a Next Step Tokamak. The central task of this phase will be to investigate impurity control in a new divertor configuration, the primary functions of which are to remove the principal source of impurities as far from the main plasma as possible, and to retain the impurities produced at the divertor target plates. The production of impurities will be minimized by reducing the divertor plasma temperature as far as possible, and by proper selection of target plate materials. Retention of the impurities in the divertor region will be enhanced by inducing a strong flow of deuterium towards the target plates.

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The general features of the proposed divertor are illustrated in Fig. 1. It is of the "open" type, with beryllium clad, water cooled (Hypervapotron) copper target plates. The in-vessel four-coil system allows both for horizontal sweeping of the plasma along the target to spread the heat load, and for vertical motion of the X-point to vary the connection length and plasma volume. The possible divertor configurations span the range from a "slim" 6 MA plasma (total divertor coil current of 1.8 MA) with a connection length in the divertor region, corresponding to a field line 1 cm away from the separatrix in the equatorial plane, of about 6.7 m, to a "fat" 6 MA plasma (divertor current of 0.74 MA) with a connection length 3.5 m.

Impurities produced at the target plates are subject to several forces, the most important of which, in steady state, are the ion thermal force, directed away from the target plates, the deuterium-impurity friction force towards the target, and the impurity pressure gradient force which arises to establish a force balance. The ion thermal force is determined principally by the ion heat flow in the scrape-off layer (SOL), while the counter-acting friction force depends on the magnitude and spatial distribution along the field (coordinate s) of the deuterium particle flux, Γ(s), and is inversely proportional
FIG. 2. Sketch of behaviour of hydrogenic flux to the target (above) and resulting impurity distribution (below) for various recirculation patterns. $S$ is the distance along a field line from the target, while $S_0$, $S_x$, and $L$ refer to effective ionization length, connection length and stagnation point locations, respectively.

to the ion temperature to the $3/2$ power. Thus the friction force can in principle be set at a level sufficient to overcome the thermal force. From global balance considerations of the SOL and the divertor channel plasma (DCP), the magnitude of the flux at the divertor target plate, $\Gamma_t$, is determined by the mid-plane SOL plasma density and the power flow in the SOL. The form of $\Gamma(s)$, denoted $\Gamma(s) = \Gamma(s)/\Gamma_t$, is controlled by the hydrogen particle source distribution in the SOL and the DCP as sketched in Fig. 2. In the "local recycling" zone, which extends approximately one effective ionization length normal to the target, $\Gamma(s)$ increases rapidly going towards the plate, due to ionization of neutrals emanating from the target plates. Some of the neutrals, however, are not locally recycled, but escape and then may re-enter the DCP further up the flux tube via reflection from the side walls, or transmission across the private flux region. We refer to this as "distant recycling". It extends the
region of significant hydrogen flux further from the target plates, thus enhancing the effectiveness of the friction force term. Case a of Fig. 2 corresponds to purely local recycling at the target plates, while Case b corresponds to roughly half of the target plate flux being uniformly "distantly recycled" between the local recycling zone and the X-point. The relative amount and distribution of the distantly recycled particles depend both on divertor plasma conditions and on divertor geometry.

It is also possible, in principle, to extract some of the incident neutralized ion flux from the targets and directly "recirculate" it, e.g. by baffles, into the X-point region, as was suggested in Ref. [1]. This would produce a \( \hat{\Gamma}(s) \) distribution as shown in Case c, Fig. 2. Finally it may be necessary in some cases, for impurity retention, to induce a moderate \( \hat{\Gamma}(s) \) in the SOL, in excess of the natural flow from the main plasma, by strong gas puffing or shallow pellet injection. This imposed flow, which we term "external recirculation", would require pumping of an equivalent neutral flux from the divertor chamber to maintain a steady state. The resulting flux distribution is shown as Case d.

The purpose of this paper is to determine the conditions under which the recycling (local plus distant) is sufficient to retain target-produced impurities in the divertor, and to calculate the amount of additional "recirculation" flow required when the "natural" flow is insufficient. We begin (Section 2) with a simple analytic model which elucidates the most important physical processes. This is followed by a discussion of the full 1-1/2 code and its results (Sect. 3). Conclusions are drawn in Section 4.

II. Impurity Retention Overview

Impurity retention in the divertor may be estimated by examining the magnitude of the different forces on the impurity ions as given by the steady state momentum equation
Here we have considered a single impurity species of charge state $Z$, mass $m_z$, density $n_z$, temperature $T_z$, pressure $p_z$, and flow speed $v_z$. The simplification to a single charge state is well justified, because the most important forces acting on the impurity all scale approximately as $Z^2$, and thus their ratio is independent of $Z$. The subscripts $e$ and $i$ refer to the electrons and the hydrogenic ions, and $\tau_{zi}$ is the hydrogenic ion-impurity collision time. The forces on the right hand side of Eq. (1) are due to the impurity pressure gradient, where $p_z = n_z T_z$, the electric field $E$, the frictional force, and the thermal forces with coefficients $\alpha_z$ for electrons and $\beta_z$ for ions.

This equation, along with the other appropriate conservation equations, are solved in detail using a 1-1/2 code, as discussed in Section 3. To gain insight into those results, however, we begin with a simplified analytic description of the impurity distribution, as deduced from Eq. (1) in the trace approximation, where the hydrogenic flow is specified $a$ priori, and is not affected by the impurities.

The electric field may be eliminated from Eq. (1) by using the electron momentum equation and taking the electron pressure to be nearly constant along $s$. In addition, we take $\frac{dp_z}{ds} = 0$ and $T_z = T_i$, so that Eq. (1) becomes

$$m_z n_z v_z \frac{dv_z}{ds} = -\frac{dp_z}{ds} + n_z Ze E + \frac{m_z n_z (v_i - v_z)}{\tau_{zi}} + n_z \alpha_z \frac{dT_e}{ds} + n_z \beta_z \frac{dT_i}{ds} \quad (1)$$

$$T_i \frac{dn_z}{n_z ds} = \frac{m_z (v_i - v_z)}{\tau_{zi}} + (\alpha_z - 0.71 Z) \frac{dT_e}{ds} + (\beta_z - 1) \frac{dT_i}{ds} \quad (2)$$

with $\alpha_z = 0.71 Z^2$, and $\beta_z = 1.76 Z^2$ for the case of Beryllium with $Z = 3$ [2].

For most cases of interest $\nabla T_e \ll \nabla T_i$, since they are both determined by classical heat conduction and $\kappa_{||e} \gg \kappa_{||i}$, so that we can neglect the $\nabla T_e$ term in Eq (2). In addition, we take $v_z = f_v v_i$ as a
further simplification, where $f_v = 0$ everywhere except in a thin impurity ionization zone near the target. Integrating the resulting equation gives

$$\frac{n_z(s)}{n_z(0)} = \exp\left\{-\int_0^s F_{fi} \, ds + \int_0^s F_{Ti} \, ds\right\}$$

(3)

with

$$F_{fi}(s) = C_f \cdot Z^2 \frac{\Gamma(s)}{T_i(s)^{5/2}} \, m^{-1}$$

(4)

and

$$F_{Ti}(s) = (\beta_z - 1) \frac{1}{T_i} \frac{dT_i}{ds} \, m^{-1}$$

(5)

Note that $F_{fi}$ and $F_{Ti}$ are inverse local scale lengths associated with the friction and thermal forces, respectively. Here and in the following, the units are MKS, except for $T_i$, which is expressed in eV. The value of $C_f$ is then $1.2 \times 10^{-20}$.

For a large enough deuterium flux, $\Gamma$, the impurity density decreases going towards the X-point until the point where $\Gamma(s)$, determined by the distribution of recirculated neutrals, becomes small enough that $F_{fi} < F_{Ti}$. Beyond that point, $n_z(s)$ begins to increase. Fig. 2b shows qualitatively the behaviour of the solution for the four source distributions of Fig. 2a. If we adopt the criterion that $n_z(L)/n_z(0)$ be less than a certain value, say $e^{-3}$, the criterion for retention of impurities becomes

$$\int_0^L F_{fi}(s) \, ds \geq \int_0^L F_{Ti}(s) \, ds + 3$$

(6)

In order to carry out the integrals, the hydrogenic ion temperature and particle flux distributions along $s$ must be specified. The latter is postulated, and then checked by a Monte-Carlo calculation, as described in section 3. The temperature is governed by both conduction and convection and is given, to close approximation, by

$$T_i(s)^{7/2} = \frac{7}{2} \left(\frac{q_{\text{eff}}}{K_{o,i}}\right) s + T_{di}^{7/2}$$

(7)
where \( q_{\text{eff}} = 0.6 q_i \), the ion heat flow per unit area, for values of recirculated \( \Gamma \) near the minimum required to entrench the impurities [3], \( T_{\text{di}} \) is the divertor ion temperature, and \( K_{0,i} \) is a constant.

Combining equations (4), (5), and (7) shows that the ratio of the ion friction and thermo-electric terms scales approximately as

\[
\frac{F_{\text{fi}}(s)}{F_{\text{Tf}}(s)} \sim \Gamma(s) \cdot (s + s_o)^{2/7}
\]

where

\[
s_o = \frac{2}{7} \left( \frac{K_{0,i}}{q_{\text{eff}}} \right)^7 T_{\text{di}}^{7/2}
\]

Consider the case \( \Gamma(s) \sim \text{constant} \) over some interval, for example Case c of Fig. 2. Because of the relatively weak variation of \( F_{\text{fi}}/F_{\text{Tf}} \) with \( s \), it can be shifted from less than one to greater than one over most of the interval by a small change in \( \Gamma \). Thus, one can go from a case of very poor impurity retention to good retention with small changes in \( \Gamma \) which extend over long distances, e.g. by distant recycling and/or external recirculation.

Incorporating equations (4), (5) and (7) into (6) yields, as the criterion for retention of impurities,

\[
\frac{C_f(1 - f_o)Z^2\Gamma_t}{\left( \frac{7}{2} q_{\text{eff}} \right)^{3/2}} \int_0^L \frac{\bar{f}(s)ds}{(s + s_o)^{3/2}} \geq \left( \frac{2}{7} \right) \left( \beta - 1 \right) \ln \left( \frac{L}{s_o} + 1 \right) + 3
\]

The target flux, \( \Gamma_t \), is determined by specifying the mid-plane or, more generally, the stagnation point separatrix plasma density \( n_b \), and the parallel heat fluxes, \( q_{||e} \) and \( q_{||i} \), using a model similar to that of Ref. [4]. For given power into the divertor and very high values of \( n_b \), \( \Gamma_t \) is sufficiently high that local plus distant recycling alone is sufficient to satisfy relation (9) for the proposed divertor geometry. As \( n_b \) is decreased for a given power, \( \Gamma_t \) decreases and \( T_{\text{di}} \) increases,
which tends to reduce the left hand side of relation (9). In order to retain impurities, \( \hat{\Gamma}(s) \) must be "extended" further up the field line, either by (single point) recirculation, e.g. at the X-point (Case c), or by injection of particles uniformly into the SOL (Case d).

For X-point injection, the flux required for impurity retention is

\[
\Gamma_x \geq \frac{2}{7} \left[ \frac{7 q_{\text{eff}}}{2 K_{o,i}} \right]^{5/7} \left[ \frac{2}{7} (\beta_2 - 1) \ln \left( 1 + \frac{L}{s_0} \right) + 3 \right] \left[ \left( s_x + s_o \right)^{2/7} - s_o^{2/7} \right]
\]

The relation (10) assumes \( \Gamma_x \) is constant along \( s \) and thus does not include the contribution of the local high recycling region, which is small for \( \Gamma_x \) non-negligible compared with \( \Gamma_t \).

Relation (10) implies that the minimum required flux can be written

\[
\Gamma_x \sim q_{\text{eff}}^{5/7} f(q_{\text{eff}}, T_{di}, s_x, L)
\]

where the variation of \( f \) with its arguments is small. For example, a change of \( q_{\text{eff}} \) by a factor of 10 changes \( f \) by only 5%. Furthermore, a factor of 4 variation in \( T_{di} \) changes \( f \) by 20%, while doubling the distance to the injection point, \( s_x \), from 4m to 8m reduces \( f \) by ~25%.

Thus, as the main result, this simple analysis predicts that for the case of single point recirculation at the X-point, the required \( \Gamma_x \) depends primarily on the ion heat flow and somewhat less strongly on the connection length. \( \Gamma_x \) scales as \( \Gamma_t^{5/7} \cdot s_x^{-2/7} \) for \( s_x^{2/7} >> s_o^{2/7} \).

We note that the electron temperature does not enter this analysis directly since the dominant forces are associated with the ions. However, the electron temperature and density in the divertor are important in that they determine the fraction of neutrals that is distantly recycled, rather than locally, and thus they control the amount of "recirculation", which must be added to retain impurities.
The integrals in Eq. (3) can be evaluated for simple \( \hat{\Gamma}(s) \) other than the step function form corresponding to relation (10), without changing the basic scaling appreciably. For example, for uniform injection into the DCP between \( s=0 \) and \( s_x \), corresponding approximately to the expected distant recycling pattern at low SOL densities, the required flow is increased by 35% relative to X-point injection, while uniform injection into the SOL between \( s_x \) and \( L \) (external recirculation) reduces \( \Gamma_x \) by about 23%.

III. Code Calculations of Impurity Retention

a) The Model
To study the requirements for effective impurity control in more detail a 1-1/2D model (EDGE1D) of the plasma boundary in the proposed pumped divertor configuration has been developed. (For a detailed description of the model see Ref. [5].) The model is based on the plasma model of the JET 2D boundary code EDGE2D [6] and the impurity model of Ref. [7]. The latter is valid for arbitrarily high impurity concentrations.

Fluid equations for the conservation of particles, momentum and energy along the magnetic field are solved for electrons, hydrogenic ions and impurity ions. Transport coefficients, friction, thermal forces and electric fields (not necessarily ambipolar) are classical [8] and allow for arbitrarily high impurity concentrations [7]. In general, the full non-coronal distribution of impurity charge states and the corresponding radiated energy losses are determined. A single impurity temperature, set equal to the hydrogen temperature, is assumed. The electron density is evaluated from quasi-neutrality.

For the boundary conditions at the target plates the Bohm condition is adopted, i.e. it is assumed that hydrogenic and impurity ions reach their sound speed at the plates. The heat fluxes at the
plates are prescribed with the appropriate energy transmission coefficients for a two component plasma.

Transport transverse to the magnetic field is replaced by specified temperature and density profiles assumed to decay exponentially outwards from the separatrix with decay lengths, $\lambda_T = 1.5 \text{ cm}$ and $\lambda_n = 1 \text{ cm}$, respectively. The width, $\delta = 3 \lambda_T$, of the SOL plasma is chosen such that the residual fluxes are negligible.

A full 2-D Monte Carlo code (NIMBUS) [9] is retained for the sources of neutral particles. Deuterium particles are recycled at the target plates as neutrals. The Monte Carlo code defines the spatial distribution of the ionisation sources in the vicinity of the target plates. To enhance the plasma flow in the divertor region, some of the recycled neutrals are removed and recirculated into the scrape-off layer as plasma ions, with prescribed spatial distribution along the separatrix.

The Monte Carlo code defines also the spatial distribution of the impurity sources due to erosion at the target plates. Impurity atoms are generated by sputtering due to plasma ions and neutrals of any species. An effective sputtering coefficient is taken into account following Ref. [10]. However, it is also possible to normalise the source of impurities to their outgoing flux onto the target plates, thus implying a prescribed impurity content.

Work is currently in progress to implement a full 2-D version of the plasma flow, including impurities at arbitrary concentrations.

b) Numerical Results
As an example, we now discuss results of calculations based on the "slim" divertor configuration. Figure 3 shows the corresponding computational mesh and the most important input parameters. As a "reference" case we choose a total input power of 40 MW, corresponding to the maximum available power in JET. This power is shared equally between the ions and electrons ($P_e = P_i$) and is
input uniformly along the separatrix into the SOL. The calculations are for beryllium impurities which are produced self-consistently by sputtering from the target plates. The magnitude of the externally recirculated flux is varied parametrically and, in most cases, is injected uniformly along the separatrix from the main plasma into the SOL. For given \( n_b, P_i, P_e \) and external recirculation fraction \( f_R \) of the target flux, all other SOL and DCP parameters are determined.

Figure 4 shows calculated profiles of the deuterium and total impurity densities, the ion and electron temperatures, and the principal forces acting on the impurity ions as functions of the distance along the field line from the outside target for runs typical

FIG. 3. Computational mesh and some input parameters used for EDGE1D code calculations.
FIG. 4. Behaviour of various plasma parameters versus $s$, for a case of poor impurity retention (left) and good impurity retention (right). $f_R$ is the fraction of target flux which is externally recirculated.
of poor (run 724) and good (run 712) impurity retention. It can be seen that for run 724, a large fraction of the impurities accumulate in the SOL adjacent to the main plasma. To quantify the degree of impurity retention we use the ratio of the number of beryllium ions retained in the two divertor regions to the total number in the SOL and DCP; we call this ratio $\eta(\%)$. The two cases displayed in Fig. 4 differ mainly in the level of external recirculation, and only slightly in the separatrix stagnation point density ($n_b = 8.3$ and $9.1 \times 10^{19} \text{m}^{-3}$, respectively). For run 724 there was no external recirculation while for run 712 10% of the target flux was extracted and reinjected uniformly into the SOL. It can be seen that this relatively small external recirculation makes a major change in the retention, increasing $\eta$ from 16% to 92%.

The second set of frames for Fig. 4 shows the electron and ion temperatures. In all our simulations the ion temperature in the SOL is higher than that of the electrons, mainly due to the lower ion thermal conductivity along field lines. The increased ion temperature near the stagnation point in run 724 arises from the high impurity density there, and the dependence of thermal conductivity on the effective charge state. The prediction that the ion temperature in the SOL should exceed the electron temperature is in agreement with experimental observations in JET [11]. Also the temperatures near the target plates are in the 10 - 20 eV range at higher edge densities and correspondingly larger hydrogen fluxes to the targets.

The third frame of Fig. 4 shows the force balance. For $f_R = 0$, the distribution of $\Gamma$, which arises from recycling alone, does not provide a large enough integrated friction force to retain the impurities, and they accumulate near the point of maximum temperature, with the impurity pressure gradient force balancing the ion thermal force. For $f_R = 0.1$, the external recirculation extends the range of the friction force sufficiently to enable it to overcome the thermal force all the way to the SOL, and the impurities are
effectively retained in the DCP. The fourth frame shows the
distribution of fluxes, on an enlarged scale, for the two cases. It can be
seen that the change in the level of the flux is small compared to the
target fluxes (which are off scale and correspond to 7.0 and 6.0 \times 10^{24}\text{m}^{-2}\text{s}^{-1}, respectively), but nevertheless the impurity distribution
is completely altered.

Relation (9), which gives the criterion for impurity retention,
can be satisfied in either of two ways. One can increase the target flux
\(\Gamma_t\) to a large value, which occurs automatically at very high SOL
densities. Alternatively, when operating at moderate SOL densities,
and hence moderate target fluxes, the integral can be increased by extending the range of $\hat{\Gamma}$ by the combined effects of distant recycling and, when necessary, external recirculation. Figure 5 shows the distribution of flux for four different code runs, corresponding to different target fluxes and distributions. Each of these runs produced retention greater than 90%, achieved by different combinations of target flux and external recirculation, as indicated on the figure.

FIG. 6. External recirculation particle flux fraction $f_R$ (a) and particle flux $\Gamma_R$ (b) required for 90% impurity retention versus SOL density, for three values of ion power flow in the SOL.
Figure 5b displays the normalized flux, $\hat{\Gamma} = \Gamma(s)/\Gamma_t$, which shows its "extension", and hence the extension of the friction force, to larger values of $s$ by the use of uniform external recirculation in the SOL. Figure 5a shows the non-normalized flux, $\Gamma(s)$, and indicates that the values of the actual local $\Gamma(s)$ over much of the divertor plasma (0 to 8m) are roughly similar, in agreement with the predictions of the analytic model.

Figure 6 summarizes the results of an extensive set of runs. Fig. 6a shows the fraction, $f_R$, of the target flux $\Gamma_t$ which must be externally recirculated in order to retain 90% of the impurities in the DCP versus SOL density, $n_b$, for three values of the total (to both targets) ion heat flow, 20 MW, 10 MW and 5 MW. A simple SOL model (e.g. that of Ref. [4]) shows that $\Gamma_t \sim n_b^2$ for $T_d > 15$ ev. Thus, at high values of $n_b$, $\Gamma_t$ is very large and the friction force integral arising from the natural recycling which occurs is sufficient to retain impurities. As the edge density is decreased, $\Gamma_t$ decreases and an increasingly large flow fraction $f_R$ must be externally recirculated.

![Graph showing integral of ionization versus computational mesh point. The distant recycling fraction increases as the SOL power decreases.](image)
The required recirculated flow itself, however, becomes relatively independent of density for low and medium $n_b$ (Fig. 6b), consistent with Eqs. (10) and (11). Under these conditions the required externally recirculated flow is fairly large and is probably not compatible with achievable steady state pumping rates.

The solid curves of Fig. 6 are results obtained by using a self-consistent beryllium sputtering model. In those simulations, the ratio of the total number of beryllium ions to deuterium ions in the edge plasma decreases from 1.5% at the lower densities to 0.1% at $n_b \equiv 1 \times 10^{20} \text{m}^{-3}$. Since these values are somewhat lower than those which have been observed in JET experiments with beryllium to date, we have repeated the simulations with a fixed relative impurity content of 1.5% and the results are shown as dashed lines in Fig. 6a. In this case, the required fractions $f_R$ are increased slightly.

The situation is somewhat improved for reduced power flows, as can be seen from the graphs in Fig. 6. Lower values of $f_R$ are required for a given SOL density, $n_b$. This arises in part because the divertor density falls substantially. This increases the distant recycling in the divertor and reduces the flow which must be provided by external recirculation. This effect is illustrated in Fig. 7, which shows the fraction of neutrals, which emanate from the plate and are reionized in the DCP, as a function of computational mesh point (as shown in Fig. 3). For high power, most of the ionization occurs very near the target, whereas at lower power the effective range of ionization extends much further along the DCP. The increased effective ionization range also increases the fraction of neutrals which are pumped by the cryopump.

Preliminary results from the code indicate that the flow requirements for the "fat" magnetic configuration are relatively unchanged from those of the "slim" magnetic configuration described above.
IV. Conclusions

The entrenchment of impurities in a forced flow of plasma towards the target plates is a candidate concept for impurity control in Next Step Tokamaks and is the focus of the Pumped Divertor planned for JET. Analysis shows that the effectiveness of the friction force in overcoming the ion thermal force is strongly reduced at high ion temperature in the SOL, which occurs at high ion power flows ($P_i$). The hydrogenic flow required for effective impurity retention increases approximately as $P_i^{5/7}$. At high SOL densities, recycling at the target plates is very high, and this "natural" recycling flux is sufficient to retain the impurities effectively. At moderate and low SOL densities, however, the "natural" recycling flux must be augmented by "external recirculation" to retain impurities. The fraction of the target flux needed for impurity retention increases as density decreases, but decreases with decreasing ion power flow. Pumping the required external recirculation flows is likely to be possible only at rather high edge densities, approaching $n_b \sim 10^{20}$ m$^{-3}$. The situation improves at lower $P_i$. Operation with clean plasmas is likely, therefore, only at high densities. This may have severe consequences, e.g. for current drive.

First steps towards validating these complex SOL models have been taken by establishing consistency with observed edge data in JET X-point operation [11]. In particular, the predicted conditions $T_i > T_e$ in the SOL, and divertor temperatures $T_{ed}$ in the range 10 - 20 eV at high $n_b$, have been observed.

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References


DISCUSSION

K.H. BURRELL: The results of your modelling depend critically on the magnitude and functional form of the thermal force. What experimental evidence do we have that the theoretical form for this term is correct?

M. KEILHACKER: JET experiments with gas puffing into the scrape-off layer (SOL) or the private flux region strongly support the view that thermal forces play an important role in impurity transport along open magnetic field lines. A more detailed assessment of the magnitude and functional form of this force will only be possible when the planned pumped divertor experiment is carried out.
P. GHENDRIH: Radiation is obviously a key issue in this model because it will increase the thermal force, but it is also necessary to radiate the power. Can you comment on this point?

M. KEILHACKER: In our model, impurities are generated self-consistently by sputtering at the target plates and their radiation is then calculated on the basis of a full non-coronal model. For the case of beryllium impurities considered here, the resulting radiation in the SOL and divertor plasma amounts to only 2 or 3 MW. We have artificially increased the radiation either by using a multiplier to the self-consistently calculated radiation or by applying a constant radiation efficiency. The results depend very much on the density and, in particular, the temperature range; but, in general, any increase in thermal force caused by reducing the temperature is largely offset by a corresponding increase in the friction force.

R.J. TAYLOR: Is your design aimed at working in the L- or H-mode or both? And how do you maintain the required high density in the SOL in the H-mode?

M. KEILHACKER: Our design is intended to work in both L- and H-mode operations. The maximum density in the SOL at which one can operate is set by the density limit, which seems to be an edge phenomenon and increases with the power flow into the SOL (for both L- and H-modes).

R.J. HAWRYLUK: Dr. Rebut pointed out in his paper the importance of impurity control as a means of extending the operating range of JET. He noted that up to 35 MW of power had been used to heat JET and that the highest neutron fluxes had been achieved at low density (~3 \times 10^{19} m^{-3}). In addition, lower hybrid RF power is being installed on JET. It appears that your study has identified a potential discrepancy between the JET requirements and the operating range in which the new divertor configurations will suppress impurities. Could you comment on this?

M. KEILHACKER: It is true that our studies indicate that impurity control at high ion power flow requires rather high edge densities and relatively flat density profiles in order to avoid large ‘external recirculation’ flows. This applies to JET as much as to any NEXT STEP device. Since in the end we do not want to work in the ‘hot ion mode’ regime but rather at $T_i \approx T_e$, this is not so much a restriction as the only way to reach high performance. In addition, in JET pumping the external recirculation flow will be helped by the strong pumping action of the beryllium walls.

S.A. COHEN: Good impurity retention in the divertor can be augmented by inclusion of a magnetic sheath in the analysis. Does your model include a magnetic sheath and if so, what fraction of your ‘good retention’ is due to the magnetic sheath and not the $\nabla T_i$ force?

M. KEILHACKER: The ‘magnetic sheath’ is not included in the code calculations. The effect of such a sheath is to extend the range of the electric field to approximately an ion gyroradius. Booth (Phys. Fluids 82 (1990) 1858) has shown that tungsten is ionized within the magnetic sheath, and thus more effectively retained, but that the effect is relatively unimportant for carbon, owing to its longer ionization mean free path. We therefore expect the effect to be relatively unimportant for beryllium, too.
EDGE AND DIVERTOR STUDIES ON ASDEX


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Abstract

EDGE AND DIVERTOR STUDIES ON ASDEX.

Recent edge and divertor studies on ASDEX, mostly under boronized wall conditions, are presented. Edge scalings derived from an experimental edge data base turn out to be consistent with classical transport along field lines. The existence of an edge density limit, of thermo-electric currents in the scrape-off layer, and of toroidal and poloidal asymmetries of the power deposition on the target plates is shown. A strong reduction of the divertor impurity retention with decreasing safety factor q is found. Specific results are presented for the quasi-stationary H-mode with frequent edge-localized modes (ELMs): while the copper sputtering at the target plates increases dramatically during an ELM, the average copper content in the bulk plasma remains low and stationary because of the combined effect of divertor retention and bulk plasma screening by the ELMs. The power deposition at the target plates shows a moderate toroidal asymmetry and the in–out asymmetry is smaller than in L-mode. Roughly half of the total deposited power is carried by the ELMs.

1. Introduction

Extensive edge and divertor investigations have been done in ASDEX over the years in practically all accessible regimes, including various external heating methods, several modifications of the divertor chambers, and different wall conditioning (e.g. [1-5]). Boronization in parallel with a stepwise, mechanical reduction of the divertor gas leakage into the main chamber has recently further improved the plasma performance [6]. The results reported below are mostly obtained under these favorable conditions. In the following we give first an overview of recent edge investigations in ohmic and L-mode discharges. In the second part we present more specific edge and divertor measurements for the quasistationary H-mode with regular ELMs in view of the possible relevance of this mode of operation for ITER.

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2. Recent edge and divertor investigations in the ohmic and L-mode regimes

An edge data base has been established recently which, together with statistical methods and analytical and numerical edge modelling, has significantly contributed to the understanding of the physics of the edge region as well as its influence on the global plasma behavior [3]. The edge scalings obtained turn out to be consistent with standard high recycling divertor scrape-off layer models assuming classical transport along field lines, provided that the power flow into the scrape-off layer and to the target plates is correctly taken into account, i.e. that radiation losses are properly accounted for. A corrected separatrix position in the midplane of about 1 to 1.5 cm (roughly one temperature decay length) outside the magnetically determined one was used for comparison of midplane and divertor data [7,8]. This separatrix position was inferred from the X-ray emission of fast electrons hitting material targets, which were mounted onto the midplane reciprocating Langmuir manipulator. The result confirms an earlier measurement, where the reciprocating Langmuir probe itself was taken as the target [2].

An edge density limit, probably correlated with the global density limit, has been predicted on the basis of the mentioned scrape-off layer models and has been qualitatively confirmed in ASDEX. This aspect has been further clarified under boronized wall conditions [9] and by use of the edge data base [10]. Plasma radiation under these conditions is, up to the maximum densities, only a small fraction of the input power and does therefore not directly cause the density limit disruption. At medium values of the safety factor q the density limit is correlated with a low divertor temperature (= 5 eV) in accordance with the edge density limit model. For q \leq 2.4, the disruption occurs at higher divertor temperature (at the location of the Langmuir probe) and the power deposition onto the target plates shows significant toroidal asymmetry. Field line tracings based on an axisymmetric equilibrium with superposed external perturbation fields [2] shows a large, partly ergodized \( m=2 \) island moving towards the more and more distorted separatrix with decreasing q. For q = 2 the \( m = 2 \) island still exists and is located right at the edge with practically no closed flux surface outside of it. This numerical result, though not self-consistent, may qualitatively explain the toroidal asymmetry of the time integrated power deposition. Because of the toroidal variation, a low divertor temperature may be again obtained at the instant of the disruption away from the toroidal position of the divertor Langmuir probe, which is located near the power deposition maximum.

The reciprocating divertor Langmuir probe can be positioned in close proximity to the target plate. Hence a probe tip electrically connected to the target plate should collect about the same current as that flowing to the target itself in the absence of the probe. Currents up to a substantial fraction of the ion saturation current are measured. These are interpreted as thermo-electric currents driven by the poloidally asymmetric power flow to the target plates as measured by calorimetry [11]. One should notice that within this model the toroidal asymmetry discussed later on indicates the possibility of a helical thermo-current pattern in the edge. Such edge currents provide an additional
FIG. 1. Argon XVI emission from the bulk plasma as function of time for ohmic discharges with safety factor q as parameter. A short argon pulse is puffed into the main chamber at $t_0 = 0.8$ s. A strong change occurs for $q < 3$. The curves are fitted by an exponential with two time constants according to a simple transport and divertor pumping model.

coupling mechanism between the divertor plasma and the main chamber edge. At the L/H transition a significant change of current profile dependent on magnetic field polarity is observed.

The divertor retention of inert gases as well as of target produced impurities has been investigated on ASDEX. Generally, a strong reduction of the retention of e.g. argon at the onset of auxiliary heating is found with moderate differences for the various modifications of the divertor configuration. The retention depends also significantly on the magnetic field or the safety factor $q$ ($I_p=\text{const}$) as seen in fig. 1, where the temporal decay of argon XVI is shown as function of $q$. The measurements have been compared with a transport and divertor pump model. Particularly for $q<3$ a drastic reduction of the divertor confinement together with a pronounced rise of the effective confinement time in the main chamber occurs. A classical diffusive impurity transport model along field lines (showing a strong dependence on charge state, ion temperature distribution etc.) may qualitatively explain some of the observations.

Poloidal variations of edge plasma parameters and power deposition onto target plates (measured by calorimetry) dependent on the configuration and confinement regime is to be expected in an axisymmetric tokamak and is clearly observed in ASDEX (see also sect.3). But even toroidal asymmetry of the power deposited onto target plates is a quite general feature in ASDEX, especially for $q<3$, where low rational surfaces and islands approach the
scrape-off layer. The strongest toroidal asymmetry is observed with Lower Hybrid Heating, where up to 50% of the input power is found in the divertor with a toroidal variation of up to a factor of five and with quite a narrow deposition peak. It is attributed to localized heat deposition by fast electron loss along perturbed magnetic field lines [12]. In contrast, a nearly sinusoidal $n=1$ toroidal variation is obtained when locked $m=2$, $n=1$ modes are present and extend out to the edge. The spatial phase of power deposition is clearly correlated with the magnetically determined one, which is always the same, i.e. locked to fixed external perturbations.

3. Edge and divertor parameters during stationary H-mode with ELMs

Stationary H-mode operation in ASDEX requires frequent ELMs, as discussed in a separate paper [13] at this conference. Since this mode of operation may be desirable also for future fusion experiments such as ITER, we have investigated the edge and divertor parameters under these conditions in more detail. Two topics of specific interest to be treated below are the power deposition at the target plates and sputtering and retention of target material.

Radial plasma profiles close to the upper outer target plate are obtained from the reciprocating Langmuir triple probe which measures two radial profiles during each in-out movement (sampling rate 10 kHz). In addition, the current to a fourth pin grounded to the target plate is recorded. The profiles obtained are similar to those measured in the old divertor in case of frequent ELMs (see fig. 34 of ref. [5] for more details). With a probe velocity of 1 m/s many ELMs occur during one transit. Since the ELMs have approximately constant shape and amplitude, the lower envelope of the experimental profile (i.e. neglecting the ELM spikes) represents the profile between ELMs, while a curve connecting the ELM maxima roughly represents the ELM profile. In fig. 2 we compare the power deposition profiles for different confinement regimes. A specific profile is shown for ohmic heating and an L-mode plasma. For the H-mode, average profiles over eight discharges are given: The dashed line is the profile between ELMs, while the crosses show the peak values of all ELMs recorded by the probe, i.e. they indicate an average ELM shape. Despite the large ELM amplitude, the energy which the ELMs carry to the target is about equal to the integral in between ELMs since the ELM spacing is much larger than the ELM width. The average profile between ELMs has roughly the same amplitude as the L-mode, but is half as wide. This is consistent with equal power flow during L-mode and a quasi-stationary H-mode with the same heating power. Comparing with the calorimetrically measured power deposition on the adjacent target (taking into account the field line inclination, the deposited fraction of divertor radiation, etc), the charged particle power flow is too low by at least a factor of two. The difference could be caused by the assumption of equal electron and ion temperature for the analysis of Langmuir probe data. In fact, a factor five higher average ion temperature could nearly account for the difference. This seems to apply to other tokamaks, too [14]. An even higher ion temperature of several hundred eV is determined in the midplane near the separatrix from low energy charge exchange neutrals. This hot ion population
FIG. 2. Power flow density along field lines as calculated from the Langmuir probe signals assuming $T_i = T_e$. Single profiles are given for an ohmic and L-mode plasma. For the H-mode the profile between ELMs is shown, averaged over eight discharges. For the same discharges the maxima of all ELMs occurring during a movement are indicated by crosses. Together they indicate an average ELM profile.

can reach the divertor and form a two temperature ion distribution there together with the cold recycling particles. There are indications from fast particle reflection in the divertor [15] as well as from the copper sputtering discussed below that at least during an ELM the ion energy profile is wider than the deposition profiles given in fig.2.
The time integrated toroidal distribution of power deposition as measured by target plate calorimetry is shown in fig.3 for an H-mode with regular ELMs compared to an L-mode. Both discharges are upshifted by 2 cm, therefore only data for the upper two target plates are given. There is a similar, but moderate toroidal variation in both cases (a stronger sinusoidal variation reported for ELMy H-mode in [5], fig.19, was at least partly caused by the presence of locked modes). The distribution of power between inner and outer target plate, however, is different. The strong asymmetry observed for L-mode is probably caused by dominant radial transport to the low-field side. It is rather insensitive to the vertical shift (double null looks rather similar). With ELMs the asymmetry is significantly reduced indicating a poloidally more uniform transport into the scrape-off layer.

An important question for future machines is whether the ELMs cause excessive target sputtering and whether a quasistationary state is possible without accumulation of target material in the bulk plasma. In ASDEX with watercooled copper target plates, such discharges are routinely obtained for several seconds. A typical temporal evolution of the $D_{\alpha}$ light emission in the upper outer divertor, the copper sputtering at the lower outer target (Cu I line intensity), and the intensity of a copper line from the bulk plasma (Cu XIX) is
FIG. 3. Toroidal variation of the power deposition on the upper inner and outer target plate as measured by cooling water calorimetry. An L-mode discharge is compared with an ELMy H-mode discharge. Both are vertically upshifted by 2 cm (single null) and have about the same safety factor \( q \approx 3 \).

shown in fig.4 (#33155, \( B = 2 \) Tesla, \( I_p = 320 \) kA, \( n_e = 3 \times 10^{19} \) m\(^{-3} \); 2 cm vertical upshift, i.e. single null, but with second outer separatrix close to the inner one). More ELM details can be seen from the divertor signals expanded in time. The \( D_\alpha \) light is recorded at 5 kHz and there is already some aliasing, since the ELM duration (as derived from \( D_\alpha \) occasionally recorded in the MHz range) is of the order of only a few 100 microseconds with a rise time shorter than that. The copper line signals are recorded with 1 kHz, but since photons are counted and summed up between two samplings, there is just a pulse shape distortion; but the integral over an ELM should be correct. We find that the total copper sputtered there during the ELM is more than a factor of two larger than the integral in between ELMs. As an estimate for the true peak value of the copper sputtering during an ELM, we have obviously to multiply the recorded peak by the ratio of the distorted pulse width to the true one, which is roughly ten. The peak value determined this way is much higher than the value between ELMs and even during L-mode. They are consistent with the strongly non-linear sputtering yield as a function of temperature. Remember that this strong ELM effect in the bottom divertor occurs despite the vertical upshift of 2 cm (1.2 cm distance between the two separatrices in the midplane). This indicates a strong increase of the ion energy beyond the second separatrix, which may partly be caused by an ion temperature profile widening during the ELM. In a double null discharge the recorded peaks are twice as large and the signal between peaks increases more than a factor five. This is expected since in this case the whole scrape-off layer is connected to both divertors. A quantitative
comparison is, however, not possible, since the ELMs are much less regular for double null and the H-mode quality is worse.

4. Conclusions

Recent edge and divertor studies, mostly under boronized wall conditions and with mechanically tight divertor domes, have provided further insight into the basic physical processes governing the edge dynamics. Under these improved plasma conditions, a more reliable comparison of empirical edge scalings, edge density limitations, thermo-currents and impurity results with theoretical models has been tried. Generally, the relevance of classical transport along field lines in competition with some anomalous cross field transport has been confirmed within the existing experimental data base covering a large variety of plasma regimes, different auxiliary heating etc. The existence and importance of poloidal and especially of toroidal asymmetries has been demonstrated.
Specific interest has been devoted recently to the stationary H-mode with regular ELMs. Power deposition and impurity production and transport have been analysed. Though quasi-stationary long-pulse discharges of this type without significant impurity problems have been routinely obtained in ASDEX, the understanding and modelling of edge and divertor under these conditions is still in its infancy.

REFERENCES


DISCUSSION

P. GHENDRIH: What ratio of $T_i/T_e$ would you expect to explain the heat load difference between calorimetry and other measured values? Also, are you interested in beryllium as a coating, since you have achieved such good results with boron?
J. NEUHAUSER: $T_i \approx 5 \langle T_e \rangle$ must be formally used to explain the difference. There are indications for high $T_i$ at the edge from a low energy neutral analyser ($E_i \approx 400$ eV during neutral injection). Analysis of these neutral spectra by a Monte Carlo code to get $T_{i,\text{edge}}$ is under way. To answer your second question, we are not interested in using beryllium for ASDEX. We are happy with boron.

S.I. ITOH: With respect to the toroidal/poloidal asymmetry of the heat flux onto the divertor, you referred to the 'critical(?)' q value, namely $q = 3$. Is this related to the connection length in the scrape-off layer?

J. NEUHAUSER: We assume that the increased asymmetry at low q is connected with external magnetic perturbations. Their effect increases strongly when the q = 2 surface approaches the edge (see also Ref. [2]).

S.I. ITOH: Do you have the parametric dependence of the heat deposition width (profile) of the divertor plate?

J. NEUHAUSER: An edge database has been established [3] and a limited amount of Langmuir probe data, including power deposition profile widths, is being added. I cannot yet give you the desired scaling, but it should be available in the future.

S.A. COHEN: Do you find kinetic effects operating in the scrape-off layer? In particular, do you find 'two temperature’ electron distributions (or tails) near the divertor plate?

J. NEUHAUSER: We have no direct measurements. We have indirect information from the fact that the electron heat conduction along field lines is consistent with the classical formula; a strong heat flux limit does not seem to be required.

R.J. GOLDSTON: Perhaps you could explain in more detail why it is you feel that argon retention is not well correlated with $L_i$ in your experiment?

J. NEUHAUSER: There is a strong variation with q for $q < 3$ in contrast to $q > 3$. A preliminary analysis shows non-monotonic behaviour of argon retention with the field line length $L_i$. This might indicate that a different mechanism is operative at $q < 3$, and it would be in line with the edge anomalies observed at low $q$ with respect to many other quantities, as shown in our paper.
COMPARISON OF BERYLLIUM AND GRAPHITE FIRST WALLS IN JET

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Abstract

COMPARISON OF BERYLLIUM AND GRAPHITE FIRST WALLS IN JET.

JET has operated with beryllium as a first wall material in 1989 and 1990. An initial period with beryllium evaporation onto the original graphite surfaces was followed by operation with beryllium belt limiter tiles. Beryllium Faraday shields for the ICRH antennae and lower X-point target tiles were installed for experiments in 1990. The use of beryllium has increased the density limit, significantly reduced deconditioning following disruptions, allowed heavy gas fuelling for impurity control, reduced the impurity influx from the ICRH antennae so that ICRH only H modes were possible for the first time and permitted hot ion plasmas on the outer limiters. The paper describes the primary effects of beryllium which led to these improvements in performance.

1. INTRODUCTION

JET has used graphite limiters, X-point target plates and vessel protection tiles since the start of operation. Although high performance plasmas have been obtained, it has become apparent that the lack of density control, the fuel dilution, radiative density limit and the "carbon bloom" all present serious obstacles to the development of the JET programme. Beryllium was proposed [1] for JET as an alternative to graphite in order to circumvent some of these difficulties. It has the obvious advantage that its nuclear charge (and so radiative cooling rate) is lower than carbon whilst having physical sputtering yields which are similar to graphite. Although beryllium melts at a very much lower temperature than graphite sublimes, this disadvantage is offset by its higher thermal conductivity and the phenomenon of "radiation enhanced sublimation" in graphite. Pilot experiments with beryllium limiter tiles were conducted on the UNITOR [2] and ISX-B [3] tokamaks.

* See Appendix to IAEA-CN-53/A-I-2, this volume.
The results of these experiments were sufficiently encouraging that the decision was made early in 1989 to conduct the beryllium experiment on JET.

There have been four experimental phases of the beryllium experiment:

(i) A graphite only phase which was used to bring the machine back into operation following the 1988/89 shut-down. Since the machine was poorly conditioned in this phase, data representative of graphite operation is taken from 1988 experiments.

(ii) The beryllium evaporation phase was intended to separate the gettering effect of beryllium from its use as a material limiter. Four evaporators, equispaced in the mid-plane of the machine, were used to deposit a layer about 150 Å thick on the inside wall and about ten times that on surfaces near to the evaporators. The evaporations were performed overnight in order not to interrupt tokamak operation.

(iii) The beryllium limiter phase saw the replacement of the graphite tiles in the two outer belt limiters by beryllium components. Evaporation continued to be used to coat the ICRH antennae, the X-point target tiles and the inner wall.

(iv) In 1990 the nickel Faraday shields of the ICRH antennae were replaced by beryllium components and a trial set of lower x-point targets were constructed out of tiles originally intended for the protective frame around the ICRH antennae.

This paper describes the primary effects of the use of beryllium and compares them with observations from the graphite only phase. The resulting high performance plasmas are described in greater detail in the other JET papers which are presented at this conference [4,5,6,7,8].

2. IMPURITY SOURCES

2.1 Graphite Only Phase
Although there was some variability, according to the kind of operation and its timing relative to vessel vents, the sources of carbon and nickel during operation with graphite surfaces were consistent with expectations based on laboratory sputtering data. The effective yield of carbon (which is defined as the ratio of the carbon to deuterium source rates) falls from about 8% at low density to 4% near the high density limit, as shown in Fig. 1. As the input power is increased, the effective yield curve moves to higher density. Calculations of the effective yield, which take account of
self sputtering and use Langmuir probe data for the boundary plasma conditions, give a satisfactory account of the experimental data [9]. Measurements of the amount of erosion found at the wettest part of the limiters are in general agreement with this picture. For the main part, the nickel source emanated from the ICRH Faraday shields. Again, when use was made of a model for the plasma in front of the antenna and the circuit formed between it and the Faraday shields, the data could be described by sputtering [10]. Sputtering of nickel by carbon is an important component of the ICRH induced nickel source.

Oxygen and chlorine behave quite differently from carbon and nickel. Both depend on the conditioning history of the machine. The sources are strongest close to the density limit. An interesting result, which became apparent after the beryllium gettering, is that a significant proportion of the carbon source is due to sputtering by oxygen (which has a yield close to unity). Thus, periods close to major vents, such as that early in 1989, are characterised by larger effective sputtering yields for carbon than when the surfaces are well conditioned.
2.2 Beryllium Evaporation
The most obvious effect of the beryllium evaporation was the almost total disappearance of the oxygen source due to gettering. For otherwise similar plasma conditions, the OII signal fell by a factor of more than 20, the carbon source was also reduced, a factor of two being typical. Whilst this could be due to the coverage of graphite surfaces with beryllium, modelling of this and the variability in the graphite phase suggests that it is the elimination of oxygen sputtering which is responsible. Once beryllium had been introduced in the vacuum vessel, chlorine became the next most important impurity instead of oxygen.

Following an evaporation, the beryllium film would be removed from the region of the limiters in contact with the plasma in a few pulses. In the course of these pulses, the carbon source increased with the exposed graphite surface area. Subsequently there would be very little development in any of the impurity sources over periods of at least 40 pulses. The coating of the ICRH Faraday shields by beryllium reduced the nickel source and would gradually deteriorate over 10-20 pulses.

2.3 Beryllium Limiter Phase
The installation of the beryllium limiters had the obvious effect, for limiter plasmas at least, of replacing the carbon source by beryllium. This further reduction in the carbon source and its consequent effect on the ICRH antennae brought about the virtual elimination of the nickel influx from the Faraday shields. As will be described later, the improvement in the operating range with ICRH was dramatic.

The beryllium effective yield behaves in a similar fashion with density and power to that of carbon, as shown in Fig. 2. However, the yield is as much as 2-3 times greater at the low density end of the curve than that anticipated from normal incidence physical sputtering. The boundary ion temperature is inferred from measurements to be much greater than that of the electrons in these plasmas. A possible result of this is that the ions arrive at the limiter surface at grazing incidence so that the effective yield is enhanced, particularly by self-sputtering.

Both carbon and beryllium limiters had hot spots at which sublimation or evaporation took place. The effective yields described above pertain to power levels and regions of the limiters which are not significantly affected by hot spots. The impurity source from the hot spots causes the "carbon (or beryllium) bloom" which has a severe effect on high performance plasmas [11].
The beryllium X-point targets have been used for plasma experiments but, at the time of writing, the results have not yet been evaluated.

3. IMPURITY CONCENTRATIONS AND PLASMA RADIATION

The impurity concentrations observed in the different phases of operation can be broadly summarised by remarking that the effective yield for all the impurity species is nearly identical with the ratio of the core density of that species to the deuteron density. This is a somewhat surprising result because the ionization lengths for each of the species involved are quite different. Nevertheless, the relationship holds for the spectroscopically derived particle sources and the density ratios $n_c/n_D$ and $n_{Be}/n_D$ as shown in the lower halves of Figs. 1 and 2.

The deuterium pumping conferred by the beryllium limiter allows heavy gas fuelling to be used to reduce impurity concentrations during high power heating. This gas fuelling has
little effect on the density of the core plasma but does raise the density at the boundary. The discharges used in Fig. 2 had just enough fuelling to sustain the line density by feedback. However it should be noted that the effective yield and the \( Z_{\text{eff}} \) value can by multi-valued at a given density. Interestingly, the relationship between the effective yield and \( n_{\text{Be}}/n_D \) is maintained during the heavy fuelling. This shows that its effect must be to reduce the overall particle confinement time. The mechanism behind this is not understood and is being investigated.

The radiated power is measured with bolometers and VUV/XUV survey spectrometers. The sum of radiation components from each species, calculated from selected line strengths, give good agreement with the bolometric total radiated power.

With graphite limiters the radiated power accounted for 30-100% of the input power even with beryllium evaporation. Carbon and oxygen contributed the bulk of the radiation prior to the introduction of beryllium. In the evaporation phase, chlorine became an important radiator, particularly near the density limit, and carbon provided the rest.

Plasmas on the beryllium limiters radiate 15-60% of the input power with the lower end of the range being typical. Only chlorine and beryllium contribute significantly. During operation, the chlorine was gradually eliminated and thus led to a considerable improvement in the density limit, which is described in the next section. Thus the lower charge of beryllium, compared with that of carbon, has conferred a definite advantage by reducing the radiated power and increasing the operating range.

4. THE DENSITY LIMIT

The density limit is determined by the loss of thermal stability when the radiated power becomes a substantial fraction of the input power. A MARFE forms, whereupon the plasma column either contracts radially until it becomes MHD unstable and disrupts or the MARFE phase which ejects a sufficient number of particles that thermal stability and poloidal symmetry are restored without a plasma disruption. In JET the former behaviour is associated with graphite limiters and the radial contraction occurs when the radiated power is 100% of the input. The softer limit is observed with beryllium limiters and is triggered when 50-60% of the power is radiated. During the MARFEing phase the radiation can exceed 100% of the input.
The Ohmic heating density limit on the beryllium limiters at \( q_v \approx 3.8 \) and \( I_p = 3 \text{MA} \) was improved by 50% over the best value achieved with the graphite limiters. Reducing the field to 1.55 T, so that \( q_v = 2.6 \) was obtained, increased the density limit by a further 20%. The line average density was \( 0.6 \times 10^{20} \text{m}^{-3} \).

In the graphite phase, NBH could be used to increase the density limit. As expected from simple power balance arguments, the density increases as the square root of the total input power. On the other hand, with ICRH the limit remained close to the Ohmic value because of the impurity influx from the antennae. However, the beryllium evaporation coated the Faraday shields sufficiently that the performance with ICRH became identical to that with NBH. When the beryllium limiter was introduced the impurity influx was so reduced that ICRH only H-modes became possible for the first time.

5. PLASMA PUMPING

Density control with graphite surfaces had always been a problem in JET because the recycling coefficient was close to unity. As shown in Fig. 3, the characteristic decay time for the deuteron density, when the gas feed was switched off, was typically 30 seconds. The density control could be improved for a few pulses by conditioning the limiter with helium plasmas.

Beryllium evaporation caused a long term change in this behaviour and the density decay time dropped to about 5-12 seconds. It can be seen in Fig. 3 that the pumping was still effective after one day's operation, although the decay time would approximately double. The beryllium limiters brought about a further improvement and the pumping time fell to 1-2 seconds. Some temporary deconditioning would occur after very high power (>30MW) heating. Helium plasmas are also pumped by the beryllium limiters but not at all by the graphite (even with beryllium evaporation). The density decay rates are fairly similar to those of deuterium.

In respect of the deuterium pumping, JET with beryllium is very similar to other all metal tokamaks. This pumping brings about the following improvements in plasma operation:

(i) Effective density control is possible,
(ii) as noted earlier, large influxes of deuterium can be used to control impurities during additional heating without a consequent density rise and
(iii) the deconditioning effect of MARFE's or disruptions has been significantly reduced.

With graphite limiters, gas desorption in pulses subsequent to a major disruption would often be sufficient to prevent plasma start-up. It often proved necessary to recondition the chamber by GDC in helium for a few hours. This loss of experimental time has largely been eliminated by the use of beryllium, either because the gas load following a disruption is reduced or because of the pumping of the following start-up. Many more high performance pulses can be obtained in a session as a result.

The pressure in the torus falls off as \( t^{-0.75\pm0.1} \) following a plasma shot with beryllium surfaces. At least 80% of the deuterium used is recovered in the pumps. From the time dependence of the \( D_2 \) and DH fractions, it is concluded that the gas release is recombination dominated.

6. CONCLUSIONS

JET has been successfully operated with beryllium evaporation onto graphite surfaces and with beryllium limiters, ICRH Faraday shields
and X-point target tiles. Although hot spots and surface melting have been observed on the beryllium components, the plasmas can be optimised so that improved performance over the graphite phase is obtained.

Beryllium has shown itself to be a viable replacement for graphite as a plasma facing material. Applications can be seen for both in the first wall of a next step tokamak. However neither material can satisfactorily match the requirements for divertor target tiles, because of the problem of blooms, unless most of the power is radiated from the target plasma. The demonstration of a divertor configuration consistent with the requirements of next step tokamaks is the goal of the proposed pumped divertor phase of JET.

REFERENCES


DISCUSSION

R.J. HAWRYLUK: Figure 3 of your paper indicates that both pulse duration and power were lower in the Be divertor than in x-point operation with C before the onset of a 'bloom'. Is this a general result?

P.R. THOMAS: The beryllium x-point target is of a provisional design, using tiles originally intended for the ICRH antennas. There are gaps and edges presented to the magnetic field lines so that the target is rather less than optimal with respect to blooms. In fact, it is rather remarkable that the H-mode lasts as long as it does.

R.J. TAYLOR: On the basis of your presentation, I cannot accept that the carbon and beryllium blooms are the same. Do you have photographic data comparing these blooms?
P.R. THOMAS: As far as we can tell the carbon and beryllium blooms are identical in form. Often, sawtooth crashes, giant ELMs or the termination of the H-mode will inflict enough power density on both kinds of target to provoke the bloom.

We have obtained photographic data of both carbon and beryllium blooms with charge couple device (CCD) TV cameras. Shortly before the bloom the light from the targets overloads the cameras so we cannot see what is going on. However, at the onset of the bloom, or rather its effect on the plasma, a MARFE-like, luminous plasma blob is seen to emerge from the target region and move to the inboard of the main plasma.
IMPURITY CONTROL AND HELIUM EXHAUST EXPERIMENT IN JT-60


Abstract

IMPURITY CONTROL AND HELIUM EXHAUST EXPERIMENT IN JT-60.

Impurity characteristics have been investigated in high power (20 MW) NB heated discharges of outer and lower X-point configurations with TiC coated molybdenum or graphite first walls. The light impurity concentrations and $Z_{eff}$ had the lowest values in the outer X-point discharges with TiC coated molybdenum first walls. In high density discharges, the heat load on the divertor plate was suppressed by radiative cooling originating from carbon and hydrogen in the divertor region. To simulate the behaviour of the helium ash in a fusion reactor, helium particles were fuelled by a 30 keV helium beam or gas puffing in NB heated hydrogen plasmas. The neutral pressures of helium and hydrogen in the divertor region were proportional to $n_e^2$. For the helium transport in the main plasma, the inward flow pinch of helium ions was slightly more pronounced than that of the hydrogen ions.
1. INTRODUCTION

In realizing a fusion reactor, impurity control and helium ash exhaust in high
temperature plasmas are important issues, because of radiation loss and fuel dilution.
These problems are closely related to the selection of first wall materials, the charac-
teristics of the boundary plasmas, and also impurity and helium transport in the main
plasma [1]. In this paper, we report on (1) the impurity characteristics in the limiter,
outer X-point and lower X-point discharges with metallic graphite first wall, (2) the
heat flux on the divertor plate and the radiation in the divertor region and (3) the
helium exhaust experiment.

2. IMPURITY CHARACTERISTICS

JT-60 experiments have been carried out with metallic (TiC coated Mo) and
graphite first wall in the limiter, outer X-point (closed divertor: with divertor cham-
ber) or lower X-point (open divertor: without divertor chamber) configuration [2].
In these experiments, the impurity characteristics were investigated [3]. For the high
power (∼20 MW) NB heated discharges, $Z_{\text{eff}}$ values are plotted in Fig. 1 against
the line averaged electron density $\bar{n}_e$. Table I summarizes the $Z_{\text{eff}}$ values, impurity
concentrations and radiation losses of the main plasmas at $\bar{n}_e = 4 \times 10^{19} \text{ m}^{-3}$ for
different wall materials and configurations.

![Graph showing $Z_{\text{eff}}$ versus $\bar{n}_e$]

FIG. 1. $Z_{\text{eff}}$ versus line averaged electron density $\bar{n}_e$. 

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Table I

<table>
<thead>
<tr>
<th>First Wall Material</th>
<th>Configuration</th>
<th>$Z_{\text{eff}}$ Values</th>
</tr>
</thead>
<tbody>
<tr>
<td>TiC-Mo</td>
<td>Outer X-point (Closed divertor)</td>
<td>$\triangle$</td>
</tr>
<tr>
<td></td>
<td>Lower X-point (Open divertor)</td>
<td>$\times$</td>
</tr>
<tr>
<td></td>
<td>Limiter</td>
<td>$\times$</td>
</tr>
</tbody>
</table>

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$Z_{\text{eff}}$ versus line averaged electron density $\bar{n}_e$. 

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TABLE I. IMPURITY CHARACTERISTICS OF JT-60 IN DISCHARGES WITH DIFFERENT CONFIGURATIONS AND FIRST WALL MATERIALS (P_{NB} \sim 20 \text{ MW}, n_e = 4 \times 10^{19} \text{ m}^{-3}).

<table>
<thead>
<tr>
<th>First wall</th>
<th>Metallic wall (TiC-Mo)</th>
<th>Graphite wall</th>
</tr>
</thead>
<tbody>
<tr>
<td>Configuration</td>
<td>Outer X-point (closed divertor)</td>
<td>Outer X-point (closed divertor)</td>
</tr>
<tr>
<td>(Z_{\text{eff}})</td>
<td>1.6</td>
<td>2.2</td>
</tr>
<tr>
<td>(n_c/n_e)</td>
<td>0.2 %</td>
<td>0.8 %</td>
</tr>
<tr>
<td>(n_0/n_e)</td>
<td>1 %</td>
<td>2 %</td>
</tr>
<tr>
<td>(n_{\text{Ti}}/n_e)</td>
<td>0.006 %</td>
<td>0.0005</td>
</tr>
<tr>
<td>Radiation loss (Main plasma)</td>
<td>10 %</td>
<td>&gt; 60 %</td>
</tr>
</tbody>
</table>

In limiter discharges with metallic first wall, the bulk of the plasma energy was lost through radiation of metallic impurity ions; therefore, no high current or high power heating discharges were achieved. On the other hand, in divertor discharges, a wide range of discharges was achieved by sweeping the separatrix lines in order to spread the heat flux widely over the divertor plates. The \(Z_{\text{eff}}\) values were about 1.5 in the range of \(n_e > 3 \times 10^{19} \text{ m}^{-3}\). The radiation losses from the main plasmas were about 10% (originating from oxygen ions) of the input power.

In discharges with graphite first wall, the main impurities were oxygen and carbon, and no significant difference in the impurity effects was observed between limiter and divertor discharges. The radiation losses from the main plasmas were about 20% of the input power.

The effect of divertor configuration and wall materials on the impurity concentrations in the main plasma becomes clear from the experiments undertaken in the various conditions described above. The lower \(Z_{\text{eff}}\) values are obtained in divertor discharges with metallic first wall rather than in those with graphite wall, owing to the lower oxygen concentration. With graphite first wall, a lower carbon concentration is obtained in closed (outer X-point) divertor discharges than in open (lower X-point) divertor discharges. In high power (~MW) NB heated discharges, the \(Z_{\text{eff}}\) values were 2 to 3 in closed divertor discharges and 3 to 5 in open discharges in the range of \(n_e = (3-4) \times 10^{19} \text{ m}^{-3}\).

As to impurity transport, anomalous diffusion was dominant, and no impurity accumulation was observed in OH and NB heated L-mode discharges [4]. On the
FIG. 2. (a) Heat flux density profiles before \( t = 4.18 \) s and during \( t = 4.5 \) s enhanced radiative cooling in the divertor region; (b) time evolution of radiation losses and impurity lines in main plasma and divertor regions; heat load inside and outside divertor plate.
other hand, impurity accumulation was observed in pellet fuelled plasmas, when peaked density profiles and sawtooth-free plasmas were obtained [5]. Fortunately, the total radiation loss in the main plasma did not increase during the accumulation, because of the very low titanium concentration \(n_T/n_e < 10^{-5} \text{ m}^{-3}\). This accumulation was well simulated with neoclassical transport and a very small anomalous diffusion coefficient in the inner region of the \(q = 1\) surface.

3. HEAT FLUX AND RADIATION IN THE DIVERTOR REGION

In outer and lower X-point discharges, the heat flux on the divertor plate decreased with increasing electron density, because of enhanced radiative cooling in the divertor region [6].

In the lower X-point discharges, the heat flux on the divertor plate and the radiation in the divertor region were measured by an infrared camera and spectrometers, respectively. A remarkable asymmetry of the heat flux in the inside and outside regions of the divertor was observed; this asymmetry was reversed by changing the ion \(\nabla B\) drift direction. In discharges with the ion \(\nabla B\) drift towards the X-point, the radiation in the divertor region increased simultaneously with the increase in the C II intensity in the divertor region, and the inside heat flux on the divertor plate decreased largely [7] in the improved divertor confinement (IDC) regime [8], as is shown in Fig. 2. These observations are consistent with Hinton and Staebler’s theory [9]. In this phase, a low temperature (\(\approx 20\) eV) and high density (\(\approx 10^{20} \text{ m}^{-3}\)) plasma was formed in the divertor region. These parameters were estimated from the heat flux and the \(H_a\) intensity in the divertor region. Using these parameters, the radiation originating from carbon and hydrogen and the charge exchange loss from the hydrogen ions were estimated, with the intensities of C II, C III, C IV and the \(H_a\) lines being measured [10, 11]. The total power resulting from these estimates was consistent with the radiation power measured bolometrically.

4. HELIUM EXHAUST EXPERIMENT

In the discharge of lower X-point (open divertor) configuration, helium particles were fuelled by a 30 keV helium beam in the hydrogen plasma, heated by hydrogen beam injection in order to simulate the behaviour of helium ash in a D–T fusion reactor. A helium gas puff fuelling experiment [12–14] was also carried out [15] to compare with beam fuelling. The radial profiles of the helium ions (\(\text{He}^{2+}\)) density were determined by measuring the intensity of the He II (\(n = 4\) to 3; 4685.7 Å) line originating from a charge exchange recombination process between \(\text{He}^{2+}\) ions and the hydrogen injection beam. The central helium ion density was also determined by measuring the helium particles neutralized by a two-electron capture process between helium ions and a 160 keV diagnostic helium beam [16]. The neutral particle pressure
of helium and hydrogen in the divertor region was measured by a residual gas mass analyser (RGA). For $B_T = 4.5$ T, $I_p = 1$ MA and $P_{NB} = 10$ MW (hydrogen heating beam), the density dependence of the hydrogen and helium neutral pressures in the divertor region were measured. For helium beam fuelling, these pressures increase proportionally to $\bar{n}_e^3$ as is shown in Fig. 3(a). The enrichment factor of the helium density in the divertor region defined by $\eta = (n_{He}/n_{H})_{DIV}/(n_{He}/n_{H})_{MAIN}$ is 0.3 to 0.5, where $(n_{He}/n_{H})_{DIV}$ and $(n_{He}/n_{H})_{MAIN}$ are the ratios of helium to hydrogen particles in the divertor region and in the main plasma, respectively. It was observed that the enrichment factor $\eta$ increased with the electron density. From these results, the design of a helium ash exhaust system in a fusion reactor to be operated in the high density region becomes a realistic task.

Helium transport in the main plasmas was investigated [17] by comparing the measured density profile of helium ions with the calculated profile by a one-dimensional impurity transport code [18]. Figure 3(b) shows the measured radial helium ion density profiles at 0.3, 1.05 and 1.95 s after the start of helium beam fuelling, and calculated profiles for peaking parameters of $C_V = 1$ (solid lines) and $C_V = 1.5$ (dotted line) for $B_T = 4.5$ T, $I_p = 1$ MA and $P_{NB} = 10$ MW. The measured profile was well simulated when the peaking parameter $C_V$ (defined in terms of the inward flow velocity $V_A$ and the anomalous diffusion coefficient $D_A$ as $V_A = -C_V D_A 2r/a^2$) was 1 to 1.5. On the other hand, the $C_V$ of the bulk hydrogen ions estimated from the density profile was 0 to 1. The inward flow pinch of the helium ions was slightly more pronounced than that of the hydrogen ions.
5. CONCLUSIONS

Impurity control was obtained successfully for discharges with metallic and graphite first walls. For high power and current discharges, it was realized that sweeping the separatrix lines is useful in reducing the heat load on the divertor plates. In high density divertor discharges, radiative cooling in the divertor region could reduce the heat flux onto the divertor plate. In lower X-point discharges, a strong asymmetry was found in the heat fluxes inside and outside the divertor. It is possible to increase the radiative cooling rate in the inside divertor region by choosing the ion $\nabla B$ drift direction towards the X-point. The helium enrichment in the divertor region was observed quantitatively. The neutral pressures of helium and hydrogen in the divertor region were proportional to $n_e^2$. The results are promising for helium exhaust in a future device. For helium transport in the main plasma, the inward flow pinch of the helium ions was slightly more pronounced than that of the hydrogen ions.

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DISCUSSION

F. WAGNER: You pointed out the low \( Z_{\text{eff}} \) and carbon concentration in the standard configuration with the outside x-point. On the other hand, you described in detail the beneficial aspects of the proper ion grad-B drift in the low x-point configuration. Which of the two principles — long connection length or proper ion \( \nabla B \) drift — is more important for performance optimization?

T. SUGIE: Both connection length and ion \( \nabla B \) drift direction are important for reducing impurity ions in the main plasma. But it is not clear which is more important from these experimental results.

R.J. HAWRYLUK: The data on the comparison of TiC coated divertor plates with carbon plates indicate that a metallic plate is superior (lower \( Z_{\text{eff}} \)). If this interpretation of your data is correct, have you employed TiC–Mo plates in the JT-60 upgrade? If not, what determined the choice of limiter material?

T. SUGIE: At first, the material used for the divertor plate was high heat conduction C/C composite graphite with a height misalignment of adjoining tiles of less than 0.5 mm. In future, we intend to use metallic divertor plates.
IMPURITY AND DEUTERIUM PELLET STUDIES ON TFTR*

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Abstract

IMPURITY AND DEUTERIUM PELLET STUDIES ON TFTR.

Medium speed lithium and carbon pellets ($v_p=0.4-0.8 \text{ km/s}$) and solid deuterium pellets ($v_p=1.3-1.5 \text{ km/s}$) have been injected into Ohmic, L-mode, and supershot TFTR plasmas. They have been used to study fueling, profile modification, particle and energy transport, and the dependence of confinement on density peakedness and heating power profile. Injection of Li and C pellets $\sim1 \text{s}$ before the neutral beams into low recycling, supershot target plasmas has resulted in a significant increase (5-20%) in the neutron rate during the beam heated portion of the discharge. In addition, injection of Li pellets and the use of Zeeman polarimetry on line emission from the Li ablation cloud have allowed measurement of internal poloidal magnetic field profiles.

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Multiple D2 pellets have been used to fuel neutral beam heated discharges to densities of $5 \times 10^{20} \text{m}^{-3}$ without disruption or indication of enhanced core particle transport. Plasma current and power scaling studies of deuterium pellet densified, $^3$He-minority ICRF heated discharges have been performed. Reduced core transport is observed at both plasma currents studied. Global energy confinement measured at the maximum of the stored energy exhibits L-mode type scaling. Interferometric measurements of the prompt ($\sim 100 \mu s$) responses of the electron density to these pellets are also reported. A long-lived, localized region of high electron density ablatant is routinely observed in the leading edge of the ablatant wake. These high density perturbations propagate at velocities of $\sim 10^4 - 10^5 \text{m/s}$ and, at least for D2 pellets, appear to travel along field lines.

**Impurity Pellet Injection Studies**

*Penetration of Medium Speed Li Pellets:* Because of the large sublimation energy of solid Li ($\Delta E_{\text{sub}} \approx 1.6 \text{ eV/Li}$) relative to solid D2, Li pellets are observed to penetrate deeper into hot TFTR plasmas than D2 pellets with a similar number of electrons. Cameras that image the pellet ablation clouds have made possible the systematic study of Li pellet penetration. The observations have been compared with the model of Parks, et al. [1]. The agreement is satisfactory, except in extremely low density plasmas, where slide-away electrons (which are not accounted for in the model) may be present. Typically, 0.75 km/s Li pellets of 1 mm radius, carrying $\sim 7 \times 10^{20}$ electrons, reach the magnetic axis of 0.8 m minor radius discharges which have $T_{e0} \lesssim 3.5 \text{ keV}$ and $n_{e0} \lesssim 2 - 3 \times 10^{19} \text{m}^{-3}$. Although the Li pellets perturb the density significantly ($\frac{n_{e,\text{post-pellet}}}{n_{e,\text{pre-pellet}}} \lesssim 9$), their injection does not result in plasma disruption.

*Li Pellet Injection into Ohmic Plasmas:* Li pellets have been injected into Ohmic discharges with plasma currents ranging from 0.8 MA to 2.0 MA. The response of a 1.4 MA discharge to a single Li pellet is shown in Fig. 1. Many of the features of this response are similar to what is observed after D2 pellet injection [2,3]. The pellets strongly peak the electron density ($n_e$) profiles; single pellets have increased the central electron density ($n_{e0}$) to $2 \times 10^{20} \text{m}^{-3}$, and peaking factors ($\frac{n_{e0}}{<n_e>_{\text{vol}}}$) of 4 to 6 which persist for $>0.5 \text{ s}$ have been produced. In discharges whose $n_e$ profiles are strongly peaked by the pellet, the $Z_{eff}$ profile in the inner quarter of the discharge becomes strongly peaked. The background,
FIG. 1. Typical response of an Ohmic plasma to a single Li pellet which has deposited a significant number of its particles in the core. Shown are the time histories of $I_p$, $n_{e,r}$, $T_{e,r}$, $E_{tor}$, the electron density peaking factor, $Z_{eff}$ at $r/a \sim 0.1$ and 0.5, the half-width at the 1/e height of the soft X-ray chord brightness profile, and the neutron rate.

intrinsic impurities (primarily C) participate in this peaking, since the changes are greater than can be accounted for by the Li alone. In addition, analysis of soft X-ray profile evolution indicates that the intrinsic, heavier impurity densities become more peaked after Li pellet injection in these discharges. Sawteeth, if present, typically cease. The neutron rate has been observed to increase by a factor of up to 3 after injection
of a Li pellet. As seen in Fig. 1, the total stored energy ($E_{\text{tot}}$) increases after injection of impurity pellets. The global energy confinement time ($\tau_E$) typically shows a marginal ($\lesssim 10\%$) or no increase.

**Impurity Pellet Injection into Supershot Target Plasmas:** Single Li and C pellets were injected into the low density, low recycling [4] plasmas which are the usual target plasmas for supershots. For pellet injection $\sim 1$ s before the start of neutral beam heating, the subsequent peak neutron rate was increased by 5-20%, and the peak $E_{\text{tot}}$ was increased by $\sim 5\%$ when compared to otherwise similar discharges with no impurity pellet injection. Li pellet injection was used in this way in the discharge for which a TFTR record neutron rate of $4.4 \times 10^{16}/s$ was measured. Fig. 2(b) shows the differences in the neutron rate, $E_{\text{tot}}$, and $T_e$ which are typically observed following the early injection of a Li pellet. (The timing is shown in Fig. 2(a)). Since most of the pellet particles have left the discharge before the onset of neutral beam heating, it is presumed that the beneficial effects of the pellet result from gettering or covering the surfaces which are the sources for carbon and (cold) D$_2$ during the remainder of the discharge. This model is supported by evidence from an experiment in which Li pellets were injected into a series of low recycling, Ohmic discharges, which are routinely used to reduce carbon and D$_2$ influx from the limiter. These limiter 'conditioning' shots, which have only a He pre-fill and no gas puffing, were produced, without Li pellet injection, until the electron density showed no significant decrease from shot to shot. Then Li pellet injection into the next $\sim 25$ shots resulted in further, gradual decreases of $\sim 10\%$ in both the electron density and in the edge carbon emission. We hypothesize that a similar effect occurs in the supershots, reducing the carbon influx and the thermal D ion density relative to the beam ion density. During the NB heating phase, edge C emission, $Z_{\text{eff}}$, and $n_e^{\text{edge}}$ are observed to decrease, while the neutron rate, $E_{\text{tot}}$, $T_e$, electron density peaking factor, and $\tau_E$ are observed to increase.

**Poloidal field measurements using Li Pellets:** Li pellet injection has also allowed measurement of the pitch angle of the internal magnetic field. The diagnostic technique relies on measuring the Zeeman-polarized line emission from the Li pellet ablation cloud; in the experimental viewing geometry, the $\pi$ components of the emission are linearly polarized parallel to $B$ [5,6]. Emission from Li$^+$ is used, and on TFTR, with $B_o=4.8T$, the emission at $\lambda = 5485$ Å is approximately 30% polarized. This is sufficient to yield pitch angle measurements and to reconstruct $q$ profiles
FIG. 2. Time histories of two supershots. Traces shown with the solid curve had a Li pellet injected 1 s before the initiation of 29 MW of NBI; traces shown with the broken curve had no Li pellet injection. Shown are (a) the total number of electrons versus time, and (b) the neutron rate, $E_{\text{meq}}$, and $T_{\text{e0}}$ versus time.

with $\sim \pm 15\%$ accuracy. As an example, internal field measurements have been compared at two different times in an Ohmic discharge. The results are shown in Fig. 3. One profile was obtained just after the flat-top current was reached, but before the onset of sawtooth activity, and another was obtained about 1 second after the sawtothing had begun. At the earlier time, $q_0 \gtrsim 1.0$, while 1.4 s later, $q \approx 1.0$ at $r/a \approx 0.30$. At this later time, the pellet did not penetrate to the magnetic axis, and $q_0$ was not measured. The $q$ profiles shown are self-consistent reconstructions of the plasma equilibria, which have included both the external magnetic measurements and the internal measurements of the field pitch angle.
FIG. 3. (a) Comparison of the magnetic field pitch angle profiles at two different times as measured by polarimetry on the Li$^+$ emission from the Li pellet ablation cloud. The curves through the two sets of data points are the respective polynomial fits, including points at the magnetic axis ($\theta = 0$). (b) Comparison of the $q$ profiles from the reconstructed equilibria at the two different times (broken line at $t = 2.1$ s, solid line at $t = 3.5$ s). Also shown is the sawtooth inversion radius (from ECE) at the later time. The error bars shown on $q(r)$ represent only the simple transformation of the uncertainties in the pitch angle measurements.

The technique has been employed for measurements of the field pitch profile in other discharge types as well [7,8].

**D$_2$ Pellet Injection Studies**

The Deuterium Pellet Injector (DPI) experimental program on TFTR employs the injection of medium speed D$_2$ pellets to explore physics issues relevant to CIT and to enhanced Q$_{DT}$ operation in TFTR, as perturbations to study core transport [8], and to study the details of the pellet fueling process. A central issue is the transport of particles and energy in peaked, high density plasmas and the effects that auxiliary heating power deposition profiles have on this transport. Current ex-
Experiments have demonstrated the feasibility of producing high density, peaked density profiles in sawtooth-free TFTR plasmas suitable for high power auxiliary heating. Heating experiments using ICRF indicate that core transport is reduced in pellet fueled TFTR plasmas as observed on JET [9]. Global energy confinement ranges from 1.25 to 1.6 times the L-mode value depending on the auxiliary heating method [10] and time in the auxiliary heating pulse. Confinement time is higher early in the pulse, up to 1.6 times L-mode, while the stored energy is rising, and then decreases slowly from its peak value down to 1.2 - 1.4 times L-mode through the period of maximum stored energy. Initial scaling experiments with plasma current from 1.4 to 2.1 MA indicate that both reduced core transport and global scaling are retained at higher current. Future work will extend these results to higher RF power (12.5 MW) and higher plasma current (3.0 MA) to study particle and energy transport within the plasma core in low collisionality, \( T_i = T_e \) discharges at densities of 4 to \( 5 \times 10^{20} \text{ m}^{-3} \).

**Pellet Fueling to High Density with Neutral Beam Injection:** Following extensive helium discharge conditioning of the inner wall bumper limiter and deposition of a boron layer within the vacuum vessel to serve as an oxygen getter, pellet fueling experiments using multiple pellet sequences were conducted to produce peaked density profiles at high density in the presence of strong auxiliary heating. The first pellets in the fueling sequence produced and sustained slightly peaked density profiles at moderate density by off-axis deposition of pellet mass. A central density of approximately \( 1 \times 10^{20} \text{ m}^{-3} \), and peaking factor of 2, were produced. Pellet penetration was to \( r/a = 0.14 \). Central electron temperatures of 2 keV were sustained during this phase of the fueling sequence by 4 MW of beam heating. Hollow density profiles were formed immediately following pellet deposition. The subsequent evolution of the density profile within the plasma core formed a peaked profile over a 100 ms period. Later pellets in the fueling sequence penetrated to and beyond the magnetic axis, increasing the central density from \( 1 \times 10^{20} \text{ m}^{-3} \) to \( 5 \times 10^{20} \text{ m}^{-3} \). Density profiles were more strongly peaked reaching a peaking factor of 3.5. The fraction of input power radiated, \( P_{rad}/P_{heat} \), was 0.4. The parameter \( \bar{n}_e R/B_T \) reached \( 12 \times 10^{19} \text{ m}^{-2} \text{T}^{-1} \) at \( q_a = 3.3 \) without disruption or any indication of strongly enhanced core particle transport and is three times the present TFTR limit in gas fueled discharges. This value is roughly twice the limit proposed by Greenwald et al. [11] and is obtained with a density at the \( q = 2 \) surface greater than \( 1 \times 10^{20} \text{ m}^{-3} \). At this density,
beam heating in TFTR is no longer an effective means of maintaining central temperature. With 14 MW injected power, the central temperature is 1.2 keV and, in the core where \( \frac{P_{\text{OH}}(W/cm^3)}{P_{\text{TR}}(W/cm^3)} = 0.7 \), Ohmic heating dominates even though \( \frac{P_{\text{OH}}(W)}{P_{\text{NBI}}(W)} = 0.09 \).

**ICRF Heating of Pellet Fueled Target Plasmas:** At high density, ICRF provides efficient core heating. D\(_2\) pellet-fueled target plasmas with \( n_{eo} = 2 \times 10^{20} \text{ m}^{-3} \) and peaking factors of 4 - 4.5, have been heated with ICRF at power levels up to 4.5 MW in both the H and \(^3\text{He}\)-minority regimes [10]. The \(^3\text{He}\)-minority regime has been used to heat both 1.4 and 2.1 MA plasmas, providing a first indication of confinement scaling with plasma current in ICRF heated, peaked density profile plasmas. Global confinement measured at the maximum of the stored energy in these discharges is found to scale with power and plasma current as 1.25 times L-mode, comparable to typical high recycling neutral beam heated discharges and high density, gas-puff-fueled, ICRH plasmas in TFTR. In Fig. 4, the peak neutron emission rate is plotted vs. RF power at 1.4 and 2.1 MA for D\(_2\) pellet-fueled target plasmas, heated by \(^3\text{He}\)-minority ICRF. The neutron emission rate is observed to peak early in the heating pulse while the density profile is peaked, and decreases in time as the profile broadens. In these discharges sawteeth are suppressed during the entire heating pulse. The increase in neutron emission early in the pulse is suggestive of reduced core transport in these \(^3\text{He}\)-minority ICRF heated discharges. Indeed, in the H-minority heated pellet-fueled target plasmas, even though global confinement is L-mode, transport (\( \chi_{\text{eff}} \)) within the plasma core (\( r/a < 0.5 \)) is reduced during the high density, peaked profile phase of the discharge by more than a factor of two, as previously observed on JET [12].

Following pellet fueling, the density decreases during the heating period. The decay time scale is found to increase with plasma current. The density and peaking factors decayed from maximum value to half maximum value, relative to the late phase of the ICRF heating period, in about 400 ms at 1.4 MA and 650 ms at 2.1 MA. Finally, the duration of the enhanced phase and time to reach peak neutron rate was greater at the higher current, 400 ms compared to 300 ms.

**Prompt (~100\(\mu\)s) Response of \( n_e \) to \( \text{D}_2, \text{Li}, \) and \( \text{C} \) Pellets**

Rapid sampling (100 kHz) interferometry [13] has been carried out during injection of each of the three pellet types. The interferometer is located about 1 m (18°) toroidally displaced from the \( \text{D}_2 \) injector.
FIG. 4. Plot of peak neutron emission rate versus ICRF power for $^3$He-minority heating of deuterium-densified plasmas at 1.4 and 2.1 MA plasma current. The insert shows the neutron emission for the $P_{RF} = 4$ MW points.

and about 10 m ($180^\circ$) from the Li/C injector. The 10 $\mu$s resolution of the interferometer facilitates the observation and characterization of the density increase resulting from the ablating pellet. The surprising experimental observation is that under conditions of mild ablation ($10^{18}$ - $10^{19}$ atoms ablated per cm travelled by the pellet, a situation which typically occurs outside of $r/a \sim 0.8$), the ablated particles appear to be transported initially through the discharge in the form of a localized, high density perturbation, which travels at a speed of $10^4$-$10^5$ m/s. As the perturbation passes toroidally across the interferometer beams, it is detected as a high density pulse in time. This initial pulse is symmetric in time (having equal rise and fall times of 150 - 250 $\mu$s). The density perturbations from the D$_2$ pellets have been most closely studied, and the ablatant appears to move along field lines. Thus, the ablatant travels
around the torus with a symmetric, high density pulse embedded in its leading edge. Related phenomena have also been observed in the soft X-ray emissivity. These observations are not in qualitative agreement with the present generation of theories describing the details of pellet ablation [14].

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SYNERGETIC EFFECTS OF COMBINED NBI-CO INJECTION WITH ICRH AND NBI-COUNTER INJECTION IN TEXTOR

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Abstract

SYNERGETIC EFFECTS OF COMBINED NBI-CO INJECTION WITH ICRH AND NBI-COUNTER INJECTION IN TEXTOR.

Important synergetic effects are observed when combining NBI-co and ICRH. These include (i) increase of incremental confinement time when each heating method is applied to a plasma preheated by the other one; (ii) enhancement of the current drive efficiency resulting in a 50% increase by the addition of an equal amount of RF power; (iii) enhancement of the density domain where sawtooth stabilization (of the monster type) is achieved; and (iv) significant enhancement of the neutron yield due to beam-target interaction. These effects generally need for their interpretation the assumption of a direct beam acceleration by the RF waves as is indeed observed by the NPA spectra. A synergetic effect is also seen in the combination NBI-co + NBI-counter leading, as for NBI-co + ICRH, to an improved confinement with respect to the Kaye-Goldston or Goldston L-mode scaling.

1. INTRODUCTION

Auxiliary heating on TEXTOR consists of 1.7 MW co-Nl, 1.7 MW counter-Nl and 4 MW ICRH. Up to 6 MW of total heating power has been presently coupled to the plasma leading to $\beta_p = 1.5$ for which $\beta_I$ corresponds to 70% of the Troyon limit [1] and to an average power density of 19 W/cm$^2$ to the first wall (a value similar to the one of JET). The results presented are obtained under the following conditions: $I_p = 340$ kA, $B_t = 2.25$ T, $a = 46$ cm, $R = 175$ cm. The plasma is limited by the toroidal belt limiter ALT-II and is circular in shape. The walls are boronized allowing a recycling coefficient below 1. The Nl-beams have the following characteristics: beam voltage = 50-55 kV, divergence less than 1 deg., D injection into D tangentially to $R = 175$ cm. ICRH is launched by two pairs of antennas on the low field side, operating with $\pi$ phasing.

The three heating methods and their combinations have been compared in Ref. [1]. It is shown there that NBI-co is able to produce at low density a hot ion mode characterized by a total energy content (tails included) exceeding the L-mode scaling, by a large energy anisotropy and non-inductively driven current and by the stabilization of the sawteeth (of the monster type). A significative enhancement of these effects and of the neutron production is observed when ICRH is added to NBI-co. The addition of NBI-counter to NBI-co leads also to an interesting operation regime characterized by a peaked density profile, large excess of the total energy with respect to the L-mode scaling, stabilized sawteeth and a large increase of the neutron yield. This paper is devoted to the characterization of the synergetic effects occurring in these two combined heating regimes.

2. NBI-CO + ICRH

2.1. Nature of synergetism

The synergetic effect between NBI-co and ICRH analyzed below can be due to two levels of interaction: 1°) ICRH produces a hotter target plasma for NBI-co and thus increases the slowing down times of the beam; conversely
the hotter plasma presented to ICRH absorbs the waves better. 2°) ICRH waves produce acceleration of the beam ions at \( \omega = 2\omega_{CD} \), again resulting in a rise of the slowing down times and in a larger e-folding tail temperature of the beam ions above its injection energy [2, 3, 4]; this also produces an additional damping for ICRH. This latter direct interaction, already seen in the charge exchange energy spectra of JET [5], JT-60 [6] and TFTR [7], is also observed on TEXTOR. We measure the slope of the tail beyond the injection energy of the D spectrum perpendicularly to the beam direction by a time of flight analyser with mass discrimination [8]. This tail temperature increases typically from 3.5 to 6.2 keV by the addition of 2 MW ICRH. The two above effects will also lead to a change in deposition profile for both heating methods.

2.2. Energy

For a set of successive discharges pertaining to the best regime of operation with NBI-co, Fig. 1a shows the evolution, versus the central chord density \( n_{eo} \), of the total energy \( E_{tot} \) and the thermal energy \( E_{kin} \) for pure NBI-co (\( P_{NI} = 1.7 \) MW), pure ICRH (\( PR_F = 2 \) MW) and their combination. The difference \( E_{tot} - E_{kin} = E_{//} + E_{\perp} \) gives the sum of the parallel and perpendicular tail energy. The way in which \( E_{tot} \), \( E_{kin} \), \( E_{\perp} \) are obtained from the diamagnetic and equilibrium (or MHD) energy measurements and from the density and temperature profiles is explained in [1]. From the analysis of the results shown in Fig. 1a one can conclude: (i) With combined heating both the total and the kinetic energy increases with respect to OH are more than additive. The synergism is also reflected in the fact that the incremental confinement time \( \tau_{inc} \) of each heating method is improved when applied to a plasma preheated by the other one: at \( n_{eo} = 3 \times 10^{13} \) this improvement is of \( \approx 25 \% \) for \( \tau_{inc,kin} \) and \( \tau_{inc,tot} \) (i.e. relative to the kinetic or to the total energy) of NBI-co when going from an OH target plasma to an ICRH heated target plasma and of \( \approx 50 \% \) for \( \tau_{inc,kin} \) and \( \tau_{inc,tot} \) of ICRH when applied to a NBI heated plasma instead of an OH plasma. (ii) For ICRH the tail contribution is mainly perpendicular, for NBI-co \( E_{\perp} = 0.5 \) to 0.6 \( E_{//} \) (values confirmed by shot simulations with the TRANSP code) and for the combined heating we have within the error bar \( (E_{//})_{NBI-co+ICRH} = (E_{//})_{NBI+co} \). The tail energy increase therefore is mainly due to the rise of \( E_{\perp} \) which can become such that an almost isotropic tail is obtained, i.e.

\[ 2(E_{//})_{NBI+ICRH} = (E_{\perp})_{NBI+ICRH} \] (iii) The improvement factor of the global confinement time for the total energy \( \tau_{tot} \) with respect to the L-mode scaling seen for pure NBI-co remains in the combined heating case: in Fig. 1b the corresponding \( \tau_{tot} \) values are shown versus the total power coupled to the plasma \( P_{tot} \) for three different densities. Two cases are considered: either \( P_{tot} \) is evaluated by its engineering value as in [1] (solid symbols) either \( P_{tot} \) is corrected for the NBI losses due to charge exchange and shinethrough
FIG. 1 (a) Evolution of $E_{\text{tot}}$, $E_{\text{kin}}$ versus $\bar{n}_e$ for ICRH, NBI–co and combined NBI–co + ICRH ($P_{\text{NI}} = 1.7$ MW, $P_{\text{RF}} = 2$ MW). The Ohmic energy $E_{\text{OH}}$ is also shown. $I_p = 340$ kA, D into D injection.

FIG. 1 (b) Global confinement time $\tau_{\text{tot}}$ versus $P_{\text{tot}}$ for improved confinement regimes: NBI–co, balanced NBI–co + counter and NBI–co + ICRH for various values of the density $\bar{n}_e$. The values are given for the engineering value of $P_{\text{tot}}$ (solid symbols) and for $P_{\text{tot}}$ corrected for NBI losses (open symbols). The Goldston (GOLD) and Kaye–Goldston (K–G) scaling laws are also indicated.
(open symbols). Note that, for ICRH, the absorbed power by the plasma $P_{\text{abs}} = \alpha P_{\text{RF}}$ has been determined only for the old High Field Side antennae and that a value $\alpha = 0.9 - 1$ was obtained; for the present LFS launch the determination of $\alpha$ is not completed and we use the uncorrected $P_{\text{RF}}$, i.e. $\alpha = 1$. The scaling law of Goldston [9] which is density independent and of Kaye-Goldston [10] for the considered density range are also shown in Fig. 1b. For the data corrected for NBI losses, the improvement factor with respect to the

![Figure 2(a)](image)

**FIG. 2 (a)** Time evolution of $V_b$, $T_{\text{e0}}$, $n_{\text{e0}}$, $\bar{n}_{\text{e0}}$, and $P_{\text{NI}}$ for an NBI-co heated discharge (1) and a combined NBI-co + ICRH heated one (2) (Nos. 37767, 37771; 340 kA, D into D injection).

![Figure 2(b)](image)

**FIG. 2 (b)** Beam driven current $I_{\text{BDC}}$ versus $n_{\text{e0}}$ for NBI-co (X) and combined NBI-co + ICRH discharges (O). The dotted line corresponds to the predicted $I_{\text{BDC}}$ due to NBI-co injected in a target plasma modified by ICRH (340 kA, D into D injection).
Goldston scaling ranges from 1.4 to 1.7 and those with respect to the Kaye-Goldston from 1.8 to 2.6. Note that, contrarily to the latter scaling, $\tau_{\text{tot}}$ is larger at lower density: this is due to the larger fraction of tail components with very good confinement in $E_{\text{tot}}$ at low density.

### 2.3. Current drive

Fig. 2a shows the comparison of a NBI-co heated shot with a combined heating discharge having the same $n_{\text{eo}} (1.8 \times 10^{13} \text{ cm}^{-3})$. The loop voltage $V_g$ (corrected for $d(L_p I_p)/dt$) drops down from $\approx 0.3 \text{ V}$ (with NBI-co) to $\approx 0.1 \text{ V}$ with the addition of ICRH ($P_{\text{RF}} = P_{\text{NI}} = 1.7 \text{ MW}$). The drop remains stationary for more than 1 s. The total non-inductively driven current $I_{\text{ID}} = I_{\text{BS}} + I_{\text{BDC}} = I_p - V_g / R_{\text{NC}}$ can be computed from $V_g$ and from the neoclassical plasma resistance $R_{\text{NC}}$ and is found to increase from 160 kA (NBI-co) to 240 kA (combined heating) with $I_p = 340$ kA. Using the bootstrap current $I_{\text{BS}}$ obtained from TRANSP simulations, the beam driven current $I_{\text{BDC}}$ is found to rise from 125 (NBI-co) to 200 kA (NBI-co + ICRH). Fig. 2b gives the results of a similar evaluation of $I_{\text{BDC}}$ for various NBI-co shots (the full line is a fit) or for combined heating shots. Also shown is a dotted curve corresponding to the mean NBI-co curve but with the enhancement factor predicted by TRANSP for the target plasma modifications due to ICRH. The direct interaction of RF waves with the beam is therefore generally needed to explain the observations since all the combined heating points lie above or on this curve.

\[\text{FIG. 3. Evolution of sawtooth period } \tau_{\text{sf}} \text{ versus } \bar{n}_{\text{eo}} \text{ for ICRH, NBI-co and combined NBI-co + ICRH} \quad (P_{\text{NI}} = 1.7 \text{ MW}, \quad P_{\text{RF}} = 1.4 \text{ MW}).\]
2.4. Sawtooth stabilization

Sawtooth stabilization for more than 2 seconds, only limited by the heating pulse duration, can be achieved with NBI-co, at sufficiently low $n_{\infty}$ ($I_p = 340$ kA). In Ref [11] this was shown to be linked to the build-up of a high enough pressure $p_{\perp,\text{hot}}$ of trapped ions inside $q=1$ which might result in the stabilization of the resistive $m=1$ mode. In Fig. 3 we show that, whereas ICRH alone can only achieve a modest stretching of the sawtooth period $\tau_{\text{st}}$, the combination ICRH + NBI-Co allows to considerably extend the parameter domain for sawtooth free operation. Considering that the power deposition profile with pure ICRH is broad in TEXTOR we are forced to conclude that some of the RF power is now capable of considerably enhancing $p_{\perp,\text{hot}}$ inside $q=1$, i.e. interacting at $\omega=2\omega_{\text{CD}}$ with the beam ions.

2.5. Neutron yield

The neutron yield observed during NBI-co injection, mainly due to beam-target interaction, is significantly increased by the addition of ICRH (up to 80 % with $P_{\text{RF}} = P_{\text{NI}}$). Up to 50 % of this increase is not explained by the changes of the target plasma and is attributed to the RF acceleration of the beam [12].

3. NBI-CO + NBI-COUNTER

When NBI-counter is added to NBI-co injection (nearly balanced injection), an interesting discharge regime, characterized by a very peaked density profile, a large energy anisotropy (at low density), high central $T_e$ ($\approx 2.7$ keV) and large increase of the neutron yield (from 4 to 6 times those of pure NBI-co) can be obtained [13]. This regime is also characterized by an enhanced particle confinement time generally leading to a density rise during the heating pulse. Complete sawtooth stabilization can also be achieved leading to an extended density domain over which sawtooth free operation is possible. NBI-counter alone is unable to stretch the sawtooth period like NBI-co.

The confinement time $\tau_{\text{tot}}$ for these discharges is as good as for NBI-co + ICRH, as shown in fig. 1b, and yields roughly the same improvement factor with respect to the scaling laws.

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H. KIMURA: Why can sawteeth not be stabilized with ICRF alone?
H. CONRADS: In Textor the ICRF deposition profile of the power is broad because of the geometrical constraints and plasma parameters. The trapped hot ion population induced by ICRF alone inside the $q = 1$ surface must be insufficient to stabilize the sawtooth oscillations.

B. COPPI: Your experimental curve of the density threshold for the suppression of sawtooth oscillations is very interesting. Do you have a model to explain it?
H. CONRADS: Our experimental data show that $\beta_{PH}$, i.e. the pressure established by the fast particles, is the parameter controlling sawtooth stabilization.
CONFINEMENT AND STABILITY PROPERTIES UNDER SLOW COMPRESSION IN MINOR RADIUS, LOW q DISCHARGES AND LHCD IN THE HT-6B TOKAMAK

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Abstract

CONFINEMENT AND STABILITY PROPERTIES UNDER SLOW COMPRESSION IN MINOR RADIUS, LOW q DISCHARGES AND LHCD IN THE HT-6B TOKAMAK.

Confinement and stability properties were investigated in different experimental conditions. With slow magnetic compression along the minor radius, the temperature at the plasma centre increased and profiles shrank. After compression the confinement showed bifurcation behaviour which was identified by different parameters in the hot plasma core and at the edge. In low q discharges, $q_e$ down to 1.75 was obtained. Data obtained from numbers of low q shots showed that $r_e$ was not distorted but kept the normal scaling law. The typical MHD instability property in low q discharges was observed. The LHCD experiment was conducted to investigate the influence of LHCD on MHD instability behaviour, which showed complicated patterns under various experimental conditions.

1. INTRODUCTION

This paper is a combination of contributions describing three sets of experimental results in the HT-6B tokamak. Slow magnetic compression along the minor radius was carried out by increasing $B_t$ in about 3 ms, which was about sevenfold the $\tau_e$ value before compression. Section 3 presents the results of an investigation of confinement in low q discharges. The investigation showed that with careful control of plasma parameters, especially $Z_{eff}$, the $\tau_e$ value remains consistent with the scaling law. Section 4 deals with MHD instabilities during LHCD. The results showed complicated patterns of MHD modes, which suggested the effect of energetic particles.
2. SLOW COMPRESSION ALONG THE MINOR RADIUS

The experiment of slow compression along the minor radius has been carried out on the HT-6B tokamak [1]. Figure 1 shows the typical waveforms of current, loop voltage, density, horizontal displacement and toroidal field. Before compression the plasma parameters were $I_p = 26$ kA, $B_t = 4$ kG, $n_e(0) = (1.0-1.5) \times 10^{13}$ cm$^{-3}$, $T_e(0) = 200$ eV, $\tau_s = 0.4$ ms. The toroidal field was increased to 6 kG in 3 ms and then kept a 10 ms flat-top. Plasma current was kept constant for the whole discharge. After compression the central electron temperature increased to 240 eV and the central ion temperature increased from 40 to 57 eV. The profile of $n_e(r)$ was peaked intensively. The $H_a$ emission decreased by 30–50% and the profiles of impurity emissions moved outward. Analysis of the impurity spectra showed that the electron temperature was increased.

Under certain discharge conditions, the confinement appeared to have two different states, that is to say, a bifurcation emerged. In different states the SX signals had very different characteristics, although the other parameters were the same. In
FIG. 2. Soft X-ray emissions (a) increase during compression and maintain the higher level after compression, and (b) increase during compression and decrease thereafter.
FIG. 3. (a) $n_{ea}$ and $T_{ea}$ drop during compression and recover slowly; (b) $n_{ea}$ drops and $T_{ea}$ increases during compression and they then recover quickly.
one state the central SX signal rose during compression and maintained the higher value after compression (Fig. 2(a)); however, in the other case the SX signals decreased after the process of compression, as shown in Fig. 2(b). The change in electron thermal conductivity of the plasma core before and after compression was very different. In the high sawtooth amplitude state it decreased by about 70% but in the low sawtooth amplitude state it decreased by only about 40%. The electron temperature and density at the plasma edge and in the scrape-off layer were different in the two states. As is shown in Fig. 3(a), $T_e$ was reduced to 70% and $n_e$ was reduced to 50% in the high state; in the low state $T_e$ increased 1.2-fold and $n_e$ was reduced to 50% (Fig. 3(b)). During compression MHD activity did not change significantly but the amplitude of the sawtooth decreased to half the original value, the $m = 1$ oscillation emerged and the region affected by the sawtooth was greatly reduced. Generally speaking, after compression the sawtooth recovered in the high state and did not reappear in the low state. So we could conclude that there was a bifurcation, one branch with improved confinement and the other with normal Ohmic heating confinement, during the time the plasma was compressed along the minor radius, with the compression time much longer than the energy confinement time.

3. CONFINEMENT IN LOW q DISCHARGES

Low $q$ ($q_a = 1.75–2.5$) discharges were obtained in the HT-6B tokamak. It was necessary, for a tokamak operating at these $q$ values without a conductive shell and using a material limiter, to make a careful adjustment of the plasma parameters and have an appropriate ramp-up rate of $I_p$ in the startup phase and low $Z_{eff}$. Figure 4 shows a set of typical waveforms of such discharges. One can see that the discharge was smooth and quiet. Large sawteeth and the 2/1 kink mode were found, which made a simple MHD mode picture. The profiles of $T_e(r)$ and current density $J(r)$ were broader than in high $q$ discharges. Figure 5 shows that the $q = 1$ surface determined by SX emission measurement expanded up to 6 cm ($a = 12.5$ cm). The $Z_{eff}$ value was usually 2, which was necessary for obtaining good low $q$ discharges. During low $q$ discharges $Z_{eff}$ did not change much. The line emission peaks from impurities were found to have moved from spectroscopic measurement, which gave the peak of O(VI) located at $r = 4.2$ cm for $q = 2.4$ and at $r = 4.0$ cm for $q = 4$, and of O(V) at $r = 7.5$ cm for $q = 2.4$ and $r = 5.5$ cm for $q = 4$. Simulation showed this to be consistent with the broad $T_e$ and $n_e$ profiles.

The confinement properties in low $q$ discharges were measured under different conditions. The $\tau_e$ values under different $I_p$, $n_e$ and $B_t$ values are plotted in Fig. 6, in which the PLT and JET scaling law [2] is drawn for comparison. It was clear that the low $q$ confinement in HT-6B was not distorted but was in good agreement with the scaling law. The $T_e$ values given in Table I also showed a smooth decline with $q$, without an abrupt drop as $q$ became lower than 2.
FIG. 4. Waveforms of a low q discharge.

FIG. 5. Profile of soft X ray emission of a low q discharge.
4. MHD BEHAVIOUR IN LHCD EXPERIMENTS

The LHCD system was assembled on the HT-6B tokamak in 1990 with parameters as follows: $P_{\text{RF}} = 100\ kW$, $f_{\text{RF}} = 2.45\ GHz$, and a $1 \times 8$ array multi-junction grill with $\Delta \phi = 90^\circ$, reflectivity $\langle N_1 \rangle = 3.22$ and $\Delta N_1 = 2.0$. Experiments were carried out in tokamak discharges with LHCD and inductive electrical field. When the loop voltage dropped to zero, $I_{\text{RF}} = 23\ kA$ was obtained. The current drive efficiency reached $\eta_{\text{cd}} = 0.2 \times 10^{19} \ m^{-2} \cdot kA \cdot kW^{-1}$. $\eta_{\text{cd}}$ decreased with the decrease of $B_c$ and the increase of $n_e$. The current drive effect could be seen clearly until $\omega = \omega_{pe}^2/\omega_{ce}^2$ up to 4 [3].

Dependence of the electron velocity distribution on the direction of launched travelling waves was observed. In parallel launch discharges ($V_{\phi}$ antiparallel to $E_{\text{Ohmic}}$), the velocity of the driven electrons was limited to lower than 40 keV,
which was equivalent to the phase velocity of the lower hybrid wave. But in the counter-launch case (by changing the phase shifts between wave guides), the electron velocities could be accelerated to much higher values. The energetic electrons with energy up to 1 MeV made the velocity distribution rather flat and often caused runaway discharges (Fig. 7).

The profile of driven current depended on the plasma density. For \( n_e < 1 \times 10^{13} \text{ cm}^{-3} \) the driven current was distributed mainly within 6 cm (Fig. 8(a)), while for \( n_e > 1.5 \times 10^{13} \text{ cm}^{-3} \), RF power was deposited in the outer region. This is shown clearly in Fig. 8(b). This meant that the highest driven efficiency existed in a proper density window, neither in the high density (cut-off) nor in the low density region. A heating effect could be seen from the rise of \( T_i \). Typically, \( T_i \) was increased from 45 to 65 eV.

The MHD instability property was affected by LHCD and also depended on plasma density. Figure 9 shows that 15 kHz Mirnov-like oscillations occurred on SX
FIG. 8. (a) Driven current exists mainly within $r < 6$ cm in the low density plasma; (b) driven current exists mainly at the plasma edge.
signals (detected by surface barrier detector arrays) when \( n_e < 1 \times 10^{13} \) cm\(^{-3}\). The oscillation frequency was lower and the magnitude was much higher than those in Ohmic plasma. The oscillation had its maximum at \( r = 3 \) cm and became very small at \( r = 6 \) cm.

When \( n_e > 1.5 \times 10^{13} \) cm\(^{-3}\) Mirnov-like oscillations were suppressed and sawteeth changed in two ways. Figure 10(a) shows that the sawtooth repetition period increased from 250 to 400 \( \mu s \) during LHCD and that the amplitudes increased, too. But the heat pulse did not propagate to the outer region \( (r > 6 \) cm\)), quite unlike the large sawteeth in Ohmic plasma and RHF experiments [4]. Figure 10(b) shows another effect, by which the sawteeth were suppressed during LHCD and recovered or were even amplified just after the lower hybrid wave stopped. These two effects occurred occasionally. According to the theory, the energetic particles from LHCD and the change of \( J(r) \) profiles are two possible reasons for these phenomena; it is difficult to identify the main one from our data. Further experiments are under way.
FIG. 10. (a) Sawteeth amplified by LHCD, (b) sawteeth suppressed by LHCD.
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DISCUSSION

T.N. TODD: I have a question concerning your graph of waveforms for the $q = 1.9$ discharge (Fig. 4). Are you really claiming that true MHD, field line $q$ is substantially less than 2.00? It is clear that there is no 'glitch' of any sort at the apparent crossing of $q = 2.00$, which is very surprising, especially given the low current ramp rate there.

Jikang XIE: I think this value of $q$ was determined by an electromagnetic measurement. It refers to a minor radius of about 0.95 of the geometric minor radius.

T.N. TODD: So you contend that that really is the actual $q$ of the field lines at the limiter.

Jikang XIE: In low $q$ discharge experiments on HT-6B it has been found that only through very careful adjustment of the growth of plasma current and density, reducing the impurity level and selection of the optimal magnetic configuration could we very quickly reduce $q_a$ and pass through the region of $q_a$ near 2 without disruption. Even in such cases, large MHD fluctuations still have been observed in this narrow region of $q_a$ near 2.

Once the discharge has passed through this narrow region with large MHD fluctuations, many new phenomena appear: the much larger sawtooth oscillation in SX signals than in the $q_a > 2$ case; strong sawtooth oscillation in the loop voltage, $H_a$ and ECE signals; antisawtooth oscillation in the HX signal; $r_{q=1}$ was increased to near half the minor radius of the plasma torus; loop voltage was 0.4–0.6 V higher than in the $q_a > 2$ case, i.e. the Ohmic heating power was 30% higher; the profile of the SX signal became very flat as the discharge entered the $q < 2$ region. All these phenomena show that, as the discharge passed through the narrow $q_a \sim 2$ region with large MHD fluctuations, the plasma entered a new state.

In our paper the listed $q_a$ values of discharges have been given by the formula: $q_a = a_L B_L / R_0 B_p$, where $B_p = \mu_0 I_p / 2 \pi a_L$, and $a_L$ is the limiter $a$. Of course, corrections of three kinds should be considered: (1) toroidal effect: the factor
$C_1 = 1 + 1.5e^2 = 1.12$ for HT-6B should be introduced; (2) correction due to magnetic surface: from Shafranov's formula, the correction factor is $C_2 = 1/(1 - v^2)^{1/2} = 1.013$ for HT-6B; (3) the shrinkage of the current channel due to plasma displacement: from different diagnostics on HT-6B it could be estimated that the displacement of the plasma centre was near 1 cm in our low q experiments but the plasma cross-section was shrunken and still kept near circular; therefore the minor radius of the plasma $a_p$ was near 11.5 cm, smaller than $a_L$. From all the aforementioned corrections, a factor $C = C_1C_2C_3$ should be introduced in our $q_a$ expression, which is near 0.96 for the HT-6B tokamak. Therefore, the given $q_a$ values in our paper can still be considered reliable, i.e. we have really entered the $q_a < 2$ region.
OPTIMISATION OF PERFORMANCE IN JET LIMITER PLASMAS

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Abstract

OPTIMISATION OF PERFORMANCE IN JET LIMITER PLASMAS.
 Improvements to the performance of JET limiter plasmas have been made by increasing the plasma current, by suppressing sawteeth, by peaking the density profile and by using beryllium evaporation and employing beryllium limiters. Central electron temperatures $T_e \sim 12$ keV, ion temperatures $T_i \sim 18$ keV, central electron densities $n_e \sim 2.3 \times 10^{20}$ m$^{-3}$ have been demonstrated in limiter plasmas with currents 3-5 MA. A 7 MA discharge with $q_0 \sim 3.2$ has been demonstrated, and heating experiments have just begun.

1. Introduction

It has already been shown [1] that the confinement of high current low q limiter plasmas in JET is in line with the expectations of Goldston or Rebut-Lallia scalings [2,3], but that the profiles of density and temperature are rather flat. Projections of the D-T fusion performance of such L mode plasmas in JET [4] suggest that $Q_{DT} \sim 0.5 - 0.8$ might be obtained at 7MA (including the beam driven yield). A more direct extrapolation to DT operation of 4.5MA H mode plasmas already run in Deuterium suggests [5] that higher fusion performance will be obtained using such divertor configurations rather than a limiter configuration despite the higher plasma current. In order to enhance the fusion performance of limiter plasmas without employing a regime of enhanced confinement, it is necessary to make the profiles of density and temperature more peaked as has already been demonstrated at low current in the pellet enhanced plasma [6]. In order to obtain peaked profiles

* See Appendix to IAEA-CN-53/A-I-2, this volume.
it is necessary to suppress sawteeth as for example in the monster sawtooth regime [7]. Furthermore, in sawtooothing discharges with centrally deposited power, it appears that the incremental confinement time no longer scales favourably with plasma current for \( q_{\text{cyl}} \leq 3 \) [8]. With the broad density profiles characteristic of high current sawtooothing discharges neutral beam penetration is poor. This reduces the central ion temperature, reduces beam refuelling of the core and reduces the beam driven fusion yield.

Thus to optimise the fusion yield in JET limiter plasmas requires more than a simple increase of plasma current. It is clearly necessary to suppress sawteeth, produce a high central density, maintain a low edge density and keep the influx of impurities to a minimum. The following sections describe experiments addressing these issues for plasma currents of 3-5MA. The progress of experiments at 7MA currently underway is briefly described.

2. Sawtooth suppression at 5MA

Sawtooth stabilisation during the current rise phase has already been demonstrated [9]. By using a faster current ramp at constant \( q_0 \sim 3.5 \) and employing Beryllium evaporation to control the density it has been possible to extend the sawtooth free periods well into the flat top at 5MA [10] as shown in fig 1. An axial electron temperature \( T_e \sim 12 \text{keV} \) has been obtained with a strongly peaked profile as shown in fig 2. The ion temperature is only \( T_i \sim 5 \text{keV} \) so the fusion performance is modest. Note that there are a few sawtooth crashes during the ICRH and during this time the soft X ray inversion radius grows from a small value to \( r/a \sim 0.3 \). Faraday rotation measurements indicate \( q_0 = 0.95 \pm 0.15 \) and constant in time during the sawtooth free periods. Since the inversion radius is large we infer that the \( q \) profile is rather flat. The fast ion slowing down time is long because the density is low and the electron temperature is high and therefore the fast ion energy content is large (35-50\% of total). Normal sawteeth resume shortly after the ICRH is switched off. It is probable that the sawtooth is being transiently stabilised by the fast ions in manner similar to that of the normal
FIG. 1. Various time traces for a 5 MA plasma where ICRF is applied during the current rise and showing sawtooth stabilisation.

monster. However since monster sawteeth do not normally occur during flat top heating at 5MA the detailed q profile shape must be important or it must be necessary to build up a high fast ion pressure to stabilise the sawtooth at high current. It is expected that a continuation of the current ramp will permit sawtooth stabilisation at higher current. The confinement time is plotted against power input for these plasmas and 5MA sawtoothing plasmas in fig 3. It can be seen that the confinement is very close to the Goldston prediction when allowance is made for the fast ion energy content.
FIG. 2. Density and temperature profiles for the pulse illustrated in fig 1.

FIG. 3. Confinement time plotted against loss power for 5 MA belt limiter plasmas. The various symbols indicate the limiter material and distinguish the sawtooth free cases described in section 2 and the pellet cases of section 3.
FIG. 4. Various time traces for a 5 MA plasma with pellet injection during the current rise followed by ICRF heating in the flat top.

FIG. 5. Density and temperature profiles for the pulse illustrated in fig 4.
3. Peaked density profiles at 5MA

The injection of a string of pellets during the current rise, as shown in fig 4, leads to high central densities \( n_e \approx 2.3 \times 10^{20} m^{-3} \) at 5MA. In order to obtain the strongly peaked density profile shown in fig 5 it is necessary that the pellets penetrate deeply and this is achieved by careful choice of pellet timing. The small inversion radius sawteeth present before the first sawteeth are suppressed by the pellets and indeed polarimetric measurements show that \( q_0 \) is raised above unity. It is important to keep \( q_{\psi} \) roughly constant during the current ramp and pellet injection since disruptions invariably occur at \( q_{\psi} \approx 4 \) if \( q_{\psi} \) is allowed to fall. 6MW of ICRF heating were applied at 5MA after the last pellet and, as shown in fig 4, a transiently enhanced D-D rate is observed. 1.5 sec after the start of ICRH, and after the peak D-D rate, the central density has decayed to \( n_e \approx 6 \times 10^{19} m^{-3} \). \( T_e \) is much lower than in the experiment of section 2, but \( T_\alpha \sim T_e \sim 5 keV \) and \( Q_{DD} \sim 5 \times 10^{-4} \). The global confinement time is enhanced transiently by \( \sim 30\% \) compared with gas fuelled discharges as shown in fig 3.

4. Density profile control in plasmas with strong ion heating

The use of Beryllium limiters has permitted operation with low \( Z_{eff} \approx 1.5 \) at moderate density [11]. However at low density, \( P_{TOT} / <n_e> \gtrsim 5 \times 10^{-19} MWm^{-3} \), the deuterium concentration \( n_D/n_e \) is typically \( \sim 0.6 \) similar to Carbon limiters conditioned by Beryllium evaporation or by extensive pulsing in Helium. The main effect of Beryllium (either limiters or evaporation) is improved density control which extends the range of \( P_{NB}/n_e \) and, as shown in fig 6, this has resulted in higher ion temperatures in 3MA belt limiter plasmas.

For the case of Carbon limiters and Beryllium evaporation a density profile peaking factor \( n_{e0}/<n_e> \sim 3 \) was obtained as a result of central beam fuelling and low edge recycling compared to \( n_{e0}/<n_e> \sim 1.5 \) for bare carbon belt limiter plasmas [12]. Unexpectedly, for low density plasmas with Beryllium limiters the density profiles were flat even with central beam fuelling. In order to produce peaked density profiles in this case it was necessary to fuel the target plasma with deeply penetrating
FIG. 6. Axial ion temperature plotted against neutral beam power normalised to the axial density for 3 MA belt limiter plasmas. The symbols indicate the limiter material and distinguish pellet fuelled target plasmas.

FIG. 7. Density profile evolution reconstructed using LIDAR data for several similar discharges where heating was applied to pellet fuelled target plasmas.
pellets. Fig 7 shows the time history of the density profile reconstructed from a series of similar shots. It can be seen that the initial central density is $n_{e0} \sim 7 \times 10^{19} m^{-3}$ with a peaking factor $n_{e0}/\langle n_e \rangle \sim 4$, but that the central density decays during the heating though the profile shape remains peaked. The ion temperature reaches $T_i \sim 18kev$ with a strongly peaked profile $T_{i0}/\langle T_i \rangle \sim 7$. This profile peaking enhances the thermonuclear performance by a factor 3 compared with 'normal' flat profiles.
in L mode, even though the global confinement is unchanged. In these discharges the ICRH suppresses sawteeth which otherwise would flatten the profiles. In addition ICRF acceleration of injected deuterons increases the driven D-D reactivity by 30-40% in such discharges [13]. The broadening of the neutron spectrum due to the ICRF is clearly visible in fig 8. In this case the D-D fusion gain was $Q_{DD} \sim 9 \times 10^{-4}$ which was slightly improved over previous inner wall plasmas at 3MA. Recently yet higher gains, $Q_{DD} \sim 1 \times 10^{-3}$, were obtained in inner wall plasmas at 4.7MA.

5. 7MA plasmas

Plasma currents of 7MA were demonstrated in 1988 [1]. These plasmas were obtained using a simultaneous ramp of toroidal field and plasma current with $q_{\psi} \sim 2.5$. The flat top was only 2 seconds limited by volts seconds. In order to pass $q_{\psi} = 3$ early in the current rise without disruption it was necessary to establish early sawteeth by strong gas puffing. This discharge is not suitable for current rise heating because of the variation in toroidal field. Recently, the fast current rise developed for the experiments described in sections 2 and 3 has been extended to 7MA. The flat top was 3 seconds long but 8 Volt-sec remain to extend this further. By making the plasma more D shaped $q_{\psi}$ was held at $\sim 3.2$ at 7MA. Sawteeth have been suppressed well into the flat top of a 6MA discharge and suppressed during the rise to 7MA by applying ICRF in the current rise. An electron temperature $T_e \sim 9keV$ was obtained in both cases. Ion heating experiments, using NBI, have begun in these high current discharges.

6. Conclusions

It has been shown that sawtooth can be suppressed and peaked density profiles formed at high current low q belt limiter plasmas. The former produces very high electron temperatures $T_e \sim 12kev$, the latter high densities $n_e \sim 2.3 \times 10^{20}m^{-3}$. The use of Beryllium to control recycling together with central fuelling by pellets has allowed peaked density profiles in 3MA beam heated plasmas with ion temperatures up to 18keV enhancing the thermonuclear reactivity over normal flat profiles. Here ICRF suppresses sawteeth and enhances the beam plasma reactivity by
the acceleration of injected deuterons. The highest value of D-D fusion gain yet obtained in a limiter plasma on JET is $Q_{DD} \sim 1 \times 10^{-3}$.

Thus the foundations have been laid for performance optimisation at the highest plasma currents in JET (up to 7MA). These promising results suggest that the performance projections for 7MA D-T operation in JET [4] might indeed be pessimistic.

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DISCUSSION

Y. KAMADA: For the 7 MA discharges, is the current penetration time classical or faster than the resistive diffusion time during $I_p$ ramp-up?

P.J. LOMAS: The current penetration in such discharges is very close to neoclassical resistivity for the constant-q part of the ramp when the initial MHD has decayed.

R.J. TAYLOR: What is the loop voltage when you go from 3.5 to 7 MA without auxiliary heating on JET?

P.J. LOMAS: The loop voltage in an Ohmic discharge at 3 MA is of the order of 0.8 V ($\pm 0.1$ V). At 7 MA the loop voltage is typically 1.1 V, but note that the flat-top is shorter than the L/R time and is, therefore, not fully in equilibrium.

G. TONON: Have you observed any ion mass dependence of the global energy confinement time in the JET limiter discharges?

P.J. LOMAS: We have not systematically investigated the dependence of confinement on ion mass. The scaling prediction for the 5 MA data shown used the Goldston (Aachen 1983) auxiliary scaling and did not have a specific mass dependence, but this was for deuterium plasmas.
SAWTEETH AND THEIR STABILIZATION IN JET

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Abstract

SAWTEETH AND THEIR STABILIZATION IN JET.

Sawtooth relaxations control the central plasma parameters in most JET discharges and, as a result, have a significant influence on fusion performance. Several techniques have, therefore, been developed for the stabilization of sawteeth in various JET operating regimes. Under some circumstances, sawtooth suppression appears to be associated with the central safety factor, $q_0$, rising above unity. However, in other cases, $q_0$ remains significantly below unity, and it is thought that stabilization is due to the influence of a population of energetic particles. Detailed measurements of the evolution of the current profile in JET during sawteeth have confirmed that $q_0$ remains below unity throughout the sawtooth cycle and thus the trigger for the sawtooth collapse remains problematic. Further diagnostic investigations of the sawtooth collapse have shown that the collapse takes the form of a convective plasma flow, but the cause of the resultant energy transport is not understood.

1. Introduction

Sawteeth are one of the fundamental mhd instabilities of tokamak plasmas. They play a key role in limiting central plasma parameters, and thereby fusion reactivity, in the near—ignition regime. Nevertheless, in spite of considerable experimental and theoretical investigations, both at JET (see e.g. [1]) and elsewhere, no consistent picture of the instability has emerged. Furthermore, while several techniques for sawtooth stabilization exist [2—4], their extension to reactor relevant plasmas is not assured in view of the limited understanding of the mechanisms involved.

Recent experiments at JET have focussed on a number of aspects of the sawtooth problem: the development of a detailed description of the sawtooth instability and its relationship to the the q—profile; the attainment of sawtooth stabilization at the highest plasma parameters; and, finally, the investigation of the interaction between the sawtooth instability and energetic particles, both in relation to fusion performance and to the stability of the m=1 mode.

* See Appendix to IAEA-CN-53/A-I-2, this volume.
2. The Sawtooth Instability

In JET, sawteeth are accompanied by a variety of mhd activity [1] which may give rise to island–like structures with poloidal and toroidal mode numbers $m=n=1$. However, these do not appear to play a role in the sawtooth instability, which occurs as a rapidly growing helical deformation of the plasma core, also with $m=n=1$, followed by a flattening of temperature and density profiles across the central plasma, the entire event having a timescale of $\sim 100\mu s$. The topology of the sawtooth collapse in JET, as first elucidated by tomographic analysis of soft X–ray (SXR) emission measurements made with orthogonal cameras [5], resembles the convective flow of plasma suggested by the ‘quasi–interchange’ model [6] rather than the behaviour proposed by full–reconnection models (e.g. [7]).

Recently, possible limitations of the SXR analysis, particularly in relation to its poloidal resolution, have been explored by simulation of the behaviour expected from different models of the collapse [8]. In these investigations, flow patterns predicted by the ‘quasi–interchange’ and full reconnection models were used to generate SXR emission profiles and the line–integrated measurements which would be obtained from the X–ray cameras. The ‘observations’ were analyzed to produce tomographically inverted profiles such as those derived experimentally. The results are illustrated in figure 1, which shows the input and output flow patterns. There are clear differences between the reconstructions which enable the topologies associated with the two models to be unambiguously distinguished. More generally, it was found that, for the case commonly studied in JET (in which the plasma core is displaced along the major radius), the analysis technique accurately reproduced the emission profiles.

Further confirmation of the topology of the sawtooth collapse was obtained from two–dimensional reconstructions of local temperatures derived from electron cyclotron emission (ECE) measurements obtained with a 12–channel grating polychromator [9]. The method exploits the rotation of mhd structures using assumptions about their helical symmetry (for sawteeth, it is assumed that reconstructed modes have $m=n$) [10]. Application of the technique is, therefore, restricted to cases where the structure changes on a timescale somewhat greater than its rotational period. For example, analysis of the sawtooth collapse requires rotational frequencies of $\sim 10–20kHz$ and is limited to discharges with high power NBI. Figure 2 compares an ECE reconstruction of a sawtooth collapse with a SXR tomographic reconstruction from a similar discharge. Both show the formation of a crescent–shaped, or ‘convective’, structure.

It has been recognized for some time (e.g. [6]) that the rapid timescale of the sawtooth collapse presents a significant challenge to theoretical models.
FIG. 1. Comparison of simulated soft X-ray analyses of the two most common models for the sawtooth collapse: full reconnection (Kadomtsev) and ‘quasi-interchange’ (convection).

More fundamentally, it is found that the rapid growth of the m=1 arises spontaneously and cannot be explained by the increase of a linear growth rate arising from the evolution of the plasma equilibrium [11]. Figure 3 shows the growth of the instability, as deduced from the displacement of the peak of tomographically inverted SXR profiles, for three sawtooth collapses in different discharges. The initial noise level is -1cm and the displacement rises out of this noise with a growth rate \(-(25\mu s)^{-1}\). Extrapolation of this growth backwards in time yields a displacement equal to the Debye length, perhaps the smallest realistic scale for the instability, on a timescale of order \(-100\mu s\). In contrast, existing models of the instability involve changes in quantities such as q or \(\beta\) which would require timescales \(-100\text{ms}\) to generate such large growth rates.
FIG. 2. Comparison of two-dimensional reconstructions of the rapid sawtooth collapse in JET as obtained by (i) soft X-ray tomography and (ii) ECE reconstructions. The reconstructions, which were derived from data obtained on similar discharges, show a convective rather than a reconnection behaviour.
3. Current and Safety Factor Profiles

The determination of the $q$—profile or, perhaps more precisely, the way in which the $q$—profile is modulated by sawtooth activity, is one of the central problems in understanding sawteeth. On JET, the principal diagnostic of the $q$—profile is a multichannel far—infrared interferometer/ polarimeter. A new analysis of the polarimetric signals [12], which combines the constraints imposed by polarimetric measurements of the internal poloidal fields with external magnetic measurements in a self—consistent equilibrium calculation, has confirmed earlier conclusions that $q_0$ remains well below unity ($0.75 \pm 0.15$) during normal sawteeth.

By Abel—inverting the difference in polarimetric signals before and after a sawtooth collapse, it has been possible to study the variation of $q_0$ during both sawteeth and sawtooth—free periods. As the sawtooth period is significantly longer than the integration time of the diagnostic electronics, sawteeth are visible in the raw data, so that coherent averaging over many sawteeth is unnecessary. Results of this analysis for a series of discharges with ohmic and ICRF heating are shown in Figure 4. It is found that
$|\Delta q_0|$ increases with $\tau_S$, the time between sawtooth collapses, at a rate of $0.05 - 0.2 \text{s}^{-1}$. Further confirmation of such modulation of the central $q$-profile has been obtained by analysis of the expansion of the sawtooth inversion radius with increasing sawtooth duration. The difference in the inversion radii, $\Delta r_i$, of successive sawtooth collapses was deduced from tomographically inverted SXR emission profiles and a relationship $\Delta r_i/\Delta r_S = 0.06 \text{ms}^{-1}$ obtained.

A major inconsistency remains, however, in our understanding of the behaviour of the $q$-profile during sawtooth activity. The most direct determination shows that $q_0$ remains well below unity during sawteeth, and does not return to unity at sawtooth collapses, a result which is supported by the observation of $m=1$ magnetic structures immediately following the sawtooth collapse. In contrast, measurements of shear at the $q=1$ surface, as inferred from observations of the 'snake' [13] and of pellet ablation [14], indicate that very low values of shear ($dq/dr \sim 5 \times 10^{-2}$) persist throughout the sawtooth cycle. Simple considerations of resistive diffusion suggest that such persistence implies the existence of a low shear region across the plasma centre, with $q_0$ within 1–2% of unity. Such a conclusion would also be consistent with the convective flow observed during the sawtooth collapse in JET.
FIG. 5. Suppression of sawtooth activity following pellet injection in JET. Polarimetric measurements indicate that the suppression is due to $q_0$ remaining above unity.

FIG. 6. Results of an ICRH resonance position scan. The longest sawtooth-free period at each resonance position is plotted against the resonance major radius. The sawtooth inversion radius ($r_{inv}$) is indicated for each case, as is the inversion radius of the longest 'monster' sawtooth.
4. Sawtooth Stabilization

Several techniques for sawtooth stabilization have been investigated in JET, and another, lower hybrid current drive, is currently being commissioned. Sawteeth have been suppressed for up to 5s following the injection of pellets, both in the current rise phase and during the flat-top. As is illustrated in figure 5, polarimetric measurements of $q_0$ indicate that sawtooth suppression is due to a broadening of the current profile, leading to $q_0 > 1$, which results from the substantial changes in electron temperature produced by pellet ablation. This is supported by observations of mhd activity with mode numbers $(m,n) = (3,2)$ and $(2,1)$ deep in the plasma core.

The ‘monster’ sawtooth regime [3], in which sawteeth are stabilized as a result of central heating, has been the subject of detailed investigations. By scanning the position of the ICRH resonance across the plasma centre, it has been found that stabilization is produced most efficiently (i.e. the longest stable periods are obtained for a given power) when the heating resonance is located at the magnetic axis. In addition, stabilization is achieved only for heating inside the sawtooth inversion radius. The results of this experiment are shown in figure 6, where the longest stable period observed is plotted against the major radius of the heating resonance. In each case, the corresponding sawtooth inversion radius is indicated. It has also been observed that stabilization is most probable at low ICRH minority concentrations, when the stored energy in the fast particle population is maximized.

In 3MA discharges with $q_{\psi} \approx 4$, stable periods of up to 5s have been obtained and previously reported observations [1] that $q_0$ is significantly below unity ($q_0 \approx 0.6-0.8\pm0.15$) have been confirmed. In addition, long sawtooth–free H–modes are now produced routinely during ICRF and ICRF/ NBI combined heating of X-point plasmas. However, stabilization by this technique becomes more difficult as the edge safety factor decreases ($q_{\psi} < 4$). To extend the regime to higher current, therefore, auxiliary heating has been applied during the current rise phase, either just before or just after sawteeth start. The high electron temperature attained is expected to slow the inward diffusion of current, but the smaller radius of the $q=1$ surface at such times may also be a significant factor. This approach has succeeded in extending stabilization to plasma currents of 5–6MA (see e.g. [15]).

5. Interaction of Sawteeth and Energetic Particles

Interactions between the sawtooth instability and energetic particle populations, produced either by auxiliary heating or by fusion reactions,
are of growing importance as fusion plasmas approach breakeven conditions. For example, considerable attention has been given to the suggestion [16] that energetic ions, accelerated by RF fields, might be responsible for sawtooth stabilization in the 'monster' sawtooth regime. This process can occur when the average fast ion energy, $E^\text{th}$, is such that the bounce–averaged magnetic drift frequency of deeply trapped fast ions, 

$$\omega_{\text{Dh}} = \frac{E^\text{th}}{Z^\text{th}eB_\varphi R_0 r_1},$$

exceeds the mode frequency in the plasma rest frame ($Z^\text{th}$ is the hot ion charge, $B_\varphi$ the toroidal magnetic field, $R_0$ the plasma major radius and $r_1$ the radius of the $q=1$ surface).

A growing body of experimental evidence is now consistent with this interpretation. Not only is it known that ICRF minority heating in JET can accelerate minority ions to energies of order 1MeV, well above the relevant energy threshold (<100keV), but several experimental observations are as expected from theory. The result that stabilization is produced most efficiently with on–axis heating is consistent with the requirement that the hot ion pressure profile be peaked. As noted previously, low minority concentrations, corresponding to high stored energy in the fast ions, are also optimal for attaining stabilization. The difficulty experienced in producing stabilization as $q_\psi$ falls below values ~4 can be associated with the prediction that the maximum stable values of $\beta_p$ and $\beta_{\text{ph}}$ scale as $(r_1/a)^{-\alpha}$, where $\alpha = 1.5$ for $\beta_p$ and $2.0 < \alpha < 2.5$ for $\beta_{\text{ph}}$. Moreover, this observation stimulated the development of the current—rise heating scenario, which takes advantage of the smaller $q=1$ radius in the initial phase of plasma development to extend the stabilization to higher currents and lower values of $q_\psi$. Perhaps the most direct evidence for the role played by fast ions is the observation that, following switch–off of RF heating, a sawtooth collapse terminates the stable period within 200ms (often far less), i.e. on a timescale which is of the order of the average fast ion slowing down time within the $q=1$ radius.

Quantitative investigation of the theory is limited by difficulties associated with the determination of the fast ion distribution in space and energy, with uncertainties in the $q$–profile and with the general uncertainty surrounding the nature of the sawtooth instability. Nevertheless, by making several simplifying assumptions, essentially that certain local quantities can be approximated by global parameters, such a quantitative comparison between theory and experiment has been performed [17]. The analysis is performed in the $(\Gamma, H)$ plane, where

$$\Gamma = \frac{\gamma_{\text{mhd}}}{\omega_{\text{Dh}}^\text{max}}, \quad H = \frac{r_1}{s_0 R_0} \frac{\omega_A}{\omega_{\text{Dh}}^\text{max}} <\beta_{\text{ph}}> a$$

(1)

$\gamma_{\text{mhd}}$ is the ideal mhd growth rate, $\omega_{\text{Dh}}^\text{max}$ the maximum fast ion precession frequency as defined previously, $\omega_A$ the Alfvén frequency, $s_0$ the shear at
FIG. 7. Analysis of a series of sawtooth-free discharges in terms of the fast-particle stabilization theory [16, 17]. \( \Gamma \) and \( H \) are defined in the text. The solid curve encloses the region predicted to be stable. The dashed lines indicate the trajectories of the discharges in the stability diagram as the \( q = 1 \) radius expands during the sawtooth-free period.

Results of this analysis are shown in figure 7, where the area enclosed by the curve represents the stable region and the dashed lines indicate the trajectories followed by discharges as the \( q = 1 \) radius expands during a sawtooth-free period. Closed circles represent discharges in a variety of conditions (\( 4 \leq q \psi \leq 9, \ 7 \leq P_{RF} \leq 10 \text{MW}, \ <n_e> = 2-3 \times 10^{19} \text{m}^{-3} \)) where stabilization is readily obtained. The closed triangle corresponds to a pulse in which stabilization was obtained at the lowest power to date (\( P_{RF} = 1.2 \text{MW} \)) and the open symbol to one of the longest sawtooth free periods (5s). Both pulses had \( q \psi = 4.5 \). As is indicated, the experimental uncertainties in the calculations are large, and significant approximations have been applied to the theory to facilitate the analysis. Nevertheless, in the light of the more detailed experimental observations outlined above, the theory provides a working hypothesis for the development of an understanding of the stabilization mechanism.

the \( q=1 \) radius and \( \langle \beta_{ph} \rangle \) the poloidal beta associated with the perpendicular fast ion energy (which, for central ICRH, is a reasonable measure of the fast ion energy within the \( q=1 \) radius).
Comparison of local neutron emission profiles measured before and after a sawtooth collapse. Although the global neutron emission falls by only 18%, the axial emission falls by a factor of 6.
In JET plasmas with ohmic and ICRF heating, neutron emission is modulated in accordance with expectations based on the redistribution of ion thermal energy at the sawtooth collapse. However, the major contribution to the fusion yield of high performance discharges is non-thermal in nature and is due to beam-plasma and beam-beam reactions. The modulation of the reactivity due to \( m=1 \) activity and, in particular, the way in which \( m=1 \) instabilities affect energetic ions is, therefore, central to the optimization of fusion reactivity. While 'fishbone'-like bursts are observed in both NBI- and ICRF-heated plasmas [18], only at the highest \( \beta_t \) values is there evidence of detectable (<10%) modulation of neutron production. However, recent studies of neutron emission profiles derived by tomographic analysis of measurements obtained from orthogonal arrays of neutron detectors has revealed a strong modulation of central neutron production due to sawteeth [19].

Figure 8 shows neutron emission profiles calculated at timeslices just before and just after a sawtooth collapse in a JET hot-ion H-mode plasma heated with 18MW of NBI. While the modulation of the global neutron yield is small, -18\%, it is found that the emission profile is changed dramatically, the FWHM broadening from 0.36m before to 1.2m after the collapse. In addition, the axial neutron emission falls by a factor of 6, from \( 1.9 \times 10^{15} \text{m}^{-3}\text{s}^{-1} \) to \( 2-3 \times 10^{14} \text{m}^{-3}\text{s}^{-1} \). Calculations show that, in this case, the emission is almost entirely non-thermal and that the ratio of beam-plasma to beam-beam reactions is 2:1. The fall in central emission can be explained, therefore, by a substantial redistribution of the fast, beam-injected ions to large minor radii, the redistribution occurring on a timescale of less than 1ms.

6. Discussion

A complete understanding of the sawtooth instability remains elusive, but several key diagnostic observations of the sawtooth collapse phase have been confirmed. In particular, the initial phase of the collapse, which resembles the convective behaviour predicted by the 'quasi-interchange' model, has now been observed in SXR and ECE reconstructions. However, the rapid timescale of the collapse and, more fundamentally, the rapid switch-on of the instability, are not understood.

Measurements of the \( q \)-profile show that \( q_0 \) remains well below unity throughout the sawtooth cycle. The persistence of a \( q=1 \) surface is further confirmed by observations of the 'snake' and of other long-lived \( m=1 \) structures. However, the topology of the sawtooth collapse implies that \( q_0 \approx 1 \), as does the observation, derived from pellet ablation measurements, that a region of very low shear exists close to the \( q=1 \) surface at all phases
of the sawtooth. More precise measurements of the evolution of the q—profile during sawteeth are required to resolve this basic inconsistency.

Sawtooth stabilization, lasting for up to 5s, is now routinely obtained in L— and H—mode plasmas and at currents of up to 6MA. Its importance has been emphasized by recent demonstrations of local modulation of fusion reactivity and the redistribution of high energy ions due to sawteeth. In cases where stabilization follows pellet injection, the dominant effect is the modification of the current profile, due to pellet ablation, which causes q₀ to rise above unity. On the other hand, in the ‘monster’ sawtooth regime q₀ is below unity and there is growing evidence that the stabilization is due to an energetic particle population.

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References


DISCUSSION

K. TOI: I have a question on sawtooth stabilization relating to monster sawteeth. It concerns the behaviour of the plasma near the edge. How does the $H_a/D_o$ emission behave, and how do the density and temperature profiles change during the monster sawtooth?

D.J. CAMPBELL: There is no significant change in the edge temperatures and density or in the edge $H_a$ emission during the sawtooth-free period.

K. McGUIRE: Does $\chi_e$ from the heat pulse propagation (hpp) agree with $\chi_e$ from the power balance, firstly at the monster sawtooth, and secondly at low $q_L$, with a high current of 5–7 MA? Does the value of $\chi_e$ from the hpp change for the monster sawtooth or as a function of $I_p$? And does $\chi_e$ from the hpp depend on $\Gamma_{inv}$?

D.J. CAMPBELL: $\chi_e$ cannot be derived from hpp measurements at the monster sawtooth collapse since the inversion radius, and hence the mixing radius, is large. As a result, the region of plasma over which the hpp analysis can be performed is small and localized to the plasma edge. No useful information has, therefore, been obtained from such analysis in the case of monster sawteeth.

Most hpp results in JET have been obtained in plasmas at moderate q, and I have no information on how \( \chi_e \) derived from hpp measurements in JET depends on \( I_p \) or \( \Gamma_{\text{inv}} \).

Y. KAMADA: With regard to the tomography of the sawtooth crash, do you find any changes in the crash characteristics as a function of the beta value and elongation of the q = 1 surface? Secondly, can you describe the characteristics of the monster sawtooth crash?

D.J. CAMPBELL: The topology of the sawtooth collapse does not appear to be affected significantly by the plasma elongation over the range of parameters investigated. At the highest \( \beta_i \) values, we have not reconstructed the sawtooth collapse as the strong coupling of modes at different rational surfaces precludes the use of two-dimensional tomographic analysis.

To answer your second question, the topology of the crash of the monster sawtooth is exactly as has been described for normal sawteeth.

R.J. GOLDSTON: Four years ago at Kyoto the JET team made the interesting observation that \( \chi_e^{\text{bpp}} \) was of about the right magnitude to fit with \( T_E^{\text{pc}} \). In the intervening period, have you attempted to correlate \( \chi_e^{\text{bpp}} \) with \( T_E^{\text{pc}} \) as a function of, for example, \( I_p \), \( B_t \), density, \( Z_{\text{eff}} \) or any other parameter?

D.J. CAMPBELL: The most systematic analysis of \( \chi_e^{\text{bpp}} \) performed on JET data revealed a dependence on the square root of \( Z_{\text{eff}} \) and average electron temperature, but no apparent dependence on power or density. The relationship between \( \chi_e^{\text{bpp}} \) and \( T_E^{\text{pc}} \) has been investigated for individual cases, but I am not sure whether a systematic correlation has been carried out. Such analysis as has been performed is already described in the literature.
ROLE OF THE EDGE ELECTRIC FIELD AND MICROTURBULENCE IN THE L-H TRANSITION *

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Abstract

ROLE OF THE EDGE ELECTRIC FIELD AND MICROTURBULENCE IN THE L-H TRANSITION.

Two clear signatures of the L-H transition in the DIII-D tokamak are that the edge poloidal rotation velocity $v_e$ abruptly increases, implying that the radial electric field $E_r$ becomes more negative, and the level of density fluctuations is abruptly suppressed in a narrow region at the plasma edge. The absolute values of $E_r$ and $v_e$ have maxima within one to two poloidal gyroradii $p_{\phi}$ of the plasma edge, indicating that large negative and positive gradients of $E_r$ and $v_e$ exist near the plasma edge in a region called the 'shear' layer. Density fluctuations are reduced, and the gradients of $T_i$, $T_e$, $n_e$, and the carbon density increase in the shear layer at the time of transition. These data qualitatively support theoretical scenarios in which changes in $E_r$ or in the gradients of $E_r$ or $v_e$ cause a suppression of microturbulence, resulting in increases in the $T_i$ and density gradients.

INTRODUCTION

Theoretical predictions that $E_r$ should change from the L- to the H-mode [1,2] were followed by the observation in the DIII-D tokamak that $E_r$ within 1–3 cm of the plasma edge was negative in the L-mode and became more negative in the H-mode [3–6]. The poloidal flow velocity of He ions was also observed to increase at the transition. A large body of observations based on emission spectroscopy of He II [3,5,7] show that these changes occur independently of the heating method used to produce the H-mode: ECH-heating only, Ohmic-heating only, co-injected NBI heating and counter-injected NBI heating. These changes also occur when the H-mode is produced in a plasma limited by a material limiter or when

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the toroidal field is reversed from its normal direction [4-6]. In all cases, the changes in $E_r$ and $v_\theta$ start to occur either simultaneously with or a few milliseconds before the drop in the $D_\alpha$ signal at the transition. Based on these experimental results, Biglari, Diamond and Terry noted that an increase in $|dv_\theta/dr|$ could suppress microturbulence [8], and Shaing and Crume extended their theory and emphasized that a more positive value of $dE_r/dr$ could reduce the amplitude of microturbulence [9], leading to improved confinement in the H-mode. Experimental data from DIII-D show that each of these requirements for suppression of fluctuations is fulfilled in the H-mode plasma [5]; furthermore, application of a criterion developed in Ref. 8 shows that $|dv_\theta/dr|$ is large enough to suppress modes with wavenumbers greater than 0.1 cm$^{-1}$, a range which covers the modes which are thought likely to be present at the plasma edge [5]. Additional evidence for the role of $E_r$ is that the transport barrier, originally discovered on ASDEX [10], is correlated with shear in $E_r$ and $v_\theta$, as shown in recent data from JFT-2M [11]. Very strong evidence for the role of $E_r$ comes from work in the CCT tokamak in which the H-mode was produced in Ohmically-heated discharges by biasing the plasma with an internal electrode [12].

Several scattering experiments have shown a general decrease in microturbulence during the H-mode [13-16]. FIR scattering data on DIII-D also shows a narrowing of the $k$ spectrum in the H-mode [15]. Reflectometer measurements, which have much better spatial localization than scattering measurements [17], have also shown a drop in microturbulence in the H-mode [18-21]. The reflectometry system on DIII-D has shown for the first time that coincident with the $D_\alpha$ drop there is a reduction of turbulence in a narrow zone which extends inwards from the separatrix for several centimeters [19-21]. Although it has been shown that the reduction of fluctuations in this “suppression” zone is correlated with the changes in $E_r$ [6], it has not been demonstrated that confinement improves throughout the suppression zone or that the suppression and shear layers are identical. This paper presents data from recent well-diagnosed DIII-D discharges for which profiles of $v_\theta$, $E_r$, $T_i$ and the density of carbon were obtained with 1 ms time resolution and reflectometer data were obtained with 50-100 microsecond resolution. The results show that, within error bars, the shear layer and suppression zone overlap and are established simultaneously; furthermore, the confinement in this layer, as indicated by the gradients of $T_i$, $T_e$, $n_e$ and carbon density, improves markedly at the onset of the H-mode. These results have been established for a small number of well-diagnosed discharges for which profiles of $E_r$ are available by the use of active CER spectroscopy. They are also consistent with the large body of data obtained from passive spectroscopy of He II [3-7].
RESULTS

Simultaneous measurements of $T_j$, $v_\theta$ and $v_\phi$ have been made with the techniques of passive emission spectroscopy [7], for which a single point measurement was possible, and with active CER spectroscopy, for which profiles are available [22]. In either case, $E_r$ is calculated from the force balance equation for a single ion species [3-7]. The contribution of the pressure gradient term is estimated to be 20%-30% of the $V \times B$ term for the single point measurement and is measured to be 10%-20% when profiles are available. Due to atomic physics effects near the plasma periphery, it is possible that the pressure gradients and therefore the poloidal rotation velocities of different species are not the same. However, the value of $E_r$ determined from the sum of the $V \times B$ and pressure gradient terms applies to all of the plasma species. The instrumentation used to measure profiles [22] has overlapping views from horizontally-viewing and vertically-viewing systems, with each system having eight chords separated in minor radius by 1.5 cm. The nominal spatial resolution is 0.5 cm. Data have been taken by measuring the C VII flow velocity, as inferred from the Doppler shift of the C VI 5290 Å line produced by charge transfer between a neutral beam and C VII ions.

Density fluctuations are studied with an O-mode reflectometer system, consisting of seven discrete homodyne channels spanning 15-75 GHz, with corresponding critical densities of $2.8 \times 10^{18}$ to $7 \times 10^{19}$ m$^{-3}$. An X-mode system, using a frequency tunable BWO source is used to obtain fluctuation data from critical densities intermediate to those available from the O-mode system [21], thus increasing the spatial resolution.

The time histories of several edge parameters during an L-H transition are illustrated in Fig. 1 for a discharge with $I_p = 1.6$ MA, $B_T = 2.1$ T, $P_b = 8$ MW, and for which D$^+$ beams were injected into a D$^+$ plasma. Figure 1 shows that for a chord about 1 cm inside the separatrix, large changes in $v_\theta$ and $E_r$ occurred in a time much less than 1 ms coincident with the transition. These results are consistent with the results obtained from the emission spectroscopy measurements, which showed that the changes in $E_r$ and $v_\theta$ were coincident with the transition, with the exception of some observations showing precursors a few ms prior to the transition [5]. Although the issue of causality has not been settled, a possible explanation for the precursors is that they were caused by motions of the He II emission shell further into the shear region in the L-mode. The location of the He II shell is determined by atomic physics and could change as the edge $n_e$ and $T_e$ profiles evolve in the L-mode. Similar effects could explain the changes in $E_r$ and $v_\theta$ after the transition, which occurred on
Fig. 1. $D_\alpha$ (a) shows that L-H transition occurred at 1859 ms, as indicated by dashed line. (b), (c), and (d) show $T_i$, $v_\theta$, and $E_r$ from $\vec{V} \times \vec{B}$ term obtained with 1 ms time resolution for location approximately 1 cm inside separatrix. (e) shows integrated power in range of 0–400 kHz from reflectometer channel viewing the same location. $E_r$ and $v_\theta$ increase and fluctuation level decreases immediately at time of transition. Increase in $T_i$ is delayed for a few ms.

the timescale for edge profile changes [5,7]. As shown in Fig. 1(e), the fluctuation level, monitored by a reflectometer channel 1–2 cm inside the separatrix, decreases coincidentally with the drop in the $D_\alpha$ signal. The decrease in fluctuation levels near the separatrix is a consistent feature of all L-H studies in DIII-D.

Figure 2 compares profiles of several quantities obtained in the time intervals 1–2 ms prior to the transition, 1–2 ms after the transition and 10–11 ms after the transition. The changes in poloidal rotation [Fig. 2(a)]
Fig. 2. Profiles of $v_\theta$, total $E_r$, $T_i$ and C VII density vs major radius 1–2 ms before, 1–2 ms after and 10–11 ms after L-H transition displayed in (a), (b), (c), and (d) respectively. $T_i$ and density gradients increase in region of sheared $E_r$ and $v_\theta$ immediately after transition and continue to increase during ELM-free phase of H-mode.
and $E_r$ [Fig. 2(b)] occur abruptly at the time of transition and persist during the H-mode. The changes are seen most clearly on the chord at 2.27 m and may extend to one chord (1.5 cm) on either side of 2.27 m. The region of shear in $E_r$ extends over 3-5 cm. The peaks in the profiles of $E_r$ and $v_\theta$ occur about 0.8 cm inside the nominal separatrix and the poloidal ion gyroradius $\rho_{i\theta}$ is about 0.7 cm at the time of the transition. The uncertainty of the separatrix location, determined from magnetics, is about 1 cm. These results are supported by the data from emission spectroscopy which showed that large values of $E_r$ and $v_\theta$ existed 1-3 cm inside the separatrix and that the $E_r$ and $v_\theta$ profiles were sheared just inside the separatrix [5,7].

The gradients of $T_i$ [Fig. 2(c)] and the carbon density [Fig. 2(d)], measured from the amplitude of the C VI signals, increase abruptly in the shear region at the onset of the H-mode, demonstrating very high correlation between the formation of the transport barrier and the increase in shear. As the H-mode progresses, the $T_i$ and carbon density gradients remain approximately constant in the shear layer, but the boundary values of $T_i$ and the carbon density increase, so that pedestals are produced on the $T_i$ and density profiles. The $T_i$ and carbon density gradients for locations more than 3-4 cm from the edge show little, if any, change from the L-mode to the H-mode, indicating that the transport barrier forms only in the shear layer. Data obtained from a Thomson scattering system with closely spaced channels near the separatrix show that $T_e$ and $n_e$ also steepen in the region of shear coincident with the transition.

As shown in Fig. 3, the suppression zone at the time of transition as determined from the reflectometry channels is approximately 4 cm wide with an uncertainty of ±1 cm. The width and location of this zone is identical to that of the shear layer within the error bars. Data obtained for a 0.8 MA discharge, for which $\rho_{i\theta}$ is twice that of the 1.6 MA discharge, show that the suppression and shear layers have about the same width as for the 1.6 MA discharge, although the error bars on the shear width are slightly larger. For both the 0.8 and 1.6 MA discharge, the suppression zone coincides with the steep $n_e$ gradient, as is expected if a reduction of fluctuations improves the local confinement.

CONCLUSIONS

A large body of data have been obtained in the previous two years from the DIII-D tokamak regarding the role of the edge $E_r$ and $v_\theta$ in the L-H transition. This work has demonstrated the existence of a shear
Fig. 3. (a) Ratio of power from reflectometer channels after the transition to power prior to transition. No reduction in signals occurs outside suppression zone with roughly order of magnitude reduction in suppression zone. (b) Suppression zone for microturbulence at onset of L-H transition is shown in shaded region, superimposed on $n_e$ profile obtained 7 ms before the L-H transition. Width of zone is determined by finding largest critical density whose signal is suppressed coincident with L-H transition, under assumption that $n_e$ profile has not changed before the instant of transition.

layer in $v_\theta$ and $E_r$ at the plasma edge and that $v_\theta$ abruptly increases and $E_r$ becomes more negative at the L-H transition. Profile data show that changes in $v_\theta$ and $E_r$ and suppression of density fluctuations occur simultaneously in the same layer at the edge of the plasma as soon as the $D_\alpha$ signal starts to drop at the L-H transition. Profiles of $T_i$, $T_e$, $n_e$ and of the carbon density show that the region of improved confinement near the edge also overlaps the shear layer. These results are consistent with recent theoretical ideas which indicate that changes in the gradients of $E_r$ or of $v_\theta$ or changes in the magnitude of $E_r$ itself cause a suppression of microturbulence which leads to confinement improvement at the L-H transition [8,9].
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DISCUSSION

K. TOI: I would like to ask about the connection between edge turbulence suppression and the depression of $H_a/D_a$ emission. In some H-mode shots, $v_e$ starts to change a few milliseconds before the transition. However, the edge fluctuations are not suppressed until the transition. This experimental evidence obviously contradicts Biglari's theory. Could you comment on this?

R.J. GROEBNER: We have never seen the edge fluctuations change before the drop in $D_a$. On the other hand, our emission spectroscopy measurements have, on occasion, shown changes before the drop in $D_a$. With the knowledge obtained from our recent active CER measurements, I am concerned that the emission data could give false indications of precursors if the He II emission region moves in the L-mode. Changes of the edge $T_e$ and $n_e$ profiles during the L-mode could move the He II emission region by perhaps 1 cm into a region of higher poloidal flow velocity.
G. FUSSMANN: You showed dramatic changes in the poloidal rotation within 1.5 ms. Undoubtedly, very large forces are needed to produce these abrupt changes, since we know, for instance from toroidal measurements in the case of NBI heating, that the plasma speeds up gradually. Do you have any explanation for the origin of these forces?

R.J. GROEBNER: We do not have a satisfactory quantitative explanation for the rate of change of poloidal rotation or for the steady state value. Simple order of magnitude estimates of thermal ion orbit losses can explain our steady state values of poloidal rotation, but a better model for the rate of ion orbit loss and the poloidal damping rate is required. Neoclassical theory predicts values that are somewhat low compared to the measured values, but there are significant uncertainties in the theory for our mixed collisionality regime. We also see dramatic changes in the toroidal rotation velocity in the shear layer; the changes are much slower for smaller radii.

A. GIBSON: In the examples you showed of the increase in ion temperature gradient at the L–H transition, the actual temperature reached was nevertheless the same at the same place further in, in both L- and H-phase. What is the reason for this?

R.J. GROEBNER: I interpret the ion temperature data as showing that the transition starts in the shear layer and then propagates inwards. The ion temperature at radii less than that of the shear layer starts to increase a few milliseconds after the transition and continues to increase to very high values, even near the plasma edge. This increase continues for about 100 ms for the data I presented.

P.R. THOMAS: As both you and Dr. Wagner have pointed out at this meeting, this is a dog and tail problem. The plasma is presumably in mechanical equilibrium, so the increased electric field is an indication of the change in pressure gradient at the L–H transition, resulting in the improved confinement. How can we tell instrumentally whether dog or tail is doing the wagging?

R.J. GROEBNER: If we can show beyond any doubt that changes in the electric field occur a few transport times before the transition, we will be able to rule out the electric field as the causal agent. If we continue to find that the changes are coincident with the transition, we will not be able to prove causality. Since transport time scales are ~10–100 ms near the edge, I doubt whether our diagnostics have the necessary time resolution. In that case, we must rely on a theory which satisfactorily integrates results from machines with spontaneous and externally stimulated transitions.

Y. MIURA: I would like to comment on the change of ion temperature after L–H transition. You measured the impurity ion temperature and found that it changes after 10 ms, but our measurements in JFT-2M by the time of flight method indicate that the change in the hydrogen ion energy distribution is faster than the drop in $\text{H}_\alpha/\text{D}_\alpha$ intensity by about 400 $\mu$s.

R.J. GROEBNER: We cannot definitively rule out the possibility that changes in the ion temperature start to occur shortly before the transition. If the changes occur 400 $\mu$s before the transition, we might not see them because of our 1 ms time resolution. It is possible that the measurements in JFT-2M and in DIII-D are sampling different parts of the ion energy distribution and that changes occur first in the high energy ions.
ENHANCED CONFINEMENT IN CCT AND PISCES-A IN THE PRESENCE OF RADIAL ELECTRIC FIELDS

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Abstract

ENHANCED CONFINEMENT IN CCT AND PISCES-A IN THE PRESENCE OF RADIAL ELECTRIC FIELDS.

Plasma edge rotation induced by radial currents injected into a tokamak is observed to produce H-mode like behavior. The particle confinement time increases by a large factor (> 5) at all densities. The central energy confinement time improves by a smaller amount (< 2) mainly as a result of the extension of the linear confinement regime to higher densities (larger Murakami parameters). Radial correlation measurements of the $E \times B$ turbulence imply an electron fluid mixing approximately on the collisionless skin depth scale, which may be responsible for the large electron heat conduction. The fluctuational particle transport can be counteracted by an externally applied DC radial electric field. The anomalous electron heat conduction cannot be eliminated but may be mitigated by indirect means, e.g. the edge plasma rotation, resulting in better MHD stability at higher densities. Giant sawteeth have been observed in these “edge rotation stabilized” ohmic discharges, with the central energy confinement time attaining large values at high densities, limited by an MHD-like hard collapse of the sawtooth and/or of the plasma channel. Radial particle transport in a sheared electric field has also been studied in the PISCES-A linear plasma device. Drastic reductions of both the fluctuation-induced transport and the steady state $E_\theta \times B_\theta$ convection are observed only when a double shear electric field layer is formed at the plasma edge.
1. INTRODUCTION

Previously, in the Macrotor tokamak, we have shown that electron injection\[1\] into the core of a tokamak plasma can change the radial sheath potential from positive to negative. It was also observed that the particle transport was modified\[2\] and mass accumulation could be induced. The resulting discharges had features common to the H-mode\[3\] discharges as could be seen from the temporal and spatial evolution.

In the past year we have returned to the edge biasing experiments in the Continuous Current Tokamak (CCT)\[4\]. Our experiments have demonstrated the connection between damping of the poloidal rotation and bifurcation\[5\]. Since the re-initiation of these experiments, discharges with radial potentials up to 5 kV have been produced. In the present paper we only present data from plasmas with low voltage negative bias. These discharges duplicate most aspects of H-mode behavior found in other tokamaks\[6\].

\[\text{Figure 1. The electrode geometry used for driving the poloidal rotation and heating the plasma edge. } r_b \text{ is the radius of the bifurcation layer. }\]

\[\text{Figures 1b-c show typical oscillograms of the tokamak plasma. Matrix arrays of Langmuir probes are used for measuring plasma radial electric fields, rotation and correlation times and lengths. The optional heater current shown in the figure generates both a reverse EMF and a reverse current in the driver circuit in accordance with MHD theory.}\]
2. EXPERIMENTAL CONFIGURATION

The experiments were carried out in CCT ($R = 1.5$ m, $a = 0.35$ m, $B = 0.25$ T, $I = 50$ kA, $n < 2 \times 10^{19}$ m$^{-3}$, $T_e = 150$ eV). The electrical circuit used for modifying the edge plasma ambipolarity and rotation is shown in Fig. 1a for a two electrode configuration.

Figures 1b-e show histograms of the two electrode experiments. First, a tokamak discharge is initiated in the usual way and the plasma position, current and density are stabilized. The plasma rotation is slowly increased until a rapid bifurcation in the rotation is achieved. The turbulent properties are evaluated at various times into the discharge. A second electrode ("heater" in Fig. 1a) may be inserted and optionally used to inject energetic electrons at an arbitrary radius.
3. E x B TURBULENCE PROFILES

Measurements of edge turbulence and particle transport have been made during the transition from CCT ohmic discharges to fully bifurcated H-mode discharges. Fluctuation levels and electrostatic turbulent driven particle fluxes taken at the outside midplane are shown in Fig. 2a-2c. Measurements were taken just before (at t_0 in Fig. 1c), during the radial current ramp, and during the bifurcated H-mode. In the scrape-off layer (r > 35 cm), the density fluctuation levels are not significantly changed (Fig. 2a), while the floating potential fluctuation levels are reduced substantially (Fig. 2b) as the rotation rate is increased. The turbulent driven particle flux, as measured by standard techniques [7], is also reduced by factors of 3-6 at the limiter radius as the rotation rate is increased (Fig. 2c). Assuming symmetry, the radial particle flux at the limiter radius implies a global particle confinement time of 5-8 msec during ohmic discharges and 20-40 msec during the CCT H-mode.

Although the above fluxes are not in contradiction with global values of \( \tau_p \), we have evidence that other transport mechanisms (inward pinch due to the DC radial E field, large scale quasi-stationary convection due to \( E_\theta \times B \) drift) can also be present. These mechanisms can locally dominate fluctuation driven transport and we are studying their effect on global particle balance.

4. POLOIDAL AND RADIAL CORRELATION PROPERTIES

Matrix-style arrays of Langmuir probes have been used to determine fluid flow and space-time correlation functions in ohmic and H-mode discharges. Turbulence with \( f > 100 \) kHz is used for the data presented below. Fig. 3a shows the cross-correlation between poloidally separated probes as a function of radius during bifurcation (t = 67 msec). During the H-mode, the radial correlation function slightly narrows on the low field side and widens on the high field side. The poloidal propagation delay reduces by a factor of 5-10 due to the strong plasma rotation \( U_\theta \) (shown in Fig. 3b). The typical correlation time in the rotating frame is unchanged, however, and is \( \tau_c \approx 3 \mu \text{sec} \). The 2-point estimation of the poloidal correlation length \( L_{\text{pol}} \), defined as \( \tau_c U_\theta \), is also plotted as a function of radius. From the data of Fig. 3a, \( L_{\text{pol}} \) varies between 0.5 cm (ohmic) and 18 cm (bifurcated H-mode), depending on poloidal rotation frequency. The fluctuations are incoherent and evanescent on the scale of the correlation length (0.5 cm), which is also a typical value of the collisionless skin depth. Due to the plasma motion, detected fluctuations undergo a doppler shift: the spectral bandwidth increases as \( U_\theta \) and the
Figure 3. (a) Pair correlation functions in the poloidal direction at 1 cm separations. (b) Summary data of the correlation lengths in the laboratory frame and the amplitude profile of the density fluctuations in the L-mode (dots) and the H-mode regimes.

low frequency amplitude diminishes as $U_\theta^{-1}$, although $L_\theta^{\text{fluid}}$ remains $< 1$ cm. These features indicate that the turbulence is largely invariant in the fluid frame, in substantial disagreement with some of the current theories[8].

Improvements in the central energy confinement time cannot be directly related to the changes in the edge plasma. However, improved central confinement at higher density is observed.

5. HIGH DENSITY EDGE STABILIZED DISCHARGES

In the presence of the rotating edge layer, the Murakami density limit is never observed. We attribute this to improved MHD stability of the plasma; large scale density and magnetic Mirnov oscillations are reduced or eliminated. A discharge without edge stabilization is shown in Fig. 4a. These discharges are density saturated and are reproducible. Two edge stabilized discharges are shown in Figures 4b and c for comparison. The density rise in these discharges is terminated by an internal disruption related to the collapse of the giant sawtooth. The collapse is not related to the radiation limit.
Figure 4. (a) Ohmic plasma without radial edge current. The fluctuations in the interferometer trace (density fringe) are due to large scale MHD activity. (b) and (c) Reduction of MHD fluctuations and appearance of giant sawteeth in edge rotation stabilized discharges. Radial bias current is 20 A.

6. SUMMARY

Large improvements in the particle confinement times can be produced by DC radial electric fields. Effects related to the applied bias can easily counteract the particle diffusion normally observed in ohmic discharges. Heat transport cannot be directly controlled by the radial fields since the intrinsic $E \times B$ fluctuations are not really suppressed throughout the temperature profile. However, a "copper shell-like" stabilizing influence of the rotating layer appears to contribute to central confinement improvements at high densities as evidenced by the formation of giant sawteeth.

7. CONCLUSIONS

Our experiments corroborate that the control of the edge properties is important for the confinement of high density plasma. While we believe that a direct control of the particle transport at the edge is both necessary and achievable for tokamak reactors, a reduction of the heat transport is neither likely nor necessary.
Part II: REDUCED EDGE TURBULENCE AND CONVECTION IN PISCES-A

Detailed studies of particle transport in the presence of an electric field layer are also carried out in the PISCES-A device (a linear reflex arc steady-state discharge)[9]. Both fluctuation induced transport and steady state convection are investigated. The helium plasma, produced by a hot LaB$_6$ cathode, flows towards a non-conducting plate with a circular aperture 3.5 cm in diameter. On the downstream side of the aperture plate, an annular biasing electrode (inner diameter 4.2 cm) is mounted coaxially with the plasma. Fast scanning probe assemblies are used to record the radial profiles of ion saturation current, floating potential, and plasma potential (measured with an emissive filament). Fluctuation spectra and the fluctuation-induced particle flux are inferred with a frequency response up to 500 kHz.

The radial profiles of radial electric field and plasma density are shown in Fig. 5a and 5b. These results were obtained at moderate collisionality, where the effect of the perpendicular ion mobility on transport is negligible. Without biasing the annular electrodes, the electric field increases abruptly at the aperture radius and is almost constant outside the aperture. A rotating boundary layer is formed as a result of the $E_r \times B$ drift. This rotating layer, intrinsically present in a reflex discharge, can destabilize Kelvin-Helmholtz modes[10] and cause large radial particle losses. Fig. 5c and 5d show the density and potential fluctuation levels without and with positive bias. The fluctuation spectrum extends to several hundred kHz, with the main portion of the spectral energy content in the range $f < 200$ kHz. The radial fluctuation-driven particle flux is shown in Fig. 5e. The particle flux is found to be proportional to the square of the $E_r \times B$ velocity in the rotating layer. This scaling, $\langle n_i \delta \rangle \propto (E_r/B_z)^2$, has been observed in particle simulations for the saturated Kelvin-Helmholtz instability[11]. A poloidal component of the steady-state electric field ($E_\theta \approx 2 - 10$ V/cm) has also been observed close to the aperture radius. The convective transport driven by this field is poloidally asymmetric and can be directed radially inwards on one side of the plasma column and outwards on the opposite side. More detailed results will be published elsewhere.

An electric field profile very similar to the spatial structure observed in the CCT plasma for low values of the applied potential (< 130 V), can be produced by biasing the annular electrode. The radial electric field increases to about 2-3 times the value for the unbiased case at the aperture radius and decreases further outwards. The resulting poloidal flow inferred from the $E_r \times B$ drift velocity resembles the profile of a
Figure 5. a) Radial electric field $E_r$ versus radial position for the single shear layer case (no bias) and the double shear layer case; $B_z = 0.17$ T; neutral pressure $p_{He} = 4 \times 10^{-4}$ torr. b) Radial plasma density profile for both cases. The aperture radius ($r = 1.7$ cm) and the position of the annular biasing electrode are indicated. c) Normalized fluctuation level $\tilde{n}/n$ versus radius. d) Normalized potential fluctuations $e\phi/kT_e$ versus radius. e) Fluctuation driven radial particle flux $\Gamma^{fl}_\perp$ versus radius.
jet of fluid in a quiescent background. In this case, both density and potential fluctuations are greatly reduced, with the maximum reduction observed outside the aperture radius. The electric field profile is radially symmetric with respect to the axis and steady-state convection is nearly eliminated ($E_\theta < 0.3 \text{ V/cm}$). The fluctuation induced transport decreases by a factor of 2-10. The edge density gradient and to some extent the central density increase as a result of enhanced radial confinement. The net radial transport inferred from the particle balance equation is reduced by a factor of 3-6 at the aperture radius, and is a factor of 2-3 below Bohm diffusion.

These experiments suggest that the magnitude of the radial electric field and the symmetry and spatial structure of the electric field layer have an important influence on particle transport. We have shown in PISCES that a large degree of poloidal symmetry, brought about by edge plasma biasing, substantially reduces radial particle transport driven by fluctuations and steady-state convection.

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K. TOI: I have a fundamental and simple question concerning discharge behaviour. Your experiment is very important for any discussion of the cause of the L–H transition, as was pointed out in the previous presentation. However, the H mode discharge in CCT does not bear the clear signature of a typical H mode. In typical H modes, the $H_\alpha/D_\alpha$ signal is suddenly changed (depressed or increased) at the L–H and the H–L transitions, whereas in CCT the $H_\alpha/D_\alpha$ signal changes only gradually. Could you comment on this?

R.J. TAYLOR: We can have both gradual and sharp changes in the $H_\alpha$, which is monitored by a horizontal chord spectrometer. In a rapidly driven bifurcation the $H_\alpha$ (the transport improvement) is fast. If we approach bifurcation slowly, then the improvement is gradual, a fact which highlights the role of small rotation.

A. HASSAM: We normally understand plasma flow as being composed of an $E \times B$ drift and a diamagnetic drift. From your data, it seems that the diamagnetic flow is almost a factor of 10 smaller than the $E \times B$ flow, after transition. Allowing for the fact that it is the shear in flow as opposed to absolute flow that counts, would you say, on the basis of your data, that we can discount the diamagnetic portion of the flow in accounting for the phenomena associated with H mode?

R.J. TAYLOR: In general, no. But at high velocity limits the barrier can be so large that you can ignore the grad P, which is equivalent to ignoring diamagnetism.
H MODE BEHAVIOUR INDUCED BY RADially OUTWARD FIELDS IMPOSED IN TEXTOR


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Abstract

H MODE BEHAVIOUR INDUCED BY RADially OUTWARD FIELDS IMPOSED IN TEXTOR.

Edge polarization experiments were performed on TEXTOR showing that H mode behaviour can be set up with both positive and negative radial electric fields. The H mode can be sustained in the presence of auxiliary heating, and improvements in energy confinement of up to about 1.4 and in particle confinement of 2 to 2.5 were obtained with positive bias. First indications are that the particle confinement might increase more with negative bias. The viscous damping of poloidal rotation is found to be in good agreement with some versions of neoclassical theory. An expression can therefore be derived for the minimum non-ambipolar radial current that must be set up in the plasma edge to provoke an L to H transition.

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1. INTRODUCTION

Experiments on edge polarization similar to the pioneering work on the Continuous Current Tokamak (CCT) [1,2] were performed on TEXTOR. The results confirm and substantiate the model according to which the H mode is induced when the plasma edge transits from a state of slow poloidal rotation \( U_p \) (low \( E_r \)) to one of fast rotation (high \( E_r \)) [3]. Whereas the rotation can be set up by an initial charge loss from the plasma induced, for example, by auxiliary heating [4], probes can also be used [1,2,5]. The bifurcation behaviour is attributed to the existence of a maximum in the viscous force opposing poloidal rotation [4], as qualitatively verified on CCT [2]. Here we present a more detailed comparison with theory. Concomitant with the transition in the edge plasma parameters which are typical for H mode behaviour, improvements in energy and particle confinement are observed. These are presently thought to be due to the generally observed [6,7] suppression of fluctuations. Since our results show H mode behaviour and confinement improvement for both signs of \( E_r \), albeit not totally symmetric, they do not contradict a possible role of flow shear as the mechanism for suppression of turbulence by decorrelation [8].

2. EXPERIMENTAL RESULTS

2.1. System description

A biasing electrode was introduced up to 8 cm (standard for the present work) beyond the toroidal belt limiter (\( a = 46 \text{ cm} \)) in full grade TEXTOR plasmas (\( D_2 \) gas, \( B_t = 2 \text{ T} \), \( I_p = 170 \text{ kA} \), \( n_e = (1-2) \times 10^{13} \text{ cm}^{-3} \), \( T_e0 = 1-1.5 \text{ keV} \)). The electrode does not itself become a limiter and is capable of sustaining the heat load from the regular discharge and the extra load connected with the biasing. The applied voltages are typically 600–900 V, the electrode current reaching about 150 A before transition and 20–40 A during the H mode. The minimum penetration length is about 5 cm.

2.2. Positive polarization

The viscous friction that opposes poloidal rotation has a maximum value [2,4]. The access to the H mode is only possible by drawing an electrode current \( I_E \) which allows this maximum to be overcome. For negative polarization a non-emissive probe is limited to the ion saturation current. We observe under most conditions that the latter is too low to exceed the said extremum. Most of our results pertain therefore to positive polarization (radially outward pointing fields).

We reported previously [5] the global characterization of the polarized Ohmic plasmas, displaying clear H transitions showing the characteristic features of density rise, profile steepening and \( D_e \) drop. The energy increases combined with an aver-
age observed loop voltage increase of 5% led to improvements in energy confinement
time $\tau_E$ of up to 40%. These results have since been extended in the presence of
auxiliary heating and the confinement properties analysed in more detail. Figure 1(a)
shows most of the features distinguishing polarized from unpolarized cases. Two dis-
charges are compared: discharge A with bias ($V_E = +850$ V) and discharge B
without bias. The initial density in A has been chosen to reach the same OH line den-
sity $\bar{n}_e$ as in B at time $t_1 = 1.13$ s. Figure 1(b) shows the electron temperature pro-
files in A and B at this time, exhibiting a slight contraction with bias. At $t = 1.15$ s
NB co-injection is applied at a level of 250 kW. The discharge stays in the H mode
until the moment the bias is removed. In A the energy is higher than in B (time $t_2$)
at the same temperature (see Fig. 1(c)) but higher density. The density profiles
remain rather similar in shape except in the region between the electrode and the
limiter, where a strong steepening is revealed by Li beam probing [5].

The stored plasma energy $E$ can double upon transition, partly due to the density
increase. Since $E$ depends on $\bar{n}_e$ in standard TEXTOR Ohmic and NBI heated
plasmas, we obtain the real gain by normalizing to constant $\bar{n}_e$. This is shown in
Fig. 2(a, b), where L and H conditions are compared for Ohmic and beam heated
discharges, respectively. For the NBI case both the thermal energy $E_{\text{kin}}$ and the total
energy including the non-thermal beam components $E_{\text{tot}}$ are given. It is seen that the
polarization brings about an increment which is roughly density independent. The
relative improvement of $\tau_E$ is then weakest at high density. At $\bar{n}_e = 1.4 \times 10^{13}$ cm$^{-3}$ it reaches about 1.2 in Ohmic and 1.3 in beam heated discharges.

The changes in the global particle confinement time $\tau_p$ inferred from the recy-
cling flux ($D_\theta$) are stronger. In a normal Ohmic discharge the recycling is divided
between the main limiter, the (unbiased) electrode and the wall in the ratios 0.55, 0.3
and 0.15. Concomitant with the density rise accompanying the transition, all three
fluxes drop. Normalized to the same density the decreases reach, for Ohmic dis-
charges, factors of about 1.75 at the limiter, 4 at the electrode and 4 at the wall. A
$\tau_p$ improvement of 2 to 2.5 results.

Radial electric field measurements were performed by means of a double probe
system in the H mode Ohmic discharges. The $E_r$ profile peaks at $r = 43.7$ cm, has
a width (FWHM) of about 1.5 cm and reaches a maximum value of about 250 V/cm
when 880 V is applied. Rotation speed measurements, using the correlation in the
fluctuations of the floating potential between poloidally spaced probes, yield poloidal
rotation speeds $U_p$ of $1.0 \times 10^4$ m/s (i.e. ten times the Ohmic unpolarized value) in
the ion diamagnetic direction, attaining about one fifth of the ion thermal velocity.
Note that $E_r/B_i$ yields the same value as this measured speed.

2.3. Negative polarization

As explained in Section 2.2, H modes with negative $E_r$ have hitherto been rare
events on TEXTOR. Recently we succeeded in an NBI assisted mode of operation
which is shown in Fig. 3. When NBI is switched off ($t = 1.45$ s), H mode behaviour
FIG. 1. (a). Time evolution of plasma energy $E$, NBI power $P_{\text{NI}}$, $H_a$ emission at the wall, electrode bias voltage $V_E$ and line density $\bar{n}_e$ for two discharges, discharge A with bias and B without bias. (b) and (c). Electron temperature profiles in A and B at $t = t_1$ and $t = t_2$, respectively.
FIG. 2. Plasma energy versus line density for (a) Ohmic and (b) beam heated discharges without (open symbols) and with (closed symbols) positive electrode biasing. For the NBI case both the thermal energy $E_{\text{kin}}$ and the total energy $E_{\text{tot}}$ are given. The large symbols pertain to the negative bias case of Fig. 3.

FIG. 3. Time evolution of same signals as in Fig. 1 for a discharge with negative bias.
is maintained in the following Ohmic phase until the bias drops below threshold. The rather sparse data set allows some tentative conclusions: At the same $|V_E|$, the effect on the density and temperature profiles is stronger and the increase in particle confinement is about two times higher than with positive fields. The energy confinement is, however, about the same (see data points from Fig. 3 in Fig. 2).

3. VISCOUS DAMPING FORCE

Poloidal rotational equilibrium arises from the balance between the viscous damping force $F_{\text{visc}}$ and the forces resulting from radial non-ambipolar loss or charging currents, capable of torquing up the plasma. The latter can be set up in the plasma edge in different ways: fast ion losses resulting from NBI [4, 9] or ICRH, enhanced thermal ion losses induced by fluctuations [10, 11] or simply by means of a probe, as reported in this paper. Only in the latter case is the driving force easily known, since it is simply given by $j_E B$, where the current density $j_E = I_e / A$ and $A$ is the area of the magnetic surface at the electrode radius.

Figure 4 shows the experimental dependence, for positive fields, of $j_E = F_{\text{visc}} / B$ on $E_r$ for a plasma with an average collisionality $v^* = 10$ (note that the

\[ X = \frac{E_r - E_{\text{crit}}}{E_{\text{crit}}}, \]

where $E_{\text{crit}} = V_{\text{thi}} / B_r$. **FIG. 4.** Experimental dependence of $j_E = F_{\text{visc}} / B$, on $E_r$ for positive fields and a plasma with an average collisionality $v^*_E = 10$. The parameter $X$ denotes $(E_r - E_{\text{crit}})/E_{\text{crit}}$, where $E_{\text{crit}} = (1/en)(dp/dr)$ and $E_{\text{crit}} = V_{\text{thi}} / B_r$. 
TABLE I. COMPARISON OF EXPERIMENTAL AND DERIVED VALUES

<table>
<thead>
<tr>
<th></th>
<th>CCT</th>
<th>TEXTOR</th>
</tr>
</thead>
<tbody>
<tr>
<td>Experiment</td>
<td>$(0.7 \pm 0.3) \times 10^{-3}$</td>
<td>$(1.5 \pm 0.5) \times 10^{-3}$</td>
</tr>
<tr>
<td>Eq. (1)</td>
<td>$(3 \pm 1) \times 10^{-5}$</td>
<td>$(1.0 \pm 0.3) \times 10^{-5}$</td>
</tr>
<tr>
<td>Eq. (2)</td>
<td>$(1 \pm 0.3) \times 10^{-3}$</td>
<td>$(0.5 \pm 0.5) \times 10^{-3}$</td>
</tr>
</tbody>
</table>

plasma conditions vary as the electrical field increases). $E_r$ is the value corresponding to the maximum field in the rotating layer. The abscissa can also be expressed in units of the neoclassical critical field $E_{\text{crit}} = V_{th}B_p$ (or critical rotation speed $U_{p,cr} = V_{th}B_p/B_i$) using $X = (E_r - E_{r0})/E_{\text{crit}}$ [4], where $E_{r0} = (1/en_i) (dp/dr)$. The dependence of $j_e$ on $X$ can be approximated by $X/(1 + X^2)$ [2]. Our data confirm the predicted maximum near $X = 1$ and the strongly decreasing damping once the rotation exceeds $U_{p,cr}$. It has already been pointed out [2, 4] that it is the very shape of the friction force in its dependence on $E_r$ and $U_p$ which creates the possibility for bifurcation in poloidal rotation. It is important to stress that this property exists irrespective of the particular way in which the rotation is set up.

The maximum value of $F_{\text{visc}}$ controls the access to high fields accompanying the H mode. Its value sets the value of the minimum loss or charging current that should be applied to trigger a bifurcation. Our experiments permit a derivation of this critical current. Indeed, the maximum current is uniquely determined by the condition $X = 1$ on the one hand, and the initial slope $dF_{\text{visc}}/dU_p(U_p = 0)$ on the other. In Table 1 we compare the experimental values of

$$\frac{dj_e}{dE_r}(0) = \frac{1}{B_i^2} \frac{dF_{\text{visc}}}{dU_p}(0) \left( \Omega^{-1} \text{m}^{-1} \right)$$

found in CCT [2] and TEXTOR with the values derived from two low velocity expressions for $F_{\text{visc}}$ found in the literature (see Ref. [2] for a discussion):

$$F_{\text{visc}} = n_i m_i v_{li} U_p$$  \hspace{1cm} (1)

and

$$F_{\text{visc}} = n_i m_i v_{li} q_i^2 U_p$$ \hspace{1cm} (2)

where $v_{li}$ is the ion–ion collision frequency.
It is clear that the data do not at all fit the often used Eq. (1) [3, 9] and that a much better agreement is obtained with Eq. (2). From condition $X = 1$ and Eq. (2) we then obtain the expression for the critical H mode current density

$$j_E = \frac{\epsilon q}{2} \frac{p_i}{\omega_{ci}} (e_n V_{th})$$

where $\epsilon$ is the aspect ratio and $\omega_{ci}$ the ion cyclotron frequency.

4. CONCLUSIONS

(a) H mode behaviour can be set up with both positive (radially outward pointing) and negative (inward pointing) $E_r$ fields.
(b) The H mode behaviour is sustained in the presence of auxiliary heating.
(c) Improvement in energy confinement of up to about 1.4 and in particle confinement of 2 to 2.5 is obtained with positive bias. First indications are that the particle confinement might increase more with negative bias.
(d) The measured viscous force opposing poloidal rotation presents those features predicted by neoclassical theory which permit bifurcations in velocity from low to high rotation speed.
(e) To set up an H mode a minimum loss (or charging) current density

$$j_E = \frac{\epsilon q}{2} \frac{p_i}{\omega_{ci}} (e_n V_{th})$$

is needed, irrespective of which mechanism accounts for the loss (or charging).

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IAEA-CN-53/A-VI-6

DISCUSSION

H.L. BERK: You have made an estimate of the amount of current that is needed to sustain an H mode. Does it appear that the orbit loss mechanism in larger machines is capable of maintaining an H mode?

R.R. WEYNANTS: I can only tell you what I recall about the DIII data (see Ref. [3] cited in my paper), where they tried to compare the current needed to overcome friction with the driven current computed by Shaing. As I showed in my presentation, they underestimated the friction by about a factor of 10, and because of this they thought, I believe, that even the thermal ion losses were sufficient. As the friction requires revision, however, I think that the fast ion losses are needed.

G. FUSSMANN: Is it just the edge region that is involved in the poloidal rotation or does the whole plasma rotate?

R.R. WEYNANTS: Our data tend to confirm the observations from other machines that the rotation is indeed an edge effect.

J. WESSON: Do I understand correctly that, once the transition is achieved, the drive can be switched off and the transition is then permanent? Is this in fact observed?

R.R. WEYNANTS: In order to attain this regime, the probe current would have to be even lower than what we observe (there is in fact still substantial friction). Our power supply does not yet allow us to go to the higher fields needed to give this interesting regime a proper test.
PLASMA ROTATION, FLUCTUATIONS AND TRANSPORT IN THE TJ-I TOKAMAK

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Abstract

RESULTS ON PLASMA ROTATION, EMPHASIZING ITS EFFECT ON PLASMA DENSITY, CURRENT AND SAFETY FACTOR, ARE PRESENTED FOR THE TJ-I TOKAMAK. MEASUREMENTS OF ELECTROSTATIC AND MAGNETIC TURBULENCE AND ITS SCALING WITH THE SAFETY FACTOR ARE GIVEN. PARTICLE CONFINEMENT TIMES ARE DISCUSSED IN THE LIGHT OF ROTATION AND FLUCTUATION DATA.

1. INTRODUCTION

Previous studies [1–3] of particle transport in the TJ-I ohmically heated tokamak (R = 30 cm, a = 10 cm, I_p ≤ 60 kA, B_T ≤ 1.5 T) had revealed that the particle diffusion rate is similar to that found in larger devices, which makes the study of its basic drive mechanisms in this small device a fusion relevant problem. On the other hand, those studies suggested that either the radial electric field or electrostatic fluctuations might play some role in explaining parametric behaviour and absolute values of τ. To obtain a deeper insight into this problem, plasma rotation and fluctuations have been measured to evaluate their possible role in the transport in TJ-I plasmas.

The paper is organized as follows:

First, results on plasma rotation are presented, emphasizing its effect on plasma density, current and safety factor. Second, measurements of electrostatic and magnetic turbulence and its scaling with the safety factor are given. Finally, the results on the particle confinement time are discussed in the light of rotation and fluctuation data.

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FIG. 1. (a) Plot of angular poloidal velocity, $\omega$, for a central and a peripheral ion versus density; (b) variation of $\omega$ with plasma current for similar density and toroidal field discharges; (c) dependence of poloidal velocity on safety factor, for two different toroidal field ranges.
2. ROTATION MEASUREMENTS

Poloidal and toroidal rotation velocities have been determined in TJ-I plasmas by measuring, with spatial resolution, the line shift of impurity lines. A 1 m monochromator with an OMA detector attached to its focal plane is used for these measurements; the plasma chord is varied, on a shot to shot basis, by means of a rotatable mirror capable of scanning the entire plasma cross-section. Thermal drift effects were taken into account in the analysis.

Poloidal plasma rotation profiles have been measured as a function of $n_e$, $I_p$ and $B_T$. In Fig. 1(a), the dependence of the angular poloidal velocity on the plasma density is shown for a central (C V) and a peripheral (C III) ion. Bulk ions rotate in the electron diamagnetic drift direction, and their angular velocity tends to decrease as the plasma density is increased. In contrast, a peripheral ion (C III) rotates in the opposite direction at low densities and in the same direction as the central ions at higher densities. The linear dependence of the angular poloidal rotation on $I_p$ can be seen in Fig. 1(b) for the case of a central peaking ion as C V; a similar behaviour has been observed for mid-peaking ions (O VI, O V). The variation of the poloidal rotation with the safety factor is illustrated in Fig. 1(c), and although the most significant changes are exhibited by the two extreme points, which correspond to the highest and the lowest plasma currents, a moderate variation with $q$ must be quoted. As is shown in this figure, the rotation tends to decrease at the highest fields. A theoretical estimation of the poloidal rotation has been carried out by using the neoclassical expressions found by Hirshman and Sigmar [4] and assuming analytical profiles which approximately match the kinetic profiles of background plasma and impurities. The theoretical values show a trend with $n_e$ and $B_T$ that is different from the experimental trends, and their ratio ($\omega_e/\omega_T$) varies most effectively between 3 and 5. This suggests the possible importance of radial electric fields.

From the radial component of the force equilibrium equation, the radial electric field can be deduced from rotation data via $E_r = (1/n_e e) dP_e/dr + v_e B_T - v_T B_T$, where these parameters have their usual meaning [5]. The toroidal rotation velocity has been measured in TJ-I plasmas, and although its absolute value is of the same order as the poloidal rotation its contribution to the former equation is $\leq 10\%$. In Fig. 2(a), plasma potential and radial electric field versus plasma radius are plotted for a low density case. The variation of the central potential ($V_0$) and the peripheral, averaged values of $E_r$ with $n_e$ as deduced from CV rotation data is shown in Fig. 2(b). Note that $V_0$ becomes less negative as the density increases. Positive plasma potentials at the edge (0.75 cm) are not seen in these data, because of the averaging within the spatial resolution, but are evident from C III rotation behaviour and Langmuir probe results. A potential change in this layer explains the C III rotation inversion as a function of density.
FIG. 2. (a) Radial plot of relative plasma potential and radial electric field, for a low density case, deduced from CV rotation data; (b) variation of central plasma potential and averaged peripheral $E_r$, deduced from CV rotation data, versus line averaged electron density.

3. FLUCTUATION MEASUREMENTS

Edge turbulence has been characterized by means of Langmuir and magnetic probes [2]. A broadband reflectometer [6] probes the density fluctuations in the bulk plasma ($r/a \leq 0.7$). By sweeping the reflectometer frequency the radial distribution of the turbulence level is obtained in a single discharge; it diminishes when the reflecting layer is more internal as is shown in Fig. 3(a). The density fluctuation level at the edge increases as the probe is moving outward, in agreement with reflectometer results, with the relative level ($\tilde{n}/n$) ranging from 0.4 to 0.8.

Frequency and wavenumber spectra have been determined, from probe data using the floating potential signal, for toroidal and poloidal propagation using a two-point correlation technique [7]. Measurements have been obtained for different radial
positions from $r/a = 1.06$ to 0.92. The turbulent broadening of the $k$-spectra ($\sigma_k$), whose inverse measures the correlation length, has been determined for $q(a)$ ranging from 3.8 to 7.5. In the proximity of the shear velocity layer, where the poloidal phase velocity of the fluctuations reverses from the ion to the electron diamagnetic drift direction, the broadening of the $k$-spectra increases with $q(a)$ and tends to peak at the shear position. The $k$-broadening in the poloidal direction is larger than in the toroidal one, as is shown in Fig. 3(b), where the results for two extreme values of the safety factor are plotted. These results are consistent with previous findings in ATF and TEXT [8]. The radial correlation length, $L_r$, of fluctuations has been estimated to be 0.5–1 cm at $r/a = 0.7$. The edge radial magnetic fluctuations decrease as $q(a)$ increases.

4. TRANSPORT STUDIES

Independent information on particle transport is obtained in TJ-I by injection of impurities by laser blow-off, $H_a$ measurements and by studying the confinement of runaway electrons. While the first two methods yield the central and global confinement times of thermal particles, respectively, the latter approach reflects the behaviour of $E \leq 3$ MeV fast particles whose confinement should be affected by electromagnetic turbulence and convective transport, but hardly by collisions.

The injection of a short pulse of impurities by laser ablation produces an enhancement in both electron density and impurity concentration; the decay of the perturbation with these two parameters is what is designated in Fig. 4(a), by ‘electrons’ and ‘impurities’, respectively. The runaway confinement times are deduced from the slope of the relevant energy spectrum. This is accounted for by a simple model in which plasma $Z_{\text{eff}}$ and loop voltage are the critical parameters. In Fig. 4,
FIG. 4. TJ-I particle confinement times versus line averaged electron density: (a) for thermal particles (electrons and impurities), deduced from the perturbation decay of their densities, produced by injecting a short pulse of heavy particles by laser ablation of an iron film; (b) electron runaway confinement times deduced from the slope of energy spectrum modelled with a one-dimensional code; (c) global particle confinement time deduced from $H_a$ measurements taking into account poloidal asymmetries.
the behaviour of central particle confinement time measured by laser ablation (Fig. 4(a)), electron runaway hard X ray emission (Fig. 4(b)) and global confinement obtained from Hα measurements, have been plotted versus density. It is worth noticing that except for the lower densities, these confinement times have similar values. In contrast to the global confinement reported in Ref. [3], the values presented here have been corrected for Hα asymmetry effects, since Hα spatially resolved measurements have revealed out/in asymmetric particle fluxes whose ratio ranges from 4 to 1.5 as the density increases in the range shown by Fig. 4(c).

Consequently, no significant differences are observed either in the confinement of thermal particles (electrons and ions) or that of fast electrons in this density scan at 1 T, with the exception of global confinement time values for n_e below \(1 \times 10^{13} \text{ cm}^{-3}\), where the interpretation of Hα fluxes is complicated by the incomplete recycling of particles [3]. These data, in conjunction with the above reported behaviour of plasma potential with density, suggest that the preferential losses of superthermal ions might be the loss channel modifying the plasma potential in a TJ-I density scan. This hypothesis, which has already been advanced for another tokamak [5], is supported, in our case by experimental data on the confinement time of thermal and fast electrons.

5. CONCLUSIONS

Significant values of the impurity poloidal rotation have been measured in the TJ-I tokamak. The rotation diminishes with the plasma density and, slightly, with the safety factor and tends to increase almost linearly with the plasma current. The main contribution comes from the radial electric field driven term. While the bulk ions, with a few exceptions, rotate in the electron diamagnetic drift direction, the peripheral ions can switch from the ion diamagnetic direction, at low densities or low currents, to the electron diamagnetic direction at higher densities. A remarkable feature is that only in the highest density range of (2-3) \(\times 10^{13} \text{ cm}^{-3}\) all the ions rotate poloidally with the same angular velocity. A similar trend with the density of thermal particles and runaway electron confinement is consistent with the changes in the plasma potential and the radial electric field deduced from plasma rotation data, which can be attributed to a reduction of fast ion losses as the density increases.

The expected dependence of turbulence spectral broadening on the safety factor has been found in the TJ-I tokamak, matching previous findings in the ATF and TEXT devices. The turbulence radial correlation length was measured, confirming reflectometry as powerful diagnostics of fluctuations in the plasma interior.

ACKNOWLEDGEMENTS

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STUDIES OF ELECTROSTATIC AND MAGNETIC FLUCTUATIONS IN CASTOR TOKAMAK

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Abstract

STUDIES OF ELECTROSTATIC AND MAGNETIC FLUCTUATIONS IN CASTOR TOKAMAK.
The electrostatic and magnetic edge turbulence of the CASTOR plasma was investigated with poloidal resolution. The effect of fluctuation suppression in the presence of a lower hybrid wave was studied. A theoretical model is presented as an attempt to understand the microscopic processes occurring in the CASTOR plasma.

1. INTRODUCTION

The improvement of global particle confinement during combined inductive/lower hybrid (LH) current drive was routinely observed in the CASTOR tokamak (R = 0.4 m, a = 0.085 m, B = 1 T, I = 12 kA) [1]. For a better understanding of this interesting effect, the experimental effort was focused on determining the poloidal dependence of both electrostatic and magnetic fluctuations and on studying this effect near the density limit for the LH current drive (Sections 2 and 3).

The aim of the theoretical studies was to explain the particle confinement in this tokamak (the density decay appearing in Ohmic discharges) and the effect of the LH wave on the electrostatic fluctuations (Section 4).

2. EXPERIMENTAL ARRANGEMENT

The experiments were performed in the CASTOR tokamak in standard Ohmic heating discharges (OH) and in regimes with combined inductive and lower hybrid current drive (OH/LHCD). In the latter case, a lower hybrid wave (f = 1.25 GHz,
FIG. 1. Experimental arrangement of triple probes and Mirnov coil poloidal arrays in CASTOR.

FIG. 2. Arrangement of three-channel analogue correlator for determining cross-field drift induced transport parameters.
P ≤ 40 kW) was launched into an OH plasma by a multijunction waveguide grill with a broad power spectrum (Nf = 1–4) and rather low directivity. The position of the plasma column inside the liner was feedback controlled.

The MHD activity was monitored by a set of Mirnov coils uniformly distributed inside the liner (Fig. 1). The signals from the coil were digitized by seven A/D converters in the 0.1–300 kHz frequency band.

The edge electrostatic turbulence was investigated by a poloidal array of triple probes (Fig. 2). The tips of each probe were arranged in a triangular shape (d = 3 mm). The two poloidally separated floating tips are used to determine the poloidal electric field fluctuations. The third, negatively biased tip is used to monitor the density fluctuations. The triple probe is connected to a three-channel analogue correlator (Fig. 2), which gives the temporal evolution of the RMS values of density and E_p fluctuations, together with their cross-correlation. The cross-correlation represents the turbulent particle flux \( \langle \delta n_e \cdot \delta n_e \rangle \), where the radial velocity is given by the cross-field drift \( v_r = \frac{E_p}{B^2} \times B \). Moreover, the quasi-stationary values of \( \langle n_e \rangle \) and \( \langle E_p \rangle \) were monitored. Consequently, the quasi-stationary particle flux \( \langle \Gamma \rangle = \langle n_e \rangle \langle E_p \rangle / B \) can be simply derived for different poloidal angles.

3. EXPERIMENTAL RESULTS

3.1. The quasi-stationary cross-field drift induced transport in the OH regime

Figure 3 presents an example for the poloidal variation of three quantities measured by the triple probe array. We see that the poloidal dependence of the edge density (and edge electron temperature) is approximately uniform. The poloidal dependence of the quasi-stationary poloidal electric field (and also the floating potential), however, exhibits a very pronounced maximum and even changes the direction. The negative sign of \( E_p \), in our case, corresponds to the inward direction of the quasi-stationary flux \( \langle \Gamma \rangle \). A similar asymmetry is observed for the turbulent component of this flux. Nevertheless, a poloidal averaged value of \( \langle E_p \rangle = 500 \text{ V/m} \) seems to yield an outward flux of an order comparable to the measured turbulent flux and the particle flux determined from global particle balance. It indicates that the quasi-stationary value of the cross-field drift induced transport should be taken into account. However, it should be noted that \( \langle E_p \rangle \), determined here as a difference of potentials of two floating tips, is a correct value of the poloidal electric field only in the case where the tips are on magnetic surfaces which are so close that \( T_i^1 - T_i^2 \ll (V_f^1 - V_f^2) \), for a hydrogen plasma. This condition may be even more severe for non-Maxwellian plasmas.
3.2. Electrostatic and magnetic fluctuations during the OH/LHCD regime

It has been shown that the level of electrostatic turbulence drops substantially during the OH/LHCD period of CASTOR discharges (Fig. 4) [1]. Simultaneously, an enhancement of global particle confinement is observed. Here, it is documented by a slowing-down of the line average density decay during OH/LHCD [2].

Figure 4(b) demonstrates that both effects are still well pronounced near the density limit of the LH current drive, when a relatively small number of super-thermal, current carrying electrons is present in the plasma column.

The MHD turbulence is similarly affected during the OH/LHCD part of the discharge, as is shown in Fig. 5. It is seen that, despite a poloidal variation of the fluctuation levels, the relative suppression of the MHD turbulence is constant and equals 0.4. The fast Fourier transform spectra of the coil signals are shown in
FIG. 4. Temporal evolution of plasma parameters for a medium density discharge (a) and for a high density discharge, near the density limit of LHCD (b).
Fig. 5. Levels of MHD turbulence for OH and OH/LHCD regimes.

Fig. 6. Frequency spectra of magnetic fluctuations.

Fig. 6. The suppression takes place in a broad spectral range. Therefore, it seems that this effect differs from the suppression of the coherent $m = 2$ mode observed during LHCD on Petula [3].

4. THEORETICAL MODEL

The results of the measurements performed in the OH regime of CASTOR have previously been used to solve the inverse problem of the plasma transport.
This consists in determining the transport coefficients from the available experimental data using numerical simulations. The specially elaborated method [4] permitted the thermal conductivity coefficient to be determined and also a realistic picture of the evolution of the plasma parameters in CASTOR to be obtained.

Because of the steep density gradient developing from the start, the theoretical explanation of these experimental and numerical results must invoke the low frequency electrostatic drift instability. The high collisionality of CASTOR discharges drives the dissipative branch unstable which is saturated by magnetic shear and ion damping. A specific feature of the low density CASTOR discharges is the fast decay of the average density, after reaching a maximum (1 ms) (Fig. 4(a)). It appears in conditions in which no fully developed turbulence is expected. This enhanced diffusion has a rather sudden onset and is too effective for the expected low level of turbulence, at that time. Indeed, estimating the linear growth rate and the level of non-linearity [5], we obtain \((k_1 v_T)^2 / (\omega_{ce} \nu_{ce}) = 20 > 1\), which shows an almost adiabatic response of the electrons and a low level of non-linear coupling. Then, the modes can have a quasi-coherent structure which is slightly perturbed by non-linearity. It is then important to investigate the contribution of non-conventional transport mechanisms which are related to the existence of large scale coherent potential structures.

We consider the potential structure of a single drift mode which, ideally, shows two-dimensional periodicity in a finite radial region. The separatrices are perturbed non-linearly and the particle motion in this finite width zone is not integrable. The particles perform \(E \times B\) oscillations in the potential cells, but those which are close to the stochasticized separatrices can be scattered by collisions and perform jumps to the neighbouring cells. A combination of such jumps allows the particles to make significant radial excursions. The transfer of particles (i.e. a flux) from one side to the other of such periodic structures is possible when the probability of the elementary process of scattering is greater than a critical value [6]. This particle loss mechanism has important peculiarities: (1) is not directly connected to the density gradient; (2) it has a threshold; (3) it is non-local as is the basic potential 2-D structure; (4) it is not accompanied by heat conduction (the proportionality between \(D\) and \(K\) is not maintained).

A similar percolation process can be identified in a regime of higher turbulence provided that the correlation time is still longer than the characteristic time of electron drift motion. In this situation, long range potential structures can randomly occur in the 2-D stochastic geometry of a superposition of electrostatic modes. These structures, having large extensions in the radial direction, provide, via the particle drift, an efficient contribution to the anomalous transport. The diffusion arising from this additional mechanism when a sufficiently high radial electric field is present (as in the CASTOR case) is [7]:

\[
D_p = \lambda^2 \omega k_1 \Phi / E_r \quad \text{for} \quad E_r / (k_1 \Phi) > (v_{ph,1} / v_d)^{(\nu + 1) / (\nu + 2)}
\]

where \(\nu = 4/3\) is the correlation length exponent in the 2-D percolation theory,
\[ \nu_d = k_\perp \Phi / B, \quad \lambda = 2\pi / k_\perp \quad \text{and} \quad \nu_{ph\perp} \] is the phase velocity of the mode. In our case, the resulting diffusion is of the order of \(1 \text{ m}^2/\text{s}\). This value is comparable with the empirical result obtained from numerical simulations.

The particles which traverse the zone with potential cells determine a modification of the density profile by increasing the density gradient in the marginal region. Such a steep boundary profile was experimentally observed in the first milliseconds of the CASTOR discharges.

As the turbulence increases, the dominant contribution to the diffusion arises from the fluid-like response of the electrons. The non-adiabatic contribution of the electron density response was estimated from the theory of collisional drift instability:

\[
\text{Im} \left[ \left( \frac{\nabla n}{n} \right) \left( \frac{e\Phi}{T_e} \right) \right] \approx 0.6
\]

The calculated levels of fluctuations (3% at the centre and 40% at the border) are close to the experimental observations [1]. This turbulent phase of the drift instability provides a satisfactory model for the plasma evolution after the fast density decay when the density continues to decrease, but at a slower rate. The drift turbulence affects the major part of the plasma cross-section.

The fact that particle losses are due to different mechanisms can be observed in the experimental determination of edge particle flux. A separated peak appears at

**FIG. 7.** Relative drop of density and poloidal electric field fluctuations and of turbulent flux as functions of frequency.
the time of maximum density decay which, according to our model, is determined by percolations; after this the turbulent flux develops (Fig. 4).

The marginal zone of the plasma is affected by the existence of a quasi-stationary radial electric field [1] which determines a sheared poloidal flow of the plasma providing a free energy source through the relaxation of the vorticity gradient. This determines a shift of the frequencies from the diamagnetic values \(\omega_c \approx 90\,\text{kHz}\) to \(\omega_E = E_x k_x / B \approx 320\,\text{kHz}\) and explains some other observed features of the edge turbulence: the reversion of the sign of the fluctuation phase velocity near the limiter and the presence of a minimum in the fluctuation level in this region.

A very important feature revealed by the measurements performed in the OH/LHCD regime is the partial suppression for the low frequency part (<100 kHz) of the edge turbulent density fluctuations (Figs 4 and 7). Since the high frequency part, which is related to the radial electric field drive is non-perturbed, the LH wave acts on the dissipative drift wave turbulence. We have considered the direct effect of the LH wave on the linear drift mode evolution. The result was that the linear growth of the mode is inhibited only for values of the LH wave electric field which are much higher than in CASTOR experiments. It may be deduced that the LH wave acts on the non-adiabatic part of the electron response, e.g. reducing the lifetime of the small scale correlations. This is also supported by the experimental observation that the poloidal field fluctuations are much less affected than the density fluctuations.

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EFFECTS OF CONDUCTING SHELL AND LOW-q DISCHARGES IN HL-1

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Abstract

EFFECTS OF CONDUCTING SHELL AND LOW-q DISCHARGES IN HL-1.

Effects of copper shell and low-q discharges have been studied intensively in the HL-1 tokamak. The stored energy is lost in two phases when the discharge is disruptively terminated, and the copper shell can reduce the vacuum vessel damage by extending the energy dissipation process. An investigation using the quasi-linear theory of the tearing mode shows that the copper shell is more effective in suppressing the $m = 2$ mode when the current density in the peripheral region of plasma is reduced. The measures taken to achieve low-q plasmas, which are directly related to the control of the current profile, are briefly summarized. Sawteeth and energy confinement were investigated for low $q_a$ and high density operation, indicating that energy confinement is correlated with sawtooth oscillation.

1. INTRODUCTION

In the past two years, HL-1 (the largest operating tokamak in China: $R = 1.02 \, \text{m}$, $B_T = 3.5 \, \text{T}$, discharge duration $= 1.6 \, \text{s}$) performance has been improved progressively. Plasma control was refined with plasma position and line averaged electron density, $n_e$, both controlled by feedback systems. By ECR or AC discharge cleaning in hydrogen-krypton mixtures and adequate titanium gettering, $Z_{\text{eff}}$ has been reduced to 1.5 and the density limit for disruption has been raised by 80%, attaining $n_e = 7.2 \times 10^{19} \, \text{m}^{-3}$ in hydrogen and $8 \times 10^{19} \, \text{m}^{-3}$ in helium. With Ohmic heating, a good confinement plasma with $\tau_E = 38 \, \text{ms}$ has been obtained, and the highest fusion product achieved so far is $n_i(0) \tau_E T_i(0) = 1.8 \times 10^{18} \, \text{m}^{-3} \cdot \text{s} \cdot \text{keV} \pm 30\%$. Recently, the effects of a thick conducting shell and low-q discharges have been studied intensively in the device.
2. EFFECTS OF CONDUCTING SHELL AND MHD PERTURBATIONS

Although the thick copper shell in HL-1 [1] is not very close to the plasma boundary \( b/a > 1.6 \), where \( b \) is the shell radius, it has a significant effect on the behaviour of the plasma disruption, because of the transient induction in the shell caused by the large current jumps with short time constants at a disruptive disturbance.

2.1. The dissipation of energy during a major disruption

In a major disruption, the stored energy is lost in two phases. In the first phase, the plasma loses the major part of its thermal energy in a very short time (50–150 \( \mu s \)) with a large negative spike (300–600 V) appearing on the loop voltage. The amplitude of the spike normalized by the plasma current \( \Delta V/I_\text{p} \) generally decreases as its width \( \Delta t \) increases (Fig. 1), implying that the negative voltage spike is related to the changing rate of the plasma inductance. Accompanying the negative spike is a short burst of HF oscillation (~100 kHz), which grows quite fast with a rise-time of 10–80 \( \mu s \) and disappears as the negative spike is terminated [2]. During the growth, the oscillation is characterized by the appearance of turbulence.

After the release of the thermal energy, the magnetic energy stored in the poloidal magnetic field partly dissipated into the plasma and was lost with the quenching of the plasma current. Bolometric studies of the energy deposition on the wall showed that the measured losses exceed the plasma thermal energy; these losses are consistent with the dissipation of a large fraction of the magnetic energy. The waveform of the current decay in major disruptions is rather variable, but the decay rates are much lower than in other tokamak devices. In the typical disruption case, the current decreases rapidly with \( I_\text{p} \) in the range between −12 MA/s and −16 MA/s from the beginning. This kind of disruption is usually a low \( q_a \) (\( q_a > 2 \)) disruption; these disruptions are only a small fraction of the major disruptions. The majority of the major disruptions have a value of \( I_\text{p} \) between −2 MA/s and −8 MA/s, while in some cases the current falls more quickly by the end of termination. This fast quenching of the plasma current is similar to the normal termination of discharges and is probably due to plasma position control having been lost. The eddy currents and the electromagnetic force in the metallic structure were examined with the EDDYTOR-6 code. A finite element stress analysis of the vacuum vessel, which is made up of 16 segments consisting of bellows and a rigid section, showed that a shear force is exerted in the radial direction (Fig. 2) and the peak value of the stress caused by disruptions was 0–01 MPa (\( B_T = 2 \) T).

In an iron core tokamak such as HL-1, the plasma behaves as an independent system with an effective inductance \( L_{\text{eff}} \) [3]. \( L_{\text{eff}} \) can be split into \( L_i \), corresponding to the energy stored in the vacuum vessel, and \( L_e \), the external inductance. During the current decay, the transient induction in the copper shell increases \( L_e \), and the rate of energy deposition on the vacuum vessel is reduced significantly. \( L_e \) was
FIG. 1. Amplitude (ΔV/Ip) versus width (Δt) of negative voltage spikes.

FIG. 2. Poloidal distribution of electromagnetic force per finite element on vacuum vessel (B_T = 2T). The forces in the radial direction are on the bellows section (A) and the rigid section (B), the forces along the symmetric axis are on the rigid section (C) and the bellows section (D).
determined experimentally; it remains in the range of 5-6 \mu h. The loop voltages measured by flux loops located at the inside and outside shell surfaces are significantly different during a disruption, which indicates an induced change of magnetic fluxes in the shell.

As the thick conducting shell in HL-1 keeps the current and the plasma position unchanged immediately after the thermal energy release, most of the disruptions in the HL-1 discharges are minor [1, 4]. For example, of about 1000 discharges generated in 1988, less than 20 terminated in a disruption. For a fairly large fraction of these events of disruptive current termination, the major disruption occurred after a sequence of minor disruptions, which consecutively deteriorate the quality of the plasma column. In particular, a relaxation phenomenon with rather pronounced disruptions has been observed in HL-1 discharges [4].

2.2. Influence of copper shell on precursory oscillations

We have investigated the role of the copper shell in suppressing precursory MHD instabilities by using a quasi-linear theory of the tearing mode. The computational results show that the position of the resonant surface is of paramount importance. The \( m = 2/n = 1 \) island can be completely stabilized by a closely fitting, conducting shell if \( q_a \) is close to 2, but in the case of \( q_a \geq 3 \) the shell has nearly no effect on the \( m = 2 \) fluctuation, even if it is close to the plasma boundary.

![FIG. 3. Dependence of island width \( W (m = 2) \) on parameter \( \mu \) (solid line: \( q_a = 2.5 \); dotted line: \( q_a = 3.5 \)), together with plasma current profiles for different \( \mu \).](image-url)
The stabilizing effect of the shell is closely related to the current density profile. In the numerical simulation, a current profile with a power dependence \( \mu \),
\[
j(x) = j(0) \left(1 + \left(\frac{q_a}{q(0)}\right)^{\mu} - 1\right) \frac{x^{2\mu}}{x^{2\mu} + \mu} \]
is used. As the inner radius of the shell is fixed at \( b = 0.32 \) m, the width \( W \) of the \( m = 2 \) island normalized to the plasma radius is plotted as a function of \( \mu \) in Fig. 3 for different plasma radii. The present plasma radius in HL-1 is in the range of 0.16–0.20 m. It needs an increase up to 0.25 m for the shell to have a significant stabilizing effect when \( \mu \) is fairly large. As is also shown in Fig. 3, the shell will be more effective in suppressing the \( m = 2 \) mode when the current density in the peripheral region is reduced with \( q_a \) kept constant. This favourable current profile may be realized with gas puffing.

3. LOW \( q_a \) DISCHARGES

With gaseous refuelling only, a Murakami parameter of \( M = 3.3 \times 10^{19} \) \( \text{m}^{-2} \cdot \text{T}^{-1} \) was reached at a low \( q_a \) value of about 1.9. The \( q_a \) limit could be reduced to \( q_a = 1.7 \) at high plasma density.

3.1. Formation and \( m = 2 \) stability

The empirical method of achieving low \( q_a \) plasmas at high density in HL-1 can be summarized as follows: (1) slow ramp-up of plasma current; (2) careful control of gas puffing during the current ramp-up phase and flat-top so as to select the discharge trajectory passing through the centre of the stable region; (3) cleaning of wall and limiter surface by ECR or AC discharge cleaning and using a moderate amount of titanium gettering; and (4) control of plasma position, with great precision, by a feedback system. These measures are directly related to the control of current density profile and metallic impurity influx [5].

The main method of suppressing the development of an \( m = 2 \) mode in low-\( q \) discharges is to program both gas puffing and plasma current waveform in such a way that the current density profile will be sufficiently peaked during the current rise to provide constant \( q \) at the current channel boundary. Both plasma current and electron density should be increased so as to keep \( I_p/n_e \) continuously in the range of \( 2.3 \times 10^{-18} \) \( \text{kA} \cdot \text{m}^3 \) and \( 4 \times 10^{-18} \) \( \text{kA} \cdot \text{m}^3 \); in this way, the \( m = 2 \) mode has been suppressed to an extent less than 0.3% (Fig. 4). An improved one-dimensional analysis of the tearing mode was made for \( j(r) \) profile optimization in the low-\( q_a \) regime. It showed that, for \( q(0) > 1 \), stable \( j(r) \) profiles can only be found when \( q_a > 2 \), but, in the case of \( 0.5 < q(0) < 1 \), a \( j(r) \) profile stable against all modes except \( m = 1/n = 1 \) exists for \( q_a < 2 \).

At small \( q_a \) values (\( q_a < 2.4 \)), a strong, saturated \( m = 2 \) mode develops already at moderate densities and increases slowly with increasing density. Near the density limit, it leads to a disruption, but it has been found that at even smaller \( q_a \) values (\( q_a \sim 2 \)), the \( m = 2 \) fluctuation gradually drops to a very low level during the flat-top phase, a result similar to the ASDEX results [6].

FIG. 5. Soft X ray signals and time evolution of soft X ray tomography in $-6.5 \text{ cm} < x < 6.5 \text{ cm}$ and $-6.5 \text{ cm} < y < 6.5 \text{ cm}$ regions. Plasma parameters: $B_T = 2.4 \text{ T}$, $I_p = 128 \text{ kA}$, $\bar{n}_e = 5.4 \times 10^{19} \text{ m}^{-3}$, $q_a = 3$. 
3.2. Sawteeth and confinement

Sawtooth oscillations were observed in the soft X ray signals with a two-dimensional array of 35 diodes. The detailed nature of the collapse has been investigated by tomography, showing that the shape of the hot core has changed from a circle to a crescent and moves towards the outside (Fig. 5). In the low-q<sub>a</sub> discharges at high density, the sawtooth oscillation increases in amplitude (ΔA/A = 20-35%) and its period becomes longer (τ<sub>s</sub> = 3-11 ms). The normal sawteeth are accompanied by a wide range of MHD activity in the central plasma region, and there are large MHD oscillations after the collapse. The duration of the crash is usually 0.2-0.5 ms, but slower crashes with crash times up to 1.1 ms were also observed. 'Giant' Ohmic sawteeth with more than twice the period and twice the amplitude of the normal sawteeth occur intermittently between the ordinary sawteeth. For sawteeth with τ<sub>s</sub> ≤ 5 ms, the period scales as τ<sub>s</sub> ∝ \( n_b/(r_s B_T) \) and the amplitude ΔA/A is roughly proportional to 1/q<sub>a</sub> but this scaling cannot fit the data from the 'giant' sawteeth.

In low-q<sub>a</sub>, high density operation, the global energy confinement time can reach 38 ms with Ohmic heating only. The scaling of the global electron energy confinement time for discharges of 1.7 < q<sub>a</sub> < 3.2 is shown to be τ<sub>E</sub> ∝ q<sup>0.75</sup>. In simple cases, the global energy confinement time is correlated with the sawtooth period; for a particular data set, the two times scale as τ<sub>ε</sub> = 15 τ<sup>0.5</sup>. The average value of the electron heat diffusivity in the central plasma region scales as \( \chi_e(0) (m^2/s) = 0.9/\tau_s \). From the theory of confinement degradation due to sawteeth developed by Chang and Callen [7], the degradation \( d_s \) due to sawtooth in HL-1 is estimated to be about 3-15%; it scales as \( d_s = 1.2 q_b^{-1.9} \). The degradation due to the presence of giant sawteeth is not more than 15%.

4. CONCLUSIONS

During a major disruption, a large amount of energy stored in the poloidal magnetic field is dissipated into the plasma and lost after thermal energy quenching, which is accompanied by high frequency oscillations appearing on the plasma current. Because of transient induction in the thick copper shell, the dissipation rate of the magnetic energy is low, and the electromagnetic force induced on the vacuum vessel is reduced significantly. As to the effect of the copper shell on suppressing precursory MHD fluctuations, the position of the resonant surface is of paramount importance, and the stabilizing effect is closely related to the current profile.

A stable low-q<sub>a</sub>, high density and good confinement plasma has been obtained in HL-1. To obtain low-q discharges at high density, careful control of both plasma current waveform and gas puffing is essential, together with a careful conditioning of the plasma chamber, thus keeping the current profile sufficiently peaked in order to suppress the development of an \( m = 2 \) mode. In low-q<sub>a</sub> and high density dis-
charges, the sawtooth oscillation increases both in amplitude and period. The global energy confinement time is correlated with the sawtooth period; for a particular data set, the two times scale as $\tau_E = 15 \tau_s^{0.5}$.

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INVESTIGATION OF IMPROVED CONFINEMENT REGIMES AND DENSITY LIMIT IN THE TUMAN-3 TOKAMAK

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Abstract

INVESTIGATION OF IMPROVED CONFINEMENT REGIMES AND DENSITY LIMIT IN THE TUMAN-3 TOKAMAK.

Two improved confinement regimes, Ohmic H mode and high density mode, were experimentally achieved on the TUMAN-3 tokamak. The spontaneous transition to the H mode took place in ohmically heated discharges in circular limiter configuration. Ohmic H mode energy confinement times were a factor of 1.2-1.3 longer than typical values in ohmically heated plasmas without H mode. Another way to achieve improved plasma confinement was minor radius magnetic compression accompanied by a second fast current rise and gas puffing — the high density mode. In this mode the highest average density was $5.2 \times 10^{13} \text{ cm}^{-3}$. This value is two times higher than the density limit in OH plasmas in TUMAN-3.

1. INTRODUCTION

Regimes with improved particle and energy confinement have been observed in a number of tokamaks [1–5]. In some of them the confinement improvement was accompanied by a considerable decrease of transport in the plasma periphery. This usually results in electron temperature and density profile broadening and periphery pedestal formation [1, 2]. These regimes are called the H mode. It should be pointed out that the transport processes in bulk plasma changed slightly after transition to the H mode. In another set of regimes (B in T-10 [3], IOC in ASDEX [4] and supershot in TFTR [5]), the significant decrease of plasma transport in bulk plasma causes the formation of peaked plasma parameter profiles. The transport in the peripheral region did not change greatly in these discharges.
FIG. 1. Time evolution of main plasma parameters in HDM.
Improved confinement regimes have been found in the experiments discussed in this paper. These regimes have revealed typical features of the H mode: a sharp drop of $D_\alpha$ radiation with a simultaneous rise of electron density and stored energy, and periphery transport reduction. However, some experimental features were in contradiction with the usual H mode behaviour. In particular, some observations indicated the electron temperature and plasma current profiles peaking after the H mode transition. It should be mentioned that in our experiments the H mode regime was found in the circular limiter configuration without auxiliary heating.

According to the well known Greenwald scaling, the density limit in tokamaks is proportional to the average plasma current density [6]. On the other hand, there is some experimental evidence of a strong influence of the peripheral plasma properties on the possibility of increasing the density [7]. In order to investigate the plasma current distribution effect on the density limit in the TUMAN-3 tokamak, experiments on increasing density in discharges with different current profiles were carried out. Such profiles were shaped by means of a second current ramp-up (broadened j(r) profile [8]) or minor radius compression (peaked j(r) profile [9]). An increase of the edge plasma current was accompanied by a decrease of the density limit compared with the $n_e$ lim in ordinary OH plasmas. Conversely, a decrease of the edge plasma current increases the density limit up to (0.8–0.85)$n_e$ lim (according to Greenwald).

2. IMPROVED CONFINEMENT IN THE HIGH DENSITY MODE

The high density mode (HDM) was found in experiments with additional gas puffing during magnetic compression accompanied by a fast second current ramp-up [10]. Twofold minor radius magnetic compression and a factor of 1.4 current rise were initiated simultaneously in a steady state of the discharge. The main plasma parameter traces are shown in Fig. 1. Toroidal magnetic field and plasma current reached their maximum values at 3.5 and 4.5 ms after their initiation. The maximum plasma current value was 140 kA. Some plasma parameters in OH and in HDM

| TABLE I. PLASMA PARAMETERS IN ORDINARY OHMIC HEATING REGIME AND IN HIGH DENSITY MODE |
|---------------------------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
|                                 | $I_p$ (kA)      | $B_t$ (kG)      | $q_{cyt}$       | $n_e$ (10^{13} cm$^{-3}$) | $T_{e0}$ (keV) | $T_{i0}$ (keV) | $\tau_p$ (ms) | $\tau_E$ (ms) | $Z_{eff}$ | $\beta_i$ max (%) | $\bar{\beta}_i$ (%) |
| OH                             | 94              | 4.5             | 2.7             | 1.4              | 0.4             | 0.11            | 1.5            | 2.5             | 2.0     | 1.9                        | 0.4                        |
| HDM                            | 120             | 6.7             | 3.2             | 5.2              | 0.5             | 0.17            | 11.1           | 15              | 2.2     | 3.5                        | 0.9                        |
FIG. 2. (a) Ion saturation current in peripheral plasma and (b) electron temperature profile in HDM. Triangles: Thomson scattering at 39.0 ms; asterisks: Thomson scattering at 44.5 ms; circles: SXR at 44.5 ms.
are listed in Table I. The $D_n$ radiation, collected along the major radius in the equatorial plane of the torus, dropped at the moment of the transition. In spite of the particle flux desorbed from the walls, the density increased until the end of the discharge. The moment of the transition took place near 40 ms, after the plasma current and toroidal field reached their quasi-stationary values. This indicates particle confinement improvement taking place mostly in the plasma periphery. Langmuir probe measurements showed steep density gradient formation near the limiter (Fig. 2(a)). One can see that the electron density decreases in the SOL but continues to increase in the main plasma after the transition to the HDM. Such experimental data make it possible to suppose that the observed transition is similar to the L–H transition in tokamaks with auxiliary heating. It is worth noting the significant decrease in fluctuation induced particle flux in the limiter region. The flux $\Gamma_p \propto \langle \tilde{n}_e \times \tilde{E} \rangle$ drops by a factor of 10 in spite of the small (factor of 1.5) decrease of the relative level of density fluctuation $\tilde{n}_e/n_e$. The central electron temperature increases during 5 ms with simultaneous peaking of the temperature profile after the HDM transition [10]. These profiles, measured just before the transition ($t = 39$ ms) and 4.5 ms after the transition, are shown in Fig. 2(b). The first profile is similar to the OH one. The second profile demonstrates the significant temperature gradient near the $r = 10$ cm surface after the HDM transition which is three times as great as before. This means that some decrease of thermal conductivity in the bulk plasma took place.

3. OHMIC H MODE

An improvement of confinement similar to that of the HDM was observed in OH plasmas without compression and current rise [11]. The Ohmic H mode transition was spontaneous and characterized by a sharp drop in $D_n$ emission and an increase of electron density and stored energy (Fig. 3). The transition also could be initiated by an MHD activity burst when the $m/n = 4$ surface was passing through the plasma boundary or by a short pulse of gas puffing. The switching on of the H mode was followed by plasma periphery transport reduction. The fast formation of a steep gradient region near the plasma boundary was observed. Figure 4 shows the time evolution of the density profile, measured by a ten channel UHF interferometer, $\lambda = 2$ mm. Note that, according to the electromagnetic diagnostic data, the plasma column was shifted towards the inner wall of the torus in these experiments. As a result of this shift, the outer vertical channel of the interferometer ($R = 76.6$ cm) looked through the SOL region. As it was measured, the line density decreased in this region after the transition but increased everywhere inside the last closed magnetic surface (Fig. 5(a)). The reduction of periphery transport correlated with the decrease of fluctuation induced particle flux (Fig. 5(b)).

The temperature in the H mode discharges seemed to be slightly higher than in OH, which was confirmed by SXR measurements. After the transition the tempera-
FIG. 3. Time evolution of main plasma parameters in Ohmic H mode, $q_{95}(a) = 2.6$. 
ture and current density profiles become more peaked as a result of a decrease of transport. Profile shrinking was followed by a corresponding decrease of $q(0)$ and the growth of sawtooth activity similar to that in HDM (Fig. 6). The current profile peaking was also confirmed by the short duration loop voltage inductive jump (Fig. 3). Impurity accumulation took place in the Ohmic H mode, and discharges finished with an increased value of the effective charge as well as radiation losses, followed by disruption. Such impurity accumulation has previously been observed in ASDEX [12]. Peripheral MHD activity (ELM) was absent in these experiments as well as in TUMAN-3.

4. IMPROVED AND ORDINARY CONFINEMENT MODES IN TUMAN-3

The discovery of HDM and Ohmic H mode gives us the possibility to extend significantly the range of densities in which confinement studies can be carried out in TUMAN-3. Therefore it was interesting to compare the confinement in these regimes with extrapolation of the ordinary OH scaling. Energy confinement times, derived from both diamagnetic and profile measurements, as a function of electron density in ordinary OH discharges are shown by the closed symbols in Fig. 7. The dotted line represents the linear dependence $\tau_E \propto n_e$, plotted for TUMAN-3 data. The broken line shows the Neo-Alcator scaling [13]. Open symbols show $\tau_E$ for HDM and Ohmic H mode. Note that in improved confinement regimes $\tau_E$ is higher than an extrapolated linear dependence $\tau_E \propto n_e$ (by a factor of 1.2–1.5) and the predictions of Neo-Alcator scaling (by a factor of 2.0).
FIG. 5. (a) Line density and $D_n$ emission and (b) fluctuation induced particle flux in Ohmic H mode.
FIG. 6. Time evolution of plasma parameters in Ohmic H mode, $q_{95}(a) = 3.4$. 
The results of our experiments allow us to draw conclusions about the possibility of the appearance of the H mode in a circular limiter discharge without auxiliary heating. The observed regime has the characteristic features of the H mode: (a) a decrease of recycling, (b) steep density gradient formation in the periphery, (c) a decrease of fluctuation driven particle flux. At the same time, it should be mentioned that the confinement improved not only in the periphery region but in the bulk plasma, too; this is confirmed by the peaking of temperature and current density profiles. The possibility of such a combination of the H mode with the peaked profiles in JET has been reported recently [14]. Bifurcation of the confinement properties can be caused by the changes in the radial distribution of the poloidal rotation velocity due to variation of the boundary plasma parameters. A corresponding mechanism was suggested in Ref. [15].

5. CURRENT DENSITY PROFILE EFFECT ON THE DENSITY LIMIT

The current density profile in a tokamak can be significantly affected by a fast second current ramp-up or by minor radius magnetic compression. In addition to the HDM scenario discussed above, density increase experiments have been carried out in discharges (a) with magnetic compression and constant plasma current, and
FIG. 8. Evolution of plasma parameters in (a) magnetic compression and (b) current ramping experiments.

(b) with a fast second current ramp-up and constant magnetic field. Corresponding
time evolutions of plasma parameters are shown in Fig. 8(a, b). In the magnetic
compression case (a) and HDM the current density profile is peaked. The peaking
of the $j(r)$ profile caused the decrease of recycling and was accompanied by transition
to the improved confinement mode. Significant increase of the plasma density has
been observed in the case of intensive gas puffing in such experimental conditions,
and the density limit was much higher than that achieved in OH plasma. Figure 9
shows a Hugill diagram for TUMAN-3. One can see that the Hugill limit for the
OH regime (open circles) is $n_e (10^{13} \text{ cm}^{-3}) R (\text{ m}) q/B_t (T) = 7.3 (0.46n_e^{\text{lim}}$
according to Greenwald). The magnitude of $n_e Rq/B_t$ increased up to the value of
13.0–14.0 ($(0.8–0.85)n_e^{\text{lim}}$) in the discharges with compression and HDM. Note that
in the Ohmic H mode where peaking of the current profile and decrease of recycling
were observed as well, the density limit turned out to be slightly higher than in OH
($n_e Rq/B_t = 10.6$, or $0.66n_e^{\text{lim}}$).

Conversely, in the experiments with the fast second current ramp-up
(Fig. 8(b)), the density profile was broader and the $D_\alpha$ trace demonstrates the
increase of recycling. In this case it was impossible to increase density. The maximum
value of the parameter $n_e Rq/B_t$ was 3.8 in these discharges (Fig. 9). Our
experimental data do not confirm the concept of density limiting by the boundary
overcooling due to intensive gas puffing. It was impossible to increase the plasma
density without a disruption, in spite of the plasma boundary heating through the
increase of the peripheral current density. Our data show the importance of recycling
for the density limiting mechanism. In the experiments described the decrease of
particle transport in the plasma edge results in the decrease of recycling.

6. CONCLUSIONS

The TUMAN-3 tokamak experiments showed:

(1) The transition to the regime similar to the H mode could appear in a tokamak
with circular cross-section, in a limiter configuration and without any auxiliary
heating.
(2) In HDM and the Ohmic H mode the confinement improves not only in the
periphery plasma but also in the $0 < r < 0.8a$ region.
(3) The peripheral plasma current decrease makes it possible to reach a higher
density using gas puffing.
(4) The intensive recycling prevents density increase.
REFERENCES

STUDIES OF ELECTRON HEAT WAVE PROPAGATION FROM ECRH AND SAWTOOTH OSCILLATIONS ON T-10

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Abstract

STUDIES OF ELECTRON HEAT WAVE PROPAGATION FROM ECRH AND SAWTOOTH OSCILLATIONS ON T-10.

Heat pulse propagation from central ECRH against the background of stationary non-central heating and from sawtooth oscillations in OH regimes and regimes with central ECRH is investigated. It is shown that the value of electron thermal diffusivity can either increase or decrease at the ECRH induced heat wave front. The velocity of the sawtooth induced heat pulse propagation increases with density, while the electron thermal diffusivity, obtained from the power balance, decreases with rising density.

1. INTRODUCTION

Heat pulse propagation (HPP) from ECRH and sawtooth oscillations on T-10 was studied earlier in Refs [1-3]. It was shown that the HPP velocity from central ECRH against the Ohmic background is enhanced in all the regimes studied. Let us assume the following: the effective electron thermal diffusivity \( \chi_e = \Gamma_e/\nabla T_e \), where \( \Gamma_e \) is the electron thermal flux; \( \chi_{PB}^e \) is the value of \( \chi_e \) defined from the power balance in the stationary stage. The dynamic electron thermal diffusivities \( \chi_{EC}^{HP} \) and \( \chi_{EC}^{st} \) are equal to \( \Gamma_e/\nabla T_e \) (\( \Gamma_e \) is the electron thermal flux perturbation, \( \nabla T_e \) is the electron temperature gradient perturbation) and are defined from analysis of HPP from ECRH and sawteeth. The value of \( \chi_{EC}^{HP} \) is 2-5 times greater than \( \chi_{OH}^{PB} \) for \( q_L = 3.8-7 \), \( I_p = 170-200 \) kA, \( \bar{n} = (2-3.5) \times 10^{19} \) m\(^{-3}\), \( B_z \approx 3 \) T and \( a_L = 28-34 \) cm. There is no strong dependence of \( \chi_{EC}^{HP} / \chi_{OH}^{PB} \) on density. The value \( \chi_{EC}^{HP} \) depends strongly on the radius in the area \( 0.5a_L < r < 0.7a_L \). The normal HPP velocity (\( \chi_{EC}^{HP} = \chi_{OH}^{PB} \)) was observed when the EC resonance was shifted 9 cm outwards. In this case the outward heat pulse caused almost no change in the relative gradient, i.e. the value \( \nabla T_e/T_e \) did not change during the process of heating.

HPP with a reduced velocity (\( \chi_{EC}^{HP} \approx \chi_{EC}^{PB}/2 \)) may exist. In this paper we shall consider in detail this effect, which was observed under central ECRH against the background of stationary non-central heating. The same results were briefly presented earlier [2].
We obtained new results in the studies of HPP from sawtooth oscillations. We found that under strong enough specific central ECRH power the sawtooth oscillations became 'saturated'. The HPP from these oscillations differs essentially from the 'classical' sawtooth oscillations. Taking into account this effect, we can investigate the dependence of $x_{lc}^{\text{EC}}$ on density. It was found that $x_{lc}^{\text{EC}}$ at least did not decrease with density rise, while $x_{lc}^{\text{PB}}$ is inversely proportional to density.

2. CONFINEMENT IMPROVEMENT AT THE ECRH INDUCED HEAT WAVE FRONT

Heat pulse propagation in the regime $\bar{n}_e \approx 4 \times 10^{19}$ m$^{-3}$, $I_p = 380$ kA, $B_t = 3.0$ T was studied in the following way. Non-central ECRH with absorbed power $P_{\text{ECRH}}^* \approx 0.9$ MW ($r_{\text{ECRH}} \approx 15$ cm) was switched on and suppressed sawtooth oscillations. Central ECRH with $P^* = 0.4$ MW was started 100 ms after the non-central ECRH. The electron temperature perturbation $\tilde{T}_e(r, t)$ dependence on time is shown in Fig. 1. Applying the method of Ref. [4] results in the following: $x_{\text{EC}}^{\text{HP}} = 0.55$ m$^2\cdot$s$^{-1}$ in the zone $r = 9$–17 cm; $0.35$ m$^2\cdot$s$^{-1}$ in the zone $r = 13$–17 cm; and $0.14 \pm 0.06$ m$^2\cdot$s$^{-1}$ in the zone $r = 15$–17 cm. It is seen from Fig. 1 that there is no power absorption (from the central ECRH) in the area $r > 15$ cm. On the one hand, $x_{\text{EC}}^{\text{HP}}$ (16 cm) = $0.14 \pm 0.06$ m$^2\cdot$s$^{-1}$; on the other hand, $x_{\text{EC}}^{\text{PB}}$ (16 cm) = $0.35 \pm 0.12$ m$^2\cdot$s$^{-1}$, i.e. $x_{\text{EC}}^{\text{HP}} = x_{\text{EC}}^{\text{PB}}/2$. It is possible to

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig1.png}
\caption{Central ECRH heat pulse propagation during the non-central ECR heated stationary state; solid lines: experiments; broken line: calculation, with $x_{\text{EC}}^{\text{HP}}$ (15 cm $\leq r \leq a_e$ cm) = 0.14 m$^2\cdot$s$^{-1}$.}
\end{figure}
understand this effect as follows. Using the results of Ref. [5], namely that $\chi_e^{HP} = \chi_e + (\partial T_e / \partial \nabla T_e) \nabla T_e$, where $\chi_e = \chi_e(\nabla T_e)$, we obtain: $(-0.2 \text{ m}^2 \cdot \text{s}^{-1})$, but with large error bars. It follows that $\chi_e$ decreases at the heat wave front.

In Fig. 2 one can see the time dependence of the relative gradient at the heat wave front, normalized by $\nabla T_e/T_e$. It is seen from Fig. 2 that the value of the relative gradient $\nabla T_e/T_e$ during the stationary stage of non-central ECRH was less than the Ohmic value. This value grew closer to the Ohmic one in 10 ms after the central ECRH was switched on. At the beginning of the central ECRH the profile $T_e(r)$ was much broader than the Ohmic one because the non-central heating started 100 ms earlier than the central heating. This time is shorter than the skin time. However, it is long enough for a considerable $j(r)$ redistribution. The density profile was constant in time. Observation of the density profile $n(r, t)$ and the SXR intensity profile showed the non-shifting of the plasma column at $t < 20 \text{ ms}$. Moreover, only shifting of the column or shifting of the absorption power profile from non-central ECRH inwards can cause such a deceleration of the HPP. But we see no physical reasons for the shifting. Thus the deceleration of the HPP can be explained only with a $\chi_e$ decrease at the heat wave front.

3. STUDIES OF HEAT PULSE PROPAGATION FROM SAWTOOTH OSCILLATIONS

To define the value of $\chi^{st}$, which characterizes the outward HPP velocity after the sawtooth crash, the formula $\chi^{st} = (r^2 - r_0^2)/8 t_p(r)$ [6] is widely used, where $r_0$
FIG. 3. Schematic diagram of $\hat{T}(r, t)$ values for saturated oscillations (solid lines) and classical sawtooth oscillations (broken lines).

is the mixing radius and $t_p(r)$ is the time at which $\hat{T}(r, t)$ reaches a maximum. This formula can be used only in the case of classical sawtooth oscillations, an example of which is shown in Fig. 3 (broken lines). The ECE measurements were used for $t_p(r)$ definition and SXR signals for additional information. To define the value of $\chi^m$ in the OH stage we used the method of Ref. [6] for classical oscillations.
First let us consider the dependence of $\chi_{\text{OH}}^n$ on the central chord density $n$ in the regime $I_p = 200$ kA, $B_c = 3.0$ T, $q_L = 3.8$, $a_L = 28$ cm. It is seen from Fig. 4 that under the condition $n = 2 \times 10^{19}$ m$^{-3}$ the values of $\chi_{\text{OH}}^n$ before ECRH are less than after ECRH. For $\bar{n} = 1.0 \times 10^{19}$ m$^{-3}$ and $1.2 \times 10^{19}$ m$^{-3}$ the values of $\chi_{\text{OH}}^n$ after ECRH are almost the same as for $\bar{n} = 2 \times 10^{19}$ m$^{-3}$. Unfortunately, we are not entirely certain about $\chi_{\text{OH}}^n$ before ECRH under the condition $n = 1.0 \times 10^{19}$ m$^{-3}$, because the oscillation amplitude is too small. As is seen from Fig. 4, the value of $\chi_{\text{OH}}^n$ before and after ECRH is the same for $n = 3 \times 10^{19}$ m$^{-3}$. All these data indicate that there is no inverse dependence of $\chi_{\text{OH}}^n$ on $n$ for $n = (1-3) \times 10^{19}$ m$^{-3}$, while values of $\chi_{\text{PB}}^{\text{HR}}$ have the usual inverse dependence on $n$.

In our opinion, it is of great importance that for $n = 2 \times 10^{19}$ m$^{-3}$ and $3 \times 10^{19}$ m$^{-3}$ the value of $\chi_{\text{OH}}^n$ and the value of $\chi_{\text{EC}}^{\text{HP}}$ [3] coincide reasonably: 0.35 m$^2$ s$^{-1}$ ($r_0 = 8$ cm, HPP was studied in the region $r = 8-14$ cm) and 0.5 m$^2$ s$^{-1}$ ($r = 14-16$ cm) at $\bar{n} = 2 \times 10^{19}$ m$^{-3}$; and at $\bar{n} = 3 \times 10^{19}$ m$^{-3}$, $\chi_{\text{OH}}^n = 0.45$ m$^2$ s$^{-1}$ and $\chi_{\text{EC}}^{\text{HP}} = 0.4$ m$^2$ s$^{-1}$ at the same points. For $\bar{n} = 2 \times 10^{19}$ m$^{-3}$ and $\bar{n} = 3 \times 10^{19}$ m$^{-3}$, $\chi_{\text{OH}}^n = 0.24$ m$^2$ s$^{-1}$ and 0.18 m$^2$ s$^{-1}$ at $r = 15$ cm. The values of $\chi_{\text{OH}}^n$ and $\chi_{\text{EC}}^{\text{HP}}$ are two times greater than $\chi_{\text{OH}}^n$.

Now let us consider HPP from sawtooth oscillations under central ECRH. Under certain circumstances saturated sawtooth oscillations are observed on T-10. The values of $\bar{T}(r, t)$ for saturated oscillations are shown in Fig. 3 by the solid curves. One can see the difference between classical and saturated oscillations: for the saturated type the value of $\bar{T}(r, t)$ increases rapidly during $t = (0.1-0.15)\tau_p$ (where $\tau_p$ is the period of the oscillations) and becomes saturated in the central part of the core. The value of $\bar{T}(r, t)$ increases first in the zone $r < r_s$, and then it drops almost to its initial level. It is a matter of fact that under $t > 0.15\tau_p$ an additional heat outflow into the region $r = r_0$ takes place. Energy content in the region $r < r_s$ stops increasing, i.e. the outward flux transports the whole power from the region $r < r_s$ into the region $r = r_0$. This means that $\chi_{\text{EC}}(r_s)$ increases. The flux from the region $r < r_s$ leads to deceleration of the $\bar{T}(r_0, t)$ decay velocity and gives $\tau_p(r)$ increasing for $r > r_0$. This is clearly seen from Fig. 3. The calculations performed (according to Ref. [4]) show that the $\tau_p(r)$ value can increase 2.5-5 times, depending on the details, in the transfer from the classical disruption to the saturated one. Let us note that the method of Ref. [4] allows one to take into consideration the additional energy flux from the core centre. But as the first step we define $\chi_{\text{EC}}^{\text{HR}}$ by the previous expression [6] and then increase it 3 times for saturated oscillations.

Figure 4 shows the values of $\chi_{\text{EC}}^{\text{HR}}$ for $I_p = 200$ kA, $P_{\text{ECR}}^* = 0.36$ MW and for $I_p = 240$ kA, $P_{\text{ECR}}^* = 1.0$ MW. One can see that oscillations with $P_{\text{ECR}}^* = 0.36$ MW are saturated at $\bar{n} = 1 \times 10^{19}$ m$^{-3}$ and become classical at $\bar{n} = 2 \times 10^{19}$ m$^{-3}$. But oscillations with $P_{\text{ECR}}^* = 0.6$ MW are again saturated at $\bar{n} = 2 \times 10^{19}$ m$^{-3}$. This means that the saturated sawtooth oscillations are observed under central ECRH with high power per particle. It is seen from Fig. 4 that the value $\chi_{\text{EC}}^{\text{HR}}$ increases with density rise, while $\chi_{\text{EC}}^{\text{PB}}$ has the usual inverse dependence on $n$, $\chi_{\text{EC}}^{\text{PB}} \sim n^{-1}$. 
4. CONCLUSIONS

The $\chi_e$ value can either increase or decrease at the heat wave front, i.e. under fast processes when only the local electron temperature profile has enough time to change significantly. This means that profile self-consistency holds. It seems that the 'optimal' relative profile $\nabla T_e/T_e$ must be close to the Ohmic one. The value of $\chi_e$ increases when the profile is not the optimal one.

The behaviour of $\chi_e$ at the heat pulse front coincides qualitatively with the behaviour of $\chi_{EC}^B$, which changes according to the scaling of Ref. [7], and depends strongly on the EC power radial profile as well.

The strong dependence of HPP velocity on the sawtooth oscillation form was proven. The saturated sawtooth oscillations are observed under central ECRH with high power per particle, where an additional heat outflux from the centre takes place in the region of $r \approx r_0$. A slow temperature perturbation decay takes place in this region. With $r > r_0$ the maximum electron temperature perturbation is observed in a time 2.5–5 times greater than under the classical oscillation.

The dynamic heat thermal diffusivity $\chi^D$ defined by the HPP velocity from the sawtooth oscillations is always greater than $\chi^{PB}$. Moreover, the values $\chi^{st}$ and $\chi^{PB}$ depend on the density in an opposite manner in the Ohmic regimes as well as under central ECRH: $\chi^D$ increases with $\bar{n}$ while $\chi^{PB} \sim n^{-1}$.

It is important that the dependence of the $(\chi^{st} - \chi^{PB})$ value on density is close to linear under central ECRH. It is this value that characterizes plasma response to deviation from the optimal profile. This means that in any transport model (linear or non-linear) the term from the HPP velocity depends on and must be proportional to density.

The HPP velocities from central ECRH as well as from sawtooth oscillations up to ECRH coincide reasonably.

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LOCKED MHD PERTURBATIONS IN T-10 TOKAMAK

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Abstract

LOCKED MHD PERTURBATIONS IN T-10 TOKAMAK

The experimental results on the 'locking on' of MHD perturbations in the T-10 tokamak as well as results of a numerical simulation of the performed experiment are given. The main reason for the m = 2 and m = 3 mode locking is related to their interaction as non-linear two-dimensional oscillators.

1. INTRODUCTION

The phenomenon of plasma MHD perturbation locking manifests itself as an interruption or strong deceleration of the rotation of the magnetic island structure. It is observed in many tokamaks [1-4] and attracts interest because of its correlation with the disruptive instability and its effect on transport processes. Mode locking can take place at any stage of a tokamak discharge, but more frequently it happens at the current ramp-up stage or accompanies disruptions.

Mode locking is usually explained by two possible reasons: first, dissipation of eddy currents induced in the discharge chamber wall due to MHD mode rotation and second, mode interaction with stationary multipolar magnetic fields (due to, for example, inaccurate positioning of the toroidal field coils) [5]. The first of these two factors cannot completely explain the locking of MHD perturbations. This factor ceases to act when the frequency is decreased and, thus, cannot result in a termination, but only in some deceleration of the rotation. The second factor can explain the stop of the rotation in definite cases, e.g. in experiments with a quadrupolar coil in T-7 tokamak [6]. This factor can, however, only be decisive when the spatially resonant component of the external perturbing field is comparable with (or exceeds) the amplitude of the spontaneously developing MHD perturbation in the plasma. In most cases, when the perturbing field is not produced by special external coils, this condition is probably not satisfied. Besides, the mechanism of mode locking by an external stationary field would result in a stop of the low amplitude modes and not prevent high amplitude mode rotation; the experiments show, however, the opposite behaviour [1, 4].

In the experiments on T-10, it was found that the stops of the rotation for the spatial m = 2 and m = 3 modes occur simultaneously. In all cases of locking recorded, these two harmonics rotated in opposite directions at comparable angular
velocities just before the stop. This fact allows us to explain the phenomenon of locking on the basis of a non-linear mechanism of frequency synchronization as suggested in Ref. [7]. The experiment that led us to this conclusion as well as the corresponding numerical model will be described below.

2. EXPERIMENTAL ARRANGEMENT

A set of 24 magnetic detectors located behind the discharge chamber uniformly along the poloidal direction was used to record the poloidal magnetic field perturbations. A spatial Fourier analysis of the MHD perturbations was performed on the T-10 tokamak with an analogue analyser [8], which separated sine and cosine Fourier components of the perturbations with poloidal m-numbers from m = 1 to m = 6. The measurement of the sine and cosine components allowed us to determine the current value of the angular velocity (including its direction) of the poloidal rotation for each mode as well as the value of the mode amplitude.

The procedure of digital processing for each signal for sine and cosine components included the following operations:

- compensation of the integrating effect of the chamber wall;
- signal filtering to eliminate induced stray signals from the tokamak power supply system;
- integration of the signal to obtain the field perturbation from its time derivative;
- calculation of the instant values of the MHD perturbation amplitude and rotation velocity from the sine and cosine components obtained.

3. EXPERIMENTAL RESULTS

In T-10, mode locking usually occurs in discharges with plasma currents I_p > 300 kA during current ramp-up. In some cases, the locking was observed at other discharge stages and, as a rule, it accompanied a disruptive instability.

Typical cases of m = 2 and m = 3 mode locking at the initial stage of the discharge and during the development of the disruption are shown in Figs 1 and 2. We see that in both cases the m = 2 and m = 3 islands were rotating in opposite directions just before their stop. This feature in the m = 2 and m = 3 modes was observed in all cases of locking recorded.

On the other hand, when the m = 2 and m = 3 modes were rotating in the same direction, synchronization of the oscillations took place, but the rotation did not stop. The behaviour of the m = 2 and m = 3 sine components in a typical case of such a synchronization is shown in Fig. 3. We see that frequency synchronization has a threshold and occurs at rather high magnitudes of the perturbations, as is noted in Ref. [7].
Figure 1. Plasma current ($I_p$), loop voltage ($U_p$), amplitude ($\tilde{B}_\theta$), and frequency, $f$, of spatial $m = 2$ and $m = 3$ modes during locking process at the stage of current ramp-up. The sign of the frequency indicates the direction of the rotation: 'plus' refers to the direction of the electron diamagnetic drift.

Figure 4, which corresponds to the discharge shown in Fig. 1, shows the relationships between the amplitudes and frequencies of the $m = 2$ and $m = 3$ spatial modes in the cases of co-rotation (open circles) and counter-rotation (solid circles). As follows from Fig. 4, in the case of counter-rotation, the observed oscillation frequencies turn out to be considerably lower and virtually reduce to zero at high amplitudes, a phenomenon that looks like locking process.

The regularity revealed — counter-rotation of $m = 2$ and $m = 3$ modes before their simultaneous stop — allowed us to consider a possible explanation for mode locking, based on a description of these two spatial harmonics as coupled non-linear
FIG. 2. Plasma current ($I_p$), loop voltage ($U_p$), sine components of poloidal magnetic field time derivative ($\dot{B}_{es}$) for $m = 2$ and $m = 3$ spatial harmonics, amplitudes ($B_e$) and frequencies ($f$) of these modes during locking process accompanying the disruption.
FIG. 3. Sine components of $m = 2$ and $m = 3$ modes under mutual synchronization in the case of co-rotation.

oscillators. The phenomenon of synchronization for such oscillators is well known in the theory of non-linear oscillations [9]. Synchronization takes place at sufficiently strong coupling between the oscillators and is usually observed at one of the natural frequencies or at an intermediate frequency between the natural frequencies.

For each MHD mode one has a two-dimensional oscillator whose state is determined by two orthogonal spatial components — sine and cosine. Therefore, the oscillations of such an oscillator are characterized not only by its amplitude and frequency, but also by the direction of the mode rotation which may be represented by the sign of frequency. For the observed MHD modes, the non-linearity is a result of the tearing mode saturation going along with magnetic island growth [10]. The coupling of different spatial harmonics is a consequence of the toroidicity and the displacements of the plasma column along the major radius of the torus and in the vertical direction. If the natural frequencies of two-dimensional oscillators have the same signs and their coupling is strong enough, the oscillation frequencies change to some intermediate frequency; this looks like a synchronization without stop. If the natural frequencies have opposite signs and their absolute values are close to each other, the intermediate frequency at which synchronization occurs can turn to be close to zero. In this case, the result of the interaction between oscillators looks like a substantial deceleration of the oscillations, i.e. locking process.
FIG. 4. Relationships between frequencies (f) and amplitudes ($\tilde{B}_g$) of $m = 2$ and $m = 3$ modes in the cases of co-rotation (open circles) and counter-rotation (solid circles) of the given spatial harmonics.

4. NUMERICAL SIMULATION

The equations usually applied to represent tearing mode interaction with an external field [11–13] and adapted to the case of non-linear interaction between two modes were used in the simulation of the given experiment:

\[
\begin{align*}
\dot{\psi}_c^2 &= \gamma_2 \psi_c^2 - 2\Omega_2 \psi_c^2 + \kappa_{32} \psi_3^2 \\
\dot{\psi}_c^2 &= \gamma_2 \psi_2^2 + 2\Omega_2 \psi_c^2 + \kappa_{32} \psi_3^2 \\
\dot{\psi}_c^2 &= \gamma_3 \psi_c^3 - 3\Omega_3 \psi_c^3 + \kappa_{23} \psi_2^3 \\
\dot{\psi}_c^3 &= \gamma_3 \psi_3^3 + 3\Omega_3 \psi_c^3 + \kappa_{23} \psi_2^3 
\end{align*}
\]
Fig. 5. Calculated dependences of the oscillation frequencies ($\tilde{f}$) for spatial $m = 2$ and $m = 3$ modes (normalized to $\Omega/2\pi$) and perturbed magnetic flux amplitudes $\tilde{\psi}_m = (\psi^2_m + \psi^4_m)^{1/2}$ (the numerical values of the $m = 2$ and $m = 3$ frequencies coincide) on the natural frequencies of the magnetic island rotation for $|Q_2| = |Q_3| = \Omega$. Solid lines correspond to co-rotation and dashed lines to counter-rotation.

where $\psi_{c2}$, $\psi_{s2}$, $\psi_{c3}$, $\psi_{s3}$ are the cosine and sine components of the perturbed helical flux for the $m = 2$ and $m = 3$ modes, $\gamma_m = 2\alpha^2 \omega_R \Delta_m / W_m$ are the parameters determining the growth rate of the magnetic island width $W_m \sim (\psi^2_{cm} + \psi^4_{cm})^{1/4}$ in the Rutherford regime, $\Omega_2$ and $\Omega_3$ are natural frequencies of mode rotation, which, in general, depend on the plasma parameters and amplitudes of these modes, and $K_{c2}$, $K_{s2}$, $K_{c3}$, $K_{s3}$ are the coupling coefficients for the cosine and sine components of the $m = 2$ and $m = 3$ modes.

We do not aim at a detailed description of the observed phenomenon of mode locking, but rather intend to carry out a fundamental verification of the idea of its
possible mechanism. So, let the quantities $\Delta'_1$ and $\Delta'_3$ depend linearly on the magnetic island sizes, $\Delta'_m = \Delta'_{m0}(1 - W_m/W_{m,\text{sat}})$. Their values at a zero island widths, $\Delta'_{m0}$, the values $(\psi_{sm}^2 + \psi_{cm}^2)^{1/2}|_{\text{sat}}$, corresponding to saturated islands, $W_{m,\text{sat}}$, and the coupling coefficients are input parameters. Since the estimates of the coefficients $\gamma$, $\Omega$ and $\kappa$, performed in the conditions of our experiment, yield the same orders of magnitude, the parameters $\kappa_{c32}$, $\kappa_{s32}$, $\kappa_{c23}$, $\kappa_{s23}$, $\gamma_2(W_2 = 0)$, $\gamma_3(W_3 = 0)$, and the values $(\psi_{sm}^2 + \psi_{cm}^2)^{1/2}|_{\text{sat}}$ will be set equal to unity.

The solution of the equations is given in Fig. 5. We see that an essential reduction of the oscillation frequencies (in comparison with the natural frequencies) takes place in the case of counter-rotating modes (dashed lines). This result corresponds to the experimental data given in Fig. 4 and confirms the suggested explanation of the mode locking mechanism.

A qualitatively similar result can be obtained without numerical simulation. Indeed, the stop of rotation corresponds to the stationary solution of the set of equations, which is only possible when the determinant of the matrix of the equation set on the right hand side vanishes. This determinant is positive definite in the absence of a coupling between the modes. The coupling corresponds to the term $\Omega_2\Omega_3(\kappa_{c32}\kappa_{s23} + \kappa_{s23}\kappa_{c32})$, in this determinant. Thus, when the natural frequencies have opposite signs, this term furthers the existence of a stationary solution, in contrast to the case where the natural frequencies have the same sign.

---

**FIG. 6. Results of a simulation of mutual synchronization of the $m = 2$ and $m = 3$ modes in the case of co-rotation. The sine and cosine components of these modes are shown versus time.**
The result of $m = 2$ and $m = 3$ mode synchronization simulation, when the modes rotate in the same direction at different frequencies, is shown in Fig. 6. To find the dependence of the synchronization on the ratio of the mode amplitudes, the saturated amplitude of the $m = 3$ mode was chosen to be stationary and equal to unity, and the $m = 2$ amplitude was raised linearly in time, from 0.3 to 3. In Fig. 6, we see that, with its amplitude rising, the $m = 2$ mode first gets out of the locked state with the help of the $m = 3$ mode, making it rotate at its own natural frequency. Note the similarity of this calculation result to Fig. 4.

5. CONCLUSIONS

The experimental result obtained — counter-rotation of the $m = 2$ and $m = 3$ modes directly before their stop — and the results of the numerical simulation indicate that the main reason for MHD mode locking under the conditions of our experiment can be attributed to the behaviour of the modes as coupled two-dimensional non-linear oscillators. It should not be excluded that other factors, such as the effect of external stationary perturbed fields, deceleration due to eddy currents in the chamber wall and a dependence of the natural frequencies on the mode amplitudes, assist in the locking process.

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ERGODIC DIVERTOR EXPERIMENTS ON TORE SUPRA

EQUIPE TORE SUPRA*

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Abstract

ERGODIC DIVERTOR EXPERIMENTS ON TORE SUPRA.

The ergodic divertor (ED) creates a magnetic perturbation which changes the transport in the edge layer. The experiments reported in the paper show that the edge temperature is reduced and that the power is channelled to the neutralizer plates. The reduced plasma density and carbon density are correlated to a reduced particle life time at the very edge of the plasma.

1 Magnetic perturbation

The magnetic perturbation of the Tore Supra ergodic divertor is induced by 6 coils equally spaced toroidally, producing a radial field $\delta B_r$ with a periodic poloidal variation over $\Delta \theta \sim 2\pi/3$. As shown in Fig.1 by field line tracing, a field line undergoes a kick as it passes in front of a coil [1]. The series of kicks builds up for $q_{\text{edge}}$ close to 3, determining the ED resonance [2]. The radial shift is then $\delta r \sim 20$ mm per poloidal rotation. As a consequence of this loss of phase ($\delta q \sim 0.4$) the field lines experience a stochastic behaviour characterized by a radial diffusion coefficient $D_QL = \langle \delta r^2 \rangle / q R$. In the stochastic layer the heat and particle transport combine the intrinsic anomalous perpendicular transport and the parallel transport weighted by the field line diffusion [3]. The latter modified transport falls off rapidly towards the plasma center. Further control is achieved by connecting the stochastic layer to the neutralizer plates implemented between the current bars of the ED coils. However, up to now the experiments have been performed with a limiter standing 30 mm ahead of the ED yielding a rather involved connection topology [1].

* The Equipe Tore Supra authors' list is to be found in the Appendix to paper IAEA-CN-53/E-II-2, these Proceedings, this volume.
2 Profile modifications

The electronic temperature determined by Thomson scattering shows that the ED leads to small modifications of the central temperature while the edge temperature is reduced from 210 to 50 eV at $\rho = r/a = 0.94$ (Fig. 2). The $1D^1$ transport analysis of the confinement domain $0.5 < \rho < 0.8$ with the code Intermede [4] shows that the
steepening of the profile is correlated to the shrinking of the current channel ($l_i$ increases from 1.22 to 1.37) and that the perpendicular heat conductivity is unchanged, $\chi_\perp \sim 3.10^{19} \text{ m}^{-1} \text{ s}^{-1}$. Theoretical calculations in the ergodic non-collisional regime $\chi_{\text{erg}} \propto D_{QL}(\rho) V_{\text{the}}(\rho)$ [3] with no radiation, exhibits a strong flattening of the edge temperature up to $\rho \sim 0.85$. This would require a localized decrease of the anomalous transport at $\rho < 0.85$. Similar temperature profiles have been reported on TEXT [5].

The density profiles, measured by reflectometry ($\circ$), Thomson scattering ($\Delta$) and interferometry, depend on the plasma species. While He-plasmas exhibit unmodified central densities [6], the H or D plasmas exhibit a factor 2 density decrease when the ED is activated, Fig. 3. The normalized density profile, $n_e(\rho)/n_e(0)$, is to first order unaffected which advocates for a transport equation in the bulk, $-D \nabla n_e + V n_e = 0$, with unperturbed values of $D/V$. The density decrease is then a signature of the different connection to the wall which leads to a transfer of particles from the plasma to the wall. This view is confirmed by the reversibility of this phenomenon which is demonstrated by recovering the density profile when the ED is turned off. A similar wall pumping occurs when the plasma is moved from the outer limiter to the inner wall [7].

3 Particle and heat load on the wall

The divertor effect of the ED is demonstrated by the calorimetric measurements (Fig. 4). By increasing the magnetic perturbation more
power is diverted from the limiter and channelled to the neutralizer plates implemented between the current bars of the ED. IR imaging [8] of the limiter shows that the heat load is reduced on average, yielding lower limiter temperatures. This redistribution corresponds to a 15 mm increase of the energy e-folding length. The hot spots which appear on the leading edges of the limiter are well explained by the calculated connection topology [1]. The particle deposition monitored by the $H\alpha$ emission is also redistributed and the remaining non-uniformity correlated to the connection properties.

4 Impurity control

The VUV spectrometer measurements calibrated on the Bremsstrahlung signal and on the soft Xrays is used to determine the densities of the major impurities (i.e. carbon and oxygen) and hence $Z_{\text{eff}}$. Figure 5 shows the carbon density (monitored by CVI-brightness / $n_e$) for different configurations, limiter (o), inner wall after the shift from the limiter (Δ), and ED (▲). A typical ED operation yields a plasma density decrease to $n_e \sim 1.1 - 1.4 \times 10^{19}$ m$^{-3}$ and a carbon density decrease. This differs from the wall pumping experiment where only a weak decrease of the carbon density is observed. In Table I, the carbon and oxygen content with ED is compared to a reference limiter plasma. Under these similar plasma densities one finds a reduction of the carbon content which ranges from 50% to 80%. The lowest $Z_{\text{eff}}$ value achieved during the ED operation, $Z_{\text{eff}} \sim 2$, is a factor 2 below the ohmic reference case.
FIG. 5. Variation of carbon density (C VI line brightness/\(\bar{n}_e\)) for three different configurations. Note that the inner wall results are obtained after a plasma swing from the limiter to the inner wall.

TABLE I. COMPARISON OF ED AND LIMITER PLASMAS

<table>
<thead>
<tr>
<th></th>
<th>(\bar{n}_e)</th>
<th>(n_e/\bar{n}_e) (%)</th>
<th>(n_\alpha/\bar{n}_e) (%)</th>
<th>(Z_{eff})</th>
<th>(D_{95}(a))</th>
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<td>3.2</td>
<td>0.9</td>
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<td>-</td>
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<td>(5 \times 10^{-5})</td>
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While the innermost impurity line emission is reduced during the ED operation, the peripheral lines exhibit an increase. This effect is well correlated to the radiated power profile [9] derived from the bolometric measurements. However this edge radiation increase does not yield a significant effect on the total radiated power which decreases linearly with the density from 900 kW to 650 kW.

Non-intrinsic impurity behaviour has been studied with gaseous nitrogen injection. A time trace of the NVII/\(\bar{n}_e\) signal for these experiments is given in Fig.6. The ED is activated after a first nitrogen injection and the time scale of the signal decrease is unaffected by the ED.
FIG. 6. Time trace of nitrogen density, NVII/n_e, for two nitrogen injections (broken line on the lower trace) with an ED 45 kA pulse (full line in lower diagram).

(With respect to the reference discharge without ED). The second nitrogen injection occurs during the ED pulse. As shown on Fig. 6, this injection is screened out and the nitrogen signal reduced by a factor 5 with respect to the reference case. This screening effect is the counterpart of a reduced life time of the incoming ions which is determined by an effective increase of the transport at the very edge of the plasma or by a SOL loss effect, i.e. the impurities ionized on an open field line have a short "parallel life time" before impinging on the wall. The latter mechanism should be reinforced by both the divertor effect (the open field lines go deeper towards the bulk) and the larger "plasma e-folding length".

A similar screening of the intrinsic impurities and of the plasma ions must occur. However the different generation mechanisms and wall content (source) will lead either to a reduced content in the plasma, as for carbon and the plasma ions during the density decrease, or to a constant content when the particle source is proportional to the particle outflux, as for the plasma ions during the ED density plateau and likely for the oxygen.

5 Conclusion

The Ergodic Divertor operation is characterized by an unaffected bulk confinement and changes at the edge with different responses for the edge temperature and edge density. The reduced edge temperature extends over 15% of the minor radius according to theoretical heat transport calculations (in qualitative agreement with the reduced current channel). The particle behaviour - both plasma ions and impurities - is
only modified at the very edge of the plasma with a reduced peripheral life time. The wall particle content and desorption mechanisms then govern the ED effects. This allows to reduce the carbon content of the discharge and to decrease the $Z_{eff}$ from 4 to about 2. The divertor effect is demonstrated both by a channelling of the power to the neutralizer plates, and an increase of the "SOL thickness" which tends to average the heat load over the limiters.

REFERENCES

EFFECTS OF AN ERGODIC MAGNETIC LIMITER ON HYBTOK-II TOKAMAK PLASMA

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Abstract

The effects of a stochastic magnetic field on edge plasma properties and impurity behaviour are investigated in a small research tokamak. The power deposition profile on the limiter was broadened by edge transport enhancement, and the edge plasma tended to be cooled favourably from the viewpoint of plasma-wall interaction, although some structural modifications were identified. It was observed that the edge radiation was increased in the visible light region. The impurity control potential demonstrated for heavy metals may possibly be found to work for all impurities in larger machines.

1. INTRODUCTION

A stochastic magnetic field has many important intrinsic and extrinsic aspects in magnetic confinement fusion devices. The transport of electrons, particularly suprathermal electrons, is sensitive to the presence of such a magnetic field [1]. The power deposition profile on a divertor or a limiter plate is influenced by a stochastic magnetic field generated by a misalignment of poloidal field coils or a discreteness of toroidal field coils [2]. An applied resonant field at the tokamak edge has demonstrated many interesting edge plasma modifications and shows potential value for improving the performance of fusion devices [3, 4].

In this paper modification of plasma properties and transport at the tokamak edge by an ergodic magnetic limiter (EML) are presented and discussed.

2. EDGE PLASMA MODIFICATIONS

A resonant helical magnetic field was applied on the HYBTOK-II tokamak ($R_0 = 0.4$ m, $a_w = 0.13$ m, $a_p = 0.11$ m, $B_t = 0.4$ T, $I_p \sim 15$ kA, $\tau_d \leq 10$ ms) plasma with six local helical coils ($m/n = 6/2$). Figure 1 shows a typical example of a puncture plot of magnetic field lines on a poloidal plane [5]. We had two operational regimes, one of which had a line averaged electron density $\bar{n}_e$ of about
FIG. 1. Puncture plot of magnetic field lines on a poloidal plane. Broken line indicates the location of the Mo limiter. \( I_p = 75 \text{ kA}, B_\perp = 0.4 \text{ T}, I_h = 0.4 \text{ kA}, \Lambda = -0.74, m/n = 6/2. \)

1 \times 10^{19} \text{ m}^{-3} and the other around 0.5 \times 10^{19} \text{ m}^{-3}. We call the former the high density mode (HDM) and the latter the low density mode (LDM).

From the observed radial distribution of electron density in the SOL shown in Fig. 2(a) for the LDM, the diffusion coefficient \( D \) across the main magnetic field was obtained, neglecting ionization in the SOL and a spatial modulation generated by the stochastic magnetic field structure. \( D \) gradually increases from an intrinsic value with the helical perturbation. Therefore, the plasma transport at the edge is indeed increased by such an ergodic magnetic layer. Taking account of a finite electron excursion length between limiters, \( 2L_e = 2\pi R_\rho \), and a poloidal \( E_r \times B \) drift due to a radial electric field in the SOL, we can obtain the characteristic diffusion step size \( \Delta r = L_e(b_h/B_\perp) \) and the characteristic time step \( \tau = (l_p/\Delta l_p) (L_e/v_e) \), where \( \Delta l_p \) is the poloidal drift distance which electrons travel during the time needed to traverse the excursion length and is given by \( (E_r/B_\perp) (2L_e/v_e) \); \( l_p \) is the poloidal distance necessary for the transition to a magnetic field line independent of the original one, or in other words, the poloidal correlation length in the magnetic field structure; and \( b_h \) is the magnitude of the helical magnetic field in the SOL. The second factor of the expression \( \tau \) comes from the time needed for an interlimiter excursion. Then \( D = (\Delta r)^2/(2\tau) \) [6]. In the experiment \( D \) is proportional to \( (b_h \propto l_h)^{1.6} \), as shown in
Fig. 2. The dependence does not differ so much from the above theoretical prediction $D \propto I_h^2$ if $I_p$ is assumed constant. In the case of $I_h = 300$ A, $b_h \approx 20$ G and $D \approx 42(\Delta l_p/l_p)m^2/s$. The experiment gives the value $D \approx 2$ m$^2/s$, which means $l_p \approx 0.1$ m, corresponding to 50° in the poloidal angle. This is a reasonable value, inferring from the magnetic field structure shown in Fig. 1.

The electron temperature profile for the LDM was also flattened, as shown in Fig. 3; we observed a slight decrease of $T_e$ at the limiter edge, while a greater reduction of $T_e$ was observed in the HDM [4], where higher recycling was expected. Such a broadening of heat flow to the limiter plate is favourable for protecting against a possible erosion of target material and impurity contamination of the
core plasma. This suggests that a stochastic magnetic field near the separatrix may be expected also to have the potential to broaden the heat deposition profile on the divertor target plate.

In the HDM experiment we have pointed out a poloidal modulation of modified plasma parameters [4]. This is related to the structure of the magnetic field close to the wall. In further study of the poloidal variation of plasma properties, the ion saturation currents at the limiter edge observed at 32 poloidal points showed that the edge density did not change or increased somewhat for a weak perturbation, while a strong helical field enhanced a spatial modulation, consistent with the puncture plot of magnetic field lines [7].

Although magnetic fluctuations picked up by Mirnov coils were not affected by the EML, changes in electrostatic potential fluctuations were detected at the edge with an array of multi-Langmuir probes. These came from a helical instability with
the mode number $m/n = 2/1$, the resonance zone of which is located inside and is not directly influenced by the present external helical field of high poloidal mode number. Nevertheless, the amplitude was found to be dependent on the poloidal position, well correlated with the magnetic field structure at the edge.

Three port computed tomography (CT) using CCD line sensors gave a time series of radiation profiles in the visible light region (420-870 nm), as shown in Fig. 4, obtained in the HDM. The EML enhances the edge radiation more than that in the central region. This is thought to come from three components: the H$_a$ line, light impurity emissions, and low ionization stage metallic lines, as will be mentioned below. The contribution of H$_a$, about 20%, came from the enhanced hydrogen recycling observed in the experiment [4]. Powerful radiation at the edge could lead to a radiative cooled edge favourable to plasma-wall interactions. To improve the spatial resolution at the edge, five port CT is desirable and is now under way.

3. IMPURITY BEHAVIOUR

Here we discuss impurity behaviour in the HDM. As shown in Fig. 5(a), some intrinsic light impurities were observed to be increased by the EML. This increase is associated with an enhancement of the following physical and chemical processes caused by the increase of the contact of the plasma with solid materials. First, physical desorption on the wall and the limiter may be increased by the enhancement of photons and energetic charge exchanged neutral atoms due to the increase of hydrogen recycling. Second, an increase of chemisorption is also responsible. Chemical reactions between incident plasma atoms or ions and surface atoms lead to the formation of molecules, for example CH$_4$, CO and H$_2$O. These molecules can then be desorbed easily.

In Fig. 5(b), changes in some spectral line intensities originating from the intrinsic metallic impurities are shown. We should note that iron and nickel are the elements of the vessel wall material, and titanium is that of gettering. The open symbols indicate atomic lines (Fe(I), Ni(I) and Ti(I)), while the other symbols show ionic lines. An opposite tendency of atomic and ionic lines is clearly seen. This tendency does not change even if the correction for spectral intensities due to the density increase is not made. The reason for the increase of metal atomic lines is considered to be as follows. First, sputtering from the wall by ion bombardment may be increased by the density increase at the edge. Second, sputtering of these impurities at the limiter due to the increase of the plasma density overcomes the effect of a reduction of the electron temperature. Finally, an enhancement of sputtering from both the wall and the limiter is thought to come from an increase of energetic charge exchanged neutral atoms due to the increase of hydrogen recycling. On the other hand, the spectral line intensities of metal ions were found to be reduced. This implies a kind of impurity control associated with the EML.
FIG. 4. Visible light radiation profiles on a poloidal plane obtained with three port CT. (a) Time series of 2-D profiles with EML. (b) 2-D profiles of emission increment due to EML applied at $t = 4$ ms with a duration of 4 ms.
FIG. 5. Changes in intensities of spectral lines at $t = 5$ ms. $n_x$ is $1.0 \times 10^{19} \text{ m}^{-3}$ without EML.
(a) Emission from light impurities. Solid symbols show the values corrected by the $n_x$ increase, open symbols are raw data: (O,●): C(II) (426.7 nm); (∆,▲): O(V) (278.7 nm); (□,■): N(II) (463.1 nm).
(b) Emissions from metallic impurities, with correction due to the $n_x$ increase. Open symbols show neutral atomic lines, other symbols show ionic lines: O,●: Fe(I), Fe(II), Fe(III); ▲,▲: Ni(I), Ni(II);
□,■,■: Ti(I), Ti(II), Ti(III), Ti(IV).

Now we would like to discuss possible reasons for the different behaviour of the light and the metallic impurities. The plasma column is divided into two regions, A and B (see Fig. 1). Region A corresponds to the core plasma region not influenced directly by the external helical perturbation. Therefore, the impurity transport there may not be modified much by the EML. Region B is an ergodic magnetic layer, where a relatively low temperature and high density edge plasma is generated by the enhancement of recycling caused by increased plasma transport with the EML. In our case the HDM has shown this rather than the LDM. The resulting cooled edge is supposed to reduce impurity generation by sputtering and to provide better screening,
FIG. 6. Penetration of H, C and Fe atoms through the edge plasma. Insert shows an assumed radial profile of electron density and temperature. Solid curves show normalized inward flux of neutral atoms, broken curves the normalized ion production rate. The energy of injected neutral atoms is assumed 2 eV.

i.e. released impurity atoms from the wall or the limiter are easily ionized in this region, and the ionized impurities return to the wall along the disturbed magnetic field lines rather than diffuse to the core plasma, because we have the relation $n_{imp}(0) = \phi \lambda_{ion}/D = \phi v/(Dn_e S_i)$ [8], where $\phi = n_{imp}(a_w)v$ is the flux of incoming neutrals, $\lambda_{ion}$ is the ionization mean free path of impurity atoms, $S_i$ is the ionization coefficient at the edge, and $n_e$ is the edge electron density. Both $D$ and $n_e$ were observed to be increased. Furthermore, an increased flow of hydrogen ions due to the recycling tends to draw impurity ions back to the wall. The positive potential with respect to the wall helps impurity ions return to the wall.

Therefore, it is important to know where impurity atoms would be ionized in order to infer the effect of the EML on the impurity transport. Figure 6 shows the normalized influx $F/F_0$ and the local ion production rate $S$ for hydrogen, carbon and iron atoms. In a descending scale of the penetration depth into the plasma, we have hydrogen (the working atom), carbon (the light impurity) and iron (the metallic impurity). Hence the light impurities (C, O, N) and the working atoms (H) arrive at the core plasma (region A). That is, they are not so much influenced by the EML. On the other hand, almost all the metallic impurities are ionized in region B.

In larger machines, almost all the light impurities are ionized in region B, mainly owing to a higher electron temperature, so their behaviour is similar to that of metallic impurities in our experiment. Therefore, the above mentioned impurity screening by EML may be expected.
ACKNOWLEDGEMENTS

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PARTIAL AND FULL RECONNECTION DURING SAWTOOTH ACTIVITY AND DISRUPTIONS

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Abstract

PARTIAL AND FULL RECONNECTION DURING SAWTOOTH ACTIVITY AND DISRUPTIONS.

Reconnection phenomena in the form of sawteeth and major and minor disruptions have been studied on the TFTR and Tokapole II tokamaks. On TFTR most sawteeth are caused by a reconnection process, although the timescale is often much faster than the Sweet–Parker reconnection time, $(r_T\gamma_T)^{1/2}$. However, on Tokapole II the central safety factor, $q(0)$, has previously been observed to remain fixed at 0.7 during the sawtooth oscillations. This result implies the absence of total reconnection to the plasma center during sawteeth on the Tokapole II tokamak. It is found that the sawtooth crash enhances the magnetic stochasticity outside the reconnection volume, complicating the analysis of heat pulse propagation. However, when the effects of the enhancement in transport are properly accounted for, it is found that the heat pulses propagate consistently with the power balance thermal diffusivity. Major disruptions at high $\beta$ are found to have some similarity to sawteeth in that they typically have $(m, n) = (1, 1)$ precursors and stochasticity may also be involved in the final disruption phase. In the Tokapole II tokamak, direct measurements of the $q$ profile have been made through repetitive minor disruptions. It is found that $q(0)$ increases from 0.75 to 1.3 during the minor disruption. Precursor oscillations are consistent with the simultaneous presence of the $(m, n) = (2, 1)$ and $(3, 2)$ modes.
TFTR RESULTS

I Sawteeth

Temperature profile reconstruction\(^1\) from local Electron Cyclotron Emission (ECE) electron temperature measurements and from soft X-ray (SXR) signals has been used to obtain information about the internal phenomenology of the sawtooth crash phase. These methods rely on the assumption of rigid rotation for reconstruction and resolve changes of the profile that are slow compared to the rotation rate of the plasma, i.e., the changes must be small during each period of the precursor oscillation \(T=2\pi R/v_\phi\) which is about 30 \(\mu\)sec with the large toroidal rotation velocities (up to \(1.4 \times 10^6\) m/sec) obtained with neutral beam injection on TFTR.

The sawtooth process is a relaxation oscillation that is caused by the growth of an \((m,n) = (1,1)\) internal mode, culminating in a reconnection of internal magnetic flux.\(^2\) The precursor to the sawtooth crash deforms the magnetic flux surfaces, forming a cold island region and a 'hot spot' (the hot, helically deformed core of the plasma). The shape of the hot spot, the amount of helical distortion and the shape of the cold island theoretically depend sensitively on details of the central pressure and \(q\) profiles. Indeed, a spectrum of sawtooth precursor activity is experimentally observed on TFTR. We limit our discussion to a example of the typical sawtooth behavior observed in TFTR and two unusual examples of sawteeth which illustrate some of the variety of observed behavior.

The typical sawteeth often have little or no precursor activity before the sawtooth crash; it is found that a mode grows rapidly during the final collapse itself. Analysis of the typical sawtooth crashes on TFTR leads to the following picture: (a) the magnetic axis or hot temperature core deforms helically and moves away from the original center as the temperature island-structure
FIG. 1. Reconstruction of the ECE to form a 2-D image of the sawtooth reconnection. The hot spot is circular prior to the final reconnection. The solid circles denote the reconnection radius.

grows; (b) as the hot core deforms helically, it remains nearly circular in cross-section; and (c) in the final stage the heat spreads poloidally around the q=1 surface and the cold region expands to form the new plasma core (see Fig. 1). The flat region of the temperature profile is interpreted as a magnetic island. Assuming monotonically increasing q(r) profiles, this sequence is what would be expected from a reconnection-type model, except for the timescale involved. From the Kadomtsev sawtooth model the reconnection time is given by $(\tau_R \tau_A)^{1/2}$, where $\tau_R$ is the resistive skin time and $\tau_A$ is the Alfvén time. For TFTR parameters, this is $\approx 5$ to 10 msec; in the experiment the crash time for the central temperature is $\approx 0.02$ to 2 msec, (however the final growth of the island takes 2 - 5 times longer). However, $\mu$-wave scattering measurements made around the q=1 surface
FIG. 2. Reconstruction of an unusual sawtooth precursor. The cold bubble is roughly circular and the hot spot is crescent shaped. The final reconnection, however, occurs with a crescent shaped cold island and circular hot spot.

during the sawtooth crash often show bursts of fluctuations when the x-point of the island passes through the scattering volume. If this turbulence were to broaden the \((m,n) = (1,1)\) tearing layer, the normal Kadomtsev sawtooth crash rate could be enhanced.

Occasionally other types of sawtooth precursor behavior are observed. An example of such behavior is shown in Figure 2. In this example, during the initial phase of the sawtooth, a "cold bubble" is seen to move inward, deforming the "hot spot" into a crescent\(^3\). The "cold bubble" then causes an off-axis reconnection, i.e., it does not affect the central temperature. Following this, the cold bubble evolves to a typical crescent-shaped cold island and the hot core moves out in minor radius, remaining circular in shape. The reconnection process normally seen in TFTR then takes place. This is the only example of a "cold bubble" observed in
TFTR. In this example, the central temperature profile was nearly flat even during the precursor phase.

In a second example of unusual sawtooth behavior, we see an incomplete sawtooth where the reconnection does not penetrate to the center of the plasma. However, unlike in the partial or subordinate reconnection, the central electron temperature does show a marked decrease. This type of sawtooth is differentiated from the typical examples in that the electron temperature remains somewhat peaked following the sawtooth crash.

II Heat Pulse Propagation

The sawtooth crash also affects the plasma outside the reconnection radius, both in the precursor and in the crash phases. Simulations with the MH3D code\textsuperscript{4} predict that finite pressure and toroidal curvature cause a toroidal bulge of the helical hot spot which extends beyond the $q=1$ surface, i.e., the amplitude of the eigenfunction is poloidally asymmetric with very low or negative amplitude on the high field side of the midplane and strongest on the outboard midplane.\textsuperscript{5} This helically asymmetric perturbation can produce stochasticity beyond the $q=1$ surface. Measurements of the ECE show both quantitative and qualitative agreement with the predictions. While this modification has little effect on the physics of the sawtooth crash, it can strongly affect the propagation of the heat pulse by creating stochastic regions outside the reconnection radius. The resulting large enhancement in the electron thermal diffusivity (to levels in excess of 150 $m^2$/sec) causes rapid transport of heat in this region for 100 - 200 $\mu$sec after the sawtooth crash (Fig. 3a). While this extended redistribution of heat has long been postulated as an explanation for the anomalous heat pulse propagation,\textsuperscript{6,7,8,9} we present direct experimental evidence and numerical simulations to support this hypothesis.

To model this effect we use a temporally and spatially varying diffusion coefficient. This extended region in which
FIG. 3. (a) Profiles of the electron temperature through the sawtooth crash (10 μsec) showing the extended deposition of heat. (b) Comparison of experimental with simulated heat pulses; the simulations use the χ from power balance and an extended perturbation.
electron heat is redistributed during the sawtooth crash substantially complicates the determination of heat transport properties from the subsequent heat pulse propagation. The commonly used methods of heat pulse analysis (time-to-peak, extended time-to-peak and Fourier methods) yield results inconsistent with the analysis of the power balance in many cases. However, by numerically simulating the evolution of the extended perturbation, it is found that the relaxation of the temperature is consistent with the power balance estimates of the local thermal diffusivity (see Fig. 3b). In cases where the extended perturbation is small, i.e., more like the classical picture of sawtooth reconnection, the heat pulse and the power balance diffusivities roughly agree.

Coupling of thermal and particle diffusivities is not likely to play a large role in the propagation of heat pulses on TFTR. Similar to the results from JET, the normalized density pulse amplitudes are roughly one third the amplitude of the heat pulse and 5 - 10 times slower, e.g., the time-to-peak at a minor radius of 0.8 m is typically about 30 msec for the density and 5 msec for the temperature pulses (see Fig. 4).

III Disruptions

Studies on TFTR in hot, high-$\beta_{pol}$ plasmas show that disruptions bear a close resemblance to the typical sawtooth crashes. They have rapidly growing $(m,n) = (1,1)$ precursors and a thermal transport phase lasting 40 $\mu$sec to 1 msec. This is in contrast to cold, high-density ohmic disruptions which tend to have the more traditional $(m,n) = (2,1)$ precursor activity. The disruptions differ from the sawteeth in that they typically cause the magnetic fields to become stochastic across the entire plasma volume. Reconstruction of the disruption precursor indicates that it is predominantly $(m,n) = (1,1)$, but in contrast to the sawtooth precursor, is kink-like rather than island-like. This $(m,n) = (1,1)$
mode typically grows very quickly, often in much less than 1 msec. Another difference between the sawteeth and the disruption is that in the final stages of the growth of the (1,1) precursor, there is a large, non-thermal burst of ECE lasting 20 - 50 \( \mu \text{sec} \), occurring 50 - 100 \( \mu \text{sec} \) before the thermal quench. The rapid loss of heat (in <50 \( \mu \text{sec} \)) during the thermal quench suggests that the magnetic field is stochastic or that a reconnection has occurred across most of the plasma volume. Following the transport phase, the positive spike in the plasma current occurs 1 - 3 msec later. Rapid temperature excursions on the limiter surface are sometimes recorded both at the time of the transport phase and at the start of the current decay. An example of a high-\( \beta \text{pol} \) disruption in TFTR is shown in figure 5.
Tokapole II Results

Both sawtooth oscillations and major disruptions have long been interpreted as arising from magnetic reconnection. However, it had been observed in the Tokapole II tokamak that the central safety factor, $q(0)$, remains fixed at 0.7 during sawtooth oscillations\textsuperscript{11}. This result implied the absence of total reconnection\textsuperscript{2} to the plasma center. Similar results were subsequently reported from other tokamaks,\textsuperscript{12-14} while in yet
q(0) was about unity during a sawtooth with insufficient resolution to discern the time variation.

We report here measurement of the safety factor during a major disruption in Tokapole II. The measurements are relatively direct in that q is obtained from a two-dimensional map of the magnetic field obtained with an internal probe. We observe that in the center of the plasma, q increases rapidly during a disruption, from less than unity to above unity, while the total plasma current remains constant. This observation is consistent with the occurrence of total magnetic reconnection.

A similar observation of a rapid increase in q(0) during a disruption has been reported in the LT-3 tokamak. The present work differs in that q is determined from two-dimensional measurements, the post-disruption state is diagnosed since the plasma survives the disruption, and the results may be directly contrasted with the absence of total reconnection observed during sawtooth oscillations in the same device.

Tokapole II is a tokamak in a poloidal divertor configuration. With scrape-off plates inserted to the separatrix, major disruptions cause sudden termination of the plasma as in most tokamaks. However, when operated without scrape-off plates major disruptions do not cause termination of the discharge. The thermal loss and the inward major radius shift occur, but without the material limiter contact. This results in major disruptions appearing as repetitive giant sawtooth oscillations.

The safety factor is obtained from a measurement of the magnetic field in a poloidal plane. A single probe, consisting of two orthogonal coils, measures the poloidal field in reproducible discharges on a grid 18 cm by 18 cm with a step size of 2 cm. From the measurement, we obtain the poloidal magnetic flux function, from which we can plot the magnetic surfaces. The surfaces include the divertor x-points and display the inward shift of a major disruption. The safety factor corresponding to a
magnetic surface is obtained by integrating the local pitch over the magnetic surface obtained experimentally. These two-dimensional measurements are necessary since the surface noncircularity renders inferences of \( q \) from single chord measurements unreliable.

The safety factor profile before and after a major disruption is plotted in figure 6. The value of \( q(0) \) before a major disruption is 0.75 and rises to 1.3 after a disruption. Since the total plasma current is unvarying, the change in \( q \) is due to a rearrangement of the current density profile. After a disruption, the \( q \) profile is flat over the inner 5 cm of the plasma (the separatrix is at about 9 cm) which is approximately the original location of the \( q=1.3 \) surface. The toroidal current density is also substantially flattened after a disruption (Fig. 7). An integration of the two-dimensional toroidal current density profile indicates that the total current does not change and is equal to the value obtained from an independent measurement of the total plasma current.

In contrast to sawtooth oscillations, these results indicate total reconnection during a major disruption. However, the details of the reconnection are less clear. Precursor oscillations
are consistent with the simultaneous presence of the \((m/n) = (2,1)\) and \((3,2)\) modes.\(^{19}\) The rise in \(q(0)\) to only 1.3 indicates that coupling to other helical modes is likely to be taking place.

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EQUILIBRIUM AND STABILITY MEASUREMENTS IN A HIGH BETA TOKAMAK*

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Abstract

EQUILIBRIUM AND STABILITY MEASUREMENTS IN A HIGH BETA TOKAMAK.
The Columbia High Beta Tokamak (HBT) utilizes fast formation techniques to produce plasmas capable of testing theoretically predicted equilibrium and stability boundaries. At very high values of poloidal beta ($\epsilon \beta_p > 1.7$), internal magnetic diagnostics show the encroachment of an inboard separatrix followed by rapid loss of confinement, suggesting the existence of an equilibrium limit. Plasmas formed at lower $\beta_p$ are observed to become unstable at values of $q^*$ and $\beta$ in accordance with the Kruskal-Shafranov and Troyon relations. Internal and external magnetic field measurements identify the modes as long wavelength external kinks that grow on an Alfvén time scale. Ideal MHD numerical simulations reproduce qualitative features of the mode structure, but predict slower growth rates for $q^* > 1$ plasmas. A description of HBT-EP is given. This tokamak is designed to test the effects of an adjustable conducting shell on the external mode stability of a longer-pulse length, higher-temperature plasma.

1. Introduction

The HBT[1] experiment at Columbia University is a small tokamak ($R_t = 24$ cm, $a = 7$ cm) that uses fast formation methods to produce plasmas with parameters that are inaccessible, or difficult to achieve, in more conventional machines. The parallel-bias start-up scheme utilizes a rapid toroidal field increase ($\approx 0.4$ kG to 4 kG in 2 $\mu$s) with intense ohmic heating to create plasmas that become unstable at values of $\beta$ ($\approx 2\mu_0 \langle p \rangle / B_0^2$) and $q^*$ ($\approx 2\pi a^2 B_\phi / \mu_0 R_0 I_p$) consistent with the predictions of ideal MHD theory[2,3]. A variation of this formation technique maintains a low value of $B_\phi$ ($\approx 0.4$ kG) with an aggressive $I_p$ ramp (1 kA/$\mu$s) and results in plasmas that violate the Kruskal-Shafranov stability boundary[4,5]. The instabilities observed in each case are identified as long wavelength external kink modes that grow on an Alfvén time scale. The highest $\beta$-values ($\approx 30\%$) are achieved in HBT by rapid $B_\phi$

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reversal (reverse-bias formation method), producing short-lived (<20 μs) plasmas that appear to be bounded by an $\epsilon\beta_p \leq 1.7$ equilibrium limit.

The plasma environment in HBT ($n_e = n_i = 10^{14}$ cm$^{-3}$, $\langle T_e \rangle = \langle T_i \rangle \sim 10$ eV, $\tau_{\text{pulse}} \leq 150$ μs) permits the use of internal probes to measure the evolution of the magnetic structure throughout the discharge, thus allowing a detailed comparison of plasma behavior at equilibrium and stability boundaries with ideal MHD theory.

2. Measurement of the equilibrium $\epsilon\beta_p$ limit

As has been long recognized, a tokamak with a given vertical field curvature has an equilibrium poloidal beta limit beyond which flux surfaces open and confinement is lost[6,7]. Using HBT’s “rapid $B_\phi$ reversal” formation technique, we have directly observed this poloidal beta limit by measuring the internal magnetic field of very high, $\langle \beta \rangle \sim 30\%$, short-lived tokamak discharges. This start-up procedure is non-flux-conserving, and the plasma pressure (determined by the magnitude of theta-pinch currents) can be generated independently from the plasma current (determined by the ohmic transformer). When a discharge is created below the equilibrium limit, the internal magnetic probe clearly shows quasi-stationary closed flux surfaces with a large, outward Shafranov shift. When the plasma pressure is increased beyond the equilibrium limit, either the plasma strikes the outside wall of the vacuum vessel or, by increasing the vertical field strength, a separatrix enters the inside of the plasma destroying confinement. Fig. 1 characterizes the parameters of high beta discharges having $\beta_p$ above and below the equilibrium limit. Fig. 2 shows the measured internal poloidal field measured for a discharge representative of HBT’s $\epsilon\beta_p \leq 1.7$ limit together with the least-squares best-fit reconstruction[8] of the free-boundary equilibrium.

3. Measurement of the internal structure of fast-growing external kink instabilities

The parallel-bias formation procedure produces the longest-lived ($\leq 150$ μs) discharges in HBT. Many shots exhibit a pronounced transition from a quiescent to an unstable state as the shot evolves. A more rapid onset of unstable activity is observed when the plasma is formed at low values of $B_\phi$ such that $q^* \leq 1$ and the Kruskal-Shafranov condition is violated. Figure 3 is an experimental stability diagram in the $\epsilon\beta_p$ vs $q^*$ plane for plasmas produced by these methods. The trajectories of four example shots, A, B, C, and D, are indicated.

The internal magnetic structure of the plasma is determined using two quartz-encapsulated coil arrays displaced toroidally by $\Delta \phi = 180^\circ$. Each probe measures $B_\parallel (t)$ and $B_z (t)$ at 6-12 positions along the plasma minor radius at the midplane, $Z = 0$. The perturbed magnetic field, $\delta B_z$, is defined as the deviation of a single probe signal from the two-probe average, that is,
FIG. 1. Operational diagram for the very high beta HBT discharges characterized by the internal structure of their poloidal field profiles. Discharges with low poloidal beta have closed flux surfaces, but those at high poloidal beta have separatrices or open flux surfaces. ($\beta_{p, \text{eff}}$ assumes that $l_i = 0.5$ for all cases.)

\[
\delta B_z(R, Z, 0, r) \equiv \frac{1}{2} [B_z(\phi = 0^\circ) - B_z(\phi = 180^\circ)] - \langle \delta B_z \rangle_{10\mu s}
\]

where $\langle \delta B_z \rangle_{10\mu s}$ represents a 10 $\mu$s time-average of the quantity prior to the observed onset of the instability. The subtraction is necessary to account for any electrical or positional offsets of the probe, and minimizes any drifts caused by slowly growing modes in the plasma. Note that this definition of $\delta B_z$ eliminates any contribution due to axisymmetric changes in the plasma equilibrium position. A toroidally distributed array of magnetic pick-up coils mounted on top of the quartz vacuum chamber at the major radius indicates the predominance of the $n = 1$ mode in the cases.
studied, where $n$ is the toroidal wave number of the instability. Other effects accompanying the magnetic perturbations include large fluctuations in the line-averaged electron density and the loop voltage, as well as increased levels of H$_\alpha$ emission.

A free-boundary numerical equilibrium best representing internal poloidal field and external coil current measurements is constructed in order to facilitate direct comparison between the measured instability
FIG. 3. HBT stability diagram displaying the evolution of example shots A, B, C, and D. Shots progress from high to low $q^*$ (right to left). Troyon scaling is indicated with $\beta_N = \beta(\%) = \beta([I_e(MA)/a(m)B_T(T)]) = 3.5$.

characteristics and the theoretical predictions of ideal MHD. The numerical equilibrium is then resolved on a flux-coordinate system (EQGRUM[9]) and subsequently analyzed by the PEST-I ideal MHD stability code[10].

Figures 4(a) and (b) illustrate the evolution of the perturbed internal magnetic field of example shot A from Fig. 3. The instability becomes evident after 179 $\mu$s, at which time $q^* < 2$. The observed growth time of the mode, $\tau_g (\text{obs})$, is 2–3 $\mu$s. The Alfvén transit time, $\tau_k = \sqrt{\mu_0 \rho q^* R/B_z}$ ($\rho$ is the mass density), for this case is 2 $\mu$s. Note that this is a large scale instability ($\delta B_z(a)/B_z(a) \geq 5\%$) which appears to saturate in 15 $\mu$s. During the growth period, the plasma column contracts and eventually terminates on the inside wall of the vacuum chamber. Figure 4(c) displays the evolution of the PEST-I computed growth times. The plasma is predicted to be unstable to the $n = 1$ mode at all times during the discharge, but with a steadily decreasing growth time. At the observed onset of the instability, PEST-I predicts a growth time, $\tau_g (\text{PEST})$, of about 15 $\mu$s. In Fig. 4(d), the measured and computed perturbed magnetic field profiles are compared. Toroidal field scaling studies indicate that the inboard null in the profile is associated with the $m=3$ poloidal component of the unstable eigenfunction as determined by PEST-I, whereas the central null is related to the $m=2$ component.
FIG. 4. Example shot A data (unstable at $q^* < 2$): (a) Evolution of $\delta B_z$, measured at six positions across the plasma radius at $Z = 0$. (b) Data from (a) superimposed to show growth time evaluation, determined from exponential fits to indicated channels weighted by measurement errors at each sampling point. (c) Evolution of growth times as determined by PEST-I. Error bars derive from uncertainty in $n$. (d) Measured and numerically calculated perturbed internal magnetic field profiles.
FIG. 5. (a) Time history of the perturbed internal magnetic field profile for $q^* < 1$ shot B. (b) PEST-1 unstable profiles for shot B showing $n = 1, 2, 3,$ and 4 modes. Comparison of measured and theoretically predicted unstable profiles for (c) example shot C and (d) example shot D.
Figure 5 displays the experimentally measured and numerically computed $\delta B_z$ profiles for example shots B, C, and D from Fig. 3, each observed to become unstable at a different value of $q^*$. The $q^* < 1$ discharge of Fig. 5(a) is predicted to be unstable to the $n = 1, n = 2, n = 3,$ and $n = 4$ modes (no higher-$n$ cases were examined) with similar growth times ($\tau_g(PEST) = 1.2 \mu s$). The measured growth time in this case was $\tau_g(\text{obs}) \approx 1-2.5 \mu s$, in closer agreement to PEST-I growth times than observed in higher-$q^*$ discharges. The measured $\delta B_z$ profile is positive across most of the plasma radius, indicating the dominance of the $n = 1$ mode. Recall that the method for measuring $\delta B_z$ excludes contributions from even-$n$ modes. A $q^* < 2$ shot, similar to example shot A, is displayed in Fig. 5(c) using a higher-resolution array of internal probes. Note that the inboard and central null points in the profile are again evident. The characteristic times for this shot are $\tau_g(\text{obs}) = 2 \mu s$, $\tau_A = 2 \mu s$, and $\tau_g(PEST) = 12 \mu s$. Figure 5(d) illustrates the perturbed field of $q^* < 5$ example shot D. Although a multi-null structure is suggested, the mode saturates at a low level and is difficult to resolve within the accuracy of the measurements. For this case, $\tau_g(\text{obs}) = 3-4.5 \mu s$, $\tau_A = 2 \mu s$, and $\tau(PEST) = 65 \mu s$.


The successor to HBT is presently under construction and will begin operation in 1991. This new device, called HBT-EP (extended pulse), with $R_0 = 0.9$ m, $a = 0.15$ m, and $B_\phi = 0.5$ to 1 T, uses the toroidal coil set formerly part of the CLEO tokamak/stellarator at Culham to produce a much longer pulse length than HBT of up to 10 msec. Two unique design features of HBT-EP are (1) an adjustable internal conducting wall which can be varied in position relative to the plasma edge and (2) the capability to produce a variety of plasma cross-sections including diverted, $\epsilon \beta_p < 1$, discharges with the poloidal null on the mid-plane near the inside wall. The principal goals of this new, larger device are (1) to extend high density ohmic plasmas to higher temperatures ($T_e \sim 100$ eV), (2) to study conducting wall stabilization of external kink modes, and (3) to study $q$-profile and shaping effects on high beta stability.

5. Summary

The internal magnetic structure of plasmas at theoretically predicted equilibrium and stability boundaries has been measured in HBT. Plasmas with $\epsilon \beta_p$ exceeding 1.7 show a separatrix at the inner surface plasma, indicating an equilibrium limit. Magnetic measurements of instabilities in plasmas formed near the Troyon and Kruskal-Shafranov boundaries show qualitative agreement with $\delta B_z$ profiles determined using the PEST-I ideal MHD stability code. These external modes grow on an Alfvén time
scale, which, for $q^* > 1$ plasmas, is faster than the numerically computed growth times. A new tokamak experiment, HBT-EP, is now under construction and will test the effect of an adjustable conducting wall on the stability of longer-lived, higher-temperature plasmas.

REFERENCES

NEGATIVE CHARGE INJECTION BY NEUTRAL BEAMS TO MODIFY EDGE CONFINEMENT

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Abstract
NEGATIVE CHARGE INJECTION BY NEUTRAL BEAMS TO MODIFY EDGE CONFINEMENT.

The injection of heavy ion neutral beams, with subsequent prompt ion losses, into the edge region of a tokamak plasma is proposed. The radial electric field resulting from this injection of negative charge can produce large poloidal rotation. Estimates indicate that the poloidal rotation may be much larger than the usual H mode regime values. Thus negative charge injection may provide a new confinement regime beyond the H mode.

1. INTRODUCTION

Spectroscopic measurements in the DIII-D tokamak indicate that the radial electric field increases dramatically at the H mode transition [1]. Recent experiments by Taylor et al. [2] suggest that a transition from L to H mode confinement in tokamaks can be achieved by an electrode in the edge region of the plasma which sets up a negative radial electric field and poloidal plasma rotation. Following these experiments, it was suggested [3] that an externally controlled negative radial electric field may be produced in hotter plasmas — where use of an electrode is not feasible — by the injection of heavy ion neutral beams with subsequent prompt ion orbit losses. In this paper it is proposed that heavy ion neutral beams may produce poloidal rotations which are much larger than those encountered in the usual H mode regime and may thereby provide a new confinement regime.

Shaing and Crume [4] have proposed that a radial electric field plays a role in all L to H mode transitions through a bifurcation in poloidal momentum balance, a balance between poloidal viscous drag and a torque arising from ion orbit losses near the edge of the plasma. In that theory the poloidal viscous drag reaches a maximum for poloidal rotations, \(X = v_p B_t / \nu_{th} B_p \sim 1\), and is then monotonically decreasing for larger poloidal rotations (\(v_p\) is the poloidal velocity, \(\nu_{th}\) is the thermal velocity, \(B_t\) is the toroidal magnetic field, and \(B_p\) is the poloidal magnetic field). The plasma ion orbit losses also decrease as the poloidal rotation increases and the H mode arises from a balance of these two processes at a poloidal rotation of \(X > 1\).

The torque due to ion orbit losses can be brought under external control by the use of heavy ion neutral beams. The beam ions play the role of the electrode in
Taylor's experiment; they carry in electrons and are then lost to the wall in a small localized region. This localized current of beam ions is analogous to the current through an electrode. Beam requirements are discussed in Section 2.

Including beam orbit losses fundamentally alters the poloidal momentum balance. Unlike the viscous drag and plasma ion orbit losses, the beam orbit losses do not decrease with increasing poloidal rotation speed. Thus to achieve momentum balance in steady state additional momentum loss terms must be identified. We find that neutral drag on the rotating ions due to charge exchange can provide momentum balance in the presence of beam orbit losses at poloidal rotations considerably larger than H mode values. This momentum balance is described in Section 3.

A brief summary appears in Section 4.

2. BEAM REQUIREMENTS

In order to maximize prompt ion losses the beam ions should be born within a gyroradius of the edge of the plasma, i.e. \( \Delta r_{\text{born}} \leq \Delta \), where \( \Delta r_{\text{born}} \) is the distance from the edge of the plasma at which the fast ion is born and \( \Delta \) is the particle gyroradius. The quantity \( \Delta r_{\text{born}} \) can be expressed in terms of beam ionization processes and the above inequality becomes

\[
\mu \left( \frac{n_e}{10^{19} \text{ m}^{-3}} \right) \left( \frac{\langle \sigma v \rangle_{\text{tot}}}{10^{-13} \text{ m}^3 \text{ s}^{-1}} \right) \left( \frac{B_i}{1 \text{ T}} \right)^{-1} \geq 96
\]

where \( \mu \) is the mass ratio \( m_{\text{beam ion}}/m_{\text{proton}} \), \( n_e \) is the electron density, \( \langle \sigma v \rangle_{\text{tot}} \) is the total reaction rate, and \( B_i \) is the toroidal magnetic field. The above expression is accurate for any injection angle into a cylindrical plasma [3]. However, for injection into a toroidal plasma, there is some variation with injection angle and beam width. Even though the above criterion is satisfied, some beam ion trapping may occur.

If the width of the edge region in which the radial current flows is required to be \( N \) plasma ion banana widths, \( \langle \Delta r_{\text{born}} \rangle \geq N\Delta r_{\theta,i} \), this establishes a minimum beam energy. In the circular cross-section approximation, this leads to

\[
E_i > \frac{112}{\mu} \left( \frac{N}{5} \right)^2 q^2 T_i
\]

where \( q \) is the safety factor and \( T_i \) is the edge ion temperature.

As a specific example, with an argon beam and \( T_e = 100 \text{ eV} \), Eq. (1) becomes

\[
\left( \frac{n_e}{10^{19} \text{ m}^{-3}} \right) \left( \frac{B_i}{1 \text{ T}} \right)^{-1} > 1.5; \text{ further choosing } N = 5 \text{ and } q = 3, \text{ Eq. (2) becomes } E_i > 25 T_i.
\]

Larger beam energies are desirable to achieve higher beam currents.
An estimate of the radial current required to achieve a normalized poloidal velocity of $X \sim 1$ and thereby overcome the peak in viscous drag can be obtained from

$$\bar{J}_r \times \bar{B}_i \approx m_i n_i \nu_{ii} \bar{v}_p$$

where $J_r$ is the radial current density, $m_i$ is the plasma ion mass, $n_i$ is the plasma ion density, and $\nu_{ii}$ is the ion-ion collision time. For deuterium in a DIII-D size machine ($R = 1.6 \text{ m}$, $a = 0.6 \text{ m}$) this becomes

$$I_r = 62 \left( \frac{n_i}{10^{15} \text{ m}^{-3}} \right)^2 \left( \frac{T_i}{100 \text{ eV}} \right)^{-1} \text{ A}$$

This current may be provided by a combination of plasma ion orbit losses and externally controlled beam ion orbit losses.

3. POLOIDAL MOMENTUM BALANCE: BEYOND THE H MODE

When edge localized neutral beams are used to produce or enhance the radial electric field and poloidal rotation, poloidal momentum balance is qualitatively altered because these beam ion losses are gyro-orbit losses and they do not decrease as the poloidal rotation increases. Since the viscous drag effects considered by Shaing and Crume decrease as $X$ increases, additional drag forces must become dominant to provide poloidal momentum balance above $X \sim 1$.

Various possibilities exist for the required source of additional drag to reach equilibrium. These include classical viscosity, electron parallel viscosity, and charge exchange with neutrals. Of these, the drag due to charge exchange with neutrals seems likely to be the largest. The poloidal viscous drag force due to parallel ion viscosity considered by Shaing and Crume is of order $m_i n_i \nu_{pp} \bar{v}_p$. By contrast the classical or perpendicular viscosity is of order $m_i n_i \nu_{ee} \left( \frac{\rho}{L} \right)^2$, where $\rho$ is the ion gyroradius and $L$ is the radial scale length of the poloidal velocity shear $(\frac{d\bar{v}_p}{dr}) v_p^{-1}$. Even if $L$ is as small as the poloidal gyroradius, the factor $(\rho/L)^2$ is small, $\sim 10^{-2}$. Similarly, in the low collisionality regime being considered, the electron parallel viscosity is $\sqrt{(m_e/m_i)}$, smaller than the parallel ion viscosity and therefore also quite small.

The drag due to charge exchange, on the other hand, can be quite large, depending upon the neutral density near the edge of the plasma. The poloidal drag is $m_i n_i \nu_{0} \bar{v}_p$, where $\nu_0$ is the effective collision frequency between rotating ions and non-rotating neutrals and is given by

$$\nu_0 = n_0 (\sigma v)_{\text{ex}}$$

(4)
where \( n_0 \) is the neutral density and \( \langle \sigma v \rangle_{\text{ex}} \) is the reaction rate for charge exchange. From Freeman and Jones [5] we find that in the range \( 100 \text{ eV} \leq T_i \leq 1000 \text{ eV} \) the reaction rate for deuterium is well represented by \( \langle \sigma v \rangle_{\text{ex}} = 3.1 \times 10^{-14} \times (T_i/100 \text{ eV})^{0.3} \text{ m}^3/\text{s} \) or

\[
\nu_0 = n_0 \langle \sigma v \rangle_{\text{ex}} = 310 \left( \frac{n_0}{10^{16} \text{ m}^{-3}} \right) \left( \frac{T_i}{100 \text{ eV}} \right)^{0.3} \text{ s}^{-1} \tag{5}
\]

Comparing this with Braginskii's formula [6] for \( \nu_{ii} \) with the Coulomb logarithm \( \lambda = 15 \), we find

\[
\frac{\nu_0}{\nu_{ii}} = 62 \left( \frac{n_0}{n_i} \right) \left( \frac{T_i}{100 \text{ eV}} \right)^{1.8}
\tag{6}
\]

indicating that neutral drag can be significant, particularly at large values of poloidal rotation where the viscous drag is greatly diminished.

From the above considerations the poloidal momentum balance equation, including beam current torque and neutral drag, is given by

\[
-e \Gamma_{\text{orbit}} \times \vec{B} \cdot \vec{B}_p + \frac{I_{\text{beam}} \times \vec{B} \cdot \vec{B}_p}{\text{area}} = (\vec{B}_p \cdot \nabla \cdot \pi) + m_i n_i \langle \sigma v \rangle_{\text{ex}} \vec{v} \cdot \vec{B}_p \tag{7}
\]

Expressions for the plasma ion orbit loss, \( e\Gamma \times \vec{B} \cdot \vec{B}_p \), and the viscous drag, \( (\vec{B}_p \cdot \nabla \cdot \pi) \), are given in the paper of Shaing and Crume [4]. A comment is necessary concerning the choice of a parameter which appears in the Shaing and Crume expressions, \( \rho_{pi}/(\Delta r G) \), where \( \rho_{pi} \) is the poloidal ion gyroradius, \( \Delta r \) is the radial width over which plasma ion orbit losses occur and \( G \) is a geometric factor appearing in Shaing and Crume's equation for the ion orbit loss rate. In their paper they have chosen \( \rho_{pi}/(\Delta r G) = 12.5 \) for their illustrative example and we use the same value below.

We present a numerical solution of Eq. (7) for DIII-D type parameters by plotting the left hand side and the right hand side on the same graph as a function of \( X \). The parameters chosen were: minor radius \( a = 0.6 \text{ m} \), major radius \( R = 1.6 \text{ m} \), toroidal field \( B_t = 1 \text{ T} \), poloidal field \( B_p = 0.1 \text{ T} \), ion density \( n_i = 10^{19} \text{ m}^{-3} \), neutral density \( n_0 = 10^{16} \text{ m}^{-3} \), ion temperature \( T_i = 100 \text{ eV} \), mass ratio = 2 (deuterium), beam current \( I_{\text{beam}} = 50 \text{ A} \).

In Fig. 1, two equilibrium cases are shown. The broken lines represent the case where the beam torque and neutral drag are neglected, resulting in the usual Shaing and Crume picture with an H mode at \( X = 2.7 \). The solid lines represent the case which includes the beam torque and the neutral drag. Momentum balance for this case occurs for \( X = 11.2 \), well beyond the H mode value.
Since the plasma ion orbit losses and the viscous drag are negligible at such large values of $X$, a good estimate of $X$ can be obtained from balancing the beam torque and the neutral drag. We have

$$\frac{I_{\text{beam}} \times \mathbf{B} \cdot \mathbf{B}_p}{\text{area}} = m_i n_i \langle \sigma v \rangle_{\text{ex}} \mathbf{v} \cdot \mathbf{B}_p$$

and using the DIII-D parameters above this becomes

$$X = 13 \left( \frac{I}{50 \text{ A}} \right) \left( \frac{n_i}{10^{19} \text{ m}^{-3}} \right)^{-1} \left( \frac{n_0}{10^{16} \text{ m}^{-3}} \right)^{-1} \left( \frac{B}{1 \text{ T}} \right) \left( \frac{T_i}{100 \text{ eV}} \right)^{0.8}$$

in good agreement with the numerical solution.

4. SUMMARY

We have shown that negative charge may be injected in the edge region of a tokamak plasma using heavy ion beams with prompt gyro-orbit losses. The resulting poloidal momentum balance between the radial current torque and the neutral drag due to charge exchange can occur at poloidal rotations much larger than typical H mode values. A beam current of 50 A in a DIII-D size device may produce a new confinement regime.

One potential difficulty with such an experiment is impurity accumulation in the edge region. Steps may be required to prevent the recycling of beam ions and the wall sputtering of other impurities by the beam ions. The choice of heavy ions may be
important. Noble gas ions do not adhere to the wall and large recycling results. On the other hand, the use of metal ions, such as titanium, will require ion source development. A preliminary experiment on DIII-D using argon was dominated by impurity radiation [7]. Beam ion trapping, as discussed in Section 2, may have been responsible.

In addition to possible confinement modifications, heavy ion neutral beam experiments with prompt ion losses can provide a simulation of the effects of alpha particle orbit losses in ignited or burning plasmas.

REFERENCES

INTERNAL MICROTURBULENCE STUDIES ON DIII-D, TEXT AND TFTR*

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Abstract

INTERNAL MICROTURBULENCE STUDIES ON DIII-D, TEXT AND TFTR.

The paper describes internal microturbulence measurements on the DIII-D, TEXT, and TFTR tokamaks utilizing heterodyne collective scattering, ECE, and beam emission spectroscopy (BES) techniques. On the TEXT and DIII-D tokamaks, emphasis is placed on results obtained during saturated Ohmic operation, where ion temperature gradient driven (ITGD) turbulence is expected to exist. Collective scattering data on TEXT strongly support the existence of internal density fluctuations propagating in the ion diamagnetic drift direction, whereas on DIII-D, the existing data base is consistent with turbulence propagating in the electron diamagnetic drift direction modified by electric field/plasma rotation effects. In the TFTR tokamak, scattering measurements in Ohmic, L-mode, and supershot regimes indicate that fluctuation amplitudes increase with decreasing k to $2 \text{ cm}^{-1}$. Estimates for $k > 2 \text{ cm}^{-1}$ give $\delta n_e/n_e \sim 0.5\%$ for L-mode and supershot discharges. Estimates of $\delta T_e/T_e$ using ECE have set upper limits on interior fluctuation levels of 0.4%. The BES system observes mode activity with $k < 2 \text{ cm}^{-1}$ in the interior of supershots with $\delta n_e/n_e \sim 1$–1.5% and less than 0.5% in L-mode discharges.

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1. INTRODUCTION

The empirical discovery of enhanced confinement schemes, such as H-mode and supershots, together with the construction of larger devices, has allowed the world magnetic fusion program to advance technically almost to the point of scientific breakeven. However, a predictive capability is highly desirable for the next phase of fusion development. The prime motivation of the work presented here is to improve our physics understanding of the role played by electrostatic turbulence in various confinement regimes. State-of-the-art, fluctuation diagnostics have therefore been integrated into the TEXT [1], TFTR [2,3] and DIII-D [1] tokamaks to characterize the properties of internal microturbulence in various confinement regimes.

2. TEXT AND DIII-D RESULTS

Energy confinement saturation, observed universally with increasing density in gas-fueled, Ohmic tokamak plasmas, is often attributed to anomalous ion transport resulting from ITGD turbulence [4]. Previous measurements on TEXT revealed the coexistence of two distinct turbulent features propagating in the ion and electron diamagnetic drift directions (see Fig.2a) [5]. The "ion" feature in the measured spectra increased in magnitude (relative to the "electron" feature) as energy confinement saturated above densities of ~ 5 x 10^{13} cm^{-3}. In addition, pellet injection, used to peak the density profile, resulted in suppression of this "ion" feature [5]. These characteristics are, in many ways, consistent with the theoretical expectations of ITGD turbulence. However, there are some valid concerns that need to be addressed.

Since scattering measurements are made in the laboratory frame of reference, the effects of radial electric fields, and associated plasma rotation, on the scattered spectra have to be considered. At low densities on TEXT (2x10^{13} cm^{-3}), the internal plasma potential has been measured to be negative (inward pointing radial electric field) using heavy ion beam probe (HIBP) techniques [6]. This field results in the measured spectra being shifted towards the electron diamagnetic drift direction. Preliminary data [7] indicate that the electric field increases and remains inward pointing as densities are increased up to 5x10^{13} cm^{-3}. HIBP data are presently unavailable at higher densities. However, the "ion" feature is apparently propagating more strongly in the ion diamagnetic drift direction than is indicated by uncorrected data. The larger corrected phase velocity of the ion feature may then be in disagreement with some theoretical expectations for ITGD turbulence, where \omega = kv_{ph} is expected to be less than the ion diamagnetic drift frequency. Careful reassessment of theory and experiment is important to resolve this issue.
Recent data suggest that the observed "ion" feature exists over a broad region of the plasma cross-section. When the scattering volume is placed at plasma center, the "ion" feature is clearly observed for high density operation. This is illustrated in Fig. 1 for densities of 2 and 8 x 10^{13}cm^{-3}. At low density (Fig.1.a), the two features in the scattered spectra correspond to fluctuations propagating in the electron diamagnetic drift direction below and above the plasma midplane. At high densities, the third feature, close to zero frequency in Fig.1b, is the "ion" feature.

It has been suggested that this "ion" feature in the spectra arises near the plasma edge and is not representative of internal fluctuations. A shear layer and outward pointing electric field are known to exist near the limiter in TEXT and could result in the generation of a false "ion" feature. These effects would only be significant if fluctuations were peaked at the plasma edge and the scattering system possessed inadequate spatial resolution. However, the spatial resolution has recently been established [8] by studying the effect of perturbations induced by an applied resonant field. At electron densities \leq 3\times 10^{13}cm^{-3}, the perturbation reverses the propagation direction of the turbulence outside a minor radius of 0.8a. This is interpreted as resulting from a change in the radial electric field direction as measured using the HIBP [7]. However, when the center of the scattering volume is located within \pm 0.25a of plasma center, no changes in propagation direction are observed. This confirms that, at low density, measurements at plasma center do not originate from, and are not affected by, fluctuations at the plasma edge. In addition, preliminary data, where scattering signals from different radial locations are cross-correlated, also support the interpretation that the "ion" feature exists over a broad region of the plasma cross-section, whereas the "electron" feature is primarily localised to the plasma edge.
Recent results have also clearly linked the above "ion" feature to the observation of long-time (~ 5τE) precursors to high-density (Hugill) limit disruptions on TEXT. Coincident with the deterioration of particle and heat transport, the fluctuation level of the high-density "ion" feature was found to significantly increase well ahead of the disruption. This is illustrated in Fig. 2. Density fluctuation spectra are shown at different times during a disrupting discharge. Note that 110 ms prior to the disruption roughly equal fluctuation levels exist for the "electron" and "ion" features. This is typical of a nondisrupting discharge. However, as the disruption approaches the "ion" feature continues to increase, whereas the "electron" feature remains essentially unmodified. This growth in fluctuation amplitude, which may be related to ITGD turbulence, provides a physical mechanism for the observed concomitant degradation in particle and heat confinement (Fig. 2d). It should be noted that these long-time precursors were not observed for disruptions where the radiation (Murakami) limit was exceeded.
A heterodyne scattering system has also recently been installed on the DIII-D tokamak[1]. Fluctuation measurements have initially concentrated on investigating saturated gas-fueled, Ohmic operation in order to compare with the TEXT results. In contrast to TEXT, confinement saturates at relatively low densities (~2.5x10^{13} cm^{-3}) on DIII-D. As illustrated in Fig.3a, in limiter plasmas, density fluctuations were found to propagate dominantly in the electron diamagnetic drift direction for all plasmas (n_e ≤ 4.5x10^{13} cm^{-3}), except the locked mode, where propagation strongly in the ion diamagnetic drift direction was observed. In contrast with TEXT data, two clearly distinguishable spectral features were not observable. The existence of a broad, single spectral feature is, however, generally consistent with measurements on most other tokamaks [9]. In divertor operation on DIII-D, propagation in either the electron or ion direction was observed, dependent on plasma conditions. This is illustrated in Fig. 3b where fluctuation data during a current ramp are displayed. At low current, propagation in the electron direction is dominant, whereas, at 1MA, fluctuations propagate almost totally in the ion direction.

It should be noted that the fluctuation spectra displayed above are obtained in the laboratory frame of reference. It can be shown [10,11] that the effect of electric fields, and associated plasma rotation, can be accounted for by applying the following correction to measured frequency spectra:

\[ \omega = \omega_0 - k_0 c E_r / B_0 \]  \hspace{1cm} (1)

The radial electric field is calculated using the steady state force balance equation

\[ \frac{1}{Z_i n_i} \frac{\partial p_i}{\partial r} = c E_r - e / c (v_\phi B_\theta - v_\theta B_\phi) \]  \hspace{1cm} (2)

where \( \omega \) is the measured frequency, \( \omega_0 \) is the "theoretical" mode frequency, \( p_i \) is the ion pressure, and \( k_0 \) is the poloidal wavenumber. The toroidal rotation is measured using CER. The poloidal rotation, which is too small to be accurately measured, is calculated using the neoclassical assumption of poloidal rotation damping given by

\[ v_\theta = k_1 (\partial T_i / \partial r) c / e B_\phi \]  \hspace{1cm} (3)

where \( k_1 \) is a neoclassical coefficient dependent on the ion collisionality and impurity strength parameter[12].

Fluctuation data taken in Ohmic plasmas have been corrected using Eq.(1). The toroidal rotation was expected to be small in these Ohmic plasmas. However, as shown in Fig.3c, large toroidal co-rotation velocities (~20km/s)
FIG. 3. (a) Heterodyne collective scattering data are shown for $k = 5.5 \text{ cm}^{-1}$. The scattering volume was centered near the outer midplane at $R=2.2\text{m}$. Normal Ohmic limiter data are illustrated in (a) together with data obtained in the presence of locked modes. Data obtained during Ohmic divertor operation are illustrated in (b) for various plasma currents. Note the reversal of "apparent" propagation direction. The spatial variation in measured toroidal plasma rotation is illustrated in (c) for low (0.5MA) and high (1MA) plasma current during double null divertor plasmas.

were measured in the vicinity of the scattering volume ($R\sim2.2\text{m}$). Note also that toroidal rotation increased for larger plasma currents. The correction factor at the center of the scattering volume in high current divertor operation was significant: $200\text{kHz} (\pm 75\text{kHz})$ toward electron propagation for $k \sim 5\text{cm}^{-1}$. This implies a strongly outward pointing electric field. In contrast, for low
current divertor operation the correction was close to zero (± 50 kHz). The correction is, therefore, large for high current operation, where the measured frequency would indicate dominantly "ion" propagation. In all cases considered so far, except the locked mode, corrected data have been consistently moved towards propagation in the electron diamagnetic drift direction. Future work will concentrate on H-mode plasmas where measured values for the larger poloidal and toroidal rotation velocity profiles can be utilized.

3. TFTR RESULTS

Density and temperature fluctuations have been measured in different heating regimes in TFTR. The k-spectra of density fluctuations in the plasma interior measured with microwave scattering follow an inverse power law \( \left< \delta n_e^2 \right> \sim 1/k_\perp^\alpha \) with \( \alpha = 3 \) for Ohmic, L-mode, and supershot regimes where \( k_\perp \rho_S \sim 0.50 - 1.0 \). The value of \( \delta n/e/n_e \) in auxiliary heated plasmas estimated from scattering data (\( 2 < k_\perp < 8 \text{ cm}^{-1} \)) and from BES (\( k_\perp < 2 \text{ cm}^{-1} \)) are approximately 1% and consistent with simple mixing length estimates. Similarly, an upper limit in all regimes on electron temperature fluctuations \( 8T_e/T_e < 0.4\% \) for \( k_\perp < 0.3 \text{ cm}^{-1} \) has been established from measurements of second harmonic electron cyclotron emission.

3.1 Electron density fluctuations

The spectra seen in scattering experiments are generally identified with electron drift modes by their spectral shift, \( \omega_e^* = k_\perp T_e/eB_0L_n \) in the electron diamagnetic drift direction (Ohmic plasmas only), the apparent width, \( \sim \omega_e^* \), and by the dominant values of \( k_\perp < 0.5/\rho_S \). On TFTR, as on other devices, the frequency shift, predicted by the simple linear dispersion relation, is actually several times smaller than the measured shift. Poloidal rotation in the electron diamagnetic drift direction or, more likely, counter toroidal rotation, may be needed to explain the difference. However, direct rotation measurements to confirm this are beyond the accuracy of current instrumentation on TFTR. In beam heated discharges toroidal plasma rotation usually dominates the form of the frequency spectrum. The spectral width is caused by instrumental averaging of the toroidal rotation profile over the scattering volume. Measurements of the rotation profile by charge exchange recombination spectroscopy show that the width of the turbulent spectrum approaches the Ohmic width (<100kHz) in balanced discharges. Thus, in the absence of strong MHD activity, the apparent broadening of the turbulent frequency spectrum during beam heating is primarily a kinematic effect due to toroidal rotation.
Scattering in the plasma interior in all regimes is dominated by values of $k_\perp$ below about 4 cm$^{-1}$ which cannot be localised well [2]. The k-spectrum falls approximately $k_\perp^{-3}$. Further, with the assumptions that the fluctuations are uncorrelated over a scale length of $\pi/k_\perp$, are isotropic perpendicular to $B_\parallel$, fall within the instrumental resolution in $k_\parallel$, and cut off at 2 cm$^{-1}$, the magnitude of $\delta n_e/n_e$ approximately agrees with the mixing length estimate $1/k_\perp<\ell$ where $\ell=(L_nL_T)^{1/2}$ and the form of the k-spectrum is consistent with results from nonlinear estimates of collisionless trapped electron mode turbulence [13].

Measurements of density fluctuations using BES [3] with a resolution of 1-2 cm are sensitive to spatial scales with $k_\perp< 2$ cm$^{-1}$ where scattering has shown the highest fluctuation activity. Simultaneous 10-channel sampling has been performed radially on the horizontal midplane between the magnetic axis and the scrape-off layer, and poloidally near the edge and q=1 radius in supershot and L-mode plasmas. In supershots fluctuations are observed both at the edge and in the interior (Fig. 4). The outer 5 cm have low frequency density fluctuations with $\delta n_e/n_e \approx 10\%$ for $f>10$ kHz (Fig. 4d), whereas the interior region has spectra with frequencies from 10 to 100 kHz at a level of 1.0 - 1.5%. Some mode structure can be seen in the frequency spectra. Near the q=1
surface ($R \sim 288\text{cm}$) activity at 40, 60-80 kHz is observed over a region 10 cm wide (Fig.4a). Similarly, activity at 10-40 kHz is observed over about 6cm in radius outside the $q=2$ surface (Fig.4c). Between these regions the activity is lower (Fig.4b). In L-mode discharges the edge fluctuation spectra are similar to supershots, but interior fluctuations above 10kHz are below the present instrumental resolution of about 0.5%.

While similar in magnitude, fluctuation activity seen in scattering and BES data may not be directly related. In particular, the BES system is sensitive to the entire range $0 < k_\perp < 2\text{cm}^{-1}$, some of which may include MHD-related activity, while scattering measurements are relatively insensitive to such phenomena. Experiments to explore low-k phenomena, microturbulence, and connections to transport are the subject of the TFTR fluctuation studies program.

3.2 Electron temperature fluctuations

An upper bound on $\delta T_e/T_e \leq 0.4\%$ is inferred from electron cyclotron emission [14] from a spot 10 cm in diameter and 3 cm thick ($\Delta R$). Observed intensity fluctuations are consistent with photon noise statistics except near the plasma edge where $\delta T_e/T_e \sim 0.5 - 3\%$ between 10 and 100 kHz. These measurements are most sensitive to values of $k_\perp < 0.3\text{cm}^{-1}$. Our knowledge of the form of the k-spectrum of these fluctuations in the vicinity of $\ell m^*$ is incomplete, but these estimates indicate that $(\delta T_e/T_e)/(\delta n_e/n_e)<1$ in supershots where $\eta_i < 1$. Linear estimates for drift waves driven unstable by the combined influence of trapped electron dynamics and ion temperature gradient effects indicate that the ratio $(\delta T_e/T_e)/(\delta n_e/n_e)\sim 1$ at the mode maximum, $k_\perp \sim 1\text{cm}^{-1}$, when $\eta_i>1$ [15] as expected for L-mode discharges.

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PELLET ABLATION STUDY IN T-10 USING A PHOTOGRAPHIC TECHNIQUE

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Abstract

PELLET ABLATION STUDY IN T-10 USING A PHOTOGRAPHIC TECHNIQUE.

Both ultrafast multiframe and single frame photographies were used to study the clouds surrounding the pellets injected into the T-10 tokamak plasma. Using the VSK-5 fast camera, up to 60 frames of clouds in the spectral range of 0.4–0.6 μm were obtained in one shot. The observations were performed from the direction of injection, which allowed the influence of poloidal drifts in the plasma on the structure of the clouds to be investigated. Substantial asymmetry in the poloidal direction was observed. Striations of the clouds for hydrogen and carbon pellets were recorded. Data on the entire cloud length in the toroidal and poloidal directions are obtained as well as the decay lengths for cloud luminosity. The toroidal decay length for carbon clouds and the entire length of the hydrogen clouds are proportional to the ablation rate, while the poloidal and toroidal decay lengths for hydrogen vary, similarly to the ionization length of the pellet atoms due to primary plasma electrons. Single photography with spectral resolution yielded information on pellet trajectory deflection and the toroidal distribution of the atom and ion line luminosities. By the multiframe technique, data on the current profile were obtained both by hydrogen and carbon pellet injections. The carbon injection seems to be a better tool for \( j(r) \) measurements, owing to the simpler cloud structure.

1. INTRODUCTION

The successful solution of fuelling and plasma diagnostic problems using pellet injection needs an adequate model for the plasma–pellet interaction. The main problem of the pellet ablation theory is the determination of the shielding mechanism (neutral, plasma, magnetic) of the electron heat fluxes flowing onto the pellet surface [1]. Measurements of the light emissivity from the ablation cloud can be useful for shielding analysis. In addition, photography of pellet clouds and trajectories can be used to determine drift motions [2], magnetic fields [3–5] and fast electrons [6, 7].
There are a few studies of pellet clouds in tokamak experiments with time resolution [8, 9]. Some of these studies show cloud instabilities which must be taken into account in developing ablation models. Photography study results for both hydrogen and impurity (C, LiH, LiF) pellets ablation are presented in this paper.

2. EXPERIMENTAL SET-UP

Pellets were injected into the tokamak in the direction of the plasma core. Possible directions of pellet injection relative to the photography technique are shown in Fig. 1. Pellet parameters and types of photography are presented in Table I. The VSK-5 camera without any filters and with an Izopanchrom-17 film recorded radiation in the spectral range of $\lambda = 400$–$600$ nm. It was possible to obtain up to 72 frames during pellet ablation. The frame frequency depended on the pellet velocity expected and was chosen such that the pellet shift during the exposure was less than 0.5 cm. Single frame photography was provided by a camera with a shutter open for a period exceeding the ablation time. In this way, the pellet trajectory and the toroidal distribution of the cloud luminosity were obtained. By using various interference filters, it was possible in this case to observe the cloud in different spectral lines of atoms and ions of the injected materials: Li I ($\lambda = 670$ nm), C I ($\lambda = 600$ nm), C II ($\lambda = 510$ nm). Image processing was performed by using an AMD-1 automatic microdensitometer with subsequent computer analysis. The samples of cloud images for multiframe and single frame photographies are shown in Fig. 2. The first column presents data for hydrogen clouds, the second column for carbon pellets, and the third column includes photos of pellet trajectories for various materials and filters. The measurements were performed in the Ohmic regimes of the T-10 tokamak with plasma currents of $I_p = 200$–$400$ kA, magnetic fields of $B_t = 2.5$–$3$ T, limiter radii of $a_l = 28$–$34$ cm and electron densities of $n_e = (1$–$5) \times 10^{13}$ cm$^{-3}$. The ablation times were equal to $\sim 0.3$ and $\sim 1$ ms for hydrogen and carbon pellets, respectively.
TABLE I. PELLET PARAMETERS AND TYPES OF PHOTOGRAPHY

<table>
<thead>
<tr>
<th>Pellet material</th>
<th>Form, size (mm)</th>
<th>Velocity (m/s)</th>
<th>Type of photography(^a)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\text{H}_2)</td>
<td>Cylinder (\phi 1.35 \times 1.35)</td>
<td>400-700</td>
<td>Multiframe, I; single frame, I</td>
</tr>
<tr>
<td>(\text{C})</td>
<td>Sphere (\phi 0.2-0.5)</td>
<td>100-150</td>
<td>Multiframe, I; single frame, (C I, C II), I</td>
</tr>
<tr>
<td>(\text{LiF})</td>
<td>Cube 0.2-0.4</td>
<td>100-150</td>
<td>Single frame (Li I), II</td>
</tr>
<tr>
<td>(\text{LiD})</td>
<td>Cube 0.3-0.6</td>
<td>100-150</td>
<td>Single frame II</td>
</tr>
</tbody>
</table>

\(^a\) C I, C II, Li I: spectrally resolved photography. I, II: different directions of pellet injection relative to the camera.

3. EXPERIMENTAL RESULTS

The data on hydrogen clouds presented in the first column of Fig. 2 correspond to shot No. 49968 with \(I_p = 208\) kA, \(B_t = 2.84\) T, \(n_e = 3 \times 10^{13}\) cm\(^{-3}\), \(a_i = 28\) cm and a pellet velocity of \(V_p = 710\) m/s. In this shot, the pellet penetration was deep, so that the centre of the plasma was reached. We may class the types of the cloud observed into four categories: symmetric (Fig. 2(a)), asymmetric (Fig. 2(b)), striated (Fig. 2(d, e)) and ‘curved’ (Fig. 2(c)). Symmetrical clouds can be observed at any stage of ablation, but usually the initial stage is preferable. A poloidally asymmetric cloud picture is the most natural state for a cloud; this type of cloud is most frequent. Asymmetry both in neo- and anti-neoclassical directions may exist in one shot.

Striated clouds occur several times during pellet penetration. Striation arises at a distance of \(~1\) cm from the main cloud and displays a shape similar to the latter. The time of existence of this striation was not longer than the time interval between two successive frames (\(\leq 7\) \(\mu\)s, in our case).

‘Curved’ clouds were observed in a limited number of cases.

Along with the poloidal cloud asymmetry described above, significant asymmetry was observed in the toroidal direction. The decay length of the cloud luminosity from the current (ion) side is greater than from the opposite side, which seems reasonable because a stronger heat flux impinges on the pellet from the electron side.
FIG. 2. Results of multiframe and single frame photographies: (a)–(e) hydrogen clouds, No. 49968, exposure time 5 µs; (f)–(j) carbon clouds, No. 43116, exposure time 40 µs; (k)–(o) single frame photos; k, l, o: visible spectral range, m: line C II, n: line C I.
Carbon clouds at various stages of ablation are shown in the second column of Fig. 2. They have a structure similar to that of the hydrogen clouds. The strong poloidal asymmetry of the carbon clouds was usually observed at the initial stage (Fig. 2(f)). The middle stage of the ablation was characterized by a symmetric type of cloud. For carbon, we did not observe 'curved' clouds. Up to four striations could be observed in the carbon case.

Significant deflection of the pellet trajectories in the direction opposite to that of the plasma current is distinctly discernible in the single frame photos (Fig. 2, column 3). The trajectory is better distinguishable for impurity pellets when interference filters are used.

The distributions of light intensity, I, in the clouds along toroidal (solid curve) and poloidal (dashed curve) directions are shown on the semilogarithmic scale of Fig. 3.

In the case of carbon pellets, the intensity decay in the toroidal direction begins at the pellet position and is almost exponential, with a slowing-down near the cloud boundary (Fig. 3(a)). For hydrogen pellets, along with the regions of exponential intensity decay at the cloud boundary, a substantial zone of weak changes in luminosity exists (Fig. 3(b)). The ion and electron sides of the cloud have different toroidal decay lengths $\ell_t$. The same effect takes place for the neoclassical and opposite sides of the poloidal distribution. The decay lengths $\ell_r$ and $\ell_i$ are determined as the lengths for which the luminosity decreases by a factor of $e$.

Radial profiles of $\ell_r$ and $\ell_i$ are shown in Fig. 4 for hydrogen and in Fig. 5 for carbon pellets. For hydrogen pellets, a slight asymmetry with a larger wing in the antineoclassical direction exists in the outer region of the plasma where the minor radius $r$ is longer than 12 cm. In the inner region, significant cloud asymmetry with
a longer tail in the neoclassical direction and striations occur. The value of $\ell_\perp$ corresponding to the antineoclassical side of the cloud decreases with pellet penetration into the plasma; it is of the order of the ionization length, $\ell_{\text{ion}} = V_T/n_e \left\langle \sigma v \right\rangle_{\text{ion}}$, due to the primary plasma electrons. ($V_T \approx 3 \times 10^5$ cm/s is the expansion velocity in the carbon or hydrogen cloud). The poloidal decay length $\ell_\perp$ for carbon clouds depends weakly on the minor radius and slightly exceeds the value of the ionization length (Fig. 5).

The behaviour of the toroidal decay length differs substantially for hydrogen and carbon pellets. The $\ell_\parallel$ value for hydrogen clouds increases with the minor radius and exceeds the $\ell_\perp$ value for the antineoclassical side by a factor of 2–3 (Fig. 4(b)).
The toroidal decay length for carbon clouds is approximately proportional to the ablation rate $\dot{N}$ (Fig. 5). This dependence can also be obtained for the entire toroidal length $L_\perp$ for hydrogen clouds, as can be seen from Fig. 6 [10].

Spectrally resolved single frame photos of carbon pellets have shown that the toroidal distribution of luminosity in the cloud has the shape of a narrow spike on a broad pedestal. A typical C II line intensity distribution is shown in Fig. 7. The $\delta_1$ scales for the spike are less than 1 mm and change slightly during pellet ablation (Fig. 8). $\delta_2$ scales for the pedestal region are up to ten times greater (3–10 mm). There was no significant difference in the toroidal distributions for $\lambda = 600 \pm 5$ nm (C I) and $\lambda = 520 \pm 5$ nm (C II) filters. There was no observable pedestal in the case of a lithium cloud with $\lambda = 670 \pm 0.5$ nm (Li I) filter.

The elongation of the hydrogen cloud in the direction of the magnetic field is characterized by the smaller $\delta_1/L_\perp$ ratio than in the case of a carbon pellet cloud.
Nevertheless, we obtained some information on the magnetic field rotational angle, $\vartheta = B_p/B_n$, by using hydrogen pictures. The dependence of $\vartheta$ and the safety factor $q$ on the minor radius are shown in Fig. 9. Very low $q$-values at the plasma centre (0.2–0.3) corresponding to a sharp plasma current profile, $j \sim (1 -(r/a_L)^2)^9$, (solid curve in Fig. 9(b)) were measured. It is necessary to emphasize that at the initial stage of pellet penetration into plasma ($r > 12$ cm) the data on $\vartheta(r)$ satisfactorily correspond to the profile shape, $j \sim (1 -(r/a_L)^2)^3$ (dashed curve). The sudden changes in the $\vartheta$ at $r < 12$ cm correlate with changes in the poloidal asymmetry of $\lambda_L$, a very strong toroidal deflection of the pellet trajectory and the appearance of striations.

4. **DISCUSSION**

The luminosity of the clouds surrounding the ablating pellets studied in these experiments is determined by numerous processes related to pellet–plasma interaction. Most important among these processes are expansion of neutral gas from the pellet surface, ionization of neutrals (and multiple ionization for impurity pellet ions), capture of ionized particles by the toroidal magnetic field and subsequent expansion of the plasma cloud in the toroidal direction, inhibition of radial expansion of the neutrals by the charged component and drifts of the secondary plasma in the poloidal (or radial) directions due to intrinsic radial or pellet induced fields. In the vicinity of the pellet, on account of the very low temperatures and high densities, recombination processes can couple the charged and neutral components, as well.

The scheme of our multiframe experiment was sensitive to poloidal motions near the pellet. This was due to the not too large angle ($30^\circ$) between injection and observation lines. Hence, we think that the transformations of the cloud observed in these experiments are related to the poloidal drifts in the plasma.
The poloidally symmetric structure of the clouds can be related to the low drift velocity. Either the smallness of the electric field in the main plasma or its decrease to small values brought about by the cloud may explain the observation of symmetric clouds. However, an asymmetric cloud shape is the most typical. To clarify whether intrinsic electrical or pellet induced fields are responsible for cloud asymmetry, independent field measurements are necessary.

The striations observed in the experiments occur on the side of the cloud corresponding to the neoclassical drift tail. The appearance of striation in the hydrogen cloud points to a coupling between neutral and charged particles in the region of luminosity. We think that the nature of the striations is similar to the nature of the striations observed in barium clouds [11] and carbon clouds in a tokamak [2]. Such striations are not connected with the pellet motion in the radial direction as was
assumed in Ref. [9]. The instability due to ion inertia analysed in Ref. [12] is a possible explanation for this effect. It should be mentioned that the striations in poloidal direction presented here are not identical to the striations that are usually observed in single frame pictures [4]. The latter are more likely to be explained by simple bursts of the ablation rate during radial pellet motion.

The appearance of 'curved' clouds has not become clear so far, but could also be a result of very fast drift motions.

It is seen that the picture of clouds in consecutive frames changes significantly. Thus, some effects of integration may take place in our data. It is clear by now that a time resolution shorter than 1 \( \mu s \) is necessary to analyse the formation of striations. A short time of cloud shape evolution points to the existence of fast instabilities in the cloud and to the necessity of developing a non-stationary ablation model [9]. Two- or even three-dimensional models are required to simulate cloud formation [13].

The different behaviour of hydrogen and carbon clouds may be due to two facts: The longitudinal tails at the boundary of the carbon cloud are determined by ion radiation. Two orders of magnitude higher ablation rates for hydrogen, as compared to carbon pellets, create the conditions for the development of a large zone with uniform luminosity near the hydrogen pellets (Fig. 3(b)). In this region, recombination is likely to dominate. Outside this zone, neutrals which have left it are ionized by primary electrons. From the toroidal decay length \( \ell_t \), the velocity of the neutrals at the cloud boundary is estimated to be \( 10^6 \text{ cm/s} \). These values of velocity, corresponding to a temperature of \( \sim 1 \text{ eV} \), are in agreement with spectral measurements of secondary plasma parameters near the pellet [4, 14]. Recently, a theory predicting cloud parameters near the hydrogen pellet has appeared [15]. The comparison of the experimental with theoretical data is, however, complicated by the numerical character of the model.

The data on the \( \phi \)-angle obtained by using hydrogen pellets demonstrate the complicated structure of the current profile and may make the angle determination doubtful. However, the decrease in the current density in the region close to 10 cm also follows from the electron temperature profile measured by using second cyclotron harmonic radiation and the ablation curve of the carbon pellet injected 60 ms after hydrogen. The multiframe technique yields better results than measurements in the single frame regime using striation [4], which occur at undefined instants of time and are concentrated within a limited portion of the ablation curve. Carbon clouds seem to be more appropriate for angle measurements because they show long wings while hydrogen clouds have abrupt edges.

5. CONCLUSIONS

The photographic technique allows a good deal of information to be obtained on the cloud structure near the ablating pellet. The clouds have a complicated shape
(three-dimensional, as a rule). Symmetric, poloidally and toroidally asymmetric as well as striated clouds are shaped during the ablation process. The cloud shape varies quickly during ablation, within a time interval of a few microseconds. The complicated structure of the clouds may be related to drift motions in the cloud plasma. There is a significant difference in carbon and hydrogen cloud behaviour, which may be explained by differences in ablation rates and ion radiation. The clouds are elongated in the magnetic field direction, which allows the use of cloud images to determine the current profile. Asymmetry in the clouds is a source of errors for magnetic field measurements. The simpler structure of the carbon clouds makes carbon pellet injection preferable for current profile measurements.

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Session E-III     (Posters)        

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REVIEW OF LOWER HYBRID EXPERIMENTS ON ASDEX

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Abstract

REVIEW OF LOWER HYBRID EXPERIMENTS ON ASDEX.

Current drive and heating with lower hybrid waves have been investigated on ASDEX with two RF systems at 1.3 and 2.45 GHz. The results on current drive efficiency agree well with theory, including the effect of a DC electric field. Global energy and particle confinement are improved in the presence of a large suprathermal electron population. Stabilization of the \( m = 1 \) mode gives rise to strong central electron heating with \( T_e \lesssim 8 \text{ keV} \) and peaking of the \( T_e(r) \) profile. Local control of the current profile is achieved by launching compound spectra from independent grill antennas. The LH power deposition profile and the modification of the current profile were modelled with a fully self-consistent 3-D Fokker–Planck code.

1. WAVE ABSORPTION

Lower hybrid (LH) waves were launched in ASDEX in a first phase at 1.3 GHz with \( P_{LH} \leq 1 \text{ MW} \) from an eight waveguide grill antenna, with wave spectra in the range of \( 1 < N_i < 4 \) and a width of \( \Delta N_i = 2 \) [1]. In a second phase with a new transmitter at 2.45 GHz, up to 2 MW were injected from two stacked antennas consisting of 24 waveguides each. The much narrower wave spectrum \( (\Delta N_i = 0.8) \) can be varied in the range of \( 1 < N_i < 4.4 \) [2].

LH waves are absorbed in the plasma either by accelerating electrons or ions perpendicular to the magnetic field. Characteristic parameters of both regimes are shown in Fig. 1 for LH current drive experiments with 2.45 GHz. The transition from electron to ion interaction occurs at \( n_e = (4-5) \times 10^{13} \text{ cm}^{-3} \), where fast electron generation and LH current drive vanish and the flux of fast ions rises strongly. The switchover was found at the same density in LH experiments on FT [3] and PLT [4] for this frequency. For 1.3 GHz on ASDEX, it occurred at \( n_e \approx 2 \times 10^{13} \text{ cm}^{-3} \) [5].

The absorption coefficient \( \alpha \) of LH power has been determined experimentally from the rate of change of the total plasma energy content at the start and at the end of LH pulses on a time-scale shorter than the energy confinement time [6, 7]. The plasma energy content is determined from the magnetic measurements of equilibrium and diamagnetic beta values, taking into account the pressure anisotropy in the presence of suprathermal electrons with \( \beta_{p,\text{tot}} = (2\beta_p^{\text{equ}} + \beta_p^{\text{dii}})/3 \). No dependence of \( \alpha \) on the power was found. At low density, \( \alpha \) increases with decreasing \( N_i \), as is shown in Fig. 2. With LH current drive opposite to the DC electric field, \( \alpha \) is by a factor of two smaller than with normal current drive. Intermediate results are obtained with symmetric spectra, both for injection with two grills, or with one grill together with full current drive from the other grill. The DC electric field has, therefore, no influence on \( \alpha \). The dependence of \( \alpha \) on the direction of the launched waves is not understood at present. At intermediate densities the value of \( \alpha \) in only weakly dependent on \( N_i \) and on the width of the spectrum. With increasing plasma density,
FIG. 1. Variation of characteristic plasma parameters with density during LH current drive at 2.45 GHz with $N_t = 2.2$, $P_{\text{LH}} = 1 \text{ MW}$, $I_p = 420 \text{ kA}$, $B_i = 2.8 \text{ T}$. $\Delta U/(LH)/U/(OH)$ = relative drop of the loop voltage due to LH current drive; $\phi_{\text{HX}}$ = photon flux from hard X ray emission at 100 keV; $\phi_{\text{CX}}$ = flux of 15 keV ions measured by charge exchange analyser; $A_{\text{PD1}}$ = relative amplitude of first satellite in parametric decay spectrum.

FIG. 2. Absorption coefficient as a function of $\bar{n}_e$ for different $N_t$ spectra. Full symbols: $I_p = 420 \text{ kA}$, $B_i = 2.8 \text{ T}$; open symbols: $I_p = 320 \text{ kA}$, $B_i = 2.17 \text{ T}$.
\( \alpha \) decreases for current drive spectra, particularly strongly at low magnetic field, as is expected from accessibility considerations. Similar results were found with 1.3 GHz [7]. For the symmetric spectrum with 00\( \pi \) at the same \( N_B \) as for 90° current drive phasing, however, no density dependence of \( \alpha \) is seen. Absorption is also found in cases where no power should reach the central plasma region, according to single pass ray tracing calculations. In the previous 1.3 GHz experiments, \( \alpha \) also clearly exceeded the fraction of accessible power at higher density. This might indicate an \( N_B \) upshift of the launched spectrum at higher \( n_e \). A possible mechanism could be scattering of the LH waves from density fluctuations at the plasma edge, as suggested by the increasing broadening of the pump wave with \( n_e \) [8].

2. ENERGY CONFINEMENT

The increase in total energy content \( W_p \), calculated from \( \beta_{p,\text{IM}} \), and in thermal electron energy content \( W_e \), determined from laser light scattering, is larger for current drive than for symmetric spectra at the same \( N_B \). The difference disappears when the different absorption coefficients are taken into account.

The scaling of \( W_p \) and \( W_e \) with total power \( P_{\text{tot}} = P_{\text{OH}} + P_{\text{LH}} \) (\( \alpha = 1 \)) is shown in Fig. 3 for LH current drive at \( n_e = 1.3 \times 10^{13} \text{ cm}^{-3} \). The total energy content \( W_p \) with OH + LH current drive is about 50% higher than for OH with the same total power input. In combined operation at low density, smaller total power

![FIG. 3. Total energy content \( W_p \) (•) and thermal electron energy content \( W_e \) (×) versus total power \( P_{\text{tot}} = P_{\text{OH}} + P_{\text{LH}} \) for LH current drive (\( \Delta \phi = 90^\circ \)), together with Ohmic values for \( W_p \) (○) and \( W_e \) (□)). \( B_i = 2.8 \text{ T} \) \( I_p = 420 \text{ kA} \).]
is required to sustain the same plasma current than with OH current drive alone. In this regime, the total energy content remains above the Ohmic value, and the confinement time is improved by up to a factor of 1.8.

The same incremental energy confinement time $\tau_{E}^{\text{inc}}$ is obtained over the whole range of powers, also in the regime $P_{\text{tot}} < P_{\text{OH}}$. Its value is smaller than the Ohmic confinement time. The global confinement time, therefore, degrades continuously with power, starting already at $P_{\text{tot}} < P_{\text{OH}}$. For the thermal energy content of the electrons $W_e$, no increase above the Ohmic value is found. Hence, an improvement in global confinement has to be attributed to a superior energy confinement of the suprathermal electron population. With increasing density and an associated decreasing fraction of LH produced fast electrons, the improvement of $\tau_{E}$ with LH disappears. The incremental energy confinement time is independent of the density if the absorbed power is rated [6]. With full LH current drive, the global energy confinement time $\tau_{E}$ becomes independent of the density while the required LH power increases with the density: $P_{\text{LH}} (I_p = \text{const}) \sim \bar{n}_e$, in contrast to OH sustained discharges where $P_{\text{OH}} \approx \text{const}(\bar{n}_e)$ and $\tau_{E} \sim \bar{n}_e$ in the linear confinement regime. The two schemes of pure OH and LH current drive are compared in Fig. 4 for a wide range of operating conditions with different $Z_{\text{eff}}$ values.

![FIG. 4. Total power input and global energy confinement time $\tau_{E}$ versus density $\bar{n}_e$ for OH and for full LH current drive. $B_t = 2.8$ T, $I_p = 420$ kA. $\Delta \varphi = 90^\circ$, $\bar{N}_t = 2.2.$](image-url)
Global confinement degrades with LH power similarly to Goldston L-mode scaling. The same increase in (diamagnetic) energy content $W_d$ with total absorbed power is obtained with LH, NI and LH + NI combined in pure deuterium plasmas as is seen in Fig. 5. With LH, $\tau_{E}^{\text{in}}$ is nearly independent of the plasma current [6]. ELM-free H-phases could be obtained by switching on LH in steady state L-phases during NI. By adding LH to NI in H-phases with high frequency ELMs and L-mode-like confinement, the ELM frequency could be reduced and $\tau_{E}$ raised.

3. CURRENT DRIVE EFFICIENCY

Plasma currents $I_p \leq 460 \, \text{kA}$ were driven up to densities of $n_e = 2.8 \times 10^{13} \, \text{cm}^{-3}$. From shots with full current drive the current drive efficiency $\eta_0$ at zero electric field is determined:

$$\eta_{0,\text{exp}} = \frac{\bar{n}_e \times I_p \times R_{\text{LHO}}}{(10^{13} \, \text{cm}^{-3} \cdot \text{A} \cdot \text{m}/\text{W}) \times R} ,$$

Figure 6 shows the density dependence of $\eta_0$. With 2.45 GHz, 70% higher efficiencies are obtained than with 1.3 GHz [9]. This is attributed to the narrower $N_I$ spectrum of the 24-waveguide grill [2].
From the theory [10], the zero electric field efficiency is derived

$$\eta_{0,\text{theor}} = \frac{1240}{\ln\Lambda(5 + Z_{\text{eff}})} \langle N_i^2 \rangle$$

where $\langle N_i^2 \rangle$ is weighted with the shape of the accessible part of the $N_i$ spectrum propagating in the direction of the initial DC electric field [11]. The accessibility limit is determined from the plane wave accessibility condition with the line averaged density $\bar{n}_e$ and the magnetic field on axis. We neglect any possible broadening or upshift of the spectrum. At low $\bar{n}_e$, all spectra with $\Delta\phi \geq 60^\circ$ satisfy this accessibility condition, and the experimental results agree well with the theory. For $\Delta\phi < 60^\circ$, the fraction of accessible power drops rapidly with decreasing phase. This should result in a strong degradation of $\eta$, which is not, however, seen in the experiment. Efficient current drive with off-axis deposition or a slight upshift of the $N_i$ spectrum could explain this discrepancy. With increasing $\bar{n}_e$, on the other hand, the experimental values of $\eta_0$ decrease more than expected from the calculations, even for fully accessible spectra. This may be caused by non-linear effects related to the observed broadening of the LH pump frequency spectrum. The $N_i$ spectrum is then expected to be broadened, and part of the launched LH power would be shifted to higher $N_i$ with reduced current drive efficiency [8].
For partial LH current drive with a co-acting DC electric field, for comparison with the theory, we use an analytic approximation for $\eta_0$ without electric field [12]. At low density the loop voltage drops strongly already at low power ($P_{\text{LH}}/P_{\text{LHO}} \leq 0.25$). This cannot be explained by bulk heating. In this range, runaway electrons are formed by the residual DC electric field. They may carry part of the plasma current which is not included in the theory applied. At high density, the agreement is quite good in the whole range of powers. Experiments with LH current drive opposite to the initial DC field or with symmetric spectra can be described on the basis of the same theory [13].

For LH current drive during neutral beam injection we found the same efficiencies as with Ohmic target plasmas.

4. CONTROL OF MHD ACTIVITY

The sawtooth repetition period $\tau_{\text{st}}$ increases with LH power for low-$N_\parallel$ spectra. This may be explained mainly by the enhanced electrical conductivity in the presence of suprathermal electrons. With LH current drive sawteeth are suppressed up to densities of $\bar{n}_e = 5.5 \times 10^{13}$ cm$^{-3}$. The threshold power for sawtooth stabilization is nearly independent of phasing for $\Delta \phi \leq 120^\circ$. It increases with density with the same scaling in the experiments at 1.3 and 2.45 GHz [14].

The $m = 1$ mode grows and saturates at large amplitude after stabilization of sawteeth with LH current drive at 2.45 GHz. The radius of the $q = 1$ surface remains essentially unchanged. This behaviour is different from the 1.3 GHz experiments at lower density where $q$ increased above 1 over the whole plasma after sawtooth stabilization. With higher power LH current drive, the $m = 1$ mode is also stabilized. In this case, the internal inductance $\xi$ decreases slightly, indicating a flattening of the current profile. At even higher powers, an $m = 2$ mode may grow, lock and lead to major disruptions. After suppression of the $m = 1$ mode, the central electron temperature rises strongly and the radial profile $T_e(r)$ is peaking. Up to densities of $\bar{n}_e = 4 \times 10^{13}$ cm$^{-3}$, stronger electron heating is then obtained in current drive mode than with symmetric spectra of same $\bar{N}_e$, as is seen from Fig. 7. The gain from $m = 1$ stabilization is highest at low density. Large changes of $T_e(r)$ are obtained under these conditions with the peaking factor rising up to $Q_{T_e} = 5.5$ from Ohmic values of $Q_{T_e} = 2.2$ for $q_a = 3.2 = \text{const}$. The thermal electron energy content, $W_e$, and the total plasma energy, $W_p$, are enhanced in MHD quiescent phases after suppression of sawteeth and $m = 1$ modes. Stabilization of the $m = 1$ mode was also achieved with LH current drive during NBI, resulting in $T_e(r)$ peaking and enhanced global confinement.
5. LOCAL CURRENT PROFILE CONTROL

Local control of the current profile was attempted in different ways with the two LH systems. At 1.3 GHz the $N_0$ spectrum was shaped by varying power and phase distribution across the waveguides of the single grill launcher. Strong modification of $j(r)$ was then achieved [15, 16]. With the 2.45 GHz system two different $N_0$ spectra were launched from the two independent couplers. $j(r)$ is broadened most efficiently with compound spectra combining low-$N_0$ current drive spectra with high-$N_0$ spectra. An example is shown in Fig. 8. First a low $N_0$ current drive spectrum ($\Delta\phi = 90^\circ$, $\bar{N}_0 = 2.2$) is applied. $P_{OH}$ drops and the full plasma current is driven by the LH. $\beta_p^{\text{eq}} + (\xi - \xi_{eq}) \times 0.5$ rises more than $\beta_p^{\text{an}}$, because of the anisotropy $\beta_p^{\beta} > \beta_p^{\beta+}$. $T_{eo}$ rises after reduction of the $m = 1$ activity. In a second phase, a high $N_0$ ($\Delta\phi = 180^\circ$, $\bar{N}_0 = 4.4$) symmetric spectrum is added. $P_{OH}$ returns after a transient overshoot to the level of the first phase. The current drive efficiency at the 90° spectrum is, therefore, not changed by the additional 180° spectrum, taking the slightly higher density during the second phase into account. $\beta_p^{\text{eq}} + (\xi - \xi_{eq}) \times 0.5$
starts decaying, indicating a decrease in $\ell_i$ and, therefore, a broadening of $j(r)$, consistent with the overshoot of $P_{\text{OH}}$. Measurements of $j(r)$ with the Li beam show a flattening in the central region in this phase. This suggests broadening of the deposition zone of the 90° current drive spectrum by the 180° spectrum. The decrease in $T_{e0}$ in this phase might then be due to a reduction in central power deposition.

Results on current profile broadening with different LH spectra are summarized in Fig. 9. $\Delta \ell_i$ remains small with LH current drive at low $N_t$ (90°) up to high power. With high $N_t$ spectra of symmetric phasing ($\Delta \phi = 180°$), $\ell_i$ decreases with $P_{\text{LH}}$, starting already at low power. This may be explained by a broad profile of LH generated suprathermal electrons carrying part of the inductively driven current.
FIG. 9. Drop of internal inductance \( \xi_i \), profile factor \( Q_{Te} \) and central electron temperature \( T_{eo} \) versus LH power for injection of compound spectra; \( \bar{n}_e = 1.35 \times 10^{13} \text{ cm}^{-3} \).

With compound low/high \( N_1 \) spectra, reductions of \( \xi_i \) larger than from both spectra separately clearly reveal a synergetic effect. At fixed phasing of the current drive spectrum and fixed total power, \( |\Delta \xi_i| \) increases with the phase of the second spectrum. At fixed phasing of both spectra, \( |\Delta \xi_i| \) rises with the fraction of power in the high \( N_1 \) spectrum. These results can only be explained by a broadening of the LH deposition profile, with off-axis shift of the low \( N_1 \) power by the superposed high \( N_1 \) spectra. At a power of \( P_{N1} = 150 \text{ kW} \), no degradation of the central electron heating has been seen so far (Fig. 9), as is observed with higher power \( P_{N1} = 180^\circ = 400 \text{ kW} \) (Fig. 8), which also results in a stronger broadening of \( j(r) \).
While \( \ell \) drops and, therefore, \( j(r) \) broadens with increasing LH power upon injection of compound spectra, the central electron temperature increases and \( T_e(r) \) peaks, even slightly more than in the normal current drive case if the power in the high \( N_i \) spectrum remains moderate (Fig. 9). Already low power levels of high \( N_i \) spectra change the deposition of low \( N_i \) current drive spectra strongly. Current and electron temperature profiles can, therefore, be largely decoupled in compound spectrum operation with LH current drive.

6. MODELLING OF POWER DEPOSITION AND CURRENT PROFILES

Effective modelling of LH power deposition and current profiles requires a detailed accounting of the balance between quasi-linear distortion of the electron distribution function by RF absorption and the collisional relaxation to a Maxwellian distribution. The CQL3D/LH code consists of a 2-D-in-momentum-space, relativistic, bounce averaged, collisional/quasi-linear Fokker-Planck equation solver run on a radial array of flux surfaces [17] combined with an LH ray tracing code [19], modelling the distortion of the electron distribution function \( f_e \) resulting from the absorption of LH wave energy. Thus, account is taken of variations in \( f_e \) as functions of radius, poloidal angle and two momentum space directions. However, radial transport effects are not included.

The steady state of \( f_e \) and the power absorption are obtained by iteration between (1) the Gaussian elimination solution of the Fokker-Planck equation for the steady state \( f_e \) at each flux surface,

\[
\frac{\partial f_e}{\partial t} (u_{l0}, u_{\perp 0}, r, t) = \left\langle \frac{eE_{DC}}{m}, \frac{\partial f}{\partial u} \right\rangle + \langle C(f) \rangle + \langle Q(f) \rangle = 0
\]

where \( u_{l0}, u_{\perp 0} \) are momenta per mass of electrons at the outer equatorial plane of each flux surface, \( \langle \\rangle \) indicates a bounce average; and (2) the RF energy transport equation integrated along a ray,

\[
\nabla v_\gamma \epsilon = - \int d\Omega (\gamma - 1) mc^2 Q(f)
\]

where \( v_\gamma \) is the LH ray group velocity and \( \epsilon \) is the energy density. In Eq. (1), \( C \) is the full collision operator linearized about a Maxwellian distribution shifted in the \( u_1 \) direction to conserve momentum in the electron-electron collision process. The temperature and density profiles of the 'background' Maxwellian distributions are fixed. The quasi-linear operator \( Q \) is the full operator for the finite gyroradius, Landau/transit time magnetic pumping (i.e. zeroth cyclotron harmonic) interaction [19] generalized to include relativity. The launched spectrum of RF energy is discre-
tized into a set of rays which are injected from the plasma periphery. Each ray has a given constant (narrow) width $\Delta N_{\phi}$ in toroidal wavenumber and is discretized into length elements in the poloidal plane short compared to the plasma radius. Each such element contributes to the operator $Q$, and the damping of the ray element is obtained self-consistently by using Eq. (2). The resulting damping summed over all the rays iterates exactly to the value obtained by integrating the total quasi-linear operator with the plasma distributions. The damping rate has been compared with standard expressions [20, 21] and at low RF powers and high mesh resolution the calculated damping reduces to the Maxwellian expression. The rays are traced for multiple passes across the plasma as they mode convert to and from fast waves and reflect at the plasma periphery until the wave energy is absorbed.

The Fokker–Planck analysis has been directed towards current profile modification utilizing compound $N_1$ spectra as is shown in Fig. 8. Using the experimental $T_e$ and $n_e$ profiles late in the $90^\circ$ and $90^\circ + 180^\circ$ phases of this shot ($T_{eo} \sim 4.5$ keV,

![Graph](https://example.com/graph.png)

**FIG. 10.** Calculated LH (a) power deposition and (b) current drive radial profiles late in the $90^\circ$ and $90^\circ + 180^\circ$ phases of Fig. 8.
n_{eo} \sim 2.0 \times 10^{13} \text{ cm}^{-3}, \text{ respectively}, \text{ the CQL3D/LH code gives RF power absorption and current profiles as shown in Figs 10(a) and (b). In the 90° phase, the main spectral lobe at \(1.85 \leq N_i \leq 2.55\) is centrally absorbed. Power absorbed at \(r/a = 0.7\) to 0.9 was injected at \(-7.0 \leq N_i \leq -6.0\), corresponding to the counter-high-\(N_i\) side lobe. In the 90° + 180° phase, additional 30% of the power is launched equally in the range \(-4.75 \leq N_i \leq -4.0\), and \(4 < N_i < 4.75\) and is absorbed at \(r/a \sim 0.5\).

Note, in Fig. 10(a), that central absorption of the 90° power has been reduced by the addition of the 180° power. The mechanism is clear from the electron distribution shown in Figs 11(a) and (b) at \(r/a = 0.53\), the 180° absorption radius. The parallel cut in Fig. 11 through the distribution shows that the lower phase velocity power raises the electron distribution function tail so that some of the 90° power is absorbed at this radius. In addition, in Fig. 10(b) the 180° power has reduced the central current drive and also leads to current density which is positive everywhere, even in the region \(r/a = 0.7-0.9\) of absorption of the high \(N_i\), counter-injected side lobe.

**FIG. 11.** (a) Contour plot of electron distribution versus \(u_t\) and \(u_{\parallel}\); (b) slices versus \(u = p/m\) through the calculated 2-D distribution function, at constant pitch angles: curves correspond to (1) parallel, and (2) antiparallel to the electron drift, (3) trapped passing boundary, and (4) perpendicular; \(u\) is normalized to 1.0 at 1500 keV.
For comparison with the experiment, the absorbed LH power ($\alpha = 0.75$) is used as input to the code. In the simulation the total LH driven current agrees, within 10%, with the experimental result of full current drive of 420 kA. The calculated $j(r)$ profiles during the 90° phase are more peaked than Ohmic, with $\xi (90^\circ) = 2.9$. During the 90° + 180° phase the simulation gives broader $j(r)$ profiles than Ohmic with $\xi (90^\circ + 180^\circ) = 1.2$. The difference is much greater than measured in the experiment ($\Delta \xi_{\exp} \approx 0.25$), indicating, possibly, the importance of current diffusion, particularly in the highly peaked deposition 90° case. The total energy of the plasma electrons is 30.4 kJ in the simulation, in good agreement with the experiment. Forty per cent of this energy is contained in the high energy tail. From the code, line integrated hard X ray temperatures are 60 to 100 keV in the photon energy range 20 to 200 keV, and spectra viewing co-electron current particles are a factor $\sim 10$ larger than the counter spectra, in general agreement with the experimental values.

For a power scan of the 180° power from 0 to 100% of the 90° power, the resulting $\xi$ values are computed with the code. Beyond about $P_{180} \sim P_{90}/2$, the broadening is saturated, corresponding to most of the 90° power being absorbed in the 180° absorption region. These results agree well with the experiment.

7. SUMMARY

The lower hybrid experiments on ASDEX have addressed a number of basic problems in fusion research where the distortion of the electron distribution function by wave absorption and non-inductive current drive gave access to new operation regimes. With the generation of suprathermal electrons by LH the electrical conductivity is increased so that steady state tokamak discharges can be maintained with less power than with Ohmic current drive. The off-set linear law for the scaling of the total plasma energy content with total power input, usually found with high power additional heating, applies also to the case of low power.

The same incremental energy confinement time is valid for the whole range of powers. As $\tau_{\text{E}}^{\text{inc}} < \tau_{\text{E}} (\text{OH})$, the degradation of confinement with power starts below the Ohmic level of power input. From the separation of thermal and suprathermal energy contents, it is seen that suprathermal electrons are better confined than thermal electrons. The confinement time, therefore, depends on the particle energy.

The comparison of Ohmic and fully LH sustained discharges ($E_{\text{DC}} = 0$) shows that the density dependence of the energy confinement time is related to the mechanism of particle acceleration in the current generation process. With resonant acceleration of high energy electrons in the wave field, the required power input is proportional to $n_e$, and $\tau_{\text{E}}$ saturates with density, because of its negative power dependence. With inductive current generation through a drifting Maxwellian distribution, the required power is independent of the density, and the confinement time increases with $n_e$ in the linear confinement regime. The difference might originate from the different links between current drive efficiency and plasma parameters. The
ohmically driven current depends essentially on $T_e$ and $Z_{eff}$, while the LH driven current depends mainly on density.

The role of a DC electric field for the non-inductive current drive efficiency has been studied. Good agreement is found between the experimental results and existing theoretical models. The zero electric field efficiency agrees with theory as long as non-linear pump depletion can be neglected.

With LH current drive the radial current profile can be modified locally. The influence of sawteeth and $m = 1$ modes on the shape of the electron temperature profile could be separated. No large changes are seen after sawtooth suppression as long as the $m = 1$ mode is still active. It is only after $m = 1$ stabilization that strong central electron heating and peaking of the profile is achieved. Profile consistency seems, therefore, to be closely related to the presence of $m = 1$ modes. With LH current drive, electron temperature and current density profiles can be completely decoupled. Peaking of $T_e(r)$ and strong broadening of $j(r)$ are achieved simultaneously with the injection of compound wave spectra. With this operation scheme, it has been demonstrated that the deposition zone of low $N_y$ spectra with high current drive efficiency can be steered with an additional high $N_y$ spectrum of low power. This allows optimum local current profile control.

Modelling of the LH profile control experiments with a 3-D Fokker–Planck code can well reproduce the broadening of power deposition and current density profiles, without any need for adjustment of free parameters. More extensive analysis of the experimental results with this code should improve the understanding of the mechanisms governing power absorption and transport during the application of lower hybrid waves.

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REFERENCES


DISCUSSION

A. BERS: You have shown that for most lower hybrid current drive (LHCD) cases, except the case of 57° phasing, the current drive efficiency can be predicted without invoking an N₁ upshift. Have you done ray tracing calculations for your experiments, and if so, do you see no upshift in N₁?

F.X. SÖLDNER: We have done ray tracing calculations with the experimental parameter input. At low density the variations in N₁ due to toroidal effects are small for a single pass. We find good agreement in absolute numbers between experimental current drive efficiencies and theoretical values calculated on the basis of the original launched spectrum for phases ≥ 60°. At smaller phase angles the calculated current drive efficiency drops rapidly owing to a decrease in accessible power. In the experiments, however, the current drive efficiency decreases only slowly. With a slight N₁ upshift of the main lobe of the spectrum we would obtain agreement between experiment and theory by raising the fraction of accessible power. However, with this small upshift in N₁ the calculated values of current drive efficiency at a larger phase angle, where all power is accessible in any case, are not changed very much.
Much larger N\textsubscript{\ensuremath{\text{i}}} upshifts would be required to fill the gap between target bulk electron distribution and the wave spectrum than are invoked here to explain the observed current drive efficiencies.

Y. KAMADA: During suppression of the \( m = 1 \) mode, is \( q(v) \) less than unity or not?

F.X. SÖLDNER: During the phase of large amplitude saturated \( m = 1 \) activity the \( q \)-profile is flat in the central region. Two \( q = 1 \) surfaces are identified by soft X ray tomography. Therefore, \( q \) is smaller than 1 only off-axis, near the location of the single \( q = 1 \) surface during the sawtoothing phase. With suppression of the \( m = 1 \) mode, the island disappears at about the original position. No shrinking to the centre is seen. After complete stabilization of the \( m = 1 \) mode, \( q \) is above 1 everywhere.

Y. KAMADA: You said that the increased sawtooth period did not affect the total energy confinement time. Did you find \( \tau_{\text{sawtooth}} \) to be longer than \( \tau_{E} \) in the second stage of LH heating?

F.X. SÖLDNER: Sawteeth were stabilized for up to 1 s, the maximum duration of the LH pulse. In these conditions, this corresponds to more than 20 times the energy confinement time.

G. TONON: From your observations, does the electron temperature seem to have any effect on the current drive efficiency?

F.X. SÖLDNER: We find a slight increase of current drive efficiency with volume-averaged electron temperature.

J. JACQUINOT: What is the evolution of the electron density profile when the \( m = 1 \) instability is suppressed?

F.X. SÖLDNER: With full LH-current drive, when \( E_{\text{dc}} \approx 0 \), the density profile is flattened compared with the Ohmic profile. Studies of the transport coefficients by the gas oscillation technique show that the inward velocity drops in this case and even reverses sign to an outward flow in the central region. When the \( m = 1 \) mode is stabilized, the density profile peaks again and returns to a shape similar to the Ohmic profile or even more peaked.
ELECTRON CYCLOTRON HEATING AND CURRENT DRIVE RESULTS FROM THE DIII-D TOKAMAK*


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Abstract

ELECTRON CYCLOTRON HEATING AND CURRENT DRIVE RESULTS FROM THE DIII-D TOKAMAK.

Auxiliary heating experiments with electron cyclotron heating have been carried out in the DIII-D tokamak. Waves at 60 GHz have been launched at power levels up to 1.4 MW from both the high-field and low-field side with the appropriate polarization for damping at the fundamental resonance (2.14 T). Confinement was studied in L-mode and H-mode plasmas for a single-null, open divertor geometry. For L-mode discharges, the energy confinement scaling agrees well with the ITER-89 power law or offset linear scaling relations. With strong off-axis heating, the electron temperature profile remains peaked, and power balance analysis indicates that the transport cannot be described by a purely diffusive model. In H-mode confinement plasmas, the magnitude and scaling of the confinement time are equal to that of plasmas heated by neutral beam injection (NBI), if the energy stored in the fast ions is removed in the NBI cases. A major issue for steady-state H-mode plasmas is control of the edge-localized mode (ELM) behavior. By moving the resonance location ±5 cm around the separatrix, the frequency of giant ELMs can be changed by a factor of three. Non-inductive current drive with electron cyclotron waves has also been investigated. Driven currents up to 70 kA have been observed, but the current drive is enhanced by the residual dc electric field. Currents aiding and opposing the Ohmic current have been measured. The magnitude of the current for co-current drive is greater than expected from modelling which includes trapped particle effects, but no electric field. Preliminary calculations including the residual dc electric field can account for the observed enhancement.

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1. INTRODUCTION

The electron cyclotron heating (ECH) experiments reported here focus on two areas critical to magnetic confinement fusion—the understanding and control of energy confinement and non-inductive current drive operation. Detailed discussions are given elsewhere of the inside launch [1] and outside launch [2] waveguide systems, so only a brief discussion will be given here. The outside launch system was modified in 1988 to launch the ordinary mode. A Gaussian pattern with a linear polarization and an 11° e-folding half width is directly launched. The eight antennas are located symmetrically above and below the horizontal midplane and are directed at the midplane for central resonance. Effective heating is observed at line-averaged densities up to $3.0 \times 10^{19} \text{ m}^{-3}$. The inside launch system launches the extraordinary (X) mode and has demonstrated central heating up to a density of $6.0 \times 10^{19} \text{ m}^{-3}$. Eight of the ten antennas have a smooth mirror which directs the power into the tokamak with a linear polarization and 12° half width at 15° with respect to the major radius. The orientation is such to drive current aiding the Ohmic current for normal operation. Two antennas which launch at ±30° with respect to the major radius have recently been installed. These antennas have grooved mirrors to generate the appropriate elliptical polarization corresponding to X mode. The mode converter has been upgraded to a single corrugated waveguide converter for $\text{TE}_{01} \rightarrow \text{HE}_{11}$. The vacuum section of the waveguide transmits power at densities greater than 700 MW/m² despite the fact that the waveguide passes through the fundamental resonance at the top of the machine. In order to prevent arcing at this point, the waveguide is split into two halves along its length and biased electrically to sweep out electrons before an avalanche can occur [2].

The launched ECH power is deposited exclusively in the electrons with a very narrow spatial deposition profile. Excellent agreement is found between experimental measurements of the deposition profile and ray tracing calculations of the absorption. In most cases, a complete accounting of the power can be made [2]. There is also good agreement between the ray tracing predictions of wave cutoff and experimental measurements [3]. Localized deposition requires control of the wave polarization and antenna pattern. At present, the location of the deposition can be changed only by changing the toroidal field of the tokamak.

2. L-MODE CONFINEMENT

The deposition profile of ECH is localized and relatively independent of plasma density and temperature, in contrast to the power deposition
of neutral beam injection (NBI). Therefore, some differences might be expected in the scaling of the global confinement with plasma parameters. However, the energy confinement for central heating agrees well with both the ITER-89 power law and linear offset scaling relations developed for the ITER design [4]:

\[
\tau_E = 0.048 I^{0.85} R^{1.2} a^{0.3} n^{0.1} B^{0.2} (A \kappa P)^{0.5}
\]

\[
\tau_E = 0.064 I^{0.8} R^{1.6} a^{0.6} n^{0.6} B^{0.35} (A \kappa)^{0.2} / P + 0.04 I^{0.5} R^{0.8} a^{0.8} A^{0.5} R^{0.6}
\]

where \(\tau_E\) is the confinement time in sec, \(I\) is the plasma current in MA, \(R\) and \(a\) are the major and minor radii in m, \(n\) is the electron density in \(10^{20}\) m\(^{-3}\), \(B\) is the toroidal field in T, \(A\) is the atomic mass number, \(\kappa\) is the elongation, and \(P\) is the total power in MW. The scaling relations were derived from NBI-heating confinement data from several machines. The confinement time data plotted against these scaling relations are shown in Fig. 1. The ECH dataset contains significant variations in \(n, I, A, B,\) and \(P\). A regression fit to a power law for these variables over the ECH dataset is not definitive, because of the small sample and a significant correlation between \(I\) and \(P\). This correlation tends to make the exponent of \(1/P\) equal to that of \(I\). The \(n\) and \(I\) scaling are in rough agreement with the ITER-89 power law, as are the scaling of \(A^{0.5} B^{0.3}\) found in previous analysis [2] (hydrogen isotopes only).

![Fig. 1. Comparison of energy confinement for central ECH with ITER L-mode scaling laws: (a) ITER-89 power law scaling, (b) ITER-89 offset linear scaling.](image-url)
One role suggested for ECH is modification of the pressure profile through localized off-axis heating. The full-width at half maximum of the deposition profile is about 0.05–0.10a for the inside launch system. In the framework of a theoretical treatment of profile modification [5], this profile can be considered a delta function at the resonance location. Assuming the transport is purely diffusive, the temperature profile should exhibit a sharp change in gradient at the heating location, as observed for off-axis ICH on JET [5]. In DIII-D, there is no evidence of this profile change, even in the extreme case of 75% of the total input power deposited at \( \rho = 0.4–0.5 \) [see Fig. 2(a)]. If purely diffusive transport is assumed, a dramatic drop in the electron heat diffusivity determined from power balance (\( \chi_{e}^{PB} \)) is required to support the observed electron temperature profile as shown in Fig. 2(b).

![Fig. 2. (a) Electron temperature profiles from ECE and Thomson scattering for Ohmic and ECH times. The calculated ECH deposition profile is also shown. (b) Profiles of electron and ion heat diffusivities determined by power balance analysis (\( \chi_{e,i}^{PB} \)) and the electron diffusivity determined by a time dependent method (\( \chi_{e}^{TD} \)). Experimental conditions: \( B = 1.8 \, \text{T}, \, I = 0.9 \, \text{MA}, \, n(0) = 4.0 \times 10^{19} \, \text{m}^{-3}, \, P_{ECH} = 1.2 \, \text{MW} \).](image)

Because this radial dependence of \( \chi_{e} \) seems unphysical, a different method [6] of determining \( \chi_{e} \) was applied to the data. When the auxiliary heating power is switched on, the electron temperature profile changes to a new stationary state. The energy balance equation during the temperature rise can be expanded assuming \( T_e(\rho, t) = T_0(\rho) + T_1(\rho, t) \), where \( T_0(\rho) \) is the equilibrium temperature profile during ECH. If the perturbation to
the energy sources and sinks other than conduction can be neglected over the time period of interest, the first order energy balance is

\[ \frac{3}{2} \frac{\partial}{\partial t} (nT_i) = \nabla \cdot (n\chi_e^{TD} \nabla T_i) \]

In the experiment, \( n \) is independent of time and \( T_i(p, t) \) is measured, so this equation can be solved for \( \chi_e^{TD}(\rho) \) [see Fig. 2(b)]. Note that there is no unusual behavior near the deposition location. Assuming \( \chi_e^{TD} \) is the true \( \chi_e \), the energy balance can be solved for the “auxiliary power” profile (i.e., the apparent input power other than Ohmic) required to support the observed electron temperature profile against the known losses. This is shown in Fig. 3(a) along with the ECH deposition profile calculated from ray tracing. The magnitude of the launched ECH power is sufficient to support the observed temperature, but the deposition must be shifted significantly inward to match the temperature profile. It is important to remember that the deposition profiles from ray tracing are well supported experimentally, and there is no cyclotron resonance inside \( \rho = 0.4 \) for this plasma. To get such an inward shift of the deposition profile by direct absorption, waves would have to propagate through a layer which is predicted to absorb most of the power, then be damped by some anomalous process in the center. This explanation is clearly unsatisfactory. If the difference between the “auxiliary power” and the rf deposition is cast in the form of a convective term \( \frac{3}{2} nTV \), the inward velocity required to complete the power balance is shown in Fig. 3(b). Because the density profile is unchanged during the ECH pulse, it is unlikely that the inward flow can be explained by a particle pinch. No theoretical explanation has been identified to explain this apparent inward transport of energy.

3. **H-Mode Confinement**

Energy confinement in the H-mode regime is accessible with ECH as the sole auxiliary heating [7]. The parameter regime in which H mode with ECH only is achieved has been extended to 2.0 T with both the inside and outside launch systems. These H modes, lasting longer than 600 msec, are limited only by the pulse length of the ECH system. The phenomenology of the H mode (\( H_\alpha \) decrease, density and stored energy rise, increase in edge poloidal rotation) is the same in all respects as the H mode generated by NBI. The variation of the power threshold with toroidal field is shown in Fig. 4(a). The relation \( P_{thres} \propto B \) was first found for NBI H modes. Also shown in Fig. 4(b) is the threshold power versus resonance location. The solid line is the expected dependence due to the changing toroidal
Fig. 3. (a) The volume-integrated "auxiliary power" profile (see text) and ECH deposition profile. (b) The convective velocity implied by the difference in the two curves in (a) when a term of the form $\frac{5}{2}nTV$ is added to the energy balance equation.

field. No evidence has been observed of a lower H mode power threshold with edge heating like that reported on the JFT-2M tokamak [8].

The energy confinement in H mode with ECH is the same as with NBI when the fraction of the stored energy contributed by the fast ions is subtracted for NBI discharges. The database is quite small, but the ECH H mode confinement appears to scale linearly with current as seen in the NBI H mode. A series of experiments to compare equal power ECH and NBI H mode transport in the same discharge has been carried out. The heat diffusivities from the power balance analysis are shown in Fig. 5. For the case of the NBI, the ion diffusivity is larger than the electron diffusivity across most of the plasma. For these discharges the NBI power is almost equally split between electrons and ions, with the ion fraction slightly larger everywhere. Despite this, the electron temperature is everywhere higher in the region analyzed, which results in the lower electron diffusivity. The diffusivities are also shown for the ECH case. As discussed above for the L mode transport, the electron diffusivity drops dramatically inside the resonance location ($\rho = 0.45$). The time dependent analysis method used to determine $\chi_e$ in the L mode case will be modified to handle the time-dependent density profiles of the H mode. The ion diffusivity is only slightly lower for the ECH case compared to the NBI case. The fact that the confinement times are roughly the same for the two cases and the power to the electrons is much higher in the ECH case
would indicate the true $\chi_e$ for the ECH case must be considerably larger than $\chi_i$ in this case. The reversal of $\chi_e$ and $\chi_i$ as the larger diffusivity between ECH and NBI could be evidence of an unfavorable temperature dependence in $\chi$, but a larger database is required to support such a conclusion.

4. INSTABILITY CONTROL

It is normally assumed that some enhancement of confinement above the L-mode level is needed to make a tokamak reactor viable. For the next generation of machines, operation with H-mode confinement plasmas appears to be a leading option for the required confinement in steady state. Long-pulse operation (10 sec) in H mode has been demonstrated in DIII-D with nearly stationary conditions [9]. The key to control of the density and impurity accumulation during these H modes is control of the edge-localized mode (ELM) frequency and amplitude. If the ELM frequency is too small, impurity accumulation leads to radiation collapse or disruption. If the ELM frequency and amplitude are too large, the confinement enhancement is lost.
On DIII-D, different types of ELMs occur. The “giant” ELMs are perhaps the best understood. These ELMs occur during H mode with high auxiliary heating power. A giant ELM occurs when the edge pressure gradient reaches the first stability limit for ballooning modes [10]. At present, the shape control capability of DIII-D is used to change the ELM behavior through changes in the edge shear. The changes in the ELM behavior are consistent with expectations based on calculations of the first stability limit [11]. It is reasonable then to expect that modification of the edge pressure profile will also affect changes in the ELM behavior. Wave absorption with the inside launch system can be as high as 25% for a single pass, even at the low temperatures and densities found at the edge. Therefore, experiments to determine the effects of edge heating were performed on NBI-heated H-mode discharges with giant ELMs. In Fig. 6, the change in ELM frequency is shown as the resonance is moved around the separatrix location. The range of ELM frequencies observed is about a factor of three. With the resonance located outside the separatrix, the ELMs occur less frequently, while placing the resonance just inside the plasma leads to more frequent ELMs. When the resonance is sufficiently far into the plasma, the effect of the ECH is the same as the addition of NBI of equal power. Some increase in stored energy is observed as the time between ELMs increases, perhaps due to the increasing density. No corresponding increase in impurity line radiation or radiated
power is observed. There is no direct measurement of edge pressure gradient modification, but the effect of ECH is consistent with the previous conclusion that the giant ELMs occur when the edge ballooning limit is reached. For future machines which do not have flexibility in plasma shaping, edge ECH may be a useful tool for controlling the ELM behavior for steady-state H-mode operation.

![Graph showing effects of edge ECH on giant ELMs]

Fig. 6. Effects of edge ECH on giant ELMs. At right is shown the ratio of the average time between ELMs with and without ECH. At left are typical time histories of the divertor $D_\alpha$ light and ECH power. Experimental conditions: $I = 1.0$ MA, $\bar{n} = 6 \times 10^{19} m^{-3}$, $P_{NBI} = 5$ MW.

5. ELECTRON CYCLOTRON CURRENT DRIVE

Fully non-inductive current drive would allow steady-state operation of a tokamak. Electron cyclotron current drive (ECCD) is one scheme which is predicted to drive current at plasma parameters projected for tokamak reactors. Experiments with ECCD have been carried out for inside-wall limited discharges in DIII-D with the inside launch system. The target conditions provide strong single-pass absorption near the magnetic axis in thermal plasmas. For these experiments, the total toroidal current is held constant. Because the available ECH power is not sufficient to drive the total current, the plasma resistivity must be calculated in order to determine the driven current. The experimental profiles of electron temperature, density, and $Z_{eff}$ are used to calculate the neoclassical resistivity and the bootstrap current. The validity of these calculations is
verified by applying it to the Ohmic phase before the ECH pulse. The loop voltage $V^*$ required to drive the entire current inductively during the heating pulse is derived assuming the voltage is fully penetrated. The difference between $V^*$ and the actual loop voltage, divided by the plasma resistance, is assumed to be the driven current $I_{rf}$. Driven currents up to 70 kA have been observed. Both co- and counter-injection experiments have been performed. The systematic variation of the data indicates that the rf-generated current is driven in the same direction as the launched wave, as expected for inside launch. An example of this is shown in Fig. 7. The loop voltage and ECH power as a function of time are shown in Fig. 7(a) for a discharge where the ECCD is in the direction of the existing current. Shown in Fig. 7(b) are the same traces for a discharge where the ECCD opposes the existing current. In both co- and counter-injection, the offset in the loop voltage relative to $V^*$ is in the expected direction.

\[
\begin{align*}
1 = 300 \text{ kA} & \quad I_{rf} = 62 \pm 16 \text{ kA} \\
1 = 300 \text{ kA} & \quad I_{rf} = 29 \pm 14 \text{ kA}
\end{align*}
\]

Fig. 7. Time history of the loop voltage and ECH power for (a) co-injection ECCD and (b) counter-injection. The quantity $V^*$ is the predicted loop voltage for purely Ohmic current based on the change in the plasma resistivity. Experimental conditions: (1) $T(0) = 3.4 \text{ keV}$, $n(0) = 1.4 \times 10^{19} \text{ m}^{-3}$, $P_{ECH} = 0.9 \text{ MW}$; (b) $T(0) = 1.8 \text{ keV}$, $n(0) = 1.4 \times 10^{19} \text{ m}^{-3}$, $P_{ECH} = 1.0 \text{ MW}$.

Several ECCD discharges have been modeled using a ray tracing code with an rf current drive model [12] and a quasi-linear Fokker-Planck calculation. The ray tracing assumes damping on a Maxwellian with the bulk temperature, while the Fokker-Planck code calculates the absorption self-consistently on the deformed distribution. Both calculations include the
effects of the trapped particle population. The experimental and modeling results are shown in Fig. 8. The driven current is expected to increase monotonically with $T/n(Z_{\text{eff}} + 5)$ at constant power [13] as seen in the figure. Notice that the experimental points are equal to or greater than the predicted current drive in the absence of an electric field. This is in sharp contrast to previous results on the CLEO tokamak [14] and the Wendelstein VII-AS stellarator [15], where the current drive was reported to be a factor of 2–4 lower than predicted. The current drive modeling for both machines showed a sensitivity of the current drive efficiency to direct losses of the current-carrying electrons [14,16]. In DIII-D, if the particle confinement time can be assumed to be greater than or equal to the energy confinement time, the ratio of confinement time to the slowing-down time for the current carriers is large. The superthermal electron transport may lead to broadening of the driven current profile, but degradation of the current drive efficiency due to losses is not expected. The current drive efficiency is known to be enhanced [13,17] by the presence of a dc electric field. At the electric field strengths seen in the experiments, the Fokker-Planck calculations show a 50–100% enhancement of the current drive efficiency. On account of this effect, an experimental determination of the current drive efficiency in the absence of an electric field is not yet possible. Further experimental analysis and more extensive modeling are in progress.

![Graph](image)

Fig. 8. Current drive by ECCD versus $T/n(Z_{\text{eff}} + 5)$. Analytic calculations without trapping or electric field corrections indicate the current drive efficiency should increase with this parametric dependence. Experimental points and Fokker-Planck modeling predictions with and without dc electric field are shown.
SUMMARY

Heating with electron cyclotron waves has demonstrated capabilities important for present and future tokamaks. These include bulk heating, instability control, and non-inductive current drive. The present DIII-D database for central ECH agrees with the empirical global confinement scaling derived from NBI-heated discharges. When ECH is applied off-axis, the electron temperature remains peaked much more than a purely diffusive transport model would predict. The power balance indicates an inward transport of the absorbed power. Control of the ELM behavior with edge ECH has been demonstrated. This capability may be important for steady-state operation in H mode for the next generation of machines. Non-inductive driven currents up to 70 kA have been observed. The driven current is greater than or equal to theoretical predictions. This enhancement is consistent with modeling which includes the residual dc electric field. Experiments in each of these areas will be extended to higher power levels with the 110 GHz system presently under construction.

REFERENCES


DISCUSSION

M.C. ZARNSTORFF: In your off-axis heating experiment, where you analysed the turn-on of the ECH, did you observe a phase delay between the heating on-axis and the heating at the resonant layer (consistent with your hypothetical convective velocity)?

T.C. LUCE: The local value of $dT_e/dt$ is greatest at the resonance position. The magnitude of the inferred convective velocity is consistent with the characteristic time for the central $T_e$ to reach equilibrium.

A. GIBSON: In the cases of anomalous inward transfer of off-axis ECH deposition, can you exclude the possibility that the ECH provokes a significant change in the Ohmic heating deposition profile?

T.C. LUCE: The central Ohmic power inferred from the electron temperature is too small by a factor of 4 to account for the measured $T_e$. The profile also exists for 0.5 s — sufficient time for relaxation of the current profile changes consistent with the observed change in internal inductance.

R.J. GOLDSTON: This is a very nice example for the ‘resiliency’ of the electron temperature profile. It is a great advantage that you could heat the electrons and use the ions as a sink term. Could you comment on the heating efficiency of ECH as a function of deposition radius?

T.C. LUCE: In this case the confinement in L-mode is equal to the ITER-89 power law value despite heating at $\rho = 0.5$.

V.E. GOLANT: Do you have any data on the influence of electron cyclotron current drive on the electron temperature profile? And what is the density dependence of the current drive efficiency?

T.C. LUCE: For electron cyclotron current drive aiding the existing plasma current, we usually see electron temperature profiles more peaked than the Ohmic
profiles. Since most of our data are derived at maximum power, we cannot separate
the dependence of current efficiency on T and n. We do see that the dependence on
T/n(Z + 5) is as expected from numerical calculations.

B. COPPI: You have justified the existence of a 'profile consistency' condition
for the electron temperature in the presence of electron cyclotron heating by postu-
lating an energy inflow term in the electron thermal energy balance equation. How
sure are you that this condition cannot be described instead by a non-linear electron
thermal conductivity?

T.C. LUCE: Our purpose in showing the inward transport of energy as a
convective flow was to show that the velocity would have a strictly linear depen-
dence, 10 (r/a)(m/s) out to the resonance location. We are not yet claiming the profile
results from an actual flow. A non-linear diffusivity (such as an unfavourable
temperature dependence of \( \chi \)) is a leading candidate to explain the profiles.
LOWER HYBRID CURRENT DRIVE AND HIGHER HARMONIC ICRF HEATING EXPERIMENTS ON JT-60


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Abstract
LOWER HYBRID CURRENT DRIVE AND HIGHER HARMONIC ICRF HEATING EXPERIMENTS ON JT-60.

The paper describes the latest results and analysis of LHCD and ICRF heating experiments and beam induced instabilities on JT-60. A new multijunction LHCD launcher recorded current drive efficiency $\eta_{CD}$ ($= \eta_{CD} = \frac{R_{p} I_{p}}{P_{LH}}$) of $3.4 \times 10^{19}$ m$^{-2}$ A$^{-1}$ W and extended current drive products ($= \eta_{CD} P_{RF}$) up to $12.5 \times 10^{19}$ m$^{-2}$ MA with $\sim 4.5$ MW LHCD power. Combination with neutral beams showed that the interaction between ions or electrons and lower hybrid waves is well explained by simple lower hybrid dispersion and showed evidence of $N_{e}$ upshift. A simplified $N_{e}$ upshift model explains experimental $\eta_{CD}$ very well and predicts the achievement of the value necessary for ITER. Experimental results on third harmonic ICRF heating in combination with NBI heating and second harmonic minority heating are presented. In both cases, strong central heating by tail ion acceleration with higher harmonic ICRF waves occurs and giant sawteeth appear. In the former case, the sawtooth period reaches...
six times the energy confinement time and the energy confinement time is enhanced by 20% owing to fast ions. In the latter case, giant sawteeth appear even in the high density and low $q$ regime ($n_e \leq 6.5 \times 10^{19} \text{ m}^{-3}$ and $q_{\text{eff}} \geq 2.6$) with relatively small threshold ICRF power (~2 MW).

1. LOWER HYBRID CURRENT DRIVE

The most advanced confinement system for fusion energy research at present is a tokamak. Pulsed operation due to inductive plasma currents diminishes the attractiveness of the tokamak. Recent progress in non-inductive current drive studies has yielded information relevant to steady state operation of a tokamak reactor. LHCD currents of 2 MA and improvement of $\eta_{\text{CD}}$ with temperature increase were demonstrated on JT-60 [1, 2]. Bootstrap current opened new possibilities for economical steady state operation [3, 4]. Key issues for the development of a steady state tokamak reactor are to attain high $\eta_{\text{CD}}$, above $5 \times 10^{19} \text{ m}^{-2} \cdot \text{A} \cdot \text{W}^{-1}$, and a large current drive product (CDP) of $\sim 500 \times 10^{19} \text{ m}^{-2} \cdot \text{MA}$. This paper presents the recent progress of LHCD experiments on JT-60.

1.1. High power and high efficiency current drive

Present LHCD experiments were done in a lower X point divertor configuration with a major radius of 3.14 m and a minor radius of $\sim 0.7$ m. The lower hybrid range of frequency system for JT-60 consists of three launchers with 8 MW source power each at around 2 GHz frequency. One of them was replaced by the multijunction LHCD launcher [5] and hence two LHCD launchers, including the conventional one, were available in the present experiment. The frequency tunability of our system from 1.74 to 2.23 GHz provided a flexible tool for investigating wave–particle interaction.

A typical example of high power LHCD is shown in Fig. 1, where 4.2 MW of LH power was injected into a 1.5 MA discharge with line average density $\bar{n}_e = 2.7 \times 10^{19} \text{ m}^{-3}$. The loop voltage $V_l \sim 0$ and the value of CDP was $12 \times 10^{19} \text{ m}^{-2} \cdot \text{MA}$. The radiation power and the power to the divertor plates increased gradually but the ratios of both powers to the total absorbed power were about 40% and kept almost constant. Power flows to both electron and ion sides and the temperature of the divertor plate are also shown in Fig. 1, and it is seen that there is no anomalous heat concentration on the divertor. The half-width of the heat flux to the divertor plates was a few centimetres and was similar to that in Ohmic or NBI heated plasma [6]. These results suggest that current carrying fast electrons accelerated by the lower hybrid wave (LHW) do not have any inherent detrimental effects on the divertor plates. Hot spots, which were often observed in limiter LHCD experiments, could be reduced with divertor operation. The maximum CDP obtained was $12.5 \times 10^{19} \text{ m}^{-2} \cdot \text{MA}$, which is $\sim 10$ times larger than previous results.
Previous experiments in JT-60 had conducted current drive efficiency \( \eta_{CD} \propto (T_e)/(Z_{eff} + 5) \), with \( \eta_{CD} = 2.6 \times 10^{19} \text{ m}^{-2} \cdot \text{A}^{-1} \cdot \text{W}^{-1} \) [2]. Further improvement is expected with the optimization of \( N_1 \) (the refractive index parallel to the magnetic field). A simple theory predicts that a sharper and lower \( N_1 \) spectrum with higher directivity gives better \( \eta_{CD} \) [7, 8]. This prediction prompted the introduction of the multijunction launcher with a 4 \( \times \) 24 wave guide array [5]. The width of the spectrum \( \Delta N_1 \sim 0.5 \) is approximately half of the conventional one and the directivity is higher than the conventional one. An improvement of \( \eta_{CD} \) by 30–40% was obtained with the multijunction launcher, as shown in Fig. 2. The qualitative dependence of \( \eta_{CD} \) on the \( N_1 \) spectrum is well explained by the theory of Fisch [5], even in the presence of what is known as spectral gap, which suggests that some fraction of LH power keeps its original \( N_1 \) almost unchanged. Studies of wave–particle interaction, hard X ray measurements and their analysis using a Fokker–Planck equation indicate that a substantial fraction of LH power diffuses to a higher \( N_1 \) wave and fills the spectral gap, as discussed later. Both \( \eta_{CD} \) and CDP approached the values required for ITER and FER.

1.2. Wave–particle interaction

It is extremely important to explain the obtained LHCD results with simple conventional theories, which would confirm the extrapolation of the results to future devices. In the case of the LHW, there are certain questions that remain to be solved,
for example the spectral gap problem. The difficulty of investigating wave physics in tokamak plasmas is also an obstacle to solving the physics of LHW propagation and absorption. We have scanned the frequency of the LHRF system, the density and the beam energy of NBI ($E_B$) to study wave propagation and absorption.

Figure 3(a) shows typical examples from the wave–particle interaction study for three different frequencies where a small NBI power of ~1 MW with $E_B = 65$ keV was injected from 4.0 to 8.0 s and the density was ramped up linearly during the same period. The second section from the top shows the intensities of 150 keV fast ions accelerated by the LHW near the central region and the third section the intensities of the ECE generated by the fast electrons ($=1.5\omega_{ce}$) corresponding to each frequency. The fast ion intensity became appreciable when the ECE signal began to decrease. This critical density, $\bar{n}_{ec}$, increased when the frequency increased. Using $\bar{n}_{ec}$ it is possible to determine the dispersion relation of the LHW experimentally. If Landau damping of the LHW is assumed, the ratio of the wavenumbers $k_x/k_1$ is estimated from $E_B$ and $T_e$ [9, 10]. The data points of $k_x/k_1$ corresponding to each frequency are plotted in Fig. 3(b). The cold electrostatic dispersion relation of the LHW is written as:

$$\omega^2 = \omega_{LH}^2 (1 + (k_1/k_x)^2 (m_p/m_e)/\gamma)$$

where $\omega_{LH}$ is the lower hybrid frequency, $m_p/m_e$ is the mass ratio of a proton to an electron and $\gamma = \Sigma (n_i/n_e) (Z_i^2/A_i)$. Equation (1) is represented in Fig. 3(b), using the corresponding critical density. The experimental data points fit the calculated dispersion curves very well. The point with $E_B = 40$ keV is also exactly on the
FIG. 3. (a) Typical time evolutions of discharges in the study of wave–particle interaction of the LHW. (b) Plots of the dispersion relation of the LHW. Solid circles indicate experimental points and the four curves correspond to the dispersion relation of the critical density of each dotted point, where $B_T = 4.5\; T$, $T_e = 3\; keV$; $\tilde{n}_c = 1.2$ (broken line), 1.7 (solid line), 2.2 (dash-dotted line) and 2.4 (dotted line) $\times 10^{19} \; m^{-3}$, and $n_{ec}(0) = 1.5\tilde{n}_c$ is assumed.

dispersion curve of the corresponding density. Thus it is found that the LHW in a tokamak plasma follows the simple linear dispersion relation.

1.3. Spectral gap

Another interesting indication from the study of wave–particle interaction (Fig. 3) is that the presence of the spectral gap does not affect the wave absorption, not only by electrons but also by ions. Actually, the change of the injected $N_B$ value from 1.2 to 3.0 did not change the critical density, as expected from Eq. (1), and even in the presence of the spectral gap almost all the LH power was absorbed by electrons or ions. Whether the ions or the electrons absorb the LH power is determined by which phase velocity is nearer to the edge of the Maxwellian distribution or the injected beam velocity. These results suggest that the gap between the wave spectrum and Maxwellian electrons or beam ions was filled with the $N_B$ upshifted wave and not with the diffused particles due to instabilities, since the spectral gap effects are exactly equal for both electrons and ions.

To further confirm the above, a low frequency wave with a lower $N_B$ value (1.74 GHz, $N_B = 1.6$--2.6) and a high frequency wave with high $N_B$ (2.23 GHz, $N_B = 2.2$--3.6) were injected simultaneously, with other conditions the same as in Fig. 3(a). A simplified model of this experiment is shown in Fig. 4(a). When the density is slightly higher than the critical density of 1.74 GHz and less than that of 2.23 GHz (at $\sim 5.8\; s$ in Fig. 4(b)), LH power is almost completely absorbed by beam
ions in the case of 1.74 GHz only (broken lines). However, when the 2.23 GHz is added in the electron spectral gap region and forms a bridge, some part of the 1.74 GHz power can couple to the electrons before coupling to the ions. This is clearly shown by the solid lines in Fig. 4(b). The fast ion flux decreases and the ECE signal increases during the period from 5 to 6 s. It is found that the N\textsubscript{i} upshift is necessary to fill the spectral gap and some fraction of the power, of course, keeps its original spectrum almost unchanged [11]. These facts can be explained by a multipath ray tracing model with an N\textsubscript{i} upshift mechanism such as the toroidal effect [12].

1.4. Discussion

Ray tracing and the calculation of the change of N\textsubscript{i} of the LHW in JT-60 are carried out using the Bonoli LH code combined with the ACCOME code [13]. Since some fraction of the power shifts to a higher N\textsubscript{i} wave through multipath propagation owing to the toroidal effect to fill the gap, an upshifted N\textsubscript{i} spectrum as shown in Fig. 5(a) is assumed for simplicity of calculation, where N\textsubscript{1} and N\textsubscript{2} are the minimum and maximum of the original N\textsubscript{i} spectrum, N\textsubscript{3} is the maximum of the upshifted N\textsubscript{i} which is determined from the power balance between the input and absorbed powers, and E\textsubscript{1}, E\textsubscript{2} and E\textsubscript{3} are the corresponding energy. T\textsubscript{SG} and T\textsubscript{tail} are the effective temperatures in the spectral gap region and the original energy region which are estimated from diffusion coefficients due to the LHW and collision. After this simplification, the absorbed power in the spectral gap P\textsubscript{SG} is approximated as:

\[ P_{SG} = A \int_{v_{i1}}^{v_{i2}} \frac{(\nu_{i}^{2}/\tau_{i}) f_{e}(\nu_{i}) dv_{i}} \]
where $A = (2\pi)^2 R_p(r) \Delta r n_e m_e$, $\tau_s$ is the collision time, $\nu_i = c/N_i$ ($i = 1-3$), and $\langle r \rangle$ and $\Delta r$ are the average radial point and radial width of the absorption layer. The absorption power in the original spectrum region $P_{OR}$ is obtained by integrating from $\nu_2$ to $\nu_1$ in Eq. (2). The currents $I_{SG}$ and $I_{OR}$ are easily estimated by the integration of $n_e f_d (\nu_i)$ for the corresponding energy ranges and the absorption layer. The calculated current drive efficiency $\eta_{CD}$ is given by $(I_{SG} + I_{OR}) \bar{n}_e R_p / P_{LH}$. The results of the calculation are shown in Fig. 5(b) with the experimental $\eta_{CD}$ versus $\langle T_e \rangle / (5 + Z_{eff})$. The calculation could well explain the temperature dependence of $\eta_{CD}$. Thus the dependence of $\eta_{CD}$ on $N_f$ and $T_e$ obtained previously [2, 5] is well explained with this simplified model, which is based on the conventional theories. If this model is applied to ITER, it predicts $\eta_{CD} \approx 5$, which is the value necessary for ITER steady state operation.

2. HIGHER HARMONIC ICRF HEATING AND BEAM ACCELERATION

ICRF heating accompanied by beam acceleration in the higher harmonic regime is being investigated in JT-60. Higher harmonic heating is attractive for producing energetic ions in a high density discharge more efficiently than fundamental resonance heating [14]. In all cases of the present experiments, hydrogen is the resonant species for ICRF heating and a hydrogen neutral beam is employed. The frequency of ICRF heating is 131 MHz and the toroidal field is selected to place the pertinent resonance layer near the plasma centre, i.e. 4.3 T for second harmonic heating and 3 T for third harmonic heating. The high $k_i$ mode ($(\pi, 0)$ mode) is chosen as the antenna phasing mode, which was found to be efficient in bulk plasma heating and beam acceleration [15].
2.1. Third harmonic heating

Experiments on third harmonic heating were carried out with the following discharge conditions: $B_t = 3$ T, $I_p = 0.7-1.4$ MA, $q_{\text{eff}} = 7-3.2$, $n_e = (1.7-4) \times 10^{19}$ m$^{-3}$. Ohmic target plasmas are found to be insufficient for absorption of the third harmonic waves even with $(\pi,0)$ mode. Therefore we added neutral beam power in order to increase absorption. Both hydrogen and helium discharges were employed. A typical time evolution of third harmonic heating at $I_p = 1$ MA and $q_{\text{eff}} = 5$ is shown in Fig. 6(a). The ICRF power, $P_c$, was 2.3 MW and NBI power was varied during the shot (i.e. absorbed NBI power 5.4 MW $\rightarrow$ 10.3 MW). The central electron temperature increases significantly and giant sawtooth oscillation is observed in the combined ICRF and higher power NBI heating phase. The period of the sawteeth reaches 410 ms, which is about six times the energy confinement time of this discharge (73 ms). The plasma stored energy also oscillates with giant sawtoothing. However, the variation of the plasma stored energy after the sawtooth crash is about 5%. The energy confinement time of this discharge is about 20% larger than that of the Goldston L mode scaling.

Remarkable beam acceleration is observed during the combined heating. Figure 6(b) shows the hydrogen energy spectra from the plasma core during the combined heating phase and the NBI heating only phase. The broken line is a fitting curve obtained from a one dimensional Fokker–Planck code using measured plasma
parameters in the centre and the NBI power density calculated by a Monte Carlo code. From this, we can estimate ICRF power density, \( p_{\text{IC}} \), and tail ion stored energy density, \( w_{\text{tail}} \). In the case of Fig. 6(b), we obtain \( p_{\text{IC}} \sim 1 \text{ MW/m}^3 \). The radius of the ICRF power deposition profile is then estimated to be \( \sim 0.2 \text{ m} \). The incremental fast ion stored energy by ICRF, \( \Delta W_{\text{tail}} \), can be calculated from

\[
\Delta W_{\text{tail}} = (w_{\text{tail}}(\text{NB + IC}) - w_{\text{tail}}(\text{NB}))P_{\text{IC}}/P_{\text{IC}}
\]

Figure 7 shows a comparison between experimental and calculated values of the incremental stored energy by ICRF as a function of the electron density. Considering possible errors of the calculated values (+30%), the results agree well. We can conclude that the enhancement of the energy confinement time during the combined heating is mainly due to fast ions accelerated by third harmonic ICRF waves.

2.2. Second harmonic hydrogen minority heating

With second harmonic hydrogen minority heating in a helium discharge, a good incremental energy confinement time (\( \sim 110 \text{ ms} \)) is obtained for Ohmic target plasmas in a wide range of density \( (2.5 \times 10^{19} \text{ m}^{-3} \leq n_e \leq 7 \times 10^{19} \text{ m}^{-3}) \) [15]. This value is considerably better than the incremental energy confinement time of the pure hydrogen second harmonic heating (\( \sim 80 \text{ ms} \)). Giant sawteeth appear during second harmonic minority heating with Ohmic plasmas even in high density and low q regimes (\( n_e \leq 6.5 \times 10^{19} \text{ m}^{-3} \) and \( q_{\text{eff}} \geq 2.6 \)). A threshold ICRF power for producing giant sawteeth is around 2 MW. These results are in contrast with the JET
results [16], where fundamental minority heating is employed and sawtooth stabilization is found at electron density \( n_e \) less than \( 3 \times 10^{19} \) m\(^{-3} \). It is found that the largest sawteeth appear with Ohmic target plasmas. A small amount of additional NBI power (~ 1 MW) stops giant sawtoothing dramatically. Sawtooth periods as a function of \( P_{NB}/n_e \) are shown in Fig. 8. A Fokker–Planck calculation indicates that energetic ions (e.g. 0.5 MeV) decrease with increasing NBI power, since NBI heating is regarded as a lower energy particle source compared with ICRF heating [17]. Thus energetic ions play an important role in sawtooth stabilization [18]. It is calculated that second harmonic minority heating (and third harmonic still more) is much more efficient in producing energetic ions in high density discharges than fundamental minority heating.

2.3. Beam induced ICRF instability during NBI heating

Ion cyclotron waves were excited during high power (~ 12 MW) NBI heating. The frequency spectrum showed sharp peaks of integer ion cyclotron harmonics from the second harmonic wave and the peak spacing scaled linearly with \( B_t \) and corresponded to \( \omega_{CH} \) at the outermost plasma edge. The number of the harmonics depended on the electron density and beam power. During wave excitation, the charge exchange spectrum showed that beam ions were accelerated perpendicularly and the preferential heating of beam ions occurred locally at the plasma edge. No remarkable changes in bulk ion temperature and stored energy were seen in either limiter or divertor plasmas. The features mentioned above are explained qualitatively by lower hybrid instability [19].

FIG. 8. Sawtooth period as a function of \( P_{NB}/n_e \) for combined second harmonic ICRF and NBI heating and NBI heating only, in helium discharges.
3. SUMMARY

Substantial progress was made in LHCD experiments on JT-60 and their analysis presented the promising prospect of non-inductive current drive in a future steady state tokamak reactor. (1) $\eta_{\text{CD}}$ of $3.4 \times 10^{19} \text{ m}^{-2} \cdot \text{A} \cdot \text{W}^{-1}$ was achieved by using the multijunction launcher and is well on the way to the value necessary for ITER ($5 \times 10^{19} \text{ m}^{-2} \cdot \text{A} \cdot \text{W}^{-1}$). (2) Progress in the CDP by one order of magnitude ($\leq 12.5 \times 10^{19} \text{ m}^{-2} \cdot \text{MA}$) compared with previous LHCD experiments was obtained. Fast electrons produced by LHCD have no inherent bad effects on the divertor and the hot spots often observed in limiter experiments could be reduced with divertor operation. (3) The interaction of the LHW with beam ions from NBI provided evidence that the behaviour of the LHW follows linear dispersion and the spectral gap is filled with $N_f$ upshifted waves through multipath propagation. (4) The results of LHCD on JT-60 are well explained by a simple model based on conventional theories that also predicts the achievement of the value of $\eta_{\text{CD}}$ required for ITER.

The effectiveness of third harmonic ICRF heating in combination with NBI heating has been demonstrated. Strong central heating accompanied by giant sawteeth is observed. With the combined heating (e.g. $P_{\text{IC}} \sim 2.3 \text{ MW}$, $P_{\text{NB}} \sim 10.3 \text{ MW}$ and $n_e \sim 2.8 \times 10^{19} \text{ m}^{-3}$) the energy confinement time is enhanced by about 20% compared with the Goldston L mode scaling, owing to fast ions. Second harmonic hydrogen minority heating in helium discharges produces giant sawteeth even in high density and low q regimes ($n_e \leq 6.5 \times 10^{19} \text{ m}^{-3}$ and $q_{\text{eff}} \geq 2.6$) with relatively small threshold ICRF power ($\sim 2 \text{ MW}$).

Beam induced ICRF instabilities and their coupling with the beam ions were observed during high power NBI heating.

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DISCUSSION

F. WAGNER: Does the enhancement of confinement by 20–30% over L mode, which you quoted, refer to a standard L mode scaling relation such as the Goldston or ITER scaling, or does it refer to the JT-60 L mode databank?

T. IMAI: We employed the Goldston L mode scaling as a standard L mode confinement.

M. PORKOLAB: I would like to comment on your plot of current drive efficiency, $\eta_{CD}$, versus temperature (Fig. 2). With the exception of the JT-60 data, I prepared the original plot of $\eta_{CD}$ versus temperature in 1984; this was published in the IEEE Journal in 1984 and in the book on RF heating edited by Granatstein and Colestock (1985). However, I used $T_e(0)$, not $\langle T_e \rangle$ as you indicated in your graph. Hence, you should use $T_e(0)$ for the JT-60 data to be fair. I also believe now that the apparent strong temperature dependence of $\eta_{CD}$ in the old experiments on smaller machines was dominated by significant electron loss. However, in larger machines this effect is reduced and JT-60 should be approaching the ideal efficiency.

T. IMAI: Some of the data, such as those from ASDEX and PLT, in addition to JT-60, were plotted as a function of $\langle T_e \rangle$ in Fig. 2. Some data were plotted using $T_e(0)$ with a large error bar, as we could not get an accurate value for $\langle T_e \rangle$ from them. I also agree to some extent that a plot of $\eta_{CD}$ versus $T_e(0)$ is a good way of visualizing the phenomena; however, the hard X ray profiles suggest that currents are driven in a wide area of the plasma cross-section in our present tokamaks. It was
for this reason that we plotted $\eta_{\text{CD}}$ as a function of $\langle T_e \rangle$. The important issue here is the experimental fact that $\eta_{\text{CD}}$ increases with electron temperature. As for the fast electron losses, data from small machines were excluded in our new plot.

W.M. NEVINS: In ITER the main problem with LH current drive lies in the accessibility of LH waves in the plasma core. This leads us to choose a relatively low phase velocity and yields hollow driven current profiles. The LH current drive figure of merit we calculate is about $0.3 \times 10^{20} \text{A} \cdot \text{m}^{-2} \cdot \text{W}^{-1}$ — similar to what has been achieved on JT-60. How can an upshift in $N_i$ improve this situation and yield an LH figure of merit of $0.5 \times 10^{20} \text{A} \cdot \text{m}^{-2} \cdot \text{W}^{-1}$?

T. IMAI: In ITER, where $\langle T_e \rangle \sim 20 \text{keV}$ is expected and LHCD will be the peripheral current drive due to strong Landau damping, there is no spectral gap and our simple model indicates that $\eta_{\text{CD}}$ can be estimated with a simple Karney–Fisch formula. Therefore, $\eta_{\text{CD}}$ depends strongly on $N_i$ at the absorption point. If the absorption point is located on the high toroidal field side, then by careful selection of $N_i$ and the launching position, $N_f^{\text{esc}}$ is reduced to a lower value and the simple formula gives $\sim 0.5 \times 10^{20} \text{m}^{-2} \cdot \text{A} \cdot \text{W}^{-1}$. If the effects of trapped particles and directivity of the $N_i$ spectrum are taken into account, $\eta_{\text{CD}}$ is reduced to some extent and might approach the value found in your modelling.

S. COHEN: Does your IR camera view the entire (toroidal) length of the divertor plate? Can you eliminate toroidally asymmetric losses of fast electrons during LHCD?

T. IMAI: We viewed only one particular portion of the divertor and so we cannot demonstrate toroidal symmetry of fast electron losses experimentally. Since fast electrons are forced to hit the divertor along the magnetic field lines and exist uniformly in each magnetic surface, it seems reasonable to expect toroidal uniformity of the fast electron loss to the divertor during LHCD. The divertor is quite effective in reducing the hot spot phenomena which were often observed in limiter LHCD discharges.

R.J. GOLDSTON: This is really a question for both you and Dr. Söldner. Magnetic turbulence theories predict fast electron loss rates $\sim (E/T_e)^{1/2}/\tau_{\text{eev}}$, so long as the fast electron drift orbits are not much larger than $\sim \rho_i$. On the basis of your results of good electron confinement, can you reject such models?

T. IMAI: No, we cannot reject the magnetic fluctuation model on the basis of our data. But as far as the confinement of fast electrons is concerned, the dominant loss mechanism is the slowing-down process in JT-60. Most of the power of the fast electrons is transmitted to the bulk plasma through collisions. The confinement time of fast electrons is well explained by the slowing-down time, the stored energy increases with power (as with other heating schemes) and the $\eta_{\text{CD}}$ value agrees well with the theoretical value.
THE LIMITER H-MODE WITH LOWER HYBRID CURRENT DRIVE

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Abstract

THE LIMITER H-MODE WITH LOWER HYBRID CURRENT DRIVE.

The H-mode has been achieved in limiter discharges on JT-60 with lower hybrid current drive. Suppression of electron density fluctuations was observed near the edge where a sharp density gradient was formed. Fast electrons (> 100 keV) which were generated further outside than half minor radius appeared to be effective in attaining the H-mode. Simultaneous application of RF power at two different frequencies, 1.74 + 2.23 GHz or 1.74 + 2.0 GHz, by which the H-mode was triggered in a reproducible manner, gave rise to brighter, hard X rays with higher photon energies than in the case of single frequency application.
1. INTRODUCTION

The H-mode has been observed in limiter discharges with lower hybrid current drive (LHCD) on JT-60 [1, 2]. The threshold LH power for the H-mode transition was as low as about 1.2 MW, in contrast to the threshold neutral beam (NB) power of about 16 MW in the outer or lower divertor configurations [3]. Other heating schemes, such as neutral beam injection (NBI) or ion cyclotron resonance heating (ICRH) using the second harmonic resonance with powers of up to 20 MW and 2 MW, respectively, never produced a limiter H-mode in JT-60.

Although the H-transition mechanism has not yet been clarified, shear in the edge radial electric field [4] and in the poloidal rotation [5] is discussed as a possible candidate with respect to H-transition and the suppression of edge fluctuations. Pronounced changes in radial electric field and poloidal rotation velocity were observed at H-transitions [6, 7]. The radial electric field has been found to become more negative, which does not contradict Shaing's model [8, 9]. Both E_r and \( dE_r / dr \) become more negative in the thermal barrier formed in JFT-2M [10]. The H-mode has been induced through poloidal rotation driven by radial currents independent of its polarity [11, 12]. Ion orbit loss which brings about non-ambipolarity is usually taken into account as a poloidal momentum source. On the other hand, our observation suggests that tail electrons generated near the edge by LHCD are responsible for the low power threshold.

2. H-MODE CHARACTERISTICS

LH power was applied via two launchers: a conventional 8 (toroidal) \( \times \) 4 (poloidal) grill and a multijunction type grill of 24 \( \times \) 4 phased array waveguides. Most experiments were conducted with LH waves at two different frequencies chosen from the values of 1.74, 2.0 and 2.23 GHz. The phase differences between adjacent waveguides were changed depending on the chosen frequencies in order to keep the peaks of the refractive index parallel to the magnetic field nearly constant at about 2.1 and 1.4, from the multijunction and the conventional launcher, respectively.

Figure 1 shows a typical time evolution of the limiter H-mode with LHCD. The characteristics of the limiter H-mode can be summarized as follows:

(a) Simultaneous application of LH powers at two different frequencies is favourable; the threshold LH powers were 1.2 and 2.0 MW for the combinations of 1.74 + 2.23 GHz and 2.0 + 2.0 GHz, respectively.

(b) There is a lower limit in the line averaged electron density of about \( 2 \times 10^{19} \text{ m}^{-3} \) to achieve the H-mode.

(c) The outer gap between the plasma edge and the limiters must be kept at values not less than 4 cm in order to produce a long H-mode phase.
FIG. 1. Time evolution of plasma stored energy evaluated by diamagnetics, injected LH power, line averaged electron density, one-turn loop voltage, Hα emission, ratio of soft X ray signals through Be foils, reflection coefficient of conventional launcher and C VI line signals along chords near the edge (r/a ~ 0.9) and through centre.

(d) A reduction in the edge density fluctuations at a density layer of n_e ~ 7 x 10^{18} m^{-3} was detected by a millimetre wave reflectometer, as is shown in Fig. 2(a).

(e) The formation of an electron density pedestal was confirmed by Thomson scattering as is illustrated in Fig. 2(b). Thomson T_e profiles remained almost the same, both in magnitude and shape, inspite of the increasing density.

(f) Sudden variation of the edge plasma parameters modified the RF coupling to the plasma. The reflection coefficient of the power from the conventional launcher increased by about 0.1, as is seen in Fig. 1.

(g) The increment in plasma stored energy by up to 30% was primarily due to the increase in electron density.

(h) An ELM-free H-phase with durations up to 3.3 s was attained while the beam heated H-phases in divertor configurations often suffered from frequent ELMs (ELM stands for ‘edge localized mode’).
FIG. 2. (a) Reduction in electron density fluctuations near the edge measured by reflectometry during the H-mode with LHCD. The fluctuations in an inner and higher density layer increase with some delay; (b) electron temperature and density profiles from Thomson scattering during the L-mode (closed circles) and H-mode (open circles). Triangles represent the $n_e$ profile at 6.5 s of Fig. 1. Horizontal dotted lines indicate the location of the reflectometry measurement layers.
(i) The $Z_{\text{eff}}$ value estimated from visible bremsstrahlung emission increased slightly from about four to around five during the later phase of the H-mode. Nevertheless, the radiation power measured by bolometer arrays remained at about 40% of the total input power.

(j) No H-mode phase has been observed with LHCD alone in divertor configurations.

Figure 2 shows that a transport barrier [13] is formed near the edge, accompanied by a quench of density fluctuations. These observations agree with the characteristic signatures of the H-mode.

3. HARD X RAY ARRAY DATA

Figure 3 shows an example shot where the H-mode phase was destroyed by (NB) injection. The trace of the electron cyclotron emission at $1.5\omega_{\text{ce}}$ multiplied by $n_e$, which is a measure of LH coupling to the fast electrons, dropped at the L-transition. Hard X ray emission measured by a three channel array as illustrated in Fig. 3(c) behaved similarly to $I_{\text{ECE}}(1.5\omega_{\text{ce}}) \times n_e$. The hard X ray array views the plasma perpendicularly to the magnetic field. The time resolution of the hard X ray data is 0.2 s since they were recorded with pulse height analysers. The traces represent the time evolution of the total photon counts in the energy range from 70 keV to 1 MeV. During the L-phase with NBI, beam acceleration by the LH waves was observed in the neutral spectra shown in Fig. 3(b) as was a drop in electron cyclotron emission (ECE) and hard X ray intensities. These facts suggest that the H-mode could not be sustained when the LH power was absorbed by the beam ions.

The hard X ray spectra are different, depending on the frequency combination. Figure 4(a) shows the channel-three photon spectra during the L-mode phases in the cases of $1.74 + 2.23$ GHz and $2.0 + 2.0$ GHz, respectively. The applied total LH power was nearly the same, i.e. around 1.4 MW. The former spectrum has a stronger intensity with increasing photon energy. The numbers written at the top right corners show the fitted photon temperature which is a measure for the slope of the photon energy spectrum [14]. Around a threshold electron density for the H-mode of $2 \times 10^{19}$ m$^{-3}$, the $1.74 + 2.23$ GHz combination produced photon temperatures of $110 \sim 140$ keV, whereas the $2.0 + 2.0$ GHz combination produced temperatures lower than about 110 keV. The $2.0 + 2.23$ GHz combination generated intermediate spectra, while the $1.74 + 2.0$ GHz combination brought about fitted temperatures higher than 150 keV. Figure 4(b) plots the hard X ray intensity, $I_0$, as a function of the photon temperature, $T_{\text{ph}}$, by fitting the spectra to $I_{\text{HX}} = I_0 \exp(-h\nu/kT_{\text{ph}})$. Open and closed symbols show the distinction in the hard X ray intensity during the L-mode phase which led to the H-mode (open symbols) from those which remained in the L-mode (closed symbols). The broken line indicates that the threshold $I_0$ drops with increasing photon temperature. Hence, the hard X ray intensity of channel
FIG. 3. (a) Time evolution of a discharge where the H-mode with LHCD was destroyed by NBI; (b) beam acceleration observed in neutral particle spectra (closed circles); no ion tail existing with LHCD alone (open circles); (c) hard X ray signals obtained from an array viewing the torus perpendicularly.
FIG. 4. (a) Examples of photon energy spectra of channel three of hard X ray array in L-mode; (b) fitted photon temperatures, varying as a function of frequency combinations: circles (1.74 + 2.23 GHz); squares (1.74 + 2.0 GHz); triangles (2.0 + 2.0 GHz); inverse triangles (2.0 + 2.23 GHz). Open symbols indicate that the H-mode was triggered afterwards. The broken line indicates that there is a threshold in the hard X ray intensity.
three appears to be crucial for obtaining the H-mode. When the LH power was below the threshold, the hard X ray intensity did not reach the threshold level.

The effectiveness of applying LH power at two different frequencies may be explained by the effective generation of suprathermal electrons (\(> 100 \text{ keV}\)). When the frequencies from the two launchers were both set at 1.74 GHz, no H-mode was observed within applied powers of up to 1.5 MW. Although the combined LH power at 2.0 GHz, higher than 2.0 MW, produced the H-mode, the H-phase did not last long owing to the abruptly enhanced carbon influx. Thus the attainment of the H-mode with low LH powers is advantageous for attaining long H-mode phases.

4. DISCUSSION

The hard X ray measurements suggest a favourable effect of the fast electrons on the H-transition. In fact, electron cyclotron heating (ECH) has been reported to be more efficient than NBI in producing the H-mode. The presence of tail electrons with LHCD may be relevant to the low threshold power. The combination of two different frequencies may be preferable to filling the spectral gap with \(N_B\) upshift \([15]\) and efficiently generating energetic electrons at positions that lie relatively outside. The effect of the current profile modification by LHCD on the H-transition has not become clear since the change in the internal inductance was small.

The reason why no H-mode with LHCD has ever been obtained in divertor configurations may be explained by the fact that there are much more pronouncedly peaked hard X ray profiles than in limiter discharges. The hard X ray intensities of channel three from divertor discharges are indicated by the hatched region of Fig. 4(b). The intensities are normalized to eliminate the chord length difference between divertor and limiter configurations. Considering that the current drive efficiency is higher in divertor discharges, we realize that LH power may drive fast electrons more inside than is the case in limiter discharges.

The connection between the H-transition and tail electrons has not yet been understood. The loss of suprathermal electrons may modify the edge \(E_r\) field to establish ambipolarity. The sheath condition at interacting limiters with fast electrons may change the edge electric potential, as well.

Doppler broadening measurements of C VI line using a CXRS system indicated that the H-mode was triggered when the edge ion temperature reached the range from 0.4 to 0.6 keV. The ion collisionality, \(\nu_i^*\), is calculated to be \(\sim 0.1\). Consequently, Shaing's model \([9]\) may not be directly applicable to the limiter H-mode with LHCD.

5. CONCLUSIONS

The H-mode with LHCD was demonstrated in limiter discharges of JT-60. These discharges exhibited the characteristics of the H-mode. Hard X ray emission
measurements indicated that the H-mode was triggered when the hard X ray intensity emitted along a chord outside half tangent minor radius exceeded a threshold, which drops with rising photon temperature. Hence, the tail electrons generated near the edge by LHCD appear to be connected with the low power threshold.

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DISCUSSION

R.R. WEYNANTS: There are several observations which might suggest that the H-mode you are producing is one with a positive (radially outward pointing) radial electric field. The first is that NBI stops the H-mode, as NBI is known to add a negative field.

Secondly, one would expect LH to produce fast electrons which leave the plasma, thus giving it a positive charge. Could you comment on this? Have you made any poloidal rotation measurements?
T. TSUJI: Your comments sound plausible. However, we do not have any data on poloidal rotation at the moment and thus cannot discuss the radial electric field from the experimental point of view.

G. TONON: Why does the application of two different frequencies produce more energetic electrons?

T. TSUJI: We do not have a definite explanation. We speculate that the application of two different frequencies may have the effect of filling the spectral gap with $N_{||}$ upshift at radially different locations.
ICRF HEATING IN SEVERAL REGIMES ON TFTR


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Abstract

ICRF HEATING IN SEVERAL REGIMES ON TFTR.

Over the past year, the two ICRF antennas on TFTR have been used to apply rf powers up to a level of 6 MW. Impurity generation at the antennas has continued to be negligible for the out-of-phase toroidal rf current strap array employed, even at the highest power levels where the average power density at one antenna has exceeded 10 MW/m². The higher operating power levels have supported explorations of all three plasma regimes of interest -- the high-recycling gas-fueled regime, the low-recycling neutral-beam-heated supershot regime, and the low-recycling pellet-fueled regime. Noteworthy results include: ICRF sawtooth stabilization has been demonstrated in all three regimes; fusion reactivity has been significantly enhanced at approximately constant Q in the supershot regime; and efficient core heating at \( n_e (0) = 2.4 \times 10^{14} \) cm\(^{-3} \) has been demonstrated in the pellet-fueled regime. Extrapolations of these results to 12.5 MW operations in the D-T phase of TFTR, which is planned to provide relatively high \( \beta_\alpha \) (~0.5%) supershots in support of \( \alpha \) physics studies and hot high density pellet-fueled regimes in support of CIT operations, are discussed briefly.

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2 Grumman Corporation, Princeton, New Jersey, USA.
1. TFTR ICRF Research Objectives and Experimental Plan

The ICRF research objectives on TFTR are to ultimately (i) increase Q and alpha production and (ii) produce CIT relevant plasmas during D-T operation. The experimental plan adopted to achieve these objectives consists of providing four finite poloidal extent antennas to assure focused core plasma heating [1] up to a power level of 12.5 MW, employing this ICRF power to provide sawtooth stabilization in combination with very high power (~35 MW) neutral beam heated supershots at high plasma current (~2.5 MA) to obtain maximum Q and especially maximum $\beta_\alpha$ to support alpha physics studies, and applying this ICRF power to high-density pellet-fueled plasmas to explore CIT relevant operating regimes especially with regard to heating via wave damping on various ion species (H, $^3$He, T, D), and with regard to the feasibility of employing sawtooth stabilization to avoid periodic degradation of the plasma core reactivity.

Leading up to the D-T phase of operations on TFTR, the ICRF system is being installed in stages during the D phase of operations to permit optimization of the antenna coupling characteristics and exploration of ICRF heating in relevant D and $^4$He regimes. Currently, two antennas are employed at a frequency of 47 MHz to provide up to 6 MW of rf power for studying coupling and heating properties in the gas-fueled, beam-fueled supershot and pellet-fueled regimes with either minority H ($B_\phi \sim 32$ kG) or $^3$He ($B_\phi \sim 50$ kG) fundamental ion cyclotron damping.

2. Antenna Coupling

The box type antennas employed on TFTR [2] are movable in major radius and this movement and selection of the plasma major radius ($R \approx 2.62$ m) have been used to reduce the plasma edge -- antenna current strap separation to ~4.5 cm. Consequently, in the high-recycling gas-fueled regime, a wave loading resistance for the larger of the two antennas ("Bay M") of ~10 $\Omega$ has been obtained for out-of-phase excitation of this antenna's toroidally displaced current straps [3]. This relatively high resistance coupled with the antenna voltage standoff capability of ~50 kV has permitted operation at powers in excess of 5 MW, and consequently at average antenna surface power densities exceeding 10 MW/m$^2$. With both antennas, power levels in excess of 6 MW have been achieved for pulse durations up to 1.5 sec and impurity generation effects have remained negligible for the out-of-phase excitation. (Titanium release from the Bay M Faraday shield is observed in the case of in-phase excitation at much lower power [2].)

The edge density produced for the $R \approx 2.62$ m plasmas with low-recycling neutral beam heating is similar to that produced in the gas-
fueled case and similar loading resistances are obtained. However, pellet injection under low-recycling conditions has very little effect on the edge density so that care must be taken to provide an adequate edge density prior to pellet injection to permit the desired power transfer.

3. Heating, Confinement, and Stability Properties
   a) High-recycling gas-fueled regime

   This regime has been explored extensively in the H minority \( (B_0 \approx 32 \text{ kG}) \) case for both D and He discharges \cite{3,4}. Important observations include: there is no apparent difference in global energy confinement time between D and \(^4\text{He} \) majority ion cases; \( \tau_E \approx 1.2-1.5 \times L \) mode scaling and exhibits the projected plasma current dependence for the range of \( I_p \approx 1-1.8 \text{ MA} \) scanned and over a range of \( \bar{n}_e \approx 2-4.5 \times 10^{13} \text{ cm}^{-3} \) (\( \tau_E \approx 0.2 \text{ sec} \) at \( I_p = 1.4 \text{ MA} \) and is only weakly dependent on \( \bar{n}_e \)). The higher \( \tau_E/\tau_{MODE} \) values are obtained at the lower densities where presumably the higher effective temperature of the energetic hydrogen population leads to better overall energy confinement; and sawtooth stabilization is observed for \( P_{RF} > 4 \text{ MW} \), 2.6 MW in helium and deuterium discharges respectively for \( \bar{n}_e \approx 2.5 \times 10^{13} \text{ cm}^{-3} \).

   The sawtooth stabilization properties are similar to those observed on JET \cite{5}. This phenomenon is particularly important for optimizing reactivity in TFTR in that it results in considerably enhanced temperatures in the core plasma as illustrated in Fig. 1. Note that the core electron temperature is held at near 5 keV for several \( \tau_E \) periods before the sawtooth terminates in a "monster" crash. The sawtooth stabilization has been ascribed to the stabilization of the \( m=1 \) ideal/resistive kink mode due to the presence of the energetic ion component \cite{6,7}. Comparisons of the JET \cite{8} and TFTR \cite{9} data to the ideal theory qualitatively support this interpretation although the experiments tend to be in the resistive limit making quantitative comparisons difficult.

   The termination of the sawtooth period appears to be linked with the onset of the \( m=1 \) kink instability as the current peaks following the peaking of the temperature profile \cite{8}. However, no \( m=1 \) precursor oscillation is observed prior to the crash. Probe measurements reveal that the energetic protons escaping the TFTR plasma are "riding" the central temperature waveform during the sawteeth and the sawtooth free period and no burst of lost ions is observed before a monster crash. However, a possible link between the energetic ion loss at the surface and a loss of the energetic ion density in the core which could lead to a monster crash has yet to be made \cite{9}.
b) Low-recycling neutral-beam-heated supershot regime

ICRF heating in the D(H) regime of the plasma core of the supershot plasma results primarily in central electron heating as shown in Fig. 2a where $P_{RF} = 2.7$ MW is applied to a $P_{NB} = 9$ MW supershot. The increase in $T_e(0)$ of ~2 keV enhances the beam-beam and beam-target reactivity and leads to an increase in the
FIG. 2. ICRF heating of a sawtooth stable supershot produced with 9 MW of balanced NBI 
\(R = 2.62 \text{ m}, I_p = 1 \text{ MA}, \bar{n}_e \approx 2.5 \times 10^{13} \text{ cm}^{-3}, \text{D(H) regime).}\)

neutron (and alpha) production which is approximately linear with the total heating power in Fig. 2b. Also, the global energy confinement for this case remains essentially unchanged with the addition of the rf power (as shown in Fig. 2c) since the rf deposition is well focused to the center of the plasma.

Reducing \(P_{NB}\) to 7 MW (4.7 MW counter and 2.3 MW co) for the conditions of Fig. 2 gives the sawtooth unstable supershot regime [10]. However, sawtooth stability is regained with the addition of \(P_{RF} = 3.5 \text{ MW [9]}.\)

Preliminary measurements in the D(H) regime for \(P_{RF} = 5 \text{ MW}, P_{NB} = 23 \text{ MW}\) reveal an increase in \(T_e(0)\) from 9 to 10.7 keV at \(\bar{n}_e = 4 \times 10^{19} \text{ cm}^{-3}\). The reactivity in this case did not increase linearly with total power indicating that second harmonic heating could be contributing to reactivity in the D(H) case. However, a substantial gain in reactivity is still obtained when the ICRF heating is used to provide sawtooth stabilization.

c) Low-recycling pellet-fueled regime
The pellet-fueled discharges evolve rapidly in time making quantitative analysis more difficult. However, very important results pertaining to higher density operation are being obtained, as shown in Fig. 3. \(T_e(0)\) with \(P_{RF} = 4 \text{ MW}\) is observed to increase linearly in time from \(-0.7 \text{ keV}\) to over 5 keV while \(n_e(0)\) decays from \(-2.5 \times 10^{14} \text{ cm}^{-3}\) to \(-0.8 \times 10^{14}\) after the pellet injection. The time evolution for a pellet case with \(P_{RF} = 0.5 \text{ MW}\) (after an initial blip at 4 MW) is shown for comparison. It is clear from the rapid rise in \(T_e(0)\) that efficient core heating is occurring in the initial phase when \(n_e(0) \approx 2.4 \times 10^{14} \text{ cm}^{-3}\) for which \(T_e = 0.2 \text{ sec}\). Hence, no limit on central
FIG. 3. Pellet-fueled density evolution and electron reheat with 4 MW and 0.5 MW of ICRF heating (D^3He), B_T = 5 T, I_p = 1.4 MA).

Power deposition is indicated to this density level. Another important result is that T_e(0) [T_i(0)] reaches ~4 keV at P_{RF} = 4 MW with n_e(0) ~ 1.2 \times 10^{14} \text{ cm}^{-3} so that projection to P_{RF} = 12.5 MW supports T ~ 6-10 keV at n_e(0) ~ 2 \times 10^{14} \text{ cm}^{-3} giving a plasma with a collisionality and a magnetic Reynolds number comparable to the values projected for CIT.

The reactivity for the case of Fig. 3 is enhanced over an order of magnitude (~5 \times 10^{12} \rightarrow ~1 \times 10^{14} \text{ n/sec}) with the simultaneous application of pellet and P_{RF} over the case with either alone.
Comparable reactivity enhancement for the hydrogen minority regime suggests that direct second harmonic deuterium heating is not playing an important role in this latter case.

Figure 4 gives the general behavior of stored energy versus total power in the pellet case with and without neutral beam heating compared to the cases of ICRF + beams and ICRF alone. These ICRF heated supershott and pellet results indicate a linear dependence comparable to that for the supershots generally and have stored energies well in excess of the L mode projection at the higher powers.

4. Extrapolation to the D-T Regime

The results obtained to date support the ICRF research objectives for the D-T phase on TFTR. The antenna has operated at 10 MW/m² so that reliable operation should be possible at the projected < 8 MW/m² planned for the four antennas of the D-T phase at $P_{RF} = 12.5$ MW. Preliminary extrapolations of the power required to sustain sawtooth stabilization in the optimized supershott regime in D-T ($P_{NB} \geq 30$ MW, $I_p \geq 2.0$ MA) give $P_{RF} \geq 10$ MW where competing second harmonic damping on the majority ion species has been taken into account [9]. The added $P_{RF}$ in this regime should result in at least ~ 1/3 more $\beta_\alpha$ in the plasma core so that exploration of $\alpha$ particle induced stability/instability effects on the $\alpha$ confinement and consequently $Q$ may be possible. Heating of high density D-T plasmas relevant to CIT is feasible and present results indicate that
$T_e,\ n(0) \sim 10 \text{ keV}, \ n(0) > 1.5 \times 10^{14} \text{ cm}^{-3}$ D-T plasmas should be produced for $P_{RF} = 12.5 \text{ MW}$. Future high density experiments on TFTR should help to resolve the power requirement for sawtooth stabilization on CIT.

Acknowledgments

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DISCUSSION

P. VANDENPLAS: I have two questions. Firstly, what was the wall conditioning during this set of experiments? And, secondly, is your power threshold for sawtooth stabilization density-dependent?

J.C. HOSEA: Boronization with diborane gas was performed periodically during the experimental period. Conditioning discharges in helium were also employed to reach the low recycling conditions.
To answer your second question, a density scan from \( n_e \sim (2-5) \times 10^{13} \text{ cm}^{-3} \) was performed in the D(H) stabilization experiments and 3 MW was sufficient to stabilize in all cases, but 2 MW was not sufficient. A fine-scale threshold investigation remains to be done, but a strong density dependence is not evident.

V.E. GOLANT: How would you interpret the difference between your results for effective heating of dense plasma and the most recent results from Alcator C?

J.C. HOSEA: The effective heating at \( n_e(0) \sim 2.4 \times 10^{14} \text{ cm}^{-3} \) was obtained with no apparent problem and suggests that the Alcator C heating problem at higher density will probably not appear at CIT densities in the range \((4-5) \times 10^{14} \text{ cm}^{-3}\). I believe that the essential difference probably lies in the dipole antenna excitation on TFTR (as opposed to monopole excitation on Alcator C), which avoids surface interaction and heating problems. It is also possible that the large ripple in Alcator C has an adverse effect on energetic ion confinement.

B. COPPI: Your proposal to investigate the onset of fishbone oscillations induced by the \( \alpha \) particles from D-T reactions is interesting. According to our theory, the relevant mode should be driven unstable by resonance with the particles that have slowed down to about 400 keV. Have you estimated whether you will have enough of these particles to excite the mode?

J.C. HOSEA: We have only looked at the total number of alphas and have found that the total concentration places TFTR in the vicinity of the fishbone regime. Quantitative projections have not as yet been made, and the effect of ICRF pumping of the slowed-down alphas with the production of fishbones, as you suggested, has yet to be evaluated for TFTR.

F. TIBONE: In your experiments with 6 MW ICRF coupled to a supershot plasma, how did the power input per particle on the electrons in the central plasma compare with that on the ions (including NBI heating and collisional equipartition)?

J.C. HOSEA: For the \( P_{RF} \sim 3 \text{ MW} \) cases which have been analysed, the ICRF heating deposited most of the RF energy into the electrons but did not overcome the dominance of the neutral beam heating in the core which is described in paper IAEA-CN-53/A-III-6 presented by S.D. Scott. The higher power \((5 \text{ MW})\) case is under analysis.
ICRF HEATING IN REACTOR GRADE PLASMAS

Performance of ICRH in JET plasmas

JET TEAM

A fast wave heating and current drive system for NET/ITER


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Abstract

ICRF HEATING IN REACTOR GRADE PLASMAS.

Impurity influxes in JET discharges due to ICRH have been reduced to insignificant levels. This has allowed high quality H-modes to be produced with ICRH alone and has enhanced the density limit which is now the same as the NBI limit. Improvement in the deuterium fuel fraction has led to the generation of 100 kW of non-thermal 3He-D fusion power. Alpha-particle simulations using MeV ions created by ICRH show classical energy loss and suggest that heating in a reactor will be highly efficient. A clear demonstration of TTMP damping of the fast wave in high beta plasmas has been achieved. A broadband ICRH system is proposed for NET/ITER which will allow fast wave current drive and central ion heating for burn control and ignition.

1. Introduction

During JET operation in late 1989 and 1990 the impurity problems associated directly with ICRF heating have been eliminated for all practical purposes. This has been due mainly to beryllium gettering of the nickel RF antenna screens and their subsequent replacement by screens with Be elements. Other important measures to reduce sputtering from the screens have been dipole antenna phasing and alignment of the screen elements with the magnetic field. These improvements have enabled ICRH on JET to make substantial advances in several important areas of plasma performance and these are described in the present paper. For the first time, long 'elm'-free H-modes have been created by minority ICRH alone and have identical characteristics with those produced by neutral

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beam injection. Similarly, the density limit no longer depends on the type of additional heating. Fast ions created by ICRH have been used to produce 100kW of fusion power from non-thermal $^3$He–D fusion reactions and to simulate the behaviour of $\alpha$–particles in a reactor. Section 5 shows the results of Be influx measurements and section 6 describes a demonstration of TTMP absorption of 48MHz fast magnetosonic waves in high beta plasmas. The final section describes a proposal for NET/ITER of a broadband ICRH system operating between 17MHz to 40MHz to give fast wave current drive, central bulk ion heating and non–thermal D–T fusion.

2. H–modes produced by ICRH alone

An example of an H–mode created by RF alone in a double–null X–point plasma with $I_p=3$MA, $B_t=2.8$T is shown in fig 1. The RF power was 7MW and dipole phasing was used. The $D\alpha$ signal shows the H–mode
to extend for 1.3s during which time the energy content given by the diamagnetic loop reaches 6MJ. Comparison with the energy obtained from magnetic analysis suggests that 1.3MJ of this is due to fast ions. The confinement time of the thermal component alone is 0.7s which is twice the Goldstone L-mode value and is similar to values obtained with neutral beam injection (NBI) for the same power level. In fig 1 the central electron temperature reached 9.5keV at the end of the H-mode when the central density had risen to 3.7x10^{19}m^{-3}. Monopole phasing (Prf≈10MW) also produces H-modes and, in some of these, Monster sawteeth occur.

These experiments were assisted by a new feedback system which keeps the coupling resistance (R_c) constant by controlling the separation (X_a in fig 1) of the last closed flux surface from the antenna. In pulse 21906 the requested 3.5 ohm coupling was maintained across the L-H transition and throughout the H-mode phase by plasma movements of the order of 0.01m. Without feedback control R_c is reduced by a factor of three during the H-mode causing generator trips and RF power loss[1].

3. $^3$He–D non-thermal fusion experiments

The fusion reaction $^3\text{He} + \text{D} \rightarrow ^4\text{He}(3.6\text{MeV}) + p(14.7\text{MeV})$ has been studied in JET using $^3$He minority ICRH. The first experiments with carbon coating of the vessel achieved 60kW of fusion power [2]. Recently, experiments with Be coating have reached 100kW as shown in fig 2a. The maximum value of Q ($=P_\text{ fus}/P_\text{ rf}$) was 1% (fig 2b). The later results were obtained with I_p = 3MA, B_t = (3.0–3.4)T, n_e(0)≈4x10^{19}m^{-3} and P rf≤14MW. As previously, the reaction rate was measured by detecting the 16.6MeV $\gamma$-rays from the weak branch $^3\text{He} + \text{D} \rightarrow ^5\text{Li} + \gamma$. In most discharges the ICRF resonance was on axis. Some data were taken with off-axis heating to reduce the power density and test for heating of the $^3$He ions beyond the energy for maximum fusion rate (~0.5MeV). The
yields were similar with on and off-axis heating implying that the $^3\text{He}$ energy was about optimum. A Stix code [2] agrees well with the yields in the carbon vessel for a fuelling ratio $n_d/n_e \approx 0.4$. To fit the data with Be gettering this model requires $n_d/n_e \approx 0.7$ which is consistent with $Z_{\text{eff}} \approx 2.5$ compared with $Z_{\text{eff}} \approx 3.5$ previously.

4. Alpha–particle simulation studies

Minority ions created by ICRH in JET can reach MeV energies. The anisotropic distribution allows their energy content to be deduced from diamagnetic loop and Shafranov shift measurements. The energy content can then be compared with theory to test for non–classical loss processes which might reduce the $\alpha$–particle heating in a reactor. Such comparisons have been made for near steady state Monster sawtooth data which avoids fast ion redistribution by sawteeth. Results are shown in figs 3a and 3b for H–minority ions. Values of $W_{\text{fast}}$ are plotted against model calculations which assume a) an RF power density of Gaussian form with a width of 0.3m, b) that 65% of $P_{\text{rf}}$ is absorbed by the minority (as deduced from modulation experiments) and c) that the fast ions form a Stix[3] distribution on each flux surface. This basic model is supplemented by corrections for the effect of the fast ion orbits on the slowing down time. The slowing down time is averaged over both the orbit and the energy distribution. Fig 3a shows $W_{\text{fast}}$ compared with the results of the basic model only which generally overestimates $W_{\text{fast}}$. However, including orbit effects gives excellent agreement (fig. 3b). A lower limit of 2s is placed on any anomalous loss time which implies that $\alpha$–particles in a reactor will yield all their energy to the plasma, at least in the absence of mhd.

**FIG. 3.** $W_{\text{fast}}$ versus theoretical values for (a) no orbit corrections and (b) with orbit corrections.
5. Beryllium influx from the antenna screens

It is well known that impurity fluxes are released from the screens of powered ICRF antennae on JET. The flux depends on antenna voltage, the plasma density near the screen, the angle of the screen bars to the magnetic field, the phasing and the screen material. The sensitivity to voltage and phasing are illustrated in fig 4 for Be coated screens. Note the threefold reduction with dipole phasing. Beryllium coating of the screens and its use as a first wall material, together with dipole phasing and alignment of the screen bars with the field, have reduced the influxes to insignificant levels. The behaviour can largely be understood in terms of sputtering due to RF field rectification in the sheaths formed where the magnetic field intercepts the screen bars[4]. With the present antenna the field can connect a) adjacent bars and b) different points on the front face of the same bar by virtue of its V—shape. The latter effect is sensitive to phasing and disappears for dipole operation. Beryllium gettering is beneficial partly because of the low sputtering coefficient (especially for self sputtering) and partly because it strongly pumps oxygen and other impurities. The influx contribution to $Z_{\text{eff}}$ is $\Delta Z_{\text{eff}} \approx 0.05 - 0.1$. 

![Graph showing Be influx sensitivity to antenna voltage and phasing.](image)

**FIG. 4.** Be influx sensitivity to antenna voltage and phasing.
6. TTMP damping of the fast wave

Directed fast waves are a strong candidate for non-inductive current drive in reactor plasmas where TTMP is the predominant direct electron damping process. So far, fast wave current drive experiments have been made in low $\beta$ plasmas and with high RF frequency[5] such that the absorption was mainly by electron Landau damping (ELD). The present experiments were carried out in $I_\rho=2\text{MA}$, $B_t=1.3\text{T}$, double-null X-point hydrogen plasmas with the ICRF (48MHz) tuned to $2\omega_{ch}$ on the high field side of the magnetic axis. No mode conversion layers existed in the plasma centre where direct electron heating could only occur through damping of the fast wave by combined TTMP and ELD. The electron $\beta$ was 1.5% so that TTMP was expected to contribute significantly. The electron heating power density was obtained from the $T_e$ (ECE) response to RF power modulation. The profile is peaked on axis as shown in fig 5a and accounts for 22±5% of the input power, the rest being absorbed at $2\omega_{ch}$. Full wave and ray tracing calculations predict more peaked profiles (fig. 5a) but agree with the fraction of the total power absorbed. Data were also obtained from soft X-ray cameras which viewed more of the minor radius than the ECE measurements. Note that the damping is almost zero near the $q=1$ surface (see fig 5b). A recent Hamiltonian treatment of the wave–particle interaction [6] predicts that quasilinear theory is most likely to break down at $q = 1$ resulting in just such a reduced absorption.

7. A broadband ICRH system for NET/ITER

An ICRH system operating between frequencies of 17MHz and 40MHz has been proposed for the NET and ITER tokamaks. At 40MHz deuterium minority heating takes place in the centre of the machine. The
8. Summary

Substantially reduced impurity influx has allowed ICRH on JET to produce high quality H-modes, an enhanced density limit and 100kW of...
non-thermal $^3$He–D fusion power. Alpha particle simulations using ICRF accelerated minority ions imply efficient $\alpha$–heating in reactors. TTMP damping of the fast wave has been demonstrated in high $\beta$ plasmas. The dual purpose ICRH system proposed for NET/ITER allows both central bulk ion heating at 40 MHz for burn control, and fast wave current drive at 17MHz with $\gamma \approx 0.3 \times 10^{20}$A/W/m$^2$.

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DISCUSSION

A. BERS: I have several questions. Have you observed any limitation on antenna power density coupling to the plasma? Have you observed any non-linear effects, e.g. ponderomotive, parametric and so on? And what is the maximum antenna power density coupled to the JET plasma?

D.F.H. START: We have not observed any limitation on the ICRF power density as such. The maximum power coupled to the plasma is usually limited by voltage limitations in the transmission lines. We have not seen any non-linear effects. The maximum power density at the antenna is $\sim 3$ MW/m$^2$.

H. KIMURA: Can you comment on the agreement between experiment and calculation for $Q$ in your D–$^3$He fusion experiment?

D.F.H. START: Theoretical calculations show good agreement for both the D–$^3$He reaction rate and the fast ion energy content.

Y. KAMADA: In your ICRF + NB + pellet experiment, you applied low NB power (2.4 MW). What is the effect of the NB? And could you comment on the difference between NB + pellet and ICRF + pellet?
D.F.H. START: The 2.5 MW, 80 keV NBI was essentially a diagnostic beam for charge exchange spectroscopy and did not penetrate to the plasma centre. A recent experiment had the usual 9 MW of 140 keV NBI and 2 MW ICRF and gave results similar to pulse 22490.

P. MOREAU: You did not say this, but I think that the calculation of the fast wave current drive efficiency quoted here is based on linear damping of the fast wave. Quasi-linear effects would further increase this efficiency. Is this correct?

D.F.H. START: Yes this is correct. The development of a tail would enhance the single pass absorption and this would reduce the $n_t$ upshift due to poloidal field effects.
PROGRESS TO LONG PULSE OPERATION IN TORE SUPRA

EQUIPE TORE SUPRA* (Presented by F. Parlange)
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Abstract

PROGRESS TO LONG PULSE OPERATION IN TORE SUPRA.

Tore Supra has been operated at its nominal current and toroidal field. A multipellet injector was successfully tested and could maintain a $4 \times 10^{19} \text{ m}^{-3}$ quasi-stationary density with up to 25 pellets. The modification of the plasma boundary conditions by a pump limiter or an ergodic divertor has been investigated. Lower hybrid waves with powers of up to 3.8 MW have been launched, driving up to 80% of the total current and, thus, extending the plateau current up to 22 s. The time constants involved in the control of steady state plasmas (heat, particle, current diffusion) are discussed.

1. INTRODUCTION

Long pulse discharges are the main purpose of Tore Supra operation [1]. They allow the investigation of phenomena which are involved in steady plasma control (heat and particle exhaust, chemical composition, current profile monitoring), in particular, current diffusion, thermal equilibrium of plasma facing components and equilibrium of plasma–wall particle exchange. Tore Supra is well suited for these studies because

— it is equipped with superconducting toroidal coils;
— the plasma facing components are actively cooled and designed for up to 25 MW additional heating (neutral beams < 11 MW, lower hybrid < 8 MW, ion cyclotron heating < 12 MW);
— particle fuelling (pellet injector) and exhaust systems (pump limiter) have been installed around the torus; and
— the current can be driven by lower hybrid waves.

Most of these subsystems have been tested and the time constants of slow phenomena have been investigated.

* See Appendix.
Tore Supra, with $R = 2.37 \text{ m}$ and $a \leq 0.8 \text{ m}$, reached its full Ohmic performance by the end of 1989. The superconducting magnet sustained a 1450 A DC current, yielding 4.5 T on axis. A 1.83 MA, 3 duration, flat-top plasma current was obtained.

In 1990, the machine was currently operated in the following parameter range:

$0.8 < I_p < 1.6 \text{ MA}, 3 < B_T < 4 \text{ T}, 1 < \bar{n}_e < 5 \times 10^{19} \text{ m}^{-3}$.

A typical 1.4 MA discharge with $\bar{n}_e = 2.5 \times 10^{19} \text{ m}^{-3}$, $B_T = 3.89 \text{ T}$, has electron and ion central temperatures of 2.2 and 2 keV, respectively, and $Z_{\text{eff}} = 1.9$. The measured energy confinement time of $\tau_E = 0.22 \text{ s}$ is in agreement with neo-Alcator scaling.

The total flux swing available in Tore Supra is about 15 Wb. Part of it is necessary to create the magnetic energy stored in the plasma. It was found that, during the current plateau, all discharges with the same $q^\phi$ end up with the same equilibrium internal inductance $L_0$ after a few seconds. In addition, for a given current, the time and the flux required to reach steady state from breakdown (3–4 s) do not depend on the current ramp rate. The flux left over at that point depends only on $q^\phi$. Ramping the current faster results in a longer current plateau but the steady state phase has the same duration (about 10 s for a 1.2 MA Ohmic discharge).

2. HEAT AND PARTICLE EXHAUST

2.1. Heat exhaust

The problem of heat removal has been treated elsewhere [2, 3]. The heat flux e-folding length in the scrape-off layer was determined by IR imaging on the outboard pump limiter (provided by Sandia National Laboratories, USA). The design value $\lambda_q = 1 \text{ cm}$ has been confirmed in 700 kA helium discharges at 2 T. Heat penetration is twice as deep when the ergodic divertor is used [3, 4], leading to a share of the convective power deposition onto the carbon first wall structures (such as inner wall, ergodic divertor structure, neutralizers).

2.2. Particle exhaust

The outboard pump limiter has been operated with titanium getters (50 to 100 m$^3$/s pumping speed). Large particle outfluxes of $\Gamma_{\text{out}} = (1-1.5) \times 10^{21} \text{ s}^{-1}$, have been obtained, corresponding to the expected values. However, the effect on the plasma density is quite weak: the apparent particle lifetime for a plasma containing $N_p = 5 \times 10^{20} \text{ particles}$ (comparable to the prefill content) decreases from 30 to 10 s, i.e. only 3% of the pumped particles actually come from the plasma [5]. This discrepancy can be understood by considering that a large number of particles, $N_w \approx (1-2) \times 10^{22}$, can easily be removed from the wall and exchanged with plasma particles. Note, for comparison, that pumping between shots removes
(1-1.5) \times 10^{21} \text{ particles and one night of helium glow discharges removes}
\approx 5 \times 10^{22} \text{ atoms. Since } N_w \gg N_p, \text{ as for all carbon wall tokamaks, a pump limiter}
is expected to become efficient on a time constant } N_w/T_{\text{out}} \text{ in excess of 10 s,}
especially if one considers additional fuelling, e.g. pellets, with fluxes in the
10^{21} \text{ s}^{-1} \text{ range.}

3. PELLET FUELLING

A centrifugal pellet injector supplied by ORNL — 100 pellets with
dimensions adjustable from 2 to 20 \times 10^{20} \text{ atoms and velocities of 600 m/s} — \text{ was}
installed in autumn 1989. The expected penetration up to about mid-radius of the
discharge has been confirmed by both H_a \text{ emission time evolution and imagery.}

Up to 25 pellets ((3-4) \times 10^{20} \text{ atoms/pellet}) were injected into the plasma, the
only limit being the pulse duration (Fig. 1). Quasi-stationary conditions were
reached, i.e. \bar{n} \approx 4 \times 10^{19} \text{ m}^{-3} \text{ and } T_{\text{e0}} = 2 \text{ keV}, \text{ without gas puffing. The peaking}
factor } n_{e0}/(n_e) \approx 1.5 \text{ remained unchanged, at least for pellet frequencies lower than}
5 \text{ Hz. The current had to be kept above 1.4 MA, since lower values led to a tempera-
ture quench after a few pellets. Pump limiters were not in use, and the wall has to}
be conditioned by helium tokamak discharges to restore its pumping capability before
the experiment. In any case, it will not be possible to inject pellets for durations
longer than } N_w/T_{\text{fuelling}} \approx 10-100 \text{ s, without active particle exhaust.}

4. CURRENT DRIVE BY LOWER HYBRID WAVES

Two multijunction grills can be fed with up to 8 MW quasi-DC at 3.7 GHz.
The n_i power spectrum has a narrow width, } \Delta n_i = 0.5, \text{ and its centre can be varied
between 1.4 and 2.3, with more than 60% of the power radiated in the direction of
the electron drift.

The scattering coefficients of the multijunction launchers have been measured
as functions of their distance from the plasma and of the scrape-off layer density. The results are in good agreement with the theoretical expectations. Global reflection
coefficients of a few per cent (3-4%) can be obtained. A broad optimum, quasi-
independent of power, even to the maximum launched power of 3.8 MW, is obtained
for a grill to plasma distance of 2-3 cm, corresponding to a plasma density of (0.8-2)
\times 10^{18} \text{ m}^{-3} \text{ at the grill mouth.}

Current drive effects have been observed in a large range of densities
(\Delta V/V = 20\% at } \bar{n}_v = 5.5 \times 10^{19} \text{ m}^{-3}). \text{ The most significant effect is the buildup of an energetic electron tail, and it can be inferred from hard X ray emission spectra that fast electrons are accelerated up to energies of 400-600 keV, depending on the
antenna phasing. As can be expected, the highest energies are observed when the LH
power spectrum is shifted towards low } n_i \text{ values.}
Detailed measurements have been made in a 1 MA, 3.9 T, $\bar{n}_e = 1.5 \times 10^{19}$ m$^{-3}$ deuterium discharge, with powers of up to 2.4 MW. The best results were obtained for central $n_I = 1.6$. The relative loop voltage drop reaches $\Delta V/V = 0.8$ (Fig. 2). Part of the RF power contributes to bulk electron heating through fast electron slowing-down. The plasma centre is preferentially heated, with $T_e$ on axis
FIG. 2. Relative loop voltage drop, central electron temperature increase, estimated RF driven current and sawtooth period as functions of RF power for 1 MA, $1.5 \times 10^{10}$ m$^{-3}$, $B = 3.83$ T, $\langle n_1 \rangle = 1.63$ deuterium discharges.
FIG. 3. Main characteristics of a 22 s long, RF driven discharge: plasma current and loop voltage, line averaged density and central electron temperature, ECE radiation, energy confinement time and radiated power are shown. In each box the heavy line refers to the first signal and to the left hand vertical scale. A 1.4 MW RF pulse was applied from 2 to 22 s.
increasing from 2.5 to 4.2 keV. The RF driven current deduced from the change in the loop voltage, after correction for temperature increase and change in $Z_{\text{eff}}$, is up to 700 kA (Fig. 2). The global energy confinement time of the plasma, as derived from diamagnetic measurements, i.e. neglecting the parallel contribution of the fast electrons, varies in accordance with the L-mode scaling law ($\propto P^{6}$). At low density, an increase in $Z_{\text{eff}}$ may be observed, with $\Delta Z_{\text{eff}}/Z_{\text{eff}}$ typically $\approx 30\%$. In some cases, spectroscopic measurements show the presence of nickel (Ni XVII line), which covers the terminal part of the grill itself, and iron (Fe XV line) probably coming from the torus wall. Low-Z impurities, carbon and oxygen, which are dominant in Ohmic discharges, are generally not much affected at high densities.

The sawtooth period increases from 20 ms in Ohmic discharges to about 50 ms with 1 MW RF (Fig. 2). A sawtooth-free discharge with a density of $n_e = 1.6 \times 10^{19} \text{ m}^{-3}$ was obtained with 1.85 MW RF.

A 20 s, 1.4 MW RF pulse has been used to extend the current plateau of a 1 MA, $\bar{n} = 2.4 \times 10^{19} \text{ m}^{-3}$ helium discharge up to 22 s. The loop voltage was 0.45 V, to be compared to 0.85 V for a similar Ohmic discharge. Steady state plasma conditions were obtained for the last 19 s (Fig. 3), with a central electron temperature of 3.2 keV. $Z_{\text{eff}}$ and the energy lifetime were constant.

5. CONCLUSIONS

Most of the systems designed to control steady state plasmas have been tested separately. Fuelling the discharge by repetitive pellets was only limited by the discharge duration and led to a quasi-stationary state for a few seconds. Particle removal during short pulses changed the plasma density by a factor smaller than expected; the pumping efficiency of the pump limiter is equal to the design value, but the large amount of weakly trapped particles in Tore Supra walls plays the dominant role in particle balance. Thus, for plasma density and composition control several tens of seconds are required.

These long discharges are feasible by using lower hybrid current drive. The multijunction grill has, so far, allowed up to 3.8 MW to be launched in the plasma; the current drive signature was observed in a large range of densities ($5.5 \times 10^{19} \text{ m}^{-3}$). Even with moderate power, flux saving during long RF pulses allowed tokamak discharges to be extended over 20 s. Combined experiments with higher power, particle fuelling and exhaust should bring new insights into wall effects.

ACKNOWLEDGEMENT

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Appendix

MEMBERS OF EQUIPE TORE SUPRA


Permanent addresses

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5 PMI, Ecole polytechnique, Palaiseau, France.
6 Sandia National Laboratories, New Mexico, USA.
7 L.P.M.I. Université de Nancy I. Vandœuvre les Nancy, France.

REFERENCES

DISCUSSION

P.E. VANDENPLAS: How long was your sawtooth-free operation and to how many energy confinement times does this correspond? How did the sawtooth-free discharge end?

F. PARLANGE: The sawtooth-free situation lasted for the entire duration of the RF pulse, 3 s in this particular case. The energy confinement time was of the order of 0.1 s, 30 times shorter than the sawtooth-free situation. Although the central electron temperature increase seems much more pronounced for sawtooth-free RF discharges, the global energy lifetime does not appear (within experimental error) to be affected by the presence or the absence of sawteeth.

F.C. SCHÜLLER: Your LH pulse of 20 s duration appears to end in a disruption. Is that due to an impurity problem?

F. PARLANGE: Because of overloading of the poloidal power supplies, the plasma progressively shifted toward the inner wall during the last 500 ms, a situation which always leads to increased radiation in Tore Supra. Finally, one of the circuit breakers opened, inducing immediate plasma disruption.
TOKAMAK T-15: STATUS AND PLANS

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Abstract

TOKAMAK T-15: STATUS AND PLANS.

An experimental campaign has been carried out on Tokamak T-15 (R₀ = 2.43 m, a = 0.7 m, B₀ = 3.5-5 T, Iₚ = 1.4-2.3 MA). The main feature of T-15 is the toroidal field superconducting winding (TFSW) made of Nb₃Sn. At a TFSW temperature of 7 K, the nominal regime with B₀ = 3.5 T (TF stored energy: 370 MJ) was obtained. The main experiments were performed at a TFSW temperature of 9-12 K, B₀ = 1.5-2.7 T, Iₚ < 700 kA, nₑ < 10¹³ cm⁻³, and a breakdown voltage of Uᵥ = 15-20 V (E = 1.0-1.3 V/m). The plasma current disruptions did not result in a TFSW quench.

1. INTRODUCTION

In the I.V. Kurchatov Institute of Atomic Energy, the experiments on Tokamak T-15 have started. The objectives of this facility and its specific design features are described in Ref. [1]. The main design and attained parameters (for February 1990) are given in Table 1.

The main feature of T-15 is the toroidal field superconducting winding (TFSW) made on the basis of a forced flow (FF) cooled Nb₃Sn conductor. On the other tokamaks with TFSW, pool boiling cooling was used: an Nb-Ti conductor on T-7 [2] and Tore Supra [3] (in the latter case, superfluid He is used for cooling) and an Nb₃Sn conductor on TRIAM-1M [4]. In future fusion reactors, an Nb₃Sn conductor with FF cooling is supposed to be used; therefore, testing of T-15 TFSW is of significant importance for fusion technology.

The poloidal field coils (PFC) are made of aluminium, cooled with liquid nitrogen, and located outside the TFSW (except for the fast vertical control field (VCF) coil, which is located between TFSW and plasma chamber). The entire magnet system is placed inside the common cryostat, which is one further peculiarity of the fusion reactors [5].

A general view of T-15 is shown in Fig. 1, and Fig. 2 shows the facility schematic cross-section.
The physical start of T-15 was realized in December 1988 [6]. Here, the results of the second campaign are presented, during which the engineering experiments with the TFSW aimed at obtaining the nominal parameters were also performed in the plasma discharges.

### TABLE I. T-15 MAIN DESIGN AND ATTAINED PARAMETERS

<table>
<thead>
<tr>
<th>Parameters</th>
<th>Nominal regime</th>
<th>Forced regime</th>
<th>Values attained</th>
</tr>
</thead>
<tbody>
<tr>
<td>Major radius, R (m)</td>
<td>2.43</td>
<td>2.43</td>
<td>2.43</td>
</tr>
<tr>
<td>Minor radius, a (m)</td>
<td>0.7</td>
<td>0.7</td>
<td>0.7</td>
</tr>
<tr>
<td>TF on plasma axis/ TFSW, B₀/Bₘ (T)</td>
<td>3.5/6.4</td>
<td>4.5-5.0/8.2-9.1</td>
<td>3.5/6.4</td>
</tr>
<tr>
<td>TFSW current (kA)</td>
<td>3.9</td>
<td>5.0-5.6</td>
<td>3.9</td>
</tr>
<tr>
<td>TFSW temperature (K)</td>
<td>4.5</td>
<td>4.5</td>
<td>7.0</td>
</tr>
<tr>
<td>Plasma current, Iₚ (MA)</td>
<td>1.4</td>
<td>2-2.3</td>
<td>0.56-0.7</td>
</tr>
<tr>
<td>Discharge duration (s)</td>
<td>5</td>
<td>5</td>
<td>1.2-0.8</td>
</tr>
</tbody>
</table>

*FIG. 1. General view of T-15 facility.*
2. TOROIDAL FIELD SUPERCONDUCTING WINDING TESTS

In the temperature range of 300–100 K, the TFSW, PFC and radiation shields were cooled by gaseous helium for about 400 hours. Further TFSW cooling down was performed by using two He liquifiers LH-800, operated in parallel, and liquid nitrogen was supplied to the PFC and radiation shields. Unfortunately, the attempts to reach the design capacity of 800 L/h failed and the blocks operated only at 260–350 L/h. Besides, the heat flux to the He supply tubes appeared to be very high (0.8–1.2 kW).

Therefore, in spite of the design heat flux directly to the TFSW (0.6–1.4 kW for a cryostat pressure variation from $1 \times 10^{-3}$ to $2 \times 10^{-4}$ torr), a stable cryostating regime was obtained in the temperature range of 9–12 K. The main experiments were carried out in this range with a TFSW current of 2.0–2.5 kA.
To obtain the nominal regime current of 3.9 kA, approximately 300 L of liquid He was stored in the bath of the current leads. By using it, an inlet He temperature of 7 K was reached. At a current of 3.9 kA, the protection system was operating (Fig. 3).

To reduce the voltage in a protective energy discharge, the TFSW was divided into four segments. The time constant for the energy discharge was taken to be
FIG. 5. Temperature dependence of TFSW current carrying capacity in assembly (o), model coils (∆) and single blocks tests.

FIG. 6. Dependence of total Ohmic power in 24 TFSW blocks on $I_B$ product: ∆ — preassembly tests; o — assembled TF system.
104 s, which corresponded to maximum voltage-to-ground 250 V at a current of 4 kA. To reduce the warming of the quenched coil sections (calculations yielded a temperature rise of up to 150 K at 4 kA), a twofold increase in the discharge resistance was used at the moment where a current equal to half maximum was attained. Typical current and voltage evolutions for a fast energy discharge are shown in Fig. 4. The He outlet temperature rise was 1.6–2.0 K, in this case. The initial temperature regime was restored within 5–10 min.

It should be noted that the operating conditions of superconductors in an assembled TF system differ significantly from those realized during the preassembly tests of the coil blocks. Therefore, as a criterion for comparing data on the current carrying capacity, the $I_cB$ product was taken ($I_c$ is the critical current, and $B$ the magnetic field), which at fixed temperature is constant over a wide range of magnetic field values. The available data on the $I_cB$ dependence upon temperature are shown in Fig. 5. We see that extrapolation to the 4.5 K temperature range permits us to consider a possible attainment of a 4.5 T magnetic field at the torus axis (forced regime). In the same figure, the results of model coil and separate working block testing are shown. The $I_cB$ values obtained for an assembled TFSW in the range of $T > 7$ K are seen to be approximately 20% higher than the minimum values obtained for the coils during the preassembly tests.

In the course of preassembly block tests, rather high losses of the $Q_{\text{res}}$ level (up to 70 W per block at $I_cB = 33 \text{kA} \cdot \text{T}$) were observed. Figure 6 shows a comparison of the results obtained on T-15 with the total heat release in blocks determined at the preassembly tests. T-15 $Q_{\text{res}}$ is seen to be, at least, not higher than its values obtained for the separate blocks.

![Plasma current disruption](image-url)

**FIG. 7.** Typical outlet He temperature rise waveforms for two pulses after 400 kA plasma current disruption with a PF variation rate of 15 T/s.
In plasma discharge experiments, the current in the TFSW was maintained stable for a rather long period of time, maximally for six hours at 2.5 kA and for two hours at 3.0 kA. In this case, the moment for coil switch-off was specified by the termination of the plasma experiments and was not associated with any faults in the TFSW.

In the course of this campaign, plasma current disruptions occurred with a time constant of 10–14 ms and a PF variation rate of up to 15 T/s. At TFSW currents up to 3 kA (I/I_c ≈ 0.85), the plasma current disruptions did not result in a coil quench. At the coil outlet, the coolant was heated up to 0.25 K (Fig. 7).

3. EXPERIMENTS WITH PLASMA DISCHARGES

In T-14 the PF system comprises the Ohmic heating coils (OH), the control coils PF1, PF2, and PF3 for plasma equilibrium and shape control, fast vertical control field coils (VCF) (B_{max} = 0.2 T, \dot{B}_{max} = 3 T/s), and the horizontal control field coils (HCF) (B_{max} = 0.018 T, \dot{B}_{max} = 0.4 T/s). The OH coils consist of a 46 turn central coil OH1 and two four turn end coils OH2 (Fig. 2). According to the design scheme, OH1 and OH2 should be connected in series.

Chamber conditioning can be performed by baking, glow and induction discharges. In the case of 'cold' TFSW, the baking temperature is limited by the nitrogen refrigerator cooling capacity, because of the part of the radiation shield installed between the chamber and the TFSW.

Before TFSW cooling down, the discharge chamber was baked up to 200–250°C and maintained at this temperature for 15 days. For TFSW working conditions, a chamber temperature of 170°C was reached at a load to the nitrogen refrigerator of 160 kW. As a result of vacuum conditioning, the following elemental composition of the chamber surface was obtained: 50% C, 40% Fe, 10% O (initial composition: C > 90%, (Fe + O) < 10%). The pressure in the chamber (1 × 10^{-7} torr) was obtained at a water vapour and hydrocarbon partial pressure of <5 × 10^{-9} torr.

The experiments with an electron beam before the discharges have shown that the stray magnetic fields resulting from the residual iron magnetization did not exceed 2–3 G and those resulting from inaccurate TFSW block assembly amounted to 10^{-3} B_0 [5].

During the experiments with the PF system, an abnormally high radial magnetic field was detected in the chamber. It was established, as a result of measurements and calculations, that the OH2 upper coil was connected oppositely to OH1 and the lower OH2 coil. Figure 8 shows the experimental and calculated dependences of the radial magnetic field component on the OH current at the point R = 1.6 m, Z = 0.

In the course of this experimental campaign, OH coil lead reswitching was impossible as cryostat opening was needed. Therefore, the radial field was compensated for by means of an HCF coil with an accuracy of B_r = 5 G at I_{OH} = 15 kA. Compensation
FIG. 8. Experimental and computed dependences of radial magnetic field component at point \( R = 1.6 \text{ m}, Z = 0 \) on OH coil current in the case of a wrong connection of the upper OH2 coil.

FIG. 9. Typical waveforms of plasma current \( I_p \), loop voltage \( U_{loop} \) and average plasma density \( \bar{n}_e \).
of the vertical field $B$ generated by the currents flowing in the OH coils and the chamber at the breakdown stage was provided by the VCF coil.

Experiments with the plasma current were performed at $B_0 = 1.5-2.7$ T and an initial pressure in the chamber of $(2-3) \times 10^{-5}$ torr. Gas breakdown (hydrogen) was realized by supplying the voltage from a thyristor converter to OH ($U_{\text{max}} \leq 1.2$ kV). The plasma loop voltage at the breakdown phase was $U_{\text{loop}} = 15-20$ V ($E = 0.8-1.3$ V/m). The current and the pulse duration of this phase were $I_p = 35$ kA and $t_p = 150$ ms, respectively.

Stably repeating regimes with plasma currents of 560 and 700 kA were obtained as a result of control system optimization. The maximum pulse duration at $I_p = 560$ kA was 1.2 s (Fig. 9) and 0.8 s at $I_p = 700$ kA. $U_{\text{loop}}$ in the quasi-stationary discharge stage at $I_p = 560$ kA did not exceed 1.5 V, and the horizontal shift of the plasma column was not larger than 1 to 2 cm. The plasma density measured by laser and microwave interferometer was $(0.6-1) \times 10^{13}$ cm$^{-3}$, $q(a) = 3-4$ ($a = 0.7$ m is the graphite limiter radius). Typical waveforms of OH, PF3, VCF and HCF currents in one of the operating regimes with $I_p = 560$ kA are shown in Fig. 10.

4. CONCLUSIONS

During a five months' experimental campaign, the T-15 equipment was tested as a whole. The main result is that the nowadays largest superconducting magnet system made of Nb$_3$Sn conductor has shown its ability to operate in the nominal regime and under real plasma experimental conditions.
As the design parameters of the cryogenic system for He supply were not attained (mainly because of the low reliability of screw He compressors), the main work was carried out at a TFSW temperature of 9–12 K. TFSW nominal parameters were obtained at a temporary temperature reduction to 7 K. Plasma current disruptions did not result in TFSW quenching. Extrapolation of the results obtained by studying the TFSW current carrying capacity to the region of a temperature of 4.5 K permits us to hope to obtain a toroidal field on the plasma axis of $B_0 = 4.5$ T.

In the course of the operation, it has been found that the OH2 upper coil was wrongly connected relative to the rest of the OH parts. However, this circumstance had no impact on the plasma discharges in the framework of this campaign.

The nearest future plans are for T-15 as follows: up to autumn 1990, the OH2 upper coil reconnection is supposed to be finished (for this purpose the cryostat was opened); in January 1991, the cryostat and chamber pumping-out will be started, in February 1991, TFSW cooling-down, and in March 1991 a new experimental campaign will begin.

In parallel, work on defect elimination in the He cryogenic system is under way; the complex tests of this system are scheduled to take place in December 1990.

Before the new campaign starts, one from two NB injectors and three from 24 gyrotrons should be prepared for the additional plasma heating experiments.

REFERENCES


DISCUSSION

D.E. POST: What is the power and energy of the neutral beam heating system in T-15?

O.G. FILATOV: During the first stage of operation the power will be 6 MW and the energy 40/80 keV, and during the second stage the power will be 9 MW at the same energy.

H. KISHIMOTO: What is the configuration of the limiters in T-15, and what are they made of?

O.G. FILATOV: The limiters in T-15 are circular in the poloidal direction and the material is graphite.
PBX-M RESEARCH PROGRESS: APPROACH TO SECOND STABILITY

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Abstract

PBX-M RESEARCH PROGRESS: APPROACH TO SECOND STABILITY.

Research on the PBX-M tokamak has focussed on pursuing access to the second stability regime by developing tools to control the plasma shape as well as the plasma pressure and current profiles. Theory indicates that a broad current profile provides enhanced shaping capability and plasma stability properties necessary for accessing second stability at values of q(0) near unity. In PBX-M, the plasma surface was closely fitted to a conducting wall by broadening the current profile utilizing a current ramp technique, resulting in the achievement of high-\(\beta\) (6.8% at high values of I/aB=2.3 MA/m-T), high confinement (\(T_{e}(0)=5.5\,\text{keV},\,\tau_{E}=1.5\,\text{\text{\varepsilon}}_{E\text{-Goldston}}\)) discharges. Although these discharges terminated in disruptions, a clear influence of the close fitting conducting shell in PBX-M was evidenced by reduced growth rates of the n=1 disruption precursors relative to those in the predecessor PBX device (without passive plates) and to those in PBX-M plasmas with large plasma-plate separations. Discharges near the threshold of second stability were achieved at medium values of I/aB=0.9 MA/m-T, with \(\beta_{i}/(I/aB)\) reaching values as high as 4.5 \%-m-T/MA in this regime. These plasmas exhibited H-transitions, with resulting H-mode confinement times of order \(2\,\text{\varepsilon}_{E\text{-Goldston}}\) and 2.5\,\text{\text{\varepsilon}}_{\text{ITER-89P}}. The \(\beta\)-limiting processes in these plasmas may be related to the onset of ideal instabilities. In both the medium- and high-I/aB regimes, the maximum \(\beta_{i}\) values were 30 to 50\% greater than those achieved in PBX, indicating an overall increase in the stability and confinement of PBX-M plasmas.

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1. Introduction

The objective of the Princeton Beta Experiment (PBX-M) is to access and study the second regime of stability to ideal ballooning modes. The potential for enhanced confinement and improved performance in that regime could naturally lead to an improvement of the tokamak concept as a basis for an advanced concept reactor design. Previous theoretical work has indicated the benefits of plasma indentation and large aspect ratio for achieving second stability at values of q(0) near 1 [1,2]. More recent work based on experimentally achieved pressure profiles has pointed to the importance of being able to control the current profile, most notably by driving edge current to broaden the current profile, to achieve second stability at q(0) values near unity [3].

The PBX-M experiment combines uniquely the features necessary to optimize both plasma stability and energy confinement in order to achieve second stability. These features include low-n mode stabilization through use of a close fitting conducting shell (passive plates), low- and high-n mode stabilization through plasma shaping (indentation up to 28%) and current and pressure profile modification, and high aspect ratio (a=0.3 m, R=1.65 m, R/a=5.5). Profile modification was accomplished to date by several techniques: up to 6 MW of $D^0$ injection into $D^+$ plasmas (with the injection divided equally in two different orientations, $R_{tan}=0.35$ m and $R_{tan}=1.30$ m), inductive current ramps, multi-pellet injection, and up to 0.9 MW of Ion Bernstein Wave (IBW) heating. Two diagnostic techniques, motional Stark effect polarimetry (MSE) [4] and tangential soft X-ray imaging, proved successful in measuring q(r,t), thus providing an essential boundary condition for plasma stability calculations, and a monitor of the efficacy of the various current drive techniques. Further details of the experiment can be found in [5]. While the predecessor device, PBX, had essentially the same mission as that of PBX-M, PBX lacked the close conducting shell, the ability to create highly indented plasmas, and sufficient profile control techniques [6].

2. Experimental Results

An overview of the volume-averaged toroidal beta, $\langle \beta_t \rangle$, achieved in PBX-M as a function of $I/aB$ is shown in the top panel of Fig. 1. The $\langle \beta_t \rangle$ shown in the figure was calculated from a simplified equilibrium code which neglects eddy currents induced in the passive plates. The PBX-M experiments were
concentrated in two regimes; the first was a high-I/aB regime (=2.1 MA/m-T) where the current ramp technique was employed to broaden the current profile by adding edge current. Plasmas in this regime exhibited $\langle \beta_T \rangle$ values up to 6.8% (with eddy current corrections) and energy confinement enhanced over L-mode. As will be discussed, the effect of the passive plates was most noticeable in this regime, in which the plasma was closest to the plates.

The second regime was at I/aB=0.9 MA/m-T in which H-mode plasmas with $\langle \beta_T \rangle$=4% and $\beta_{pol}$$\leq$2.4 were determined to be near the threshold of second stability. Used in the equilibrium and stability calculations for these cases were the actual measured values of $q(0)$.

The bottom panel of Fig. 1 shows the envelope of $\langle \beta_T \rangle$ vs I/aB values for both PBX-M (thick solid curve) and the predecessor device, PBX (thick dashed

FIG. 1. (a) Volume-averaged toroidal beta, $\langle \beta_T \rangle$, in per cent plotted as a function of I/aB for PBX-M; (b) envelope of the maximum $\langle \beta_T \rangle$ versus I/aB for PBX-M and PBX.
curve). The PBX-M plasmas achieved normalized beta values, $\beta_n = <\beta_t> / (I/a_B)$, of up to 3.5 %-$m$-T/MA in the high-$I/a_B$ regime, while at medium $I/a_B$, $\beta_n$ reached values as high as 4.5 %-$m$-T/MA. The other important feature to note is the marked improvement in performance (higher $<\beta_t>$) of PBX-M plasmas over those in PBX for all $I/a_B$, going from approximately a 30% improvement at high-$I/a_B$ to a 50% improvement at medium $I/a_B$.

2.1. High-$<\beta_t>$ Regime with Broad Current Profile

The goal of the inductive current ramp experiment was to study the effects of a broad current profile produced by driving current in the region $r/a > 0.7$. Because total plasma current increased with time using this technique, it led, by its very nature, to high values of $I/a_B (=2$ MA/$m$-T) and plasma conditions which were transient. Nonetheless, the effect of the broader current profile was studied in anticipation of other, more steady-state current drive techniques, such as lower hybrid current drive, which can drive current at the edge for long periods of time without leading to an ever-increasing total plasma current.

The current ramp technique in the predecessor PBX device produced plasmas with $<\beta_t>$ values up to 5.5% that were terminated by disruptions due to growing $n=1$ external kink oscillations [6]. PBX-M plasmas in this high-$I/a_B$ regime achieved values of $<\beta_t>$ as high as 6.8%. An equilibrium flux plot of a high-$<\beta_t>$ plasma is shown in Fig. 2. To obtain this plasma, an $I_p$ ramp rate of 2.4 MA/sec was employed. The resulting broad current profile allowed operation with the plasma surface close to (within a few cm) and conforming to the shape of the passive stabilizers, as seen in the figure. Other parameters for this discharge were: $I_p=570$ kA, $B_t=1.1$ T, $P_{\text{heat}}=5.7$ MW, $q_{95}=3.2$, indentation=28%, elongation=2.2, $n_e=5 \times 10^{19}$ m$^{-3}$, $T_e(0)=2.2$ keV, $T_i(0)=5.4$ keV, and $\tau_E=53$ msec (which is a factor of 1.5 greater than the Kaye-Goldston L-mode value [7] for those discharge conditions). The MHD stability analysis of the high-$<\beta_t>$ plasmas indicated that the $n=1$ ideal kink mode was stable, although growing $n=1$ modes, leading to a discharge-terminating disruption, were observed. The stability calculation, however, may not have represented accurately the discharge conditions, as it was based on an incomplete set of kinetic profiles, and assumptions that the passive plates were ideal conductors and that $q(0)$ was 0.9.

The passive plates were successful in modifying the mode responsible for the disruption, and they led to increased edge confinement. In both PBX and
PBX-M, the current disruption was often preceded by a thermal collapse of the plasma, as inferred from the change in the soft X-ray emissivity profile, with the thermal collapse preceded by a growing n=1 mode. In the final stage of the n=1 mode growth (0.5 to 2 msec prior to the thermal collapse), higher-n components appeared and grew non-linearly. The PBX thermal collapse occurred simultaneously across the entire plasma column, while in PBX-M the collapse initially occurred only in the core; the plasma edge remained stable and actually showed a transient energy increase. Furthermore, the amplitude of the n=1 precursor just prior to the thermal collapse was approximately a factor of three to five smaller in PBX-M than in PBX, possibly due to a shift to a more internal mode structure. In PBX, the e-folding growth times of the growing n=1 mode were from tens to hundreds of μsec. In PBX-M non-indentated plasmas, with the plasma surface far from the plates, the growth times were comparable to those in PBX. However, for the PBX-M discharges in which the plasma surface was close to the plates, the precursors grew on a time scale of hundreds of μsec to several msec, a factor of ten greater growth time, and none
of the growth times were longer than the L/R time of the passive plates (20 msec). These results suggest that it was the presence of the plates, and not just different plasma shapes, that was responsible for slowing the n=1 growth.

A further measure of the effect of the passive plates was a comparison of the disruption characteristics of similar discharges with two different passive plate configurations. In the most recent operational period, the effective surface area for low-n suppression was increased by strapping another set of passive plates in with the original set, increasing the poloidal extent of the plates by 50%, and the L/R time from 20 to 50 msec. The skin time of the plates is 1 to 1.5 msec. The stabilizing effect of the plates was analyzed by monitoring the eddy currents on the passive stabilizers as measured by Rogowski coils placed at the electrical junctions between the plates at ten toroidal locations. The eddy current pattern prior to the disruption was decomposed to identify locked and oscillatory modes. In the initial configuration, the locked and oscillatory modes grew over a period of 5 to 10 and 15 to 25 msec respectively. With the extended passive plate configuration, the duration of growth increased to 10 to 15 msec for the locked modes and 30 to 50 msec for the oscillatory modes.

The above results indicate that the passive plates were effective in stabilizing the n=1 locked and oscillatory modes up to the L/R time scale, and that the extended shell was advantageous. The thermal collapse occurred 0.5 to 2 msec after the non-linear growth of coupled n=1 and 2 modes in both configurations, with the maximum amplitude of these modes being 30 to 50 kG (0.7 to 1.0 kA) at the plate surface. Because the passive plates provide a means by which the linear growth times of the n=1 locked and oscillatory modes can be increased, the opportunity exists for actively controlling profiles to avoid the strong, deleterious non-linear phenomena.

A major concern for the design of large, elongated devices such as CIT and ITER is the effect that eddy current related forces due to fast vertical disruptions will have on in-vessel hardware components. Although vertical motion is controllable by a combination of passive stabilization and active feedback systems, discharges often disrupt vertically when the limits of the control system are exceeded. Serious damage to the support structures during the initial phase of PBX-M current ramp operation, when the plates were electrically connected to the vacuum vessel, stimulated a re-examination of the model for the vertical disruption with the actual hardware arrangement. The TSC code was modified to simulate the actual hardware configuration and calculate the external current flow from the plasma surface to the plates and through the vacuum vessel via the halo plasma near the plasma boundary. For this
calculation, the electron temperature of the halo was assumed to be 10 eV. The calculation predicted plasma motion along the surface of the passive stabilizers towards the separatrix at a velocity of 0.3 to 1 m/msec, with a resulting poloidal voltage between the plates of approximately 0.3 to 0.8 kV. These values are consistent with the velocities obtained from the soft X-ray wave array and poloidal voltage measurements (through $\mathbf{E} \times \mathbf{B}$). Eliminating the halo plasma from the calculations gave estimated velocities and voltages three orders of magnitude larger (approaching ideal MHD values), inconsistent with experimental observations. Insulating the plates from the vessel eliminated the occurrence of the currents through the vacuum vessel, and, therefore, the structural damage from vertical disruptions in PBX-M.

2.2. High-$\beta_{\text{pol}}$ ($\beta_n$) Regime

In the medium-$I/aB$ regime (0.9 MA/m-T), values of $\beta_{\text{pol}} \leq 2.4$ and $\beta_n \leq 4.5$ were achieved with $\leq 5.5$ MW of neutral beam injection (NBI). Typical discharge parameters for these plasmas were: $I_p = 330-370$ kA, $B_t = 1.3$ T, $q_{95} = 5$, indentation=16%, elongation=1.9, $T_e(0) = 2.2$ keV, and $T_i(0) = 4.5$ keV.

![FIG. 3. Evolution of various plasma parameters for a high-$\beta_{\text{pol}}$ discharge.](image)
The temporal evolution of such a discharge is shown in Fig. 3. As can be seen in the various traces in the figure, the plasma current remained steady near 350 kA while the neutral beams were turned on sequentially until a total injected power of 5 MW was reached. A sharp drop in the $D_{\alpha}$ signals at 375 msec on both the midplane and in the divertor signify a transition into the H-mode. This was followed by a period free of giant ELMs in which the stored energy increased and then saturated.

The $\beta$-saturation in PBX-M, similar to that in PBX [8], was caused by sawteeth, continuous $n=1$ to 3 mode activity, and isolated or grassy ELMs. A major difference between this period in PBX and PBX-M, however, was that no fishbone oscillations or associated particle losses were observed in PBX-M. The reason for this difference is not known, although it could be related to the higher indentation of the PBX-M high-$\beta_{\text{pol}}$ plasmas relative to those in PBX, and/or different current or pressure profiles. The amount of energy loss associated with each type of MHD activity during the saturation phase of PBX-M discharges was also not known since the high sensitivity diamagnetic loop was not available. Consequently, the source of the 50% improvement in performance of PBX-M over PBX has not yet been identified unambiguously.

As seen in Fig. 3, the stored energy decrease near 500 msec was coincident with a giant ELM that brought the plasma back into the L-mode. The giant ELM could cause up to a 20% loss of plasma stored energy. Although not seen in the figure, the $n=1$ continuous mode "locked" during this collapse phase, and was coincident with increased enhanced energy loss. The deleterious effect of the giant ELM was much greater in PBX-M than in PBX.

The energy confinement times of these high-$\beta_{\text{pol}}$ discharges showed up to a factor of two improvement going from the L- to the H-phase, with the H-phase confinement times being up to a factor of 2 times Kaye-Goldston and 2.7 times ITER-89P L-mode [9] values. The transition from the L- to the H-phase resulted in a reduction in the local transport, and a transport analysis of a full set of kinetic profiles (including $T_i(r)$) indicated the reduction in transport to be primarily in the ion channel. This is seen in Fig. 4, where the electron and ion thermal diffusivities at $r/a=0.5$ going from the 3-beam L-phase to the 4-beam H-phase of the discharge shown in Fig. 3, are plotted. $\chi_i$ is reduced by over a factor of three from the L- to the H-phase, going from 4.5 m$^2$/sec to 1.3 m$^2$/sec, with essentially no change in $\chi_e$. Since the heating power increased across the L- to H-transition, it can be argued that there was an effective reduction in $\chi_e$ and an even more dramatic reduction in $\chi_i$, relative to L-phase values, taking into account the tendency for transport to increase with increasing power.
FIG. 4. Time evolution of $\chi_i$ and $\chi_e$ at $r/a = 0.5$ as calculated from a transport analysis of kinetic profiles obtained from a documentation of similar discharges. For this transport run, the L–H transition occurred at 380 msec.

The stability analysis of these high-$\beta_{pol}$ plasmas to high-$n$ ($n \geq 25$) ideal ballooning modes was performed using the CAMINO code. For this calculation, a fit to the measured thermal and calculated fast ion contributions to the total pressure profile and the actual $q(r)$ profile as measured by MSE were used. The measured $q(0)$ value was 0.7. Shown in Fig. 5 are the $S-$\(\alpha\) plots at two locations in the plasma, a core location ($q=1.8$) and a location near $r/a=0.7$ ($q=2.9$). While the central region of the plasma was well within the first stability regime, the region farther out was at the transition to the second stable regime. The analysis also indicates the discharge to be marginally unstable to ballooning modes near $q=2.4$. These results, which indicate the plasmas to be approaching second stability, offer encouragement that finer profile control, particularly with edge current as planned with LHCD, can make second stability attainable.

The separation between the plasma and the plates for these high-$\beta_{pol}$ plasmas was much larger than the few cm separation at high-$l/aB$ (approximately 5 cm at the midplane going to 20 cm near the top), which diminished the stabilizing effect of the passive plates. An example of this is the result of the stability analysis of plasma profiles at the time of a giant-ELM, a
major $\beta$-limiting mechanism in these discharges. Just prior to (1 msec) the ELM at 490 msec in Fig. 3, the stability analysis results with the wall at 14 cm indicate that the plasma edge was unstable to the pressure-driven, $n=1-3$ external kink. The calculations showed further for the pre-ELM case that the presence of the plates modifies the poloidal structure of the unstable modes in a manner consistent with observation [10]. These results, coupled with the ballooning stability results discussed above, indicate the $\beta$-limiting ELM to be more consistently explained as an external kink, rather than ballooning, mode. One objective of the upcoming LHCD campaign is to broaden the current enough to increase the plasma shaping and decrease the plasma-plate separation at low $I/a_B$ in order to stabilize this edge mode.

2.3. Profile Control

Control of plasma pressure and current profiles was attempted with counter-NBI, pellet injection, and IBW heating. Counter beam injection was used as a technique to raise $q(0)$ above 1, a condition that would facilitate access to second stability. For these experiments, $q(r)$ was measured by the MSE and $q(0)$ by the soft X-ray pinhole camera diagnostics. The value of the measured $q(0)$ varied consistently with the existence of sawteeth, and its time evolution
was well related to the change in the magnitude of the beam driven current, as calculated in the TRANSP code. Shortly after the onset of NBI, and late in the discharge, when sawteeth were not observed and the MHD changed character, the time-averaged q(0) was greater than unity, while in the ohmic phase and in the middle of the NBI phase, when sawteeth were present, the time-averaged q(0) was less than unity. This technique of q(0)-control, along with advanced wall conditioning methods (e.g., boronization) to reduce the impurity buildup during counter-injection, will be re-visited during future operation.

Multiple (up to three) pellet injection into NBI discharges succeeded in peaking the density profile and maintaining the peakedness for approximately 100 msec. Strong n=1 activity, present prior to the pellet injection, was suppressed during this 100 msec period. Pellets injected into H-mode plasmas, however, caused the plasma to transition back into the L-mode, although \( \tau_E \) remained nearly constant.

Up to 0.66 MW of IBW power (41 MHz) was injected for 150 msec into discharges with 2 MW of NBI, resulting in broader and higher density profiles. A 20% increase in stored energy was observed, with this increase due mostly to the increase in density. Other encouraging results from these IBW experiments were the absence of parametric instabilities, no increase in high-Z impurity levels, no significant change in coupling to the plasma across an H-mode transition, and an apparent reduction in sawtooth activity during the IBW pulse. Up to 2 MW of IBW power will be available during the next operational period.

3. Summary and Future Plans

The first phase of PBX-M operations focussed on the development of the tools for second stability access: plasma shape control (indentation), \( j(r) \) and \( p(r) \) profile modification, and surface mode stabilization through a close fitting conducting shell. The ideal surface mode stabilization was shown to be effective for plasmas with broad current profiles which allowed better surface shaping capability to bring the plasma surface and passive plates to within a few cm of each other. To date, broad current profiles needed to achieve this shaping have been produced by the \( I_p \) ramp technique (at high-I/aB), resulting in plasmas with high-\(<\beta_t>\) and enhanced confinement. The addition of edge current to broaden the current profile still needs to be explored in the plasmas closest to second stability (medium-I/aB) where low-n ideal surface modes may be responsible in part for limiting the achievable \(<\beta_t>\). Nonetheless, in both plasma regimes, PBX-M performance (i.e., \(<\beta_t>\) was 30 to 50% improved
over that of PBX. This enhanced performance in PBX-M indicates the apparent success of the various techniques for improving stability both at the edge and in the interior of the plasma.

A most promising method for edge current control is a rapid phase-controlled 2 MW LHCD (4.6 GHz) system presently being installed on PBX-M. For pressure profile control and as an extra heating source, the IBW is being upgraded to 2 MW, with 7 MW possible in the future. Results from the TSC code predict access to second stability across the entire plasma volume, with $\langle \beta_p \rangle = 15\%$ and $\beta_n = 9$, using 7 MW of NBI, 6 MW of IBW, 2 MW of LHCD (driving 150 kA of current near the plasma edge), and with an indentation of 35%.

In order to determine whether access to second stability has been achieved, detailed current profile and ballooning mode measurements are essential. To this end, the MSE diagnostic is being upgraded to a multi-channel system able to measure the q-profile within a single discharge. A vertical soft X-ray array to supplement the existing horizontal array, a densely packed array of Mirnov coils, a tangential phase contrast imaging system, and a third harmonic ECE system as part of a collaborative effort with MIT are being developed and installed. To optimize edge confinement further, which may also be important for access to second stability, passive plate biasing and edge diagnosis with a fast reciprocating probe will be carried out in conjunction with a group from UCLA.

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DISCUSSION

P. SMEULDERS: Why did you not get an H mode in the high $\beta$, discharges? Was this not related to the $\beta$ limit? And what were the precise modes that led to the high $\beta$, disruption?

N. SAUTHOFF: We did achieve an H mode in some high $\beta$ shots in PBX-M. The enhancement factor was approximately 1.5 compared with Kaye-Goldston scaling. The present highest $\beta$ (≈ 6.8%) was obtained with a relatively high toroidal field ($B_t = 11$ kG). Further experiments with lower $B_t$ will be needed to see whether the H mode probability is correlated with the $\beta$ in PBX-M.

Many major disruptions were preceded by a large internal disruption like a giant sawtooth crash. The precursor was dominated by the $m/n = 3/1$ mode, which occurred 10–30 ms prior to the thermal collapse.

F.X. SÖLDNER: Is it possible that impurity accumulation, starting from the beginning of counter-NBI, contributes to the observed sawtooth stabilization?

N. SAUTHOFF: During the initial period of disappearance of sawtooths just after the NBI (200–250 ms), we did not observe any change in the impurity behaviour and do not believe that impurities contributed to the raising of $q(0)$. However, in the later phase (400 ms), we observed the increase of impurity density and the increase of the radiation loss due to higher plasma density which led to a decrease in $T_e(0)$. Our interpretation is that during this later period the $q(0)$ was raised by the combination of higher neutral beam driven current and the increased resistivity near the axis.

K. KIKUCHI: What is your maximum bootstrap fraction in PBX-M?

N. SAUTHOFF: The bootstrap fraction estimated by TRANSP analysis for beam discharges with four NBI beams is up to 35% with $\beta_p \sim 2$. There have been other discharges in which the bootstrap fraction was up to 55%.

K. KIKUCHI: In a steady state reactor the frequency of the hard disruption would be important. Did you see major disruptions in your high $\beta_p$ discharges?

N. SAUTHOFF: Major disruptions leading to current collapse were not normally observed in the high $\beta_p$ discharges.

Y.I. KOLESNICHENKO: Could a plasma cross-section similar in shape to that of PBX-M be used in reactor conditions, for instance in ITER?

N. SAUTHOFF: One difficulty in adapting the beam shaping to a reactor design is the production of an extremely distorted external field pattern (which is produced by the internal indentation coil in the present PBX-M). However, this type of external pattern can be produced by coils located outside the TF coil, depending upon the elongation. Some preliminary design work for reactor applications has been carried out by the MIT group (D. Cohn et al.).

K. TOI: With regard to IBW heating, what mechanism do you think caused the reduction of sawtooth activity? Could it be an impurity accumulation effect, or some other direct $j(r)$ profile change by RF?
N. SAUTHOFF: The visible bremsstrahlung radiation profile became flat with IBW, indicating that the density as well as the electron temperature profile were affected by IBW. At present, we do not know what caused the reduction in sawtooth activity; the q(r,t) measurement will be carried out when we resume operations.
THEORETICAL AND EXPERIMENTAL INVESTIGATION OF A MECHANISM FOR IMPURITY PRODUCTION BY ICRF FIELDS

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Abstract

THEORETICAL AND EXPERIMENTAL INVESTIGATION OF A MECHANISM FOR IMPURITY PRODUCTION BY ICRF FIELDS.

Two complementary theoretical approaches — simulations of the scrape-off plasma with a self-consistent single particle code and investigation of surface fields with a full wave antenna code — are used to provide a picture of impurity production during ICRF heating of the ASDEX tokamak plasma that is consistent with the experimental observations.

1. NUMERICAL SIMULATION OF HF SHEATH FORMATION

The most plausible mechanism for impurity production during ICRF heating is sputtering by ions accelerated in the non-linear HF sheath [1, 2], which forms where a wall cuts magnetic field lines on which the alternating flux from the antenna maintains an oscillating potential drop $\Phi_0$. In the sheath the ions gain an energy of the order of $e\Phi_0$ in a time $v_i/\lambda_D \ll \omega^{-1}$, where $\lambda_D$ is the Debye length, $v_i$ the ion velocity and $\omega = \mathcal{O}(\Omega_i)$ the applied frequency. This can greatly enhance the sputtering yield of the ions hitting the wall. To model this effect, we consider a plasma slab between two metallic plates immersed in a uniform static magnetic field $B_0$, making an angle $\psi$ with the x direction normal to the plates. An irrotational electric field, $E = -\nabla \Phi(x, t)$, is assumed to be maintained by an oscillating potential drop $\Delta \Phi = \Phi_0 \cos \omega t$ between the plates. The plasma behaviour in the combined electric and static magnetic fields is simulated with a particle-in-cell Monte Carlo code [2–4]. The code is one-dimensional in real space, but calculates the 3-D orbits of ions and electrons in velocity space together with the self-consistent electric field; it can thus describe sheath formation and ion acceleration at the plates consistently. Particles hitting the wall are assumed to be lost; the plasma is replenished by ambipolar sources of ions and electrons. The assumption $| \nabla \times E | \ll | \nabla \cdot E |$ is justified by the fact that the length scale of the longitudinal part of the field is of the order of the Debye length, while that of the transverse part is of the order of the antenna dimensions. Neglecting other space dimensions requires $\omega \ll \omega_{pi}$, the ion plasma frequency.
In the runs the particle sources $S$ were assumed constant in space and time, with a strength satisfying $S \cdot L / \cos \psi = 1.2 \times 10^{23} \text{ m}^{-2} \cdot \text{s}^{-2}$; this creates a time averaged particle flux of $\Gamma_0 = 6 \times 10^{22} \cos \psi \text{ m}^{-2} \cdot \text{s}^{-1}$ to each plate. The temperature of the source particles was taken as $T_0 = 10 \text{ eV}$. The static magnetic field was $B_0 = 2 \text{T}$; the angle $\psi$ with respect to the target normal was varied between $0^\circ$ and $85^\circ$. The potential difference between the two plates was assumed to oscillate at 1.3 times the proton cyclotron frequency, $f = 1.3f_{ci} = 39.7 \text{ MHz}$ (the HF sheath rectification mechanism is insensitive to the exact value of the frequency), with an amplitude of $\Phi_0 = 10 T_0/e = 100 \text{ V}$. Since the ion acceleration is localized in the sheaths, the total length $L$ of the system is not a relevant parameter, provided it is large compared to the sheath thickness. On the other hand, the collision mean free path was assumed comparable to $L$.

Figures 1(a) and (b) show the spatial distribution of the potential $\Phi$ and density $n$ in the plasma at a time when the right hand side plate is at its lowest potential.
The two figures differ in the orientation of $B_0$, which is, respectively, perpendicular ($\psi = 0^\circ$) and grazing to the plates ($\psi = 85^\circ$). At perpendicular incidence, the potential and density profiles are rather flat, except for the abrupt sheath in front of the negative plate, whose thickness $\lambda_{es}$ agrees well with that given by the Child-Langmuir formula, $\lambda_{es} \approx (e\Phi_0/T_{es})^{3/4} \lambda_D$, where $T_{es}$ is the electron temperature at the sheath edge. For grazing incidence, on the other hand, the potential and density profiles are more trapezoidal, the non-neutral sheath being preceded by a neutral magnetic sheath in which the ion flow adjusts from being sonic along the magnetic field to having a sonic component perpendicular to the target. The thickness of the presheath is $\lambda_{ms} \approx (E_x/B_0\Omega_i) \sin \psi$; the condition $\lambda_{ms} \geq \lambda_{es}$ gives an estimate of the angle at which this effect appears [4]. It will be noticed that because of the enhanced ion acceleration in the magnetic presheath the density at the Debye sheath edge is about four times smaller than in the case $\psi = 0^\circ$. In agreement with the Child-Langmuir formula, the non-neutral sheath thickness at $\psi = 85^\circ$ is thus about twice as large as in the previous case. In both cases, outside the sheaths the density stays practically constant in time.

Figure 2 shows the time behaviour of the energy flux $Q_x$ divided by the particle flux $\Gamma_x$ (i.e. the mean energy per particle hitting the target) to one plate for ions and electrons. At the potential minimum the ions carry the full potential drop $e\Phi_0$, plus an amount due to their energy at the sheath edge and to the slightly positive plasma potential with respect to the other target. The electrons, on the other hand, always feel a repulsive potential and carry only an energy of $Q_x/\Gamma_x = 2T_{es}$.

Figure 3 shows the flux of ions hitting a plate at the instant of minimum potential, resolved in energy $W$ and angle of incidence $\theta$ with respect to the normal. For $\psi = 0^\circ$, the velocity perpendicular to target is $\approx (2e\Phi_0/m_i)^{1/2}$, compared to a tangential velocity $(2T_i/m_i)^{1/2}$; thus the angular distribution peaks at $\theta \approx 18^\circ$ (Fig. 3(a)). When the magnetic field is oblique, its component $B_0 \sin \psi$ perpendicular to $E_x$ tends to bend the ion motion. The force due to the electric field is
larger than the Lorentz force, however, and the ions are predominantly accelerated perpendicularly to the wall, as long as $\omega_{pi}/\Omega_{ci} > (e\Phi_0/T_{ei})^{1/4} \sin \psi$. This inequality is satisfied for the parameters of the run at $\psi = 85^\circ$; accordingly (Fig. 3(b)), in this case the ions hit the wall at an angle that is only slightly larger than in the $\psi = 0^\circ$ case. By contrast, in the absence of an HF field, the average angle of impact is about $65^\circ$ [5].

2. NUMERICAL INVESTIGATION OF EDGE HF FIELDS

The most obvious region where a parallel HF field can be induced is along magnetic field lines in front of the ICRH antenna. In ASDEX, however, the main impurity production was not localized at the Faraday shield or the protection limiters of the antenna. Enhanced recycling and a strong influx of metal ions were observed from a region of the bottom divertor entrance plates, peaked toroidally at the antenna position and one to two metres wide: these domains are close to the antenna, but not linked to it by magnetic field lines [6].
To understand these experimental findings, we have made extensive numerical simulations of ICRH wave excitation with the FELICE code [7], which solves the finite Larmor radius wave equations in a plane stratified model of the tokamak. After imposing the matching conditions at the antenna and the wall for each partial wave with specified wavevectors \( k_y \) and \( k_z \) in the ignorable directions (poloidal and toroidal, respectively), Fourier synthesis allows the field to be reconstructed at a number of positions in the plasma and its surface; knowing the loading resistance of the antenna, we can also evaluate the field intensity as a function of the coupled power. Particularly useful in these investigations is the fact that the code allows for an arbitrary orientation of the antenna and the Faraday shield blades with respect to the static magnetic field.
Figures 4(a) and (b) show the amplitude of the parallel field $E_z$ at the plasma edge (the surface tangent to the Faraday screen) as a function of the central density at four poloidal positions between 0 and 75 cm from the equatorial plane (the half-length of the antenna is 49 cm, the plasma radius 40 cm). The frequency is 67 MHz and the magnetic field on axis 2.36 T, locating the second harmonic of hydrogen near the plasma centre. The density and temperature profiles were taken to be those measured at a central density of $n(0) = 5.1 \times 10^{13} \text{ cm}^{-3}$. In the first scan the Faraday shield blades were assumed aligned with the static magnetic field, in the second they make an angle of $\theta_F = 3.8^\circ$, corresponding to a cylindrical safety factor of $q_a = 3.5$. The two cases show a markedly different behaviour. At $\theta_F = 0^\circ$, the parallel field is strongly localized in front of the antenna: the largest value occurs at $h = 25 \text{ cm}$, half of the antenna length while, at $h = 75 \text{ cm}$, $E_z$ is almost one order of magnitude weaker. In contrast, at finite screen inclination the values of $E_z$ at all poloidal positions are comparable. The angular scan at $n(0) = 5.1 \times 10^{13}$ (Fig. 5) shows that this is true for angles larger than about 2°. The difference is clearly illustrated by the toroidal pattern of $E_z$ along the corresponding magnetic field lines at 0° and 3.4° (Figs 6(a) and (b)).

In Fig. 4(b), moreover, an appreciable enhancement of $E_z$ is seen in the range $4.5 \times 10^{13} \leq n(0) \leq 5.3 \times 10^{13}$; similar peaks are seen in Fig. 5. They are connected with the excitation of radially evanescent torsional-like surface waves with low toroidal and poloidal wavenumbers (typically, $n_\phi = \pm 4$, $m_\theta = \pm 1$, corresponding to wavelengths $\lambda$ somewhat shorter than $c/f$); the penetration depth of $E_z$ is a few centimetres. These high parallel fields give rise to strong collisional dissipation.

![FIG. 5. $E_z$ at the plasma edge versus Faraday shield inclination.](image-url)
FIG. 6. $\text{Re} E_z$ along magnetic field lines at various distances from the equatorial plane: (a) $\theta_F = 0^\circ$; (b) $\theta_F = 3.4^\circ$. 
in the scrape-off plasma (up to 20\% of the launched power) and to a parasitic contribution to the antenna loading, dominated by a single partial wave; otherwise, the field distribution and the power deposition profiles in the plasma core are not affected. The nature of these surface modes is further discussed in Ref. [8].

The strength of the resonances is likely to be overestimated by the simplified slab geometry adopted in the code. However, an enhancement in $E_z$ is present over domains of density and angles which are much broader than the sharp peaks observed in the loading resistance. Disregarding, therefore, the peaks, we find that the maximum potential drop along a field line, estimated to be $V_m \sim E_z \lambda/4$, is nevertheless several hundred volts at a coupled power of 1 MW, even 25 cm beyond the antenna tip. These results suggest that under these conditions the HF sheath acceleration mechanism can be effective also along magnetic field lines not passing directly in front of the antenna. The poloidal inhomogeneity of the tokamak is likely to restrict the region of strong $E_z$ to the outer half of the torus, hence to field lines which, even if not transiting in front of the antenna, are not too far from it in the toroidal direction. This is consistent with the experimental findings. On the other hand, because of the almost normal incidence of the accelerated ions on the wall suggested by the Monte Carlo code, it seems necessary to invoke self-sputtering to explain the observed impurity fluxes [6]. We might also mention that in the experiments feedback control of the density during the HF pulse becomes possible only above about the density at which the code predicts that the peaks in $E_z$ disappear.

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DISCUSSION

K. ITOH: The $z$-component of the induced electric field shows a peak at a particular density. Does this mechanism work for a particular density regime?

M. BRAMBILLA: Yes, we see 'resonant' excitation of these surface waves preferentially at low densities (less than $5 \times 10^{19}$ in the centre). However, the
strength of the resonances is likely to be overestimated by the slab geometry of the code. In the estimates we have used the field values away from the peaks, which are already sufficiently large.

L.J. PERKINS: An obvious suggestion for alleviating the problem of sputtering by Faraday screen elements is to align the elements along the magnetic field. How precise must this alignment be to reduce sputtering appreciably?

M. BRAMBILLA: If the surface waves seen in our code are responsible, the transition from the ‘aligned’ to the ‘misaligned’ situation occurs at an angle of about two degrees.

R. KOCH: In your model, do you take into account the finiteness of the screen blades?

M. BRAMBILLA: No, the screen is assumed to be ideal.
EXPERIMENTS ON STEADY-STATE TOKAMAK DISCHARGE BY LHCD IN TRIAM-1M

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Abstract

EXPERIMENTS ON STEADY-STATE TOKAMAK DISCHARGE BY LHCD IN TRIAM-1M.

Steady-state LHCD operation has been extended to pulse lengths beyond one hour in TRIAM-1M. Hall-effect sensors were used in the feedback loop for plasma position control. In the long duration discharge, a plasma current of 22 - 25 kA and a line-average electron density of about $2 \times 10^{12}$ cm$^{-3}$ were maintained nearly constant and no significant impurity accumulation was observed. High density ($\bar{n}_e \approx 1.0 \times 10^{13}$ cm$^{-3}$) current drive discharges were obtained by improving the accessibility of the plasma to LH waves. The current drive efficiency was improved by increasing the plasma density. The bulk plasma has better confinement properties than a low density OH plasma.

1. INTRODUCTION

During the past two and one-half years, experiments on the superconducting tokamak TRIAM-1M (R = 0.8 m, a x b = 0.12 m x 0.18 m and $B_t = 8$ T) have concentrated on exploiting the new operational regime of steady-state tokamak discharges and high density current drive with LH waves (2.45 GHz, 50 kW). The plasma discharge duration was extended up to 3 minutes by using a conventional method with magnetic coils and analog integrators for plasma position control [1,2]. In order to obtain a longer duration discharge, we had to develop a new plasma position control system in which the poloidal magnetic field can be detected directly rather than through the integrators. We have incorporated Hall-effect sensors (Hall generators) into the position control system and achieved tokamak discharges of one hour's duration by LHCD. In the steady-state tokamak discharges, it was very important to maintain the plasma density constant. This plasma density

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FIG. 1. Block diagram of plasma position control system with Hall generators. The plasma position is controlled with the feedback loop including the eight Hall generators, the digital and analog computer system and the vertical and horizontal field coil system.
control was carried out by the gas feed control system which regulated the neutral pressure including the recycling effect [2]. The main impurities in the plasma were molybdenum from the limiters and oxygen. The accumulation of these impurities was not observed during the long duration discharge.

High density current drive discharges have also been obtained through increasing the gas puffing rate. In these discharges, an abrupt drop of the plasma current and a flattening of the current profile were observed at a critical plasma density, which is closely related to wave accessibility [3]. By increasing the toroidal magnetic field, the accessibility of the plasma to LH waves was improved and the plasma density was raised up to $n_e = 1.0 \times 10^{13}$ cm$^{-3}$. The current drive efficiency increases with the plasma density and reaches up to 0.9 (10$^{19}$ m$^{-2}$ kA/kW). The stored energy of the bulk plasma also increases with the plasma density. This density dependence may be attributed to the improvement of the fast electron confinement.

In this paper, the tokamak discharge with the longest duration (one hour) in the world and high density current drive with plasma density of the order of $10^{13}$ cm$^{-3}$ are described. Finally, preliminary results on high frequency (8.2 GHz) LHCD experiments are also described.

2. STEADY-STATE CURRENT DRIVE OPERATION

In order to demonstrate steady-state tokamak operation on TRIAM-1M, we have made efforts to extend the discharge duration as long as possible and already succeeded in producing 3-minute discharges by controlling precisely the plasma position with magnetic coils and low drift analog integrators [1,2]. However, the conventional control method has an essential problem in that the offset voltage in the integrator output signal increases with time, and the fatal defect was an obstacle to achieving a longer duration discharge. Therefore, we attempted a new plasma position control method based on Hall generators, which can measure a magnetic field directly.

2.1 Plasma position control with Hall generators

Figure 1 shows a block diagram of the plasma position control system based on Hall generators. Four pairs of sensors are set up near the upper and the lower vertical ports to pick up the poloidal magnetic field produced by the plasma current. Here, the two sensors of each pair are fixed on a mounting arm so as to measure the magnetic components $B_\rho$ and $B_\theta$ in the two directions of minor radius and poloidal angle, respectively. The poloidal magnetic field is given by $\vec{B}_p = B_\rho \hat{e}_\rho + B_\theta \hat{e}_\theta$ with the unit vectors $\hat{e}_\rho$ and $\hat{e}_\theta$. Since the output signal of each sensor includes the toroidal magnetic component $\hat{e}_t$ according to
FIG. 2. Comparison of (a) plasma horizontal position $R_p$ and (b) plasma vertical position $Z_p$ detected by magnetic coils with those detected by Hall generators. These data were obtained through the horizontal and vertical scans of the plasma position, which was accurately controlled by use of magnetic coils in 10-second discharges.

the accuracy of the mounting angle and the poloidal magnetic component required for the plasma equilibrium, we compensated these components with the monitor signals of the toroidal and the poloidal coil currents, and extracted the component $B_p(I_p)$ due to the plasma current alone. By this measuring system, we can detect the current and the position of a noncircular cross-sectional plasma.

In 10-second discharges, we compared the plasma position calculated from the Hall generator signals with that estimated through the conventional method with magnetic coils. The results were in good agreement within 3 mm for the horizontal and the vertical positions, as shown in Fig. 2. Therefore, we can expect a precise plasma position control without any restriction on the discharge duration by use of this new control system.

2.2 One-hour tokamak discharge

The above-mentioned control system based on Hall generators allowed us to maintain a plasma current for more than one hour. This current drive experiment was performed in the low density region of $n_e < 2 \times 10^{12} \text{ cm}^{-3}$ under a low RF power (~ 15 kW) injection because of the restraint of the temperature at the plasma chamber. Figure 3 shows the waveforms of the longest duration (1 hour 10 minutes) discharge. The plasma current of about 22 kA is maintained nearly constant during the discharge by the RF power alone.
The plasma position is well controlled at $R_p = 0.825$ m and $Z_p = 0$ m within the displacement of 5 mm. Here, at the initial phase ($t \leq 3$ s) of the discharge, the plasma position was controlled with the magnetic coil system in order to obtain a fast response of the feedback loop, and afterwards the Hall generator system was used in plasma position control. This discharge was stopped by the suspension of the RF power injection because of the temperature rise in the ceramic break, which is a component of the plasma chamber.

The plasma properties in 1-hour discharges were studied mainly with respect to the behaviour of the impurity and the high energy electrons. The impurity behaviour was investigated by VUV spectroscopic measurements. The main impurities in the plasma were molybdenum from the limiters and oxygen resulting from recycling and desorption. Figure 4 shows the time evolution of Mo XIV and O VI line intensities. These intensities do not change significantly with time. This indicates that the content of the impurities in the plasma core remains at a quasi-stationary level without too large an adverse effect on the discharges. The concentrations of oxygen and molybdenum in the long duration discharge plasmas were about 2 % and 0.2 % of the electron density, respectively. The information on high energy electrons which carry the plasma current was obtained through hard X-ray emission measurements.
FIG. 4. Time variation of (a) molybdenum and (b) oxygen line intensities during stable discharges of more than one hour. These impurity lines are radiated mainly in the inner plasma region.

FIG. 5. Time evolution of hard X-ray emission profile for a photon energy of 100 keV during a 1-hour discharge. The centre of the chord position r corresponds to the major radius of 0.84 m.
Figure 5 shows the time evolution of the emission profiles in the long duration discharge. The hard X-ray emission profile is more peaked than the parabolic profile and remains stationary throughout the discharge. Therefore, the plasma current can be considered to have a relatively peaked and stationary distribution.

3. **HIGH DENSITY CURRENT DRIVE OPERATION**

The plasma density in the steady-state tokamak discharge is about $2 \times 10^{12} \text{ cm}^{-3}$ and lower than that in the normal Ohmic discharge in other tokamaks. In order to study the plasma properties of steady-state tokamak discharges in the next generation tokamaks (ITER, FER etc.), it is necessary to operate at higher plasma density of the order of $10^{13} \text{ cm}^{-3}$. Therefore, we have carried out experiments on high density current drive operation.

3.1 **Extension of operational regime**

The high density current drive discharge was attempted by increasing the gas puffing rate. In the early stage, an abrupt drop of the plasma current or a suspension of the discharge was always observed in the high density discharges and the maximum plasma density was limited to $n_e \approx 5 \times 10^{12} \text{ cm}^{-3}$. This density limit is due to the deterioration of the accessibility of the plasma to LH waves with low values of the parallel wavenumber [3]. In order to improve the accessibility of the high density plasma, the current drive discharge was carried out at higher toroidal magnetic field and consequently a high density discharge was obtained, as shown in Fig.6. The plasma density increases with increasing gas puffing rate and remains at the level of $n_e = 1.0 \times 10^{13} \text{ cm}^{-3}$. In this discharge, the abrupt drop of the plasma current and the remarkable flattening of the current and density profiles, which were observed in the case of the lower toroidal magnetic field, do not appear. However, when the plasma density exceeds $n_e = 1.1 \times 10^{13} \text{ cm}^{-3}$, such phenomena are always observed. Figure 7 shows the dependence of the maximum central electron density on the toroidal magnetic field and the accessibility condition at the plasma centre. From this figure, we can see that the operational density regime is extended by improving the accessibility of the central plasma to LH waves.

3.2 **Improvement of current drive efficiency**

In steady-state tokamak discharges with low density, the current drive efficiency is very low ($\eta = 0.2 - 0.3$). However, the efficiency is improved by increasing the plasma density, as seen in Fig.6. Figure 8 shows the dependence of the efficiency on the electron density. The current drive
FIG. 6. Time evolution of injection RF power $P_{RF}$, plasma current $I_p$, line averaged plasma density $\bar{n}$, and current drive efficiency $\eta$ for a high density discharge. The efficiency was calculated in consideration of the reflection power.

FIG. 7. Dependence of the maximum central electron density on the toroidal magnetic field in high density discharges. Broken line indicates the accessible $n_e$ of 1.2 at the plasma centre.
efficiency increases with the plasma density in spite of the deterioration of wave accessibility and comes to around 0.9, which is about one third of the efficiency estimated from the quasi-linear theory. The saturation level of the efficiency increases with the toroidal magnetic field. This density dependence seems to be due to the improvement of the confinement of high energy electrons. The direct loss of high energy electrons decreases with increasing plasma density because of the reduction of the collisional slowing-down time. As predicted by theory [4], the current drive efficiency increases with the confinement time of fast particles. Therefore, the efficiency is improved by increasing the plasma density. The saturation of the efficiency seems to be caused by the deterioration of fast electron confinement due to the poor accessibility condition of the high density plasma. In fact, the profile broadening of the high energy electrons was observed in the high density region.

3.3 Confinement of bulk plasma

In order to investigate the confinement of the bulk plasma in the high density current drive discharge, the electron and ion temperatures were measured by a multichannel Thomson scattering system and a neutral energy analyzer, respectively. Figure 9 shows the dependence of the central electron temperature on the electron density. The bulk electron temperature decreases with increasing plasma density and $T_e \sim 300$ eV in the high density region. The profile of the electron temperature changes from a peaked one to a broad one with increasing plasma density, as shown in Fig.10, and the electron density profile changes in the same way. Since the profile of
FIG. 9. Central electron temperature versus line averaged electron density for long duration LHCD discharges.

FIG. 10. Electron temperature profiles from Thomson scattering for long duration LHCD discharges. Profile becomes peaked with decreasing plasma density.

the high energy electrons shows a similar tendency, these profile changes seem to be caused by that of the power deposition from high energy electrons to bulk electrons. On the other hand, the bulk ions have a high energy tail component and it is difficult to determine the ion temperature from the ion energy spectrum. However, there is no doubt that
4. HIGH FREQUENCY LHCD EXPERIMENT

We have installed a new high frequency (8.2 GHz) LHCD system to carry out long pulse operation at higher plasma density of the order of \(10^{13} \text{ cm}^{-3}\). The RF system has eight klystrons with an output power of 25 kW CW and the total RF power is 200 kW. The launcher is composed of 8 x 2 waveguides, which have a cross-section of 4 mm x 35 mm. The high frequency LH waves cannot couple well to the low density target plasma which is produced in the present long pulse operation. Therefore, the experiments at 8.2 GHz were commenced by applying the LH waves to OH plasmas. Figure 11 shows a typical
combined OH-LHCD discharge. When the RF power was injected into a target plasma, the increase of the plasma current and the decrease of the loop voltage were observed as seen in current drive experiments. This current drive effect was supported by the significant increase in the hard X-ray emission. We estimated the RF-driven current by using an equivalent circuit for the combined OH-LHCD discharge and it was found that the driven current increased with the RF input power, as shown in Fig. 12. From this figure, the current drive efficiency can be estimated to be about 2, which is close to the theoretical one. If we can produce a high density target plasma with no negative loop voltage, we can expect to achieve long pulse operation with higher plasma density (~5 x 10^{13} cm^{-3}) and with higher current drive efficiency by using the 8.2 GHz RF system.

5. CONCLUSIONS

During the past few years, substantial progress has been made in expanding the operating range of long duration tokamak discharges of TRIAM-1M. A new method to detect the plasma current and position over a long time has been developed with
Hall-effect sensors and consequently 1-hour tokamak discharges by LHCD (2.45 GHz) have been achieved by precise plasma positioning and gas feed control. No significant impurity accumulation has been observed and the impurity concentration remains at a quasi-stationary level without an adverse effect on the discharge. High density current drive discharges with $n_e \approx 1.0 \times 10^{13} \text{ cm}^{-3}$ have been achieved by improving the accessibility of the core plasma to LH waves. At the same time, the improvement of current drive efficiency has been observed with increasing plasma density. High frequency (8.2 GHz) LHCD experiments have exhibited the effectiveness of high frequency LH waves in the higher density region. By using the 8.2 GHz RF system, it will be possible to achieve a long duration discharge with plasma density of the order of $10^{13} \text{ cm}^{-3}$.

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DISCUSSION

F.X. SÖLDNER: The current drive efficiency for the high frequency experiments was determined from a discharge where only a small fraction of the plasma was driven by LH. What would be the current drive efficiency for the lower frequency if you determined it from a discharge with the same partial current drive, i.e. $\Delta U_{\text{LH}}/U_{\text{LH}}^0$ equal in both cases?

S. ITOH: We also carried out experiments on a combined OH-LHCD discharge with 2.45 GHz RF waves. When the RF power is superimposed on an OH tokamak plasma, the current drive efficiency can be estimated to be about $1 \times 10^{-19} \text{ m}^2 \cdot \text{A} \cdot \text{W}^{-1}$. Thus, the efficiency in combined OH-LHCD discharges with 8.2 GHz RF waves seems to be better than that with the 2.45 GHz RF waves.
We think that this is due to the narrow power spectrum for the high frequency launcher with eight wave guides (the lower frequency launcher has only four wave guides).

B. COPPI: Another important upgrading of your experiment would be to utilize fully the toroidal magnet cavity and obtain high plasma currents with record values of the poloidal field. When do you plan to do this?

S. ITOH: We are now proceeding to design the upgraded TRIAM-1M machine. When the design study is completed, we will put forward a plan for upgrading the TRIAM-1M tokamak. Exactly when this can be done depends on the budget.
ICRF IMPURITY AND ANTENNA STUDIES*

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Abstract

ICRF IMPURITY AND ANTENNA STUDIES.

The paper summarizes recent work on three related topics: (i) Recent experiments and theory of metal impurity influx from the Faraday screens (FS) in JET are compared and show good agreement in both magnitude and scalings with key parameters. The theoretical model describes the acceleration of plasma ions by induced RF sheaths on the screen and the resulting sputtering of the screen material. It is shown that metal impurities can be reduced to negligible levels by proper antenna design and the use of a low atomic mass screen coating. (ii) Experiments on RF sheath formation at the FS of a fast wave ICRF antenna in Phaedrus are described. It is shown that the radial feeders of a typical two-strap fast wave antenna induce a large voltage on the FS with respect to an adjacent limiter. The magnitude of the voltage varies with strap phasing. Experiments with a low power single strap antenna mounted in the Phaedrus-B central cell demonstrate RF self-biases at the FS. (iii) A computational investigation of the vacuum properties of ICRF antennas is reported on. The numerical simulation code, ARGUS, is described briefly and results from a study of the PPPL TFTR ICRF antenna are given.

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1. IMPURITY PRODUCTION BY ICRF ANTENNAS IN JET

The release of impurities during ICRF heating can be attributed to two kinds of mechanism: *global* production caused by additional power flow in the scrape-off layer (SOL) and *local* release from the vicinity of the antenna. The latter process is particularly important in understanding the origin of metal impurities in ICRF experiments with good single pass absorption. Here we describe a comparison of recent experiments [1] and theory [2] on the metal impurity influx from the JET Faraday screens (FS), which are the major source of metal impurities in JET ICRF experiments. The main conclusions from this study are that (i) metal impurities can be reduced to a negligible level by proper antenna design, and (ii) an RF sheath based impurity model provides a framework for understanding the local impurity release and for designing optimal antenna systems in ICRF experiments.

Different antenna designs and FS coatings have been used on JET. Substantial local impurity influxes, seen in early operation, have been reduced to negligible levels in the latest series of experiments. The experimental and modelling results suggest that RF sheaths are responsible for the FS impurity release. These RF sheaths [3] are caused by RF flux linkage through circuits that include magnetic field lines in the plasma which intersect the conducting surfaces of the FS gaps and front face. The sheaths non-linearly rectify the induced RF voltage, and the rectified (dc) potential accelerates ions into the FS with sufficient energy (typically, keV) to cause physical sputtering of impurities. The impurity release model [2] couples SOL, RF sheath and sputtering physics self-consistently. The model includes sputtering by the majority species, one species of limiter impurity, and self-sputtering by the FS material.

The measured metal influx on JET depends on the antenna voltage, the edge density, the magnetic field angle, the phasing of adjacent current straps, the shape of the screen and the FS coating, as well as the extent of limiter conditioning and the impurity content of the SOL plasma. Many of the observed dependences are straightforward consequences of the model:

(a) For FS materials with high atomic mass, the model predicts a large increase of the metal influx at high voltage due to high energy self-sputtering and sputtering by heavy impurities. This prediction is in qualitative agreement with data [4] from a Cr-plated FS, which indicates an influx that is two to six times larger than that obtained in the recent Be evaporation experiments [1].

(b) In dipole phasing, the model predicts that sheath potentials on the front face of the A₁ (V-shaped) FS are greatly reduced because of symmetry, which annihilates the net RF flux driving the sheaths, leaving only a smaller influx induced by gap sheaths, consistent with the observations.

(c) The calculated influx in monopole phasing increases with FS angle (relative to the equilibrium B field) because more RF flux is captured in the circuit containing the sheaths. This is in qualitative agreement with reversed field experiments on JET [5].
The model calculations suggest that self-sputtering contributed substantially to the observed influxes on JET with Cr and Ni screens.

The impurity release model also shows good quantitative agreement with recent Be evaporation experiments [1], where the FS was coated with Be (Fig. 1). The Be influx from one screen was measured spectroscopically, along a line of sight intercepting the screen, by means of the BeI (4573 Å) line intensity. These measurements were obtained for monopole phasing with C limiters and with Be limiters. An absolute calibration of the Be influx $\Phi_{Be}^{FS}$ for the C limiter case was obtained from erosion rate estimates [1] ($(1-2.5) \times 10^{19}$ atoms $\cdot$ MW$^{-1}$ $\cdot$ s$^{-1}$) and from Be concentration measurements ($(2-4) \times 10^{19}$ atoms $\cdot$ MW$^{-1}$ $\cdot$ s$^{-1}$). The spectroscopic Be influx data is illustrated in Fig. 1 for the normalization $2 \times 10^{19}$ atoms $\cdot$ MW$^{-1}$ $\cdot$ s$^{-1}$, which represents a good compromise between these two experimental methods, with a power level of 1 MW per antenna corresponding to an antenna line voltage of 18 kV. The experimental data illustrate the important voltage and (indirectly) density dependences of the impurity influx.

The theoretical model [2] and the spectroscopic data both yield the neutral impurity influx at the FS before screening of the impurities by ionization occurs. The solid curves in Fig. 1 are the result of detailed calculations in which we have used the data and model together to infer relationships among the antenna voltage, the edge density and the SOL diffusion coefficient. The inference from this procedure is that the voltage dependence of the experimental influx is primarily due to the increase in

![Figure 1](image-url)

*FIG. 1. Be influx (measured by the intensity $I_{Be}$ at the FS) versus antenna line voltage $V_1$ for the two cases of C (•) and Be-belt (○) limiters with monopole phasing. The data are normalized to the estimated flux $\Phi_{Be}^{FS} = 2 \times 10^{19}$ atoms $\cdot$ MW$^{-1}$ $\cdot$ s$^{-1}$ and the data for the C limiter case are corrected for erosion decay. The solid lines represent the theoretical model calculation.*
edge density with RF power (which has been determined from probe data) and to the increase in the local diffusion coefficient with RF voltage (assumed here to be linear). The sputtering yield of Be is insensitive to energy (and hence to antenna voltage) in the range corresponding to Fig. 1. The lower influx with Be limiters was modelled by assuming a reduced rise of edge density with RF power and a reduced fraction of limiter material available to sputter the screen (motivated by the observed gettering of O and pumping of D by the Be limiters).

2. ICRF ANTENNA STUDIES IN THE PHAEDRUS PROGRAMME AND NUMERICAL MODELLING OF TFTR ANTENNAS

2.1. Phaedrus studies

RF sheaths at the Faraday shield are now suspected to be responsible for impurity production during ICRF heating experiments. Attention first focused on sheaths in the gaps between FS elements [3]. It has recently been shown for JET that more important sheaths occur along field lines which traverse the face of the JET FS (see the preceding section of this paper and Ref. [2]). Here, we discuss an additional mechanism of sheath generation. We show that the azimuthal component of the B-dot fields generated by the radial feeders of a fast wave antenna can generate inductive RF voltages, and therefore rectified potentials (self-biases) between the FS and the nearest limiter.

The mechanism is illustrated in Fig. 2(a). An inductive voltage is generated along the path ABCD, owing to the radial feeder fields. The net flux through ABCD is clearly a function of antenna strap phasing. An in-air demonstration of this mechanism has been performed by using the two-strap antenna constructed for ICRF experiments in the Phaedrus-T tokamak. A capacitive probe is used as the voltage sensing element, and a conducting path is chosen to mimic the FS-wall-limiter geometry illustrated in Fig. 2(a). Results for three antenna phasings are shown in Fig. 2(b). Note that the inductive voltage is reduced by a factor of four as phasing is varied from 0° to 180°. The voltage peaks at the ends of the antenna straps, where the radial feeders are located. This antenna–limiter geometry is similar to the TFTR Bay M antenna–RF limiter geometry and is a possible explanation for the discharge phenomena observed on the Bay M antenna during in-phase operation of that antenna [6].

We have also measured the plasma floating potential near the FS of a low power (≤3 kW) single strap model antenna using the central cell of Phaedrus-B ($n_e(0) \leq 5 \times 10^{12} \text{ cm}^{-3}$; $T_i, T_e \sim 20-40 \text{ eV}$) as a source plasma. The antenna operates at 2.7 $\Omega_{ci}$ or 3.5 MHz. The FS bars of the model antenna are tilted 3° with respect to the static magnetic field to allow probe measurements on field lines which terminate on the bars (see Fig. 3(a)). The shield is in the limiter shadow; field lines originating on the FS terminate on the limiter. In Fig. 3(b), we show a plot of the increase in
floating potential when the model antenna is excited, as a function of probe position as the probe is scanned across the FS gap region. The peak increase in floating potential occurs when the probe is located on a field line which terminates at the centre of the FS. Only small changes in the floating potential are seen when the probe is on a field line which transits the FS gap without terminating on the shield structure.

The increase in floating potential produced by the RF scales linearly with the antenna strap current. The (grounded) FS collects electrons during antenna operation from the biased scrape-off layer plasma, as measured with a current transformer. A frequency spectrum of the collected current shows strong non-linear generation of harmonics, with the harmonic at $2\omega_{RF} < 3$ dB down from the fundamental. Strong harmonic generation has also been seen in FS experiments and capacitive probe measurements in the SOL of TEXTOR [7] during ICRF heating.
2.2. TFTR modelling

The successful application of RF heating to future fusion plasmas generated in CIT and ITER size devices will require power levels that are much higher than those in use today and, consequently, the design of suitable antennas will require the ability to predict the antenna performance accurately. We report here on progress made in developing a modelling and simulation capability for the vacuum characteristics of ICRF antennas that will play a role in future antenna design and evaluation.

The SAIC ARGUS system of codes [8] used for the ICRF antenna studies presented here is modular and contains both electromagnetic and electrostatic field solvers in the frequency and time domain, a PIC module and a fluid plasma module
that self-consistently contribute to the field solver, as well as many other features that model realistic physics of devices such as secondary emission from surfaces. ARGUS allows fairly arbitrary and complex geometrical structures to be simulated and, by means of block decomposition techniques, large data structures resulting from finely gridded computer models can be accommodated.

Experimental benchmarking of ARGUS applied to ICRF antenna configurations has been carried out in collaboration with the University of Wisconsin. The experimentally determined vacuum characteristics of a model antenna geometrically similar to the TFTR design were compared with the results from an ARGUS simulation of the laboratory test apparatus. The agreement between the computed and measured results was excellent. It was found that the currents flowing in the antenna strap and other parts of the antenna configuration (such as the backplane on which the test apparatus was mounted) should be calculated self-consistently for the actual antenna radiation pattern to be captured. In fact, because of the finite three-dimensionality of the experimental test set-up, certain features of the experimentally observed results could not be understood without the simulation model results.
FIG. 5. (a) Toroidal RF field ($B_{\text{RF}}^T$) profile in front of the current strap at midplane ($0-\pi$ phasing); (b) toroidal RF field ($B_{\text{RF}}^T$) profile in front of the Faraday shield at midplane ($0-\pi$ phasing).
More recently, ARGUS has been used to model the full PPPL antenna design for TFTR. Figure 4 shows a view of a slice through the upper right quadrant of 1/4 of the double strap antenna. Shown are the current strap, the power lead and the bottom row of Faraday bars. Also shown are the end and side walls. Not shown, but included in the model, is the centre septum. Simulations of this antenna model have shown a strong similarity to the results of the UW model antenna. The double strap antenna has been simulated for both symmetric and antisymmetric current phasing. Noticeable differences between the two phasings are present in the simulation results, but it is not known at present how these affect antenna performance. An example of the results from the present study is shown in Fig. 5. Figure 5(a) shows the toroidal RF magnetic field measured in the toroidal direction at the front of the strap at the midplane for antisymmetric phasing; Fig. 5(b) shows the same information measured in front of the FS. The effect of the actual antenna geometry is clearly displayed in the curves shown in Fig. 5. Other information available from the simulations includes all RF electric and magnetic fields generated, strap impedances, coupling coefficients, power spectra as a function of toroidal and poloidal wavenumber on the same surface, strap current distribution both along and around the strap, and all these quantities as functions of input parameters such as port voltage and frequency. ARGUS can model most types of antennas including those used for IBW heating.

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MACROSCOPIC AND MICROSCOPIC EFFECTS OF ELECTRON CYCLOTRON HEATING IN TEXT

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Abstract

MACROSCOPIC AND MICROSCOPIC EFFECTS OF ELECTRON CYCLOTRON HEATING IN TEXT.

Changes in the electron temperature, density and density fluctuation profiles were measured during electron cyclotron heating, at the fundamental resonance, in the Texas Experimental Tokamak (TEXT). Density fluctuations appeared to be correlated with edge conditions and independent from local changes in the electron temperature. The modification of the electron distribution function was obtained by solving the kinetic equation, and the corresponding X ray and electron cyclotron emission spectra were computed and compared with the experimental data. Both calculations and measurements, in the case of central heating, show the electron distribution to be non-Maxwellian near the resonance.

1. MACROSCOPIC MEASUREMENTS AND DENSITY FLUCTUATIONS

1.1. Experimental conditions

ECH heating experiments were undertaken in the TEXT tokamak ($R_0 = 1 \text{ m}$, $a = 0.26 \text{ m}$). A single 200 kW, 60 GHz gyrotron was used as the microwave source. The gyrotron was connected to one of two transmission lines, which transmitted the TE\textsubscript{01} mode to the tokamak vicinity. One line was connected to an unfocused antenna which launched a TE\textsubscript{11} mode with near O-mode polarization at an angle of 100° from the toroidal field. The other line was connected to a focused launcher which could be scanned in the vertical (poloidal) direction. The vacuum spot size ($1/e$) was approximately 9 cm in radius for the unfocused and 3 cm for the focused launcher.

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The focused launcher was oriented 80° from the toroidal field. The electron drift direction due to the Ohmic heating electric field was antiparallel to the toroidal field direction. For the results reported in this paper, the plasma current before ECH was typically 120 kA, the line averaged density was $1.5 \times 10^{13}$ cm$^{-3}$ and the central $T_e$ was 700 eV. The toroidal field was 21.6 kG in the centre of the vacuum vessel, which placed the $\omega_{ce}$ resonance on the magnetic axis. The limiter $q$ was 5.5, and the plasma exhibited very weak sawtooth oscillations.

1.2. $T_e$, $n_e$, and $\delta n_e/n_e$ profile modification by ECH

The gyrotron system delivers approximately 150 kW of power to the plasma with the O-mode polarization. The application of this amount of ECH power sharply peaks the electron temperature profile when the resonance is located on the magnetic axis. The central temperature nearly doubles while, outside $r/a = 0.3$, the increase is typically <30%. Figures 1(a), (b), and (c) show the density, temperature and density fluctuation profiles before and during the ECH pulse. The 'density pumpout'
which normally occurs during ECH is counteracted by increasing the gas puff level during the ECH pulse. The density profile is only slightly flattened compared to the pure Ohmic profile. The relative density fluctuations $\delta n/n$ are measured by the heavy ion beam probe (HIBP). The $\delta n/n$ profile increases uniformly by 25–30% across the centre half of the plasma. This is contrasted with the dramatic sharpening of the $T_e$ profile which occurs over the same region of space and suggests that the fluctuations as measured by the HIBP are not correlated with the local changes in the electron temperature.

Further evidence of the decoupling between the HIBP measured $\delta n/n$ and the local $T_e$ was obtained by varying the ECH heating location by changing the toroidal field. Fluctuation profiles were obtained for discharges with $B_t = 21.6, 20.3, 19.3, 17.2$ and $15.2$ kG, which placed the $\omega_e$ resonance at 0.5, 10, 20 and 30 cm from the magnetic axis on the high field side of the plasma. At 15.2 kG the resonance is well outside of the limiter radius, and no changes in $T_e(0)$, edge $T_e$, or $\delta n/n$ were measured for this case. Figure 1(d) shows the measured changes in the central density fluctuations $\delta n/n$ and the $T_e(0)$ increase as measured by ECE as a function of $B_t$. Note that $\Delta T_e/T_e(0)$ decreases rapidly as the toroidal field is decreased and the resonance is moved away from magnetic axis. In contrast, the increase in central $\delta n/n$ reaches a local maximum at $B_t = 19.2$ kG and then decreases as the resonance is moved further away from the axis. The relative changes in the edge temperature and floating potential fluctuation levels are shown in Fig. 1(e) as functions of $B_t$. It is clear from Figs 1(d) and (e) that the changes in the HIBP measured $\delta n/n$ scale more closely with the changes in edge temperature and fluctuations than with the changes in central temperature. During modulated ECH experiments the time evolution of the HIBP $\delta n/n$ signal during the short (5 ms) ECH pulses follows the time evolution of the edge $T_e$ more closely than $T_e(0)$. From these experiments we conclude that the central fluctuations as measured by the HIBP are not driven by the changes in the centre temperature, but appear to be correlated to changes in the plasma edge conditions.

2. DETERMINATION OF THE ELECTRON DISTRIBUTION

2.1. Solution of the kinetic equation

The electron distribution function is computed numerically for resonance with a finite amplitude electron cyclotron wave packet. A self-consistent Fokker-Planck quasi-linear code is employed, computing the coupled evolution of the distribution function and of the wave spectrum, as obtained from the relativistic wave damping, the time dependent distribution and the ray equations. The code is 2-D in momentum space and in wave vector space, 1-D in real space. It includes the following terms: the Beliaev and Budker collision term for a non-relativistic bulk; the relativistic quasi-linear term; the convective Ohmic field term; a loss term allowing relaxation
FIG. 2. Momenta of electron distribution function at different radial positions: (a) parallel function; (b) perpendicular temperature; (c) fraction of suprathermal electrons with energy larger than E; (d) fraction of suprathermal electrons versus radial position.

In Fig. 2, we present a typical distribution function computed for co-injection of the microwaves with the direction of the electron drift, and resonance on axis. Non-thermal electrons are generated by the perpendicular diffusion in momentum space induced by the electron cyclotron waves, and by the parallel acceleration due to the Ohmic field, which acts more effectively on the higher perpendicular energy, less collisional electrons. The collisional drag creates a steady state in a few milliseconds. The distribution function has a complex two-dimensional form in momentum space and is a function of the radial position. It is easier to visualize the structure of the distribution by considering its momenta. In Fig. 2(a), the parallel function $F_{\parallel}$, i.e. the distribution integrated over the perpendicular momentum, is plotted versus the normalized parallel momentum, $u_{\parallel} = p_{\parallel}/(m_e T_e)^{1/2}$. The parallel tail appears enhanced, especially in the direction of the electron drift, owing to the effect of the Ohmic field and to the pitch angle diffusion of the perpendicular energy. In Fig. 2(b), the average perpendicular energy is plotted. The perpendicular energy is higher for electrons with negative momentum since a larger part of the microwave...
power is coupled to these electrons, and because part of the fast parallel tail energy is converted to perpendicular energy by pitch angle scattering. The energy distribution of the non-Maxwellian part of the distribution function is plotted in Fig. 2(c); \( N(E) \) represents the number of suprathermal electrons with energy higher than \( E \). Since the non-Maxwellian part of the distribution has zero density \( N(0) = 0 \), the maximum of \( N(E) \) can be used as definition of the number of suprathermal electrons (Fig. 2(d)).

2.2. Emission spectra

Non-Maxwellian distribution functions are well characterized by their emission spectra. The form of the spectrum can be related to the modification of the distribution induced by the resonant interaction with the cyclotron waves and used to investigate the underlying kinetic process. Here we summarize the main results.

![Image of computed emission spectra](image)

**Figure 3.** Computed emission spectra for the conditions of Fig. 2: (a) electron cyclotron emission from the low field side; (b) from the high field side, for wall reflection coefficients \( R = 0.85, 0.90, 0.95 \). Thomson scattering temperature profile is also shown \( (T_e) \); (c) average electron temperature for electrons of energy between 10 and 18 keV, as obtained from the computed X ray spectra, versus radial position, as a function of line averaged density. Observed emission spectra; (d) high field side electron cyclotron emission (filled circles), low field side (open circles), and Thomson scattering profile; (e) suprathermal temperature measured by soft X ray near the plasma axis, for electrons of energy between 10 and 18 keV, versus line averaged density.
2.2.1. Electron cyclotron emission

The radiation from the X-mode at the second harmonic cyclotron frequency was measured with a ten channel grating polychromator from both the high and low field sides. Non-thermal spectra were observed: the high field radiative temperature was 100 eV (unfocused antenna case) to 200 eV (focused antenna case), higher than the low field temperature, for frequencies corresponding to locations from 0 to 10 cm on the low field side of the magnetic axis. Both spectra were non-thermal for frequencies corresponding to the outer edge of the plasma.

2.2.2. X ray emission

Soft X ray spectroscopy was also used to investigate the shape of the electron distribution function during heating with the unfocused antenna. The detector, a HPGe crystal of 5 mm thickness, is positioned to view the plasma along a vertical chord. Tail temperatures of the order of 6 keV and 3.5 keV were observed for co- and counter-injection of the microwave with the electron drift, in plasmas with a higher impurity content \( V_{\text{loop}} = 1.5 \text{ V}, I_p = 120 \text{ kA} \). Tail temperatures of the order of 2.5 keV were measured in cleaner plasmas \( V_{\text{loop}} = 1 \text{ V}, I_p = 120 \text{ kA} \), during co-injection. A density scan revealed a \( 1/n_e \) dependence of the suprathermal temperature.

2.2.3. Discussion

Agreement (typically, within the 10 to 15% experimental error bars) was found between the computed and measured spectra, supporting the interpretation of these results as due to a non-Maxwellian population located near the plasma axis, with the general characteristics described in Fig. 2. The conclusions are exemplified in Fig. 3, where some computed and measured spectra are shown.
APPLICATION OF INTERMEDIATE FREQUENCY RANGE FAST WAVE TO JIPP T-IIU AND HT-2 PLASMAS

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Abstract
APPLICATION OF INTERMEDIATE FREQUENCY RANGE FAST WAVE TO JIPP T-IIU AND HT-2 PLASMAS.

Fast waves well above the fundamental ion cyclotron frequency were applied to JIPP T-IIU and HT-2 tokamak plasmas to study current drive, electron heating and third harmonics ion cyclotron heating. In HT-2, up to 10 kA of plasma current was driven by a 100 MHz fast wave in the density range of \((0.3-3) \times 10^{18} \text{ m}^{-3}\), which corresponds to one to ten times above the lower hybrid current drive density limit. In JIPP T-IIU, an efficient electron temperature increase was observed on application of 130 MHz fast waves in the density range of \((2-3) \times 10^{19} \text{ m}^{-3}\), two orders of magnitude higher than the lower hybrid density limit. In a further, still higher density range of above \(3 \times 10^{19} \text{ m}^{-3}\), the third harmonics cyclotron heating of the additionally heated target plasma yields a heating efficiency that is as high as that of fundamental harmonic heating.

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1. INTRODUCTION

The intermediate frequency range between third cyclotron harmonic frequency and the lower hybrid frequency is relevant to thermonuclear fusion research in several aspects. First, fast wave current drive in this frequency range has a potential for bulk current drive in hot and dense reactor plasmas which lower hybrid current drive cannot provide [1]. Second, third cyclotron heating can increase Q through selective wave damping on high energy tritium ions. Third, the coupling of the waves to alpha particles could offer a method of alpha particle exhaust control. Fourth, this frequency range allows the use of waveguide launchers.

Research work done on this intermediate frequency range has been reported. As to the fast wave current drive applied to standard tokamak plasmas [2-4], electron Landau damping of the fast wave was confirmed in a plasma with a density well above the lower hybrid density limit although a concrete demonstration would require further investigation. As to third harmonics heating, few experimental results are available, because of the weak damping. This heating scheme is, however, expected to work better as the plasma parameters are improved and may become the major scheme applied to reactor grade plasmas. In the present paper, the results of fast wave current drive and third harmonics ion cyclotron heating experiments in JIPP T-IIU and HT-2 tokamak are reported.

2. FAST WAVE CURRENT DRIVE

HT-2 is an ion core tokamak with a major radius of 0.45 m, a minor radius of 0.12 m and a toroidal field of 2 T, at maximum. The experiment is conducted in a deuterium slide-away plasma. Three Faraday shielded loop antennas are used to generate travelling waves with a parallel refractive index $n_i = 5$, by keeping the phase difference between adjacent antennas at 120°. Extensive titanium flashing is applied to suppress impurity release from the antenna surface. The time evolution of the plasma parameters in two typical discharges is shown in Figs 1(a) and (b). The former describes the evolution at the relatively low density of $0.3 \times 10^{18} \text{ m}^{-3}$ ($n_e \sim n_{cem}$) and the latter refers to the high density of $1.5 \times 10^{18} \text{ m}^{-3}$ ($n_e \sim 5n_{cem}$), where $n_{cem}$ is the lower hybrid current drive density limit [5]. From case (a), it is seen that, with 25 kW of RF power, the loop voltage drops below zero and the transformer current stops increasing, on the application of RF power at 30 ms. In case (b), the loop voltage also drops to zero on application of 100 kW RF power. However, the density rise due to the impurity release from the antenna surface, as seen by the visible light signal of the C-III line collimated on the Faraday shield, limits the time duration of the loop voltage drop to only 2 ms. The hard X ray signal integrated over the energy range above 30 keV using the NaI scintillator placed on the horizontal port increases substantially on application of RF power, indicating high energy electron generation.
FIG. 1. Time evolution of plasma parameters for two typical shots in HT-2: (a) electron density $n_e = 0.3 \times 10^{18} \text{ m}^{-3}$; (b) $n_e = 1.5 \times 10^{18} \text{ m}^{-3}$; (c) density dependence of generated current per unity RF power, $I_{\text{gen}}/P_{\text{RF}}$. 
The dependence of the current generated per unit RF power, $I_{RF}/P_{RF}$, on the average electron density is shown in Fig. 1(c). The data points closely fit the theoretical solid line corresponding to $I_{RF}/P_{RF} \sim 1/n_e$. The current drive efficiency $I_{RF}n_eR/P_{RF}$ is found to be 0.01 A·W$^{-1}$·m$^{-2}$. According to the empirical database [5], the density limit $n_{crit}$ at 100 MHz is $0.3 \times 10^{18}$ m$^{-3}$. The efficiency does not degrade up to a density of $3 \times 10^{18}$ m$^{-3}$, which is ten times above $n_{crit}$. Beyond this density, the loop voltage drop could not be observed because of the strong impurity release from the antenna surface due to the increased RF power.

In JIPP T-IIU, a frequency of 130 MHz was chosen to assess the absorption mechanism to electrons and ions. The toroidal field is around 3 T, which places the layer $\omega = 3\omega_{CH}$ close to the plasma centre. Deuterium is used as working gas, with an admixture of 10% hydrogen. A four element toroidal array antenna is placed on the inboard (high field) side. In a plasma density range of up to $(2-3) \times 10^{19}$ m$^{-3}$, electron heating by fast waves has been observed. Figure 2 displays the typical time evolution of the plasma parameters. It is observed that the electron temperature rises.
on application of 200 kW fast waves. The density increases owing to enhancement in recycling from carbonized materials, degrading the heating efficiency in the density range above $3 \times 10^{13} \text{ m}^{-3}$. The soft X ray spectra indicate the formation of a high energy electron tail above 10 keV.

This density regime is relevant to fast wave current drive; the wave interacts with electrons at densities that are two orders of magnitude higher than the density limit $n_{\text{crit}}$ predicted for lower hybrid current drive [5]. The electron heating in this regime has been identified with ECRH heated plasma in previous work [4]. The present experiment revises the ratio of $n_e/n_{\text{crit}}$ by an order of magnitude. It is

![Graphs showing energy evolution and energy distribution](image)

**FIG. 3.** Results of third harmonic ion cyclotron heating in JIPP T-IIU: (a) time evolution of stored energy; $P_{\text{NBI}} = 300 \text{ kW}$, $P_{\text{40 MHz}} = 350 \text{ kW}$ and $P_{\text{130 MHz}} = 400 \text{ kW}$; (b) energy distribution of hydrogen ions; (c) power dependence of stored energy.
assumed that this is possible because of the relatively high Ohmic electron temperature in JIPP T-IIU, as a result of the relatively high toroidal field which allows a high plasma current density.

3. THIRD HARMONICS CYCLOTRON HEATING

Electron heating dies out as the plasma density is increased and another regime appears. In the density range above \((2-3) \times 10^{13} \text{ cm}^{-3}\), third cyclotron harmonics damping dominates. This regime is most clearly established when 130 MHz are superposed on NBI and ICRF (40 MHz; two ion hybrid heating regime).

The time history of a typical shot is shown in Fig. 3(a), in which the stored energy increases further as 130 MHz are applied, demonstrating that third cyclotron heating really works in this combined heating. In Fig. 3(b), the ion energy distribution function (with and without 130 MHz) obtained from a fast neutral analyser is shown. It is seen that a high energy hydrogen tail above 7 keV is produced by the application of 130 MHz. Without preheating with NBI and 40 MHz, a 130 MHz fast wave could not cause any appreciable increase in the stored energy. This improved heating efficiency, as it is combined with other heating methods, is interpreted as one aspect of the finite Larmor radius effect. In Fig. 3(c), the stored energy dependence on the total power is compared between the combined 130 MHz case and the 40 MHz only heating case. It is found that the heating efficiency of the combined heating is comparable to that of the established 40 MHz ICRF heating. We note that, since wave absorption in the 130 MHz case is substantial, heating on the basis of the third cyclotron harmonics may be viable as a new heating regime. In this case, Q-enhancement through tritium ion tail production and the potential use of a waveguide antenna may be regarded as merits.

REFERENCES

SIMULATION OF NON-LOCAL EFFECTS FOR RF PLASMA HEATING IN AN OPEN TRAP

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Abstract

SIMULATION OF NON-LOCAL EFFECTS FOR RF PLASMA HEATING IN AN OPEN TRAP.

The paper is concerned with numerical modelling of ion cyclotron heating of the open trap plasma. The heating scheme investigated makes use of fast magnetosonic wave (FMSW) global resonances in deuterium plasma (the calculations were made for the TE$_{102}$ mode). The FMSW energy absorption was due to cyclotron resonance for hydrogen minority ions. The minority ion cyclotron zone was situated at the RF field bunches for the FMSW global resonance. The comparison is made between the FMSW absorption when the non-local effects due to minority ion thermal motion are taken into account and the absorption calculated in the local ('cold') approximation. The relative global resonance width is used as a quantitative characteristic of the FMSW absorption value. The resonance width dependence on minority concentration is studied. Considerable differences in this dependence and in field pattern between local and non-local approximations are discovered. This is due to suppression of local Alfvén resonance in the vicinity of the hydrogen minority cyclotron zone when the non-local effects are taken into account.

Reference [1] considered the possibility of RF heating of plasma in an open trap using global resonances of the fast magnetosonic wave (FMSW). The FMSW damping and the finite width of global resonances are caused by the cyclotron absorption of RF field energy by minority ions.

It is necessary to note that in homogeneous plasma the FMSW cyclotron damping is absent for the cold minority ions because the direction of the FMSW electric field vector rotation is opposite to the cyclotron rotation ($E_+ = 0$). The damping appears when one takes account of the thermal motion of particles, owing to the fact that the cyclotron resonance condition $\omega = \omega_{ci} + k_{\parallel} v_t$ is fulfilled only for a definite group of ions [2]. In this case the material equation $D_+ = D_+ (E_+)$ does not contain the pole at $\omega = \omega_{ci}$ ($D_+$ is the left hand polarized component of the electric induction vector).

In inhomogeneous plasma the FMSW cyclotron minority absorption may take place also in the cold plasma approximation. In Ref. [1] global FMSW resonances
in an open trap have been studied and considerable 'cold' cyclotron damping for minority concentrations of 1-3% has been discovered. This paper is concerned with a similar problem taking into account the thermal motion of minority ions. The thermal motion effect on the FMSW absorption and the global resonance width has been investigated.

Taking account of the ion thermal motion leads to the non-local coupling between the electric induction vector $\vec{D}$ and the electric field $\vec{E}$. The ion minority contribution in $D_\pm$ components may be presented in the form

$$D_\pm = -2\pi i \frac{\omega_{pl}^2}{\omega^2} \int_0^\infty dv_\perp v_\perp^2 \frac{\partial f_\perp(v_\perp)}{\partial v_\perp} \int_0^\infty \frac{dv_\parallel}{v_\parallel} f_\parallel(v_\parallel) \left\{ \int_s^L ds' E_\pm(s') \right\}$$

$$\times \exp \left[ i \frac{v_{IT}}{v_1} \Psi_\pm(s,s') + \int_s^L ds' E_\pm \exp \left[ -i \frac{v_{IT}}{v_1} \Phi_\pm(s,s') \right] \right\}$$

(1)

where

$$\Phi_\pm = \frac{1}{v_{IT}} \int_s^{s'} \left[ \omega_\parallel(s'') \mp \omega \right] ds''$$

(2)

Here $v_{IT}$ is the longitudinal thermal velocity of minority ions, $F = n_0 f_\perp(v_\perp)f_\parallel(v_\parallel)$ is the distribution function of minority ions, $v_\perp$, $v_1$ are the perpendicular and the longitudinal components of the ion velocity with respect to the confining magnetic field and $s$ is the co-ordinate along the magnetic field line ($s = \pm L$ corresponds to the end points at the field line). Formula (1) has been obtained by integrating the linearized kinetic equation along the unperturbed trajectory. The particle unperturbed trajectory is determined by a sum of two motions: a uniform motion along the field line and the rotation in the transverse direction with the local cyclotron frequency. It has been assumed that the cyclotron radius is small compared with the spatial scale of RF field distribution and when the particle is reflected off the trap ends the rotation phase is randomized. Similar formulas have been used in analytical calculations in Refs [3, 4].

The non-local coupling between $D_\pm$ and $E_\pm$ in (1) is determined by the behaviour of phase integrals $\Phi_\pm(s,s')$. The dependence of $\Phi_\pm(s,s')$ on $|s - s'|$ is always monotonic. The phase $\Phi_\pm(s,s')$ is non-monotonic and has a maximum (minimum) at the cyclotron resonance point $s_\circ$, where $\omega_\parallel(s_\circ) = \omega$. Strong non-local coupling between $\vec{D}$ and $\vec{E}$ takes place if

$$|\Phi_\pm(s,s')| \approx \pi$$

(3)

For $\Phi_+$ and $\Phi_-$ this condition is fulfilled near the observation point $s$ on the distances $|s - s'| \approx \Delta_\pm = \pi v_{IT}/(\omega_\parallel \mp \omega)$, determining the width of the non-locality region.
If at the field line there is a cyclotron resonance point, then the condition (3) for $s'$ may be also fulfilled in the vicinity of the 'mirror' point $s_m$: $|s_m - s'| \leq \Delta_s$. The appearance of the mirror point is connected with the fact that the equation $\Phi_+ = 0$, apart from the root $s' = s$, may have the root $s' = s_m$ because of the non-monotonic behaviour of $\Phi_+$. Coupling between $D_+$ and $E_+$ in the non-locality regions $s' = s$ and $s' = s_m$ shows considerable differences. This is connected with the fact that the non-local effects in the $s' = s$ vicinity are derived from particles travelling from the right as well as from the left, whereas the non-local effects in the $s' = s_m$ vicinity are due only to particles crossing the cyclotron zone. When the observation point approaches the cyclotron point the widths of the non-locality regions increase, and when the condition $|s - s_c| \leq \Delta_+$ is fulfilled the two regions merge. When the observation point departs from the cyclotron point or the ion thermal velocity decreases, the non-locality regions narrow; the contribution from the mirror point becomes insignificant and the coupling between $D$ and $E$ approaches the local one.

The expression (1) for the electric induction vector $D_+$ has been employed for numerical solution of the problem of RF heating of the minority ions in the open trap by making use of global FMSW resonances. The problem of searching for global resonance conditions for the FMSW is a linear eigenvalue problem for the Maxwell equations

$$\text{rot} \ \text{rot} \ \mathbf{E} - \frac{\omega^2}{c^2} \mathbf{E} = \frac{\lambda}{n_0} \frac{\omega^2}{c^2} (\mathbf{D}(n_0) - \mathbf{E})$$

(4)

where $n_0$ is the normalizing constant of the density ($\langle \mathbf{D} - \mathbf{E} \rangle/n_0$ does not depend on $n_0$), $\lambda$ is the eigenvalue having the dimension of the plasma density. In the expression for $\mathbf{D}(\mathbf{E})$ the non-local effects, according to formula (1), are taken into account only for the minority ion contribution in $D_+$. For the component $D_z$ and the non-resonant ion contributions the local expression $\mathbf{D} = \varepsilon \mathbf{E}$ has been used, where $\varepsilon$ is the dielectric permeability tensor of the cold plasma, in which $\varepsilon_{33}$ is assumed to be infinitely large.

The PLFEM code has been used in the numerical calculations. Calculation parameters like those in Ref. [1] have been chosen for the gasdynamic trap KP-2M. The magnetic field at the trap centre has been assumed to be $B = 2.7$ kG, the mirror ratio was $R = 50$ and the plasma column radius in the central cross-section was $r_p = 10$ cm. The plasma density distribution was assumed parabolic, $n(r) = n_0 \times (1 - r^2/r_p^2)$. The plasma density along the field lines of the confining magnetic field was constant up to the point where the magnetic field lines diverge ($z = 245$ cm), beyond which it decayed exponentially. The RF field frequency was $\omega = 2.75 \times 10^7$ s$^{-1}$. First the calculations for deuterium plasma without minority ions were performed. The real eigenvalue $\lambda = 3.926 \times 10^{13}$ cm$^{-3}$ found is the plasma density value at which the global resonance of the TE$_{102}$ mode occurs. This resonance has a zero width and there is no absorption in it because $\lambda$ is real. Note that the electromagnetic field pattern and the $\lambda$ value remain unchanged when the
local value for $\overline{D}$ is used and account is taken of the non-local effects of the bulk gas ions (deuterium). This proves the applicability of the local approximation for $D_-$ and $D_+$ when the non-locality region width $\Delta_\pm$ is much less than the wavelength.

Calculations have been performed of global FMSW resonances of the $TE_{102}$ mode for deuterium plasma with a hydrogen minority of 0.4, 1.5, 2, 3 and 5% for two values of hydrogen ion thermal velocity, $2 \times 10^6$ and $6 \times 10^6$ cm/s. These and other parameters were chosen such that one can make a comparison with the results of calculations performed in Ref. [1] for a cold plasma.

Figure 1 gives the relative width of the resonance $|\text{Im}\lambda|/\text{Re}\lambda$ versus the minority concentration. It should be noted that taking the thermal motion into account leads to the shift of the curve maximum to larger concentrations compared with calculations using the local approximation, and the maximum absorption value is diminished more than twofold. The considerable absorption in the cold plasma with the minority concentration $C = 1.5\%$ was connected, perhaps, with the local Alfvén resonance excitation near the cyclotron zone for the minority where $\epsilon_{11} > 0$ (see Fig. 2).

With the non-local effects taken into account, the electromagnetic field pattern is changed considerably for such concentrations. The distributions of the azimuthal component $E_\phi$ and the component $E_n$ normal to the magnetic surfaces are smoothed;
FIG. 2. Contours of $E_n$ and $E_\phi$ fields for the global FMSW $TE_{102}$ resonance in deuterium plasma with 1.5% hydrogen minority in the cold plasma approximation. (Numbers on lines correspond to relative field amplitudes.) The eigenvalue $\lambda = 4.397 \times 10^{13} - i(5.232 \times 10^{12}) \text{ cm}^{-1}$.

the local increase of the field characteristic of the Alfvén resonance disappears (see Fig. 3). This allows one to suggest that the Alfvén resonance is absent when the non-local effects are taken into account. This explains the decrease in the absorption value and the different dependence of the absorption on the minority concentration. At small (0.4%) and large (5%) minority concentrations the local resonance was not
excited in the cold case and taking account of the thermal motion did not lead to considerable differences in the electromagnetic field pattern or in the resonance width values. This is explained by the fact that the resonance TE_{102} mode studied has a long wavelength and the correction to the cyclotron resonance condition $k_{1}v_{T}$ is small compared with $\omega_{ci}$. In such a situation the cold approximation is sufficiently accurate.
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ELECTRON CYCLOTRON LOWER HYBRID CURRENT DRIVE AND HEATING ON THE WT-3 TOKAMAK

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Abstract

ELECTRON CYCLOTRON LOWER HYBRID CURRENT DRIVE AND HEATING ON THE WT-3 TOKAMAK.

A 70 kA \((q_m = 5)\) flat-top discharge and a 120 kA \((q_m = 3.6)\) flat-top discharge with zero loop voltage are realized by 2H\(_{\text{e}}\) electron cyclotron current drive (ECCD) and lower hybrid current drive (LHCD), respectively. The figure of merit of current drive is 0.03–0.06 \(10^{19} \text{A-W}^{-1} \cdot \text{m}^{-2}\) for ECCD and 0.3–0.6 for LHCD. 2H\(_{\text{e}}\) ECH at the \(q(r) = 1\) surface on the high field side of the midplane is more effective in stabilizing sawteeth than on the low field side. LHCD can stabilize sawteeth, \(m = 1\) and \(m = 2\) MHD activities. In the full LHCD discharge, there are no sawtooth, \(m = 1\) and \(m = 2\) MHD activities.

1. INTRODUCTION

There has been much interest in non-inductive current drive and local heating of tokamaks aiming at the improvement of the stability of tokamak discharges. Experiments have been carried out on the WT-3 tokamak \((R_0 = 65 \text{ cm}, a = 20 \text{ cm and } B_t(0)_{\text{max}} = 1.75 \text{ T})\). A 56 GHz, 200 kW, 100 ms pulse gyrotron was used as an electron cyclotron (EC) power source. A Vlasov antenna with a parabolic or an elliptical reflector was used to convert the TE\(_{02}\) mode to a linearly polarized plane wave and to launch the X-mode radiation from the low field side into the plasma at an arbitrary angle of incidence with respect to the midplane.

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toroidal field [1]. A 2 GHz, 350 kW, 100 ms klystron amplifier was used as an LH power source. Two four-waveguide arrays were used as a lower hybrid wave launchers.

2. SECOND HARMONIC ELECTRON CYCLOTRON CURRENT DRIVE

A target plasma was generated by a slideaway discharge to avoid the cancelling effect of electron cyclotron current drive (ECCD) in the tokamak. A 100 ms EC pulse was launched into the target plasma at an angle of incidence of 60°. EC power selectively heated up the high energy tail electrons in the target plasma, and the second harmonic (2\omega_e) EC driven current was generated [1-3]. A 70 kA flat-top discharge was realized with zero loop voltage during a 144 kW EC pulse. The safety factor q_a at the limiter is about 5. The ECCD plasma had two electron components: bulk electrons with T_e0 = 500 eV and n_e = 2 \times 10^{12} \text{cm}^{-3} and current carrying high energy tail electrons with T_t = 100 keV and n_t = 0.02 \times n_e. Flat-top or ramp-up discharges are obtained when 2\omega_e / \omega = 1-1.15; there are no remarkable differences in the direction of wave injection. The figure of merit for the 2\omega_e ECCD, \eta = \eta_0 I_e R / \text{Pr}_{\omega_e}, is obtained by using 20-73 kA flat-top discharges; it is 0.03-0.06 (10^{19} \text{A} \cdot \text{W}^{-1} \cdot \text{m}^{-2}), which is one order of magnitude smaller than the theoretical value. To study this discrepancy, the velocity distribution of the current carrying high energy electrons was estimated from the hard X-ray energy spectra emitted at various angles to the toroidal field [4,5]. Using the induced distribution function, the energy stored by these electrons was estimated, the EC power deposition profile was calculated with a ray tracing code, and the energy confinement time was estimated to be as about 6 ms, which was much smaller than the slowing down time with bulk electrons. If the loss of fast electrons is taken into account in the calculation of the current drive efficiency, the calculated efficiency approaches the experimental value.

3. LOWER HYBRID CURRENT DRIVE

By applying the lower hybrid (LH) wave after Ohmic power shutdown, a constant plasma current of up to 120 kA was sustained with zero loop voltage (full LHCD discharge) during a 340 kW LH pulse. The safety factor q_a was 3.6 and the line averaged density \bar{n}_e = 9 \times 10^{12} \text{cm}^{-3}. In the full LHCD discharge, the maximum line averaged density obtained was 1.5 \times 10^{13} \text{cm}^{-3}, while the so-called density limit was 2.3 \times 10^{13} \text{cm}^{-3}; the maximum central electron temperature Te(0) was 410 eV. LH driven current is carried by high energy tail electrons, whose distribution function is estimated by electron
FIG. 1. (a) Ray trajectory, location of 2Ω_e layers ((b), (c), (d)) and q = 1 surface (dotted circle) in the poloidal section. Waveform of soft X ray signal from the central chord at $n_e = 0.5 \times 10^{13}$ cm$^{-3}$; (b) for 2Ω_e layer (b), $q_a = 3.4$, $B_{70} = 0.93$ T; (c) for 2Ω_e layer (c), $q_a = 3.7$, $B_{70} = 1.02$ T; (d) for 2Ω_e layer (d), $q_a = 3.4$, $B_{70} = 1.08$ T.

cyclotron emission [6] and hard X-ray spectra. The average energy of the current carrying high energy tail electrons is about 80 keV, and the density ratio of tail to bulk electrons is about 1%. The LHCD efficiency increases with plasma current [7] and electron density, and is $\eta = 0.3-0.6 \times 10^{-9}$ A·W$^{-1}$·m$^{-2}$, which is one half of the theoretical value.

4. SUPPRESSION OF SAWTEETH BY 2Ω_e ECH

The effect of 2Ω_e ECH on sawteeth was studied for different locations of the 2Ω_e ECR layer and for different plasma currents. When 2Ω_e ECH was applied, the repetition period and the amplitude of the sawteeth increased and reached maxima when the 2Ω_e ECR layer was located at the $q(r) = 1$ surface in the midplane [4]. Sometimes, we observed that the sawteeth could be completely suppressed. Such an effect was very sensitive to the location of the 2Ω_e ECR layer. In Fig. 1 three locations ((b), (c), (d)) of the 2Ω_e ECR layer are selected in the midplane by changing the toroidal field: (b) $q(r) = 1$ on the high field side, (c) magnetic axis, (d) $q(r) = 1$ on the low field side. The safety factor $q_s$ was kept constant (~3.4) in Fig. 1. The typical sawteeth
FIG. 2. Sawtooth repetition periods versus EC power normalized by OH power. Curves (b), (c) and (d) correspond to the locations of 2Ω, layers (b), (c) and (d) in Fig. 1, respectively.

appearing in the Ohmic phase have a repetition period of about 0.5 ms and precursors. In the ECH phase the repetition period and amplitude of the sawteeth increase, and a fast crash without precursor occurs. In cases (b) and (d), the amplitude of the sawteeth becomes saturated within several milliseconds and crashes abruptly (Fig. 1 (d)). In case (b), we observe complete suppression of sawteeth during the EC pulse (Fig. 1 (b)). The temperature profile obtained with Thomson scattering shows local heating near the q = 1 surface. The sawtooth repetition is plotted versus the injected EC power in Fig. 2. In case (b), the sawteeth are completely suppressed during the EC pulse when $P_{EC} / P_{OH} > 3.5$, but in case (d), no complete suppression is attained at the available EC power levels. In low current discharges ($q_a > 4.7$), we observe complete sawtooth suppression in case (d).

5. CONTROL OF MHD ACTIVITIES BY LHCD

The effect of LHCD on sawteeth [3,8] and the m = 1 and m = 2 modes was studied. When LHCD was superimposed ($V_L > 0$) on an Ohmic discharge with sawteeth, the sawtooth repetition period gradually increased, and then the sawteeth disappeared. With increasing LH power, the LH driven current increased, and, first, the sawteeth and then the m = 1 mode were stabilized. The diameter of the q = 1 surface estimated by the location of the maximum amplitude of the m = 1 oscillation was almost
constant when LHCD was superimposed on the Ohmic discharge. This means that the $q = 1$ surface still exists after sawtooth suppression and the current profile changes locally. When the opposite LHCD was superimposed on the Ohmic discharge, the soft X-ray signal increased by the same amount as in the normal LHCD but there was no effect on the sawteeth. Suprathermal electrons have no effect on the suppression of sawteeth.

Figure 3 shows typical waveforms of the full current drive discharge ($V_L = 0$). The LH power is turned on after the Ohmic discharge has been turned off. When the LH power is turned on, the $m = 2$ mode is completely stabilized (Fig. 3(c)). The sawtooth repetition period increases gradually and, finally, the sawteeth disappear completely (Fig. 3(d)).

6. CONCLUSIONS

Flat-top discharges of $q_a = 5$ and 3.6 were realized with zero loop voltage by 2$\Omega$e ECCD and LHCD, respectively. The figure of merit of the current drive is $0.03-0.06 \times 10^{19} \text{A} \cdot \text{W}^{-1} \cdot \text{m}^{-2}$ for ECCD and $0.3-0.6$ for LHCD.

Local heating near the $q = 1$ surface by 2$\Omega$e ECH effectively suppresses the sawteeth. 2$\Omega$e ECH on the $q = 1$ surface on the high field side is more effective in suppressing the sawteeth than on the low field side.
Sawteeth and $m = 1$ MHD activities in an Ohmic discharge can be stabilized by superposition of LHCD. After stabilization of the sawteeth, the $q = 1$ surface still exists in the plasma column.

In the full LHCD discharge, there are no sawtooth, $m = 1$ and $m = 2$ MHD activities.

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REFERENCES

FREE ELECTRON LASER EXPERIMENTS IN THE MICROWAVE TOKAMAK EXPERIMENT*


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Abstract

FREE ELECTRON LASER EXPERIMENTS IN THE MICROWAVE TOKAMAK EXPERIMENT.
Microwave pulses have been injected from a free electron laser into the Microwave Tokamak Experiment at up to 0.2 GW at 140 GHz in short pulses (10 ns duration) with O-mode polarization. The power transmitted through the plasma was measured in a first experimental study of high power pulse propagation in the plasma; no non-linear effects were found at this power level. Calculations indicate that non-linear effects may be found at the higher power densities expected in future experiments.

1. INTRODUCTION

The primary goal of the Microwave Tokamak Experiment (MTX) [1] is to study electron cyclotron resonance heating and control of a high density tokamak (Alcator C) by microwaves from a free electron laser (FEL) or a gyrotron. Here, we report the first measurements of high peaked power microwave absorption in the plasma. Microwaves were generated by the FEL, and their transmission and absorption in MTX were experimentally evaluated. In agreement with theory, non-linear effects [2] were not found at these power levels (0.2 GW) but are predicted at higher levels.

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FIG. 1. Measured gain curve for non-tapered FEL operation at 139.9 GHz. The electron beam parameters were 6 MeV, 2 kA, and 50 ns pulse length. The input drive power was 20 W from an extended interaction oscillator (EIO) source, injected into the wiggler waveguide, co-linear with the electron beam.

2. FREE ELECTRON LASER OPERATIONS: MICROWAVE GENERATION AND TRANSMISSION TO MTX

For the microwave transmission and absorption experiments, the ELF-II wiggler (9.8 cm wiggler wavelength and 4 m total length) was used [3]. With non-tapered operation of the wiggler, peak power levels up to 0.2 GW and pulse lengths up to 10 ns were achieved at 140 GHz. Figure 1 shows an experimental gain curve for the wiggler, obtained by varying the wiggler length (number of energized magnets). At the available drive power and electron beam current, the FEL interaction just reached saturation. Operation with the tapered wiggler achieved 0.4 GW, but this power level was not used during tokamak experiments.

The FEL microwave pulse was injected into the windowless, quasi-optical transport system [4] for transmission to the tokamak. This system consists of an evacuated, stainless steel pipe of 50 cm diameter; six aluminium mirrors transmit the beam 33.3 m. Measurements outside the tokamak port demonstrated that the transmitted microwave beam envelope was well approximated by a Gaussian curve, in good agreement with theory.
3. ABSORPTION PHYSICS EXPERIMENTS

FEL produced microwave pulses were injected into the tokamak for a variety of plasma parameters. Local measurement of the transmitted fraction of injected power was made by using a microwave horn located just below the midplane on the inside wall of the tokamak. At sufficiently high (6.5 T) toroidal magnetic fields that there was no electron cyclotron resonance in the plasma, the transmitted power showed the expected drop to negligible levels as the peak density approached cut-off. Ray tracing calculations [5] based on measured density profiles qualitatively reproduced the observed fall-off in transmitted power versus density, indicating that refraction was strong enough to be the dominant mechanism for the measured result.

Experiments were also performed with a central magnetic field of 5.0 T so that the cyclotron resonance passed through the plasma centre. Figure 2 shows the data from these experiments compared to ray tracing calculations of refraction effects and of the additional effects of linear absorption. Analysis indicates that the FEL power levels were below that at which non-linear absorption [2], Brillouin backscatter instabilities [6], and ponderomotive self-focusing [7, 8] are important.

The transmitted power fraction data from the FEL experiments are shown by solid circles. For comparison, the upper hatched region shows the calculated effect of refraction alone. This region spans the range of plasma density profiles \( \alpha_n = 0.5-1.0 \), where \( n(r) = n(0) [1 - (r/a)^2]^{\alpha_n} \) was measured in the experiments. The lower hatched region shows the combined effects of refraction and linear absorption. This band spans the range of density profiles and electron temperatures.
(0.5–1.0 keV) observed. Within the variation, there is no evidence that non-linear saturation of the absorption plays a role in these experiments, in agreement with theory [2].

Data from 5.0 T experiments at low power (1 W) and long pulselength (1–2 μs) are shown as squares in Fig. 2. The bars on the data points represent the root mean square variation in the amplitude of the microwave horn signal during the pulse. These variations also occurred in the high power experiments. Ray tracing calculations with edge density fluctuations predicted scattering of 10–15% of the power through angles of about 20°, insufficient to explain the observed signal variations. The time-scale of the variations is also much shorter than for edge turbulence periods. Multiple spatial modes (generated by wall interaction in the narrow MTX access duct), together with time varying, multipath transmission of power through the plasma (due to edge turbulence), could produce a rapidly varying interference pattern on the tokamak inside wall. This might account for the variations in the observed signals. Experiments are planned to quantify these effects.

4. INTENSE WAVE ISSUES

4.1. Ponderomotive self-focusing

Ponderomotive self-focusing calculations are relevant for future high power microwave propagation experiments in MTX. Numerical solutions of a scaled, paraxial self-focusing equation for ordinary modes perpendicularly incident to an applied magnetic field show qualitative differences from theories in which the laser beam remains Gaussian: self-focusing occurs at a shorter distance, and the beam does not remain Gaussian [8]. In the limit where ion inertia can be neglected, calculations for \( T_e = T_i = 1 \) keV, 3 cm beam radius, \( \omega^2 / \omega_e^2 = 0.5 \), \( f = 140 \) GHz, and \( \nu_0 / \nu_e = 0.24 \) (0.1 GW/cm\(^2\) at 1 keV core temperature) show that self-focusing would occur near the magnetic axis in 3–5 ns. The numerical studies further indicate that the self-focusing time and distance increase approximately as \( (\nu_0 / \nu_e)^{-1} \) for powers above threshold for self-focusing in agreement with a self-similar analysis [7]. Thus, self-focusing is less severe for higher plasma temperatures and larger beam cross-sections, which will make FEL heating applications in CIT less susceptible to self-focusing than in MTX.

Our numerical calculations overestimate the degree of self-focusing because they do not include several effects that significantly delay or prevent the focusing: ion inertia [7], scattering by edge fluctuations that increase the effective beam divergence, and absorption. Analytical and numerical calculations including ion inertia [7] and additional numerical calculations including absorption, plasma profiles, and beam divergence with ion inertia omitted [8] have shown that these effects greatly reduce the amount of self-focusing expected in MTX. These calcu-
lations also indicate how self-focusing can be avoided, e.g. by increasing the beam divergence in the plasma.

4.2. Microinstability

Microinstability of the FEL heated electrons could rapidly relax the distribution function, thus affecting non-linear physics. We have developed a new computer code that solves the electromagnetic linear dispersion for a relativistic plasma with an arbitrary distribution [9], and we have analysed distributions predicted by a particle orbit code for whistler, upper hybrid loss cone (UHLC) and cyclotron maser instabilities in an infinite homogeneous geometry. For a representative distribution we chose the one generated by injecting into an MTX plasma (at 5 T) 2 GW of power at 140 HHz in a beam with cross-section of $6 \times 8$ cm. At a typical density given by $\omega_{ce}/\omega_{pe} = 0.6$, the whistler mode has a maximum growth rate of $\gamma/\omega_{ce} = 7 \times 10^{-3}$, the UHLC mode has $\gamma/\omega_{ce} = 5.6 \times 10^{-3}$, and the cyclotron maser mode is stable. The UHLC mode is unstable for a range of wavenumbers $0.3 < k_c/\omega_{ce} < 1.5$ and $12 < k^c/\omega_{ce} < 22$. The cyclotron maser mode becomes unstable at a lower density, $\omega_{pe}/\omega_{ce} = 0.3$. The maximum growth rates for the whistler and UHLC modes increase monotonically with the density in the ranges of interest. Thus, high frequency linear microinstabilities should be observable in the operating regime of MTX and could enhance the relaxation of the heated distribution.

REFERENCES

Abstract

FLUCTUATION INDUCED TRANSPORT WITH ADDITIONAL HEATING AND A POLOIDAL FIELD DIVERTER.

Local estimates, in the DITE edge plasma, of the cross-field particle flux due to fluctuations, show that it can account for the particle losses to the limiters, over a wide range of density and ECRH power. The calculated cross-field convected heat flux agrees with that from power balance, after subtracting radiation losses, at low densities ($n_e = 2 \times 10^{19} \text{ m}^{-3}$) but underestimates the power loss at higher densities where heat conduction is more important because of the observed large temperature fluctuations ($T_e/T_i = (0.4-1)n^7n$). Magnetic fluctuations are generally too small to explain particle or heat fluxes, even with strong additional heating. Significant correlation between local fluctuations in $B_r$ and $I_{sa}$ is observed near the limiter radius which indicates that the magnetic fluctuations are probably caused by pressure fluctuations due to the basically electrostatic turbulence. COMPASS-C experiments have been successful in producing stable inboard X-point configurations in Ohmic and ECH assisted discharges where a loop voltage of 0.3 V is obtained.

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1. INTRODUCTION

In these experiments we have attempted to identify the processes responsible for particle and heat losses at the edge of tokamaks with various forms of plasma heating, with and without magnetic separatrices. Much of the work has used the circular cross-section DITE tokamak ($R = 1.19$ m, $a = 0.23$ m) with additional heating/current drive by ECR and LH waves. More recently the new COMPASS-C device ($R = 0.56$ m, $a = 0.20$ m) has come into operation. This is equipped with a very flexible poloidal field system with which a variety of plasma shapes can be produced. Initial work has focused on separatrix bounded plasmas with an inboard X-point and with ECRH.

2. TRANSPORT IN THE DITE PLASMA BOUNDARY

The edge plasma response to lower hybrid current drive (LHCD) and electron cyclotron heating (ECH) has been studied in DITE with input powers of up to 250 kW and 300 kW, respectively [1]. The edge is diagnosed by using limiter probes giving good poloidal coverage of $T_e$, $I_{sat}$ and $n_e$ in the scrape-off-layer (SOL), whilst reciprocating probes are configured to monitor fluctuations in poloidal electric field, electron density and temperature, and magnetic field.

LHCD experiments (1.3 GHz) were carried out primarily in hydrogen with $B_e \sim 0.9-2.5$ T, $\bar{n}_e \sim (0.5-1.5) \times 10^{19}$ m$^{-3}$ and $I_p \sim 75-125$ kA. During LHCD, $\bar{n}_e$ is usually observed to fall by 20-30% in the absence of gas puffing. All probes in the SOL consistently show a 50% to 70% reduction in $n_e$ (edge) and a 25% to 30% decrease in both $T_e$ (edge) and $T_i$ (edge). The total particle flux to the pumped limiter, evaluated by interpolating between probes distributed over the limiter surfaces, decreases by a factor of three to four. Coupled with an observed decrease in the wall $H_a$ signal there is strong evidence for a doubling of the particle confinement time. The fluctuation levels $\bar{E}_e$ and $\bar{E}_i$ also decrease during LHCD, and it appears that the improved confinement can be attributed to a reduction in the fluctuation driven flux. (However, the absence of simultaneous measurements of $\bar{E}_e$ and $\bar{n}_e$, in these discharges, has prevented a detailed analysis.) These observations are in complete contrast to those with ECH (60 GHz) where $\bar{n}_e$, $n_e$ (edge), $T_e$ (edge), $\bar{E}_e$, $\bar{n}_e$, $H_a$ emission and the particle flux to the limiter all increase. During ECH the particle confinement time decreases by up to a factor of two, and detailed investigations have shown that this can be accounted for by an increase in the electrostatic fluctuation driven flux [2].

Global power balance has been studied by comparing the power conducted and convected to the edge, as deduced from Langmuir probes in the limiter. Power accounting within $\pm 20\%$ is observed in Ohmic and ECRH phases but, during LHCD, there is a shortfall in the power losses corresponding to essentially all of the LH input power. This is consistent with a very local power loss and/or power loss
via a high energy component in the electron distribution which would not be registered in the Langmuir probe bias range (−100−−3 V).

The cross-field particle flux, $\Gamma_0^f (= \langle nE \rangle / B)$, was compared to the parallel flow to the limiter as deduced from $I_{sat}$ profiles and from the probes embedded in the limiter. Typical DITE cross-field fluxes are $\sim 3.4 \times 10^{20} \text{m}^{-2}\cdot\text{s}^{-1}$ at $\bar{n}_e = 3.2 \times 10^{19} \text{m}^{-3}$. Good agreement was found in the Ohmic and ECH regimes. In addition, the delay of the increase in $\Gamma_0^f$ with respect to ECH initiation, measured by modulating the ECH power, was found to be of the order of $\tau_0$ [3].

The convective power loss, $P_c$, associated with the fluctuation-derived flux, $\Gamma_0^f$, (i.e. $P_c = 3T_e\Gamma_0^f$, assuming $n_eT_e = n_iT_i$) is compared in Fig. 1 with the total power loss, after subtracting the radiation losses. It is clear that at low densities all the power can be accounted for, but at high densities there is a substantial discrepancy. Although $T_i$ of up to 120 eV is observed in the same set of discharges, the accompanying dilution factor, $n_i/n_e \sim 0.4$ [4], means that the ion and electron energy density are the same (within 10%), as assumed in Fig. 1. At high densities, there is a substantial discrepancy. Part of this may be caused by the increased heat conduction due to the observed increase in temperature fluctuations [5] ($\bar{T}_e/T_e = (0.4-1)\bar{n}/n$).

In any case, magnetic fluctuations are too small to explain either particle or heat losses, even with strong additional heating, contributing less than 2% to either of these losses. However, significant correlation is found between the electrostatic ($\bar{E}_{eo}$)
and the magnetic ($\tilde{B}_r$) fluctuations near the limiter radius, in the frequency range of maximum electrostatic activity, suggesting that at least 50% of the broadband magnetic fluctuation power at the limiter radius is associated with the local electrostatic turbulence [5]. This indicates that the magnetic fluctuations are probably caused by pressure fluctuations due to the basically electrostatic turbulence.

3. X-POINT OPERATION

Initial experiments with X-points, in the inboard midplane, have been explored in COMPASS-C, with $I_p \sim 70-100$ kA, $B_T = 1.04-1.12$ T and $n_e \sim (0.5-1.5) \times 10^{19}$ m$^{-3}$. Operation is in hydrogen with a boronized vessel. Additional heating is provided by up to $\approx 0.7$ MW of 60 GHz, second harmonic ECRH (X-mode, low field launch). Part of the programme is to investigate whether or not an L- to H-mode transition can be induced in this essentially untried configuration featuring the X-point in the good curvature region.

The configuration is formed during the plasma current plateau by changing the distribution of currents in the poloidal field windings, whilst the electron density is allowed to fall from $\sim 2.5 \times 10^{19}$ m$^{-3}$ to $\sim 1.0 \times 10^{19}$ m$^{-3}$. The horizontal and vertical equilibrium are maintained by feedback control. Plasma shaping takes place in the period from $t \sim 50$ ms to $t \sim 100-200$ ms and is followed by ECRH within 20-50 ms. The displacement, $\Delta_{\text{inside}}$, of the X-point from the limiter radius has been varied from $\sim 0$ cm to 2 cm whilst the last closed surface is detached by $\Delta_{\text{outside}} \approx 0$ cm to 3 cm. A typical discharge is depicted in Fig. 2, showing the temporal behaviour of the plasma current, loop voltage, shaping coil current and electron density. In this case, centrally resonant ECRH ($\approx 0.5$ MW; $B_T = 1.07$ T) is applied from $t = 200$ ms until the plasma termination time at 300 ms. The electron temperature, as measured by a Si(Li) detector, is $\sim 0.5$ keV throughout the shaping period and increases to $\sim 1.2$ keV during ECRH, whereas the loop voltage drops from $\sim 1.5$ V to $\sim 0.5$ V and the SXR intensity increases by about an order of magnitude. $V_{\text{loop}}$, $H_a$ and $T_e$ are almost indistinguishable from reference discharges of the same $I_p$ and with circular cross-sections similar to the X-point plasmas and resting at the inside limiter.

The largest decreases in $V_{\text{loop}}$, with ECRH, are obtained with $B_T(0) = 1.07$ T (central ECRH resonance), $\Delta_{\text{inside}} \approx 2$ cm, and $\Delta_{\text{outside}} \approx 3$ cm. The lowest $V_{\text{loop}}$ obtained was $\sim 0.3$ V with $I_p = 90$ kA and ECRH power $\sim 0.7$ MW.

During ECRH, the outer magnetic flux contours, as obtained from a current filament reconstruction model and predicted by a full equilibrium code, can be seen, in Fig. 3, to be reasonably matched by the VUV/SXR emissivity contours provided by an SXR TV camera with a tangential view of the plasma [6].

The edge behaviour is monitored with visible spectrometers, Langmuir probes and arrays of magnetic coils. The Balmer alpha ($H_a$) light intensity from the limiter region adjacent to the X-point shows no sudden decrease during the Ohmic or ECRH
phases, nor is there any ELM-type signature. Signals from spectral lines of neutral or low ionization states from various locations also register no particular structure, apart from a typical two- to threefold increase in $n_e$ when applying ECRH.

Langmuir probe measurements, during ECRH, for discharges similar to that shown in Figs 2 and 3, yield relatively high edge values of $n_e$ ($\tilde{n}_e/4$) and $T_e$ ($\sim 40$ eV) and reveal a region of shallow gradients in density and temperature, extending through the separatrix to the flux surface in contact with the limiter at the outboard midplane. At the separatrix there is a change in the electron density gradient ($\lambda_n$ decreasing from $\geq 12$ cm, outside the separatrix to $\leq 3$ cm, inside), but no obvious change in the temperature gradient. The floating potential exhibits a sharp turnover in the region of the separatrix. In this region the temperature profile is essentially flat so that the reversal of the floating potential gradient implies a reversal of the $E \times B$ velocity from the ion (outside) to the electron (inside) diamagnetic drift direction.

Preliminary analyses of results from a reciprocating Langmuir probe indicate that $\Gamma_0'$ is directed outwards in the region outside the separatrix. At 3 mm outside the separatrix, $\Gamma_0'$ is $\sim 3.3 \times 10^{20}$ m$^{-2} \cdot $s$^{-1}$, implying a particle confinement time of
FIG. 3. Flux surfaces, at $t = 225\,\text{ms}$ ($B_T = 1.07\,\text{T}, I_p \sim 70\,\text{kA}$), after transition to the X-point configuration and during ECRH ($=0.5\,\text{MW}$), obtained from measurements of the poloidal field by a set of magnetic pick-up coils. The dashed circle is defined by the limiter radius. The last closed flux surface is indicated by a thicker line and shows the X-point, $\sim 2\,\text{cm}$ displaced inwards from the limiter radius. Also shown are the relative emissivity contours with numbers assigned from the tangential SXR camera.

$\sim 3.5\,\text{ms}$. $\Gamma_{0}^{f}$ rises to $\sim 1.2 \times 10^{21}\,\text{m}^{-2}\cdot\text{s}^{-1}$ at 17 mm outside the separatrix. This increase can be attributed to strong recycling in the open field line region. During ECRH the $\Gamma_{0}^{f}$, at the limiter radius, increases threefold to $3 \times 10^{20}\,\text{m}^{-2}\cdot\text{s}^{-1}$, consistent with the simultaneous increase in the $H_{\alpha}$ intensity.

4. CONCLUSIONS

The cross-field particle flux, derived from density and electrostatic fluctuations, corresponds well with that obtained at the limiter. The associated convected heat flux agrees with that obtained from power balance at low density; at high density the increased temperature fluctuations, $(\dot{\mathcal{T}}_e/T_e \approx (0.4-1)n/n)$, suggest that heat conduction may be more important. Magnetic fluctuations do not appear to explain the particle or heat transport. Correlation between electrostatic and magnetic fluctuations
imply that the latter are caused by pressure fluctuations driven by electrostatic turbulence. Stable inboard X-point configurations have been produced in Ohmic and ECRH assisted plasmas and, when the separatrix is almost equally displaced from the inside and outside limiter positions, the lowest loop voltages (~0.3 V) are obtained.

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THE INFLUENCE OF ECRH, LHCD AND RESONANT HELICAL FIELDS ON TOKAMAK INSTABILITIES


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Abstract

THE INFLUENCE OF ECRH, LHCD AND RESONANT HELICAL FIELDS ON TOKAMAK INSTABILITIES.

The control of MHD instabilities and disruptions remains a problem for future tokamaks, and a number of experiments have been performed on the DITE and COMPASS-C tokamaks to develop control techniques. A variety of tools are available on these devices: LHCD, ECRH, and external resonant magnetic perturbations (RMP) inside and outside the vessels. The last two have been operated in fast feedback loops following rotating instabilities, as well as quasi-DC. Suppression of density limit disruptions with magnetic feedback has been demonstrated, as well as control of internal modes (m = 1, n = 1 and the sawtooth) using RF or magnetic feedback techniques. An extensive study of the effect of static RMPs on MHD activity has been performed which, together with a new theoretical model, sheds light on possible mechanisms for both mode stabilization and stimulation of disruptions with external field perturbations.

1. THE EFFECT OF RESONANT MAGNETIC PERTURBATIONS ON MHD ACTIVITY

Fast magnetic feedback experiments have been performed on DITE [1, 2, 3] using driven saddle coils inside the vessel, with sensitive Mirnov coils to provide an m = 2, n = 1 feedback signal, and controlling the gain and phase of the feedback loop with fast hybrid analogue–digital electronics. This has allowed extension of the density limit by up to ~25% [3], and, perhaps more importantly, disruption-free operation at or somewhat above the usual density limit (Fig. 1). When a disruption

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does occur it is caused by a sudden increase in the m = 2 amplitude, triggered at a sawtooth crash (i.e. not by the heatpulse) [4], which exceeds the range of the feedback system. Such an effect of the sawtooth is not observed on all tokamaks, but noise of some form will always be a limit. In addition, when LHCD is used to remove the sawtooth at low density, the resultant coupled m = 1, n = 1 and m = 2, n = 1 mode can be stabilized (factor ~8 reduction in $B_\phi$ and also a drop in $I_{sxr}$) with the feedback, with an associated improvement in central confinement [5] ($I_{sxr}$ (0) rises by ~10\%, indicating that such feedback systems have a wider application than just disruption avoidance. The m = 1 stabilization is probably due to the global nature
TABLE I. EFFECT OF RMP ON COMPASS-C.
‘stab’: reduction of background m = 2 activity; ‘stim’: stimulated disruptions, and
‘nil’: no observable effect on discharge evolution, m = 1, m = 2 modes and
sawtooth. $\tilde{B}_\theta$ is the rms value at $r = 22$ cm for the unperturbed plasma. $B_r$
($r = 20$ cm) $\sim 10$ G, $\sim 1\% B_\theta(a)$ for the dominant applied perturbation: $I_{saddle}$
$= 500$ A–1 kA (single conductor).

<table>
<thead>
<tr>
<th>Type</th>
<th>RMP configuration</th>
</tr>
</thead>
<tbody>
<tr>
<td>1,1</td>
<td>stab, stim</td>
</tr>
<tr>
<td>2,1</td>
<td>nil</td>
</tr>
<tr>
<td>3,1</td>
<td>weak stab</td>
</tr>
<tr>
<td>3,2</td>
<td>$\tau_{E,p} l$</td>
</tr>
</tbody>
</table>

TABLE II. PARAMETERS FOR THE DIFFERENT DISCHARGE TYPES USED
IN THE RMP STUDIES ON COMPASS-C.
$I_p \approx 100$ kA for all these OH discharges, with $B_\theta$ being varied to change $q(a)$.

<table>
<thead>
<tr>
<th>Type</th>
<th>$q(a)$</th>
<th>$f_{m=2}$ (kHz)</th>
<th>$\tilde{B}_p^{m=2}$ (G)</th>
<th>$\bar{n}_e$ ($10^{19}$ m$^{-3}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>$2 \leq q \leq 3$</td>
<td>13–15</td>
<td>0.5–1.5</td>
<td>$\sim 0.6$</td>
</tr>
<tr>
<td>2</td>
<td>$2 \leq q \leq 3$</td>
<td>10–12</td>
<td>1.5–3</td>
<td>3–5</td>
</tr>
<tr>
<td>3</td>
<td>$3 \leq q \leq 4$</td>
<td>14–16</td>
<td>0.4–0.8</td>
<td>1.3–2.5</td>
</tr>
<tr>
<td>4</td>
<td>$3 \leq q \leq 4$</td>
<td>12–16</td>
<td>1–3</td>
<td>3–5</td>
</tr>
</tbody>
</table>

of this mode: the stability can be affected by changing $\Delta'$ at $q = 2$ [6]. In OH
discharges, the $m = 1$ sawtooth precursor is at a lower frequency than the continuous $m = 2$ mode ($f_{m=1} \sim 4$ kHz, $f_{m=2} \sim 10$ kHz), and is not affected by the feedback.
Experiments to measure dynamic mode locking and the phase instability are
described in Ref. [3].

On COMPASS-C ($T = 0.557$ m, $a_{lim} = 0.196$ m, $I_p \leq 200$ kA,
$B_\phi \leq 1.75$ T) the emphasis has been on quasi-static perturbations using a complex
set of external saddle coils [7]. In each quadrant of the tokamak there are 10 toroidal
bars, 70° long, with poloidal links. Feeder errors have been minimized, as have errors in the main poloidal field system — error fields have been reduced to below 0.5 G (n = 1) for each PF coil by magnetic alignment [8]. Most configurations have sideband levels below 10% (calculated over a uniform $\theta, \phi$ grid). Tables I and II show the range of experiments performed and summarize the effects of the RMP. It is seen that there are five main cases: (i) no effect; (ii) reduction of the existing rotating mode; (iii) stimulated disruptions; (iv) reduction in the confinement and (v) increase in the density limit. Attempts to achieve $q_\rho(a) < 2$ using RMPs have not been successful to date. RMPs also have little effect on the behaviour of runaway electrons.

Apart from a sequence at $q(a) \sim 2$, these experiments were mostly performed before boronization of COMPASS-C. Most RMP configurations had no net helicity (i.e. $\tilde{B}(m,n) = \tilde{B}(m,-n)$); the main change when an RMP with resonant helicity is used is a reduction in the saddle current required to see any effect.

The results will be reviewed in the light of the theory of Refs [9, 10]. Little driven reconnection is predicted when $f_{\text{applied}}$ is more than a few per cent from $f_{\text{natural}}$, where $f_{\text{natural}}$ is the drift tearing frequency at which tearing modes propagate. Non-linear modelling shows that there is also a stabilizing effect due to the applied RMP slowing the mode from $f_{\text{natural}}$, until $I_{saddle}$ exceeds a threshold $I_{s,crit}$ when $f_{\text{natural}} - f_{\text{applied}}$ and full driven reconnection then occurs.

The 'nil' entries are thus interpreted in terms of minimal driven tearing and little stabilization due to the applied RMP being too small for the strength of the instability in question (enhanced power supplies are in construction).

Using a 2,1 RMP, $m = 2$ stabilization is seen at low density for $q(a) \leq 3$ (Fig. 2), usually accompanied by a reduction of $1 \sim 3$ kHz in the mode frequency before the mode signal sinks into the noise, as expected from the theory. At high density the natural $m = 2$ level is higher, and the stabilizing term is smaller. It is possible, however, to obtain a 15% enhancement in the density limit at intermediate $q(a)$ [7] — there is no strong stabilization of the $m = 2$ mode in this case, and differences in the $H_\alpha$ signal indicate that the edge conditions (which are known to affect the density limit) are changed with the RMP, and this may be the dominant effect here. As $q(a)$ is raised the effect of the 2,1 RMP is weakened, probably simply because of the greater distance between the saddle coils and the rational-$q$ surface.

At high values of $I_{saddle}$, disruptions are stimulated for 2,1 RMPs, as usual in such experiments, with $I_{saddle} \geq I_{s,crit}$. A drop in central confinement is also seen just before disruption [10], consistent with substantial tearing. These disruptions do not have conventional rotating precursors, but the final explosive growth ($\tau \sim 50 \mu s$) and energy quench ($\sim 100 \mu s$) are similar to normal low-$q(a)$ disruptions. Adding a 3,1 RMP (in quadrature, toroidally) with $I_{2,1}$ below the disruption threshold also leads to disruptions when the calculated islands touch ($I_{2,1} \approx 900 \ A$, $I_{3,1} \approx 700 \ A$) [7]. Stimulated and low $q(a)$ disruptions are generally faster than density limit disruptions on COMPASS-C ($dL_p/dt = 2 \times 10^8 \ A/s$ for all types of disruptions).
FIG. 2. Reduction of $m = 2$ instability by stationary 2,1 RMP on COMPASS-C at low $\bar{n}_e$. A variety of $I_{\text{saddle}}$ waveforms shows penetration to be fast. $I_p = 100$ kA, $q_0(a) = 3$.

Turning now to 1,1 and 3,2 RMPs, there are two main results. Firstly, there is no apparent effect on the sawtooth period, amplitude and $m = 1$ frequency, or on the $m = 2$ and secondly some confinement degradation (reduced $d\bar{n}_e/dt$ for constant fuelling, $10 \sim 20\%$ drop in $I_{m1}$ ($r = 10$ cm)) is seen with both RMPs. These results, with changes seen in $H_a$ signals, suggest there is little or no driven tearing (for the 3,2 case $W_{\text{vac}}/a \sim 8\%$ using the vacuum perturbed fields, and $I_{\text{saddle}} < I_{\text{crit}}$) and that the confinement change is due to changes at the plasma boundary, consistent with changes in this behaviour after boronization.

2. INSTABILITY CONTROL BY RF INJECTION

LHCD has been used on DITE [5] to suppress the sawtooth instability for $P_{LHCD} \approx 80-100$ kW ($P_{OH} \sim 100-200$ kW before injection). The increase in period and inversion radius before suppression, and the reliable appearance of continuous $m = 1$ activity after sawtooth removal, following a period of $\sim 5$ ms with no clear activity, both indicate that there is a $q = 1$ surface present throughout, close to the sawtooth inversion radius. It is not possible to determine $r_{\text{inv}}$ immediately after LHCD ends, because of strong Parail–Pogutse instabilities at that time. Increasing $P_{LHCD}$ to above $160 \sim 200$ kW removes the $m = 1$ activity (as on other devices), as does applying $m = 2$ magnetic feedback. No effect on the sawtooth is seen with LH heating.
### TABLE III. MHD EFFECTS OF $2\omega_{ce}$ ECRH ON COMPASS-C (PREBORONIZATION).
The entries are described in the text. $I_p = 100$ kA, $n_e = 1.3 \times 10^{19}$ m$^{-3}$, and $P_{ECRH} = 300-450$ kW.

<table>
<thead>
<tr>
<th>$B_e$ (T)</th>
<th>$r_{res}$ (cm)</th>
<th>Sawtooth removal</th>
<th>$m = 2$ mode</th>
<th>$m = 1$ mode</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.95</td>
<td>$-8.3$</td>
<td>Sometimes</td>
<td>Yes</td>
<td>No</td>
</tr>
<tr>
<td>1.0</td>
<td>$-5.6$</td>
<td>Sometimes</td>
<td>Yes</td>
<td>Yes (weak)</td>
</tr>
<tr>
<td>1.05</td>
<td>$-3.0$</td>
<td>At onset only</td>
<td>No</td>
<td>Yes (weak)</td>
</tr>
<tr>
<td>1.1</td>
<td>$-0.4$</td>
<td>At onset or none</td>
<td>No</td>
<td>Yes</td>
</tr>
<tr>
<td>1.2</td>
<td>$4.8$</td>
<td>Yes</td>
<td>Yes</td>
<td>No</td>
</tr>
</tbody>
</table>

On COMPASS localized 60 GHz ECRH [7] at the second harmonic launched from low field side X-mode antennas ($85^\circ$ from $B$) with several tangency radii has been used to remove the sawtooth for up to 60 ms (the ECRH pulse length used). The results are summarized in Table III. Suppression occurs in $\leq 1$ sawtooth periods ($T_{ST} \sim 0.4$ ms here) and may last for the whole ECRH pulse. After the ECRH pulse the sawteeth often appear with smaller $r_{inv}$, but $r_{inv} - r_{inv\text{, before ECRH}} \sim -4$ cm in a few periods. Suppression is obtained for a range of resonance positions ($-8$ cm $\leq r_{res} \leq +5$ cm) but is most reliable for $r_{res} = +5$ cm. With ECRH a high frequency internal $m = 1$ mode at $\sim 25-30$ kHz can appear inside $r \sim 6$ cm in addition to the sawtooth. Enhanced $m = 2$, $n = 1$ activity ($f = 15-17$ kHz) is sometimes seen later in the sawtooth-free period, with SXR diodes indicating a small component of $m = 1$ as well. This enhanced activity is not observed simultaneously with sawteeth with large inversion radii, so a picture where $q(0)$ rises significantly above unity and hence destabilizes the $m = 2$ mode may be consistent with these data, but may not explain the rapid sawtooth removal at the onset of ECRH.

A brief series of experiments has been performed on DITE using ECRH modulated at $f_{m\_2}$ to attempt to control steady $m = 2$ oscillations by heating the magnetic islands [11-13]. The same fast analogue–digital electronics as in the magnetic feedback experiment was used to provide a signal to modulate $P_{ECRH}$ (the deposition is predicted to be well localized poloidally and toroidally), allowing the deposition zone to be slowly swept with respect to the island X- and O-points. The modulation was applied after a period of DC heating, so as to allow separation of the usual rapid ($\sim 1-2$ ms) stabilizing effect of the mean ECRH power for $|r_{res} - r_{q=2}| \leq 2.5$ cm for $I_p = 85$ kA, $q_{opt}(a) = 2.5$. The results show that there is a definite sensitivity of the $m = 2$ amplitude to the phase in the feedback loop, when DC stabilization is seen, and stabilization/destabilization effects comparable to the DC effects are seen.
3. CONCLUSIONS

A variety of MHD stabilization methods has been attempted on the DITE and COMPASS tokamaks and sustainment of discharges above the usual density limit with magnetic feedback has been demonstrated. A range of effects from mode stabilization to stimulated disruptions has been observed with resonant magnetic perturbations which appear to be consistent with a single theory, described in a companion paper.

Sawtooth suppression has been demonstrated with both LHCD and ECRH. Suppression is rapid and is apparently not usually simply due to removal of the q = 1 surface. Reduction of m = 2 activity with ECRH is seen on DITE, with the stabilization again being rapid. The speed of response indicates that local changes in J(r) at mode rational surfaces are responsible.

Initial experiments to control m = 2 activity with modulated ECRH in a fast feedback loop indicate sensitivity to the phase in the loop, broadly supporting theoretical predictions.

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The authors are grateful for the capable assistance of the COMPASS DITE and ECRH operations teams. The m = 2 magnetic feedback work on DITE was performed under Article 14 contract No. JE8/9006 for JET.

REFERENCES

INTENSE LOWER-HYBRID WAVE PENETRATION AND CURRENT DRIVE IN REACTOR-GRADE PLASMAS

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Abstract

INTENSE LOWER-HYBRID WAVE PENETRATION AND CURRENT DRIVE IN REACTOR-GRADE PLASMAS.

Applying lower-hybrid power in short, intense pulses can overcome Landau damping, allowing penetration into the core of reactor-grade plasmas. A theoretical description of the absorption and parametric stability of the pulses is presented and results of ray-tracing calculations are shown which include the absorption calculation. Consideration of the absorption and potential source availability leads to the consideration of 5–10 GW peak power, 30–100 μs pulses for ITER, and ~2 MW, 20 μs pulses for a proof-of-principle experiment in the Microwave Tokamak Experiment (MTX).

Recent numerical studies have confirmed that current drive by conventional CW lower-hybrid power in reactor-grade tokamak plasmas is restricted to the outer periphery [1]. However, application of lower-hybrid power in short, intense pulses enhances penetration by saturating the power absorbed through Landau damping [2]. We envision multi-GW peak, 100–200 MW average power at 8–10 GHz driving 10–20 MA of d.c. current. Ref. 2 gave a partial theoretical description of the absorption and parametric stability of very short (10’s of ns) pulses, such as might be obtained from an induction-linac-driven relativistic klystron (“regime A”). We have since found that there are significant advantages to operating with longer pulses (“regime B”), with durations τ_p comparable to the collisional relaxation time τ_r of the nonlinear plateau (of width Δu in parallel velocity u). We present here a theory of the nonlinear absorption and ray-tracing results valid for this regime, and an improved description of parametric instabilities valid in either regime.

The absorption physics depends on τ_p. For τ_p << τ_r, the absorbed power P_a is dominated by the energy E_0 required to initially create the plateau in the absence of collisions. For the multi-ns regime described in Ref. 1, τ_p is
less than a transit time, so that $E_0 \propto \tau_p$ and $\langle P \rangle$, the absorbed power averaged over the pulse duration and extent, is independent of $\tau_p$. For $\tau_p$ long enough that all electrons on a flux surface pass through the beam, $\langle P \rangle \propto \tau_p^{-1}$. For still-longer times, collisional dissipation during the pulse is important. For small $\Delta u/u$, the dissipation can be calculated by a boundary-layer analysis of a one-dimensional Fokker-Planck equation (obtained by assuming that the perpendicular variation is Maxwellian) which can be written in the form [3]:

$$\frac{\partial F}{\partial t} = D + \nu(2+Z)v_t^2 \frac{\partial}{\partial u} \left[ \frac{T_e}{m_e u^3} \frac{\partial}{\partial u} \right] F$$

where $D$ is the r.f. operator (not necessarily diffusive), $F$ is the electron distribution function integrated over perpendicular velocity, $v_t = (2T_e/m_e)^{1/2}$, and $\nu = 4\pi n_e e^4 \lambda/m_e^2 v_t^2$ with $\lambda$ the Coulomb logarithm, and $Z$ is the effective charge state. The boundary-layer analysis proceeds as follows: divide $u$ space into subresonant, resonant, and super-resonant regions, $I$, $II$, and $III$ respectively; consider $D$ to be dominant in region $II$, so that $\partial F/\partial t |_{II} \approx$ constant. The absorbed power can be expressed in terms of $F_{II}$, $\partial F_{II}/\partial t$, and $\partial F_{I,III}/\partial u$ evaluated at the boundary with region $II$. These in turn can be obtained by changing variables from $u$ to parallel energy $\epsilon_\parallel$, introducing $g = F \exp(\epsilon_\parallel/T_e)$, treating as constant the coefficients of the $\epsilon_\parallel$ derivatives of $g$, and solving via Laplace transforms. Defining the absorption coefficient $\alpha = (S^{-1} dS/ds)$, where $S$ is the r.f. energy flux and $s$ denotes length along a ray, we find, for $\tau_p \ll \tau_r$,

$$\alpha = \frac{n A_\psi}{A_b S} \left[ \frac{T_e v_t H e^{-\epsilon_1}}{2 \pi^{1/2} u_r \tau_p} + \frac{\nu \Delta \epsilon_\parallel \Delta e^{-\epsilon}}{\pi^{3/4} (2\nu \tau_p)^{1/2}} + \nu T_e v_t \ln(u_2/u_1) \frac{\Delta(\text{erf} e^{1/2})}{\Delta u} \right]$$

(1)

where $\Delta x = x_2 - x_1$, with 1 and 2 denoting values at either side of region $II$, $\nu = \nu(2+Z)$, $\epsilon = \epsilon_\parallel/T_e$, $A_\psi$ and $A_b$ are the areas of the flux surface and beam, respectively, $H = e^{\epsilon_1}(u_r/3v_t) [2K \Delta(\bar{u}^3) - 3K \Delta \bar{u} + 3\Delta(\bar{u} e^{-\epsilon})]$, $\bar{u} = u/v_t$, and $K = \pi^{1/2} (2\Delta \bar{u})^{-1} \Delta(\text{erf} \ e^{1/2})$. The terms in Eq. (1) are associated with, respectively, the energy to establish the nonlinear collisionless plateau ("NL"), the transient dissipation of the boundary layers at either end of the plateau ("BL"), and steady collisional dissipation of the plateau ("SS"). For narrow plateaus, these terms scale as the third, second, and first powers of the plateau width. Note that the BL and SS terms become equal for $\tau_p \approx (2/\pi)^{1/2} (2 + Z)^{-1} e^{\epsilon_2/2} (\Delta u/u)^2$, which is a fraction of the plateau relaxation time, while the NL term is small. This choice of $\tau_p$ is a good compromise between penetration and source technology considerations. A generalization of Eq. (1) to relativistic parallel dynamics is straightforward and will be described in a forthcoming paper; the most important relativistic corrections are captured by using the relativistic energy in Eq. (1), which we do in the ray-tracing calculations described below.

A sample comparison between the analytic expression for the power absorbed, $P_a = \alpha A_b S$, and the numerical result from the Fokker-Planck code CQL[4] is shown in Fig. 1. We compare a nonrelativistic case with the following parameters: $n_e = 1 \times 10^{20}$ m$^{-3}$, $Z = 1$, $T_e = 25$ keV, $\epsilon_1 = 3.21$, and $\epsilon_2 = 3.92$ corresponding to $\Delta u/u = 0.1$ and parallel index of refraction $N_{\parallel} = 1.7$. A large, constant diffusion coefficient in the parallel-velocity direction is applied between $\epsilon_1$ and $\epsilon_2$. Flattening of the initially Maxwellian distribution takes place in about $10^{-7}$ s. After this time, the power is determined by the three terms in Eq. (1). These terms each dominate in a different range of $\tau_p$ as indi-
FIG. 1. Comparison of Fokker-Planck code and analytic results for pulse averaged absorbed power versus pulse duration.

cated in the figure. The agreement between the theory and code is good except for long times where the theory overestimates the absorption by about 30%, and where the 1-D and boundary-layer approximations used in the theory are invalid. The overestimate may actually be somewhat larger as the code result tends to decrease somewhat as the number of grid points increases; Fig. 1 was obtained using 160 points in both the speed and pitch-angle coordinates.

The plateau width is a function of both the power density and the spectral width \( \delta N \), and is well-approximated by \( [(\delta N N)^2 + (\Delta u/u)^2]^{1/2} \), where \( (\Delta u) \) is the trapping width for a single \( N \) wave. We obtain \( (\Delta u) \) by calculating the action enclosed by the separatrix from the electrostatic Hamilton, and relating the parallel electric field to \( S \) from the polarization and dispersion relations. The result is

\[
(\Delta u/u) = 1.4 [S_{\perp} (GW/m^2)/N]^{1/4} A^{-1/4} (G/E)^{1/2} (10 GHz/f)^{1/2}
\]

where \( f \) is the wave frequency, \( A \) is the ion mass in AMU, \( G = y \) \((1 - y^2)^{-1/4} (1 + z^2)^{-1/2}, x = \omega_{ce}/\omega_{ce}, y = \omega/\omega_{ih}, \) and \( \omega, \omega_{pe}, \omega_{ce} \) and \( \omega_{ih} \) are, respectively, the wave, electron-plasma, electron-cyclotron, and lower-hybrid frequencies. We have verified Eq. (2) with orbit-code calculations. This width is significant, of order 10%, even for the application of continuous lower-hybrid power at the 50–100 MW level for driving current in the edge of an ITER-like plasma.

The opacity for both regimes A and B and the nonlinear plateau width have been incorporated into the ACCOME ray-tracing and current-drive code which was developed within the framework of LLNL-MIT-JAERI collaboration [5]. The code evaluates power deposition as a function of flux surface for model temperature and density profiles and numerically computed equilibria. Additionally, the code provides essential information on the trajectory of rays and
the degree of spreading of the $N_{\parallel}$ spectrum as the rays propagate. We find that, with ITER-like parameters [$n_e(0) = 8 \times 10^{19} \text{ m}^{-3}, T_e(0) = T_i(0) = 34 \text{ keV},$ and $B_0 = 4.85 \text{ T}$], rays from plausible launcher locations can pass near the center with acceptable $N_{\parallel}$, and that, while the mean $N_{\parallel}$ shifts somewhat, the spectral width remains narrow and nearly constant until the rays pass their distance of closest approach to the plasma center and subsequently broadens. Figure 2 shows a comparison of absorption profiles for a regime B scenario (9 GW, 85 $\mu$s, 8 GHz pulses launched with $\delta N_{\parallel}/N_{\parallel} = .05$ about $N_{\parallel} = 1.8$ from 18 m$^2$ of launcher area; the ray shown is launched from 2 m above the midplane) and a conventional scenario (100 MW continuous power from 2 m$^2$ launcher area; other parameters the same). The pulsed power penetrates to the core (in fact, to somewhat beyond the center for the chosen parameters) while the continuous power does not. Comparable penetration can also be achieved for regime A [2], but requires an order-of-magnitude higher power density (possibly conflicting with parametric-instability requirements) and long, narrow launchers which follow field lines around the torus (a difficult engineering feat even for a single, known magnetic equilibrium). The absorption in regime B is dependent only on the beam area, not on the shape or number of launcher arrays. We view regime B as preferable, even though it requires a source development program for reactor applications. Note that, even though the pulses are much longer and more energetic than those for regime A [2], the energy per pulse is a small fraction of the plasma energy content, and maintenance of a d.c. current $I \sim 10-20$ MA would require a repetition period short compared to energy-confinement or current-penetration times. Hence fluctuations in the temperature and current will be small.

Because the absorption is single-pass and in the high-temperature core of the plasma, the current-drive efficiency should be as given by the narrow-spectrum results of Karney and Fisch (Fig. 3 of Ref. 6). In particular, the momentum-conserving corrections to the test-particle efficiency should be appreciable, yielding efficiencies $\eta = n_e(10^{20}/\text{m}^3)I(A)R(\text{m})/P_e(\text{W}) \gtrsim 0.5$, where $R$ is the major radius.

A detailed theoretical investigation of parametric decay instabilities (PDI) has been carried out. The growth rates have been computed numerically using the hot plasma PDI dispersion relationship, valid in the large amplitude regime [7]. The dominant instability is decay into lower-hybrid sideband modes and ion-cyclotron quasi-modes. Convective stabilization can be achieved by restricting the launcher dimensions in the toroidal and/or poloidal directions. We have evaluated the thermal fluctuation level, and obtained the conditions for pump wave depletion. We find that the usual convective threshold power may be increased by a factor of four before pump depletion sets in. Even including this effect, regime A requires a narrow launcher height ($\lesssim 20 \text{ cm}$) or short pulse length ($\lesssim 20 \text{ ns}$). In regime B the launcher height is increased to 40-60 cm and the width $L_z$ can be $\sim 1 \text{ m}$. PDI is most important near the plasma periphery where $T \lesssim 1 \text{ keV}$ and $n \lesssim 2 \times 10^{19} \text{ m}^{-3}$. Here, the growth rates are largest for modes driven by the parallel electric field of the pump wave, and decay waves are stabilized by Landau damping beyond the separatrix. Considering the finite radial distance between the separatrix and the launcher ($\sim 10 \text{ cm}$ in ITER), we find that pump-wave depletion is avoided for power densities $\lesssim 0.5 \text{ GW/m}^2$. Furthermore, in at least regime B we estimate that the PDI modes will locally generate a non-Maxwellian distribution, which should heat
FIG. 2. Comparison of absorption for 9 GW pulsed and 100 MW continuous power ITER cases: (a) ray penetration for pulsed case; each tick mark corresponds to absorption of 10% of the incident power; (b) same for continuous case; (c) power absorbed versus normalized toroidal magnetic flux $\Psi$ for pulsed (solid curve) and continuous (dashed curve) cases.
edge electrons and tend to stabilize the modes by Landau damping. Hence, even without total convective stabilization, we expect a saturated PDI spectrum without causing pump depletion. Experimental results which support this thesis can be found in Ref. 8. We have examined ponderomotive density depletion effects. These should be small ($\Delta n/n \lesssim 0.1$ in regime B). We also consider the scattering of the incoming lower hybrid wave by low frequency background fluctuations. Present estimates indicate acceptably small scattering for the expected fluctuation levels.

Finally, we note that a proof-of-principle test could be done on the Microwave Tokamak Experiment MTX. Using the modified ACCOME code, we find that 2 MW, 20 $\mu$s, $N_\parallel = 5$, narrow-spectrum pulses at 4.6 GHz from a 640 cm$^2$ launcher (which would be mounted between ports) can penetrate to the core, while shorter or lower-amplitude pulses, or pulses from a smaller launcher, do not. The experiment would attempt to confirm these predictions. Typical parameters for such an experiment would be $n_e(0) \approx 2\mathbf{-3}\times10^{19}$ m$^{-3}$, $T_e(0) \lesssim 5$ keV, and $B_0 = 5$ T.

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ANALYSIS OF FAST WAVE CURRENT DRIVE IN REACTOR SCALE TOKAMAKS THROUGH HAMILTONIAN THEORY

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Abstract

ANALYSIS OF FAST WAVE CURRENT DRIVE IN REACTOR SCALE TOKAMAKS THROUGH HAMILTONIAN THEORY.

The Hamiltonian formalism is used to analyse the direct resonant interaction between the fast magnetosonic wave and passing electrons in a tokamak. The regions in phase space where quasi-linear diffusion applies are deduced from a stochasticity criterion. A singular behaviour of rational-q surfaces is found, and the implications on the FWCD scenario in a reactor are discussed.

1. INTRODUCTION

The fast magnetosonic wave, in the frequency range of up to a few ωci, is a major candidate for driving non-inductive currents in reactor-like plasmas. Electrons interact with the wave through a parallel electric force eδEz, leading to electron Landau damping (ELD), and through the transit time magnetic pumping (TTMP) force, μνV(δBz). The interaction between the wave and the drift motion of the electron becomes significant at very high parallel energy and will be neglected here. The fast wave does not suffer, as does the lower hybrid wave, from accessibility conditions and will propagate in high density flat profile reactor plasmas. Fast wave direct electron heating has recently been observed on JFT-2M [1] and JET [2], and several fast wave current drive (FWCD) experiments are planned (DIII-D, JET, JT-60, TORE SUPRA), aiming at the implementation of an ICRF heating–current drive system on NET/ITER.

The efficiency of the proposed scenario depends on the coupling between wave and electrons. The existing theories a priori assume a quasi-linear damping in the velocity space domain where resonant velocities exist. The quasi-linear theory demands, however, a validation in terms of stochasticity and collisional effects. A Hamiltonian approach to the interaction is, therefore, necessary in collisionless regimes in order to determine the domains, self-consistently depending on field amplitude and spectrum, where trajectories become stochastic and quasi-linear diffusion occurs.
2. THE RESONANT WAVE-PARTICLE INTERACTION

The interaction is analysed for passing electrons in the tokamak geometry. The electromagnetic fast wave perturbation is a discrete sum of Fourier modes \((N_2, N_3)\), where \(N_2\) and \(N_3\) are the poloidal and toroidal mode numbers, respectively. Resonances are identified in the \((v_1, v_\perp, q)\) space:

\[
v_{\text{res}} = -\frac{\omega R}{\left(\frac{N_2}{q} + N_3\right)}
\]

where \(v_1\) and \(v_\perp\) are, respectively, the unperturbed parallel and perpendicular velocities in the equatorial plane, \(q\) is the safety factor, \(\omega\) the wave frequency and \(R\) the major radius of the tokamak. The width of the associated trapped domains is then determined.

The overlapping of trapped domains surrounding each resonant velocity leads to a stochastic trajectory and thus provides the necessary phase decorrelation for an irreversible transfer of energy between wave and particle to take place. Extrinsic decorrelation processes such as collisions will be analysed in Section 4. The existence of intrinsic chaos can be quantified through the Chirikov criterion, \(S \gtrsim 1\), where \(S\) depends on the particle velocity and on the amplitude and spectrum of the electromagnetic field. The parallel electric field is consistently derived from the magnetic perturbation through a linear analysis, and a numerical computation of the fields is under way through the ALCYON [3] full wave code.

A critical perpendicular energy \(W_{\perp\text{crit}}\) appears [4] for which the electric and magnetic forces cancel each other:

\[
W_{\perp\text{crit}} = kT_e + m_e c^2 \frac{\omega^2}{\omega^2_{\parallel\text{crit}}}
\]

The TTMP effect is dominant for electron perpendicular energy \(W_{\perp} > W_{\perp\text{crit}}\). FWCD scenarios with frequencies \(\omega \approx \omega_{\parallel}\) have \(W_{\perp\text{crit}} \approx kT_e\), leading to a significant interaction only for subthermal (ELD) and suprathermal (TTMP) perpendicular velocities. For given field amplitude and spectrum, the Chirikov parameter \(S\) depends on the velocities as \(|v_3|^{-2} |W_{\perp} - W_{\perp\text{crit}}|^{1/2}\).

3. ANALYSIS OF HAMILTONIAN CHAOS IN PHASE SPACE \((v_1, v_\perp, q)\)

In the following, we shall apply the theory to a typical reactor scale tokamak, such as NET/ITER and choose the wave frequency so as to remove the tritium
fundamental cyclotron resonance from the plasma ($R = 6 \text{ m}$, $f = 19 \text{ MHz}$, $n_e \approx 5 \times 10^{19} \text{ m}^{-3}$, $T_e \approx 25 \text{ keV}$, $B_0 \approx 5 \text{ T}$).

Assuming that the safety factor $q$ is constant along the unperturbed trajectory, Fig. 1 shows the $S = 1$ frontier in velocity space for a given field spectrum. Within the regions labelled TTMP and ELD, $S$ is larger than 1, the perturbed trajectories are stochastic, and, therefore, quasi-linear diffusion applies. The spiky nature of this frontier is very sensitive to the $q$ value, which directly acts upon the distance between consecutive resonances (Eq. 1). In particular, rational $q$ values lead to a degeneracy, reducing the effective number of resonances and, therefore, the extension of chaos. The effect is strongest near the integer-$q$ surfaces.

To display this phenomenon over a significant plasma volume, we have, in Fig. 2(a), represented a constant-$v_\perp$ cross-section of the phase space where the dark regions contain adiabatic trajectories and quasi-linear diffusion is restricted to the white domains. As expected, the vicinity of rational $q$ values shows up in exhibiting large adiabatic 'diamond shaped' structures, which, in the absence of any extrinsic phenomenon, would prevent an electron from sitting on such a flux surface to reach high parallel velocities. However, outside these singular zones and for $v_\perp = 0.5 \text{ c}$, stochasticity is effective up to $v_\parallel = 0.5 \text{ c}$. Increasing the perpendicular velocity (Fig. 2(b), $v_\perp = 0.7 \text{ c}$) pushes the adiabatic barrier to $v_\parallel = 0.6 \text{ c}$, i.e. $v \approx c$.

4. IMPLICATIONS ON THE KINETIC THEORY OF FWCD

With regard now to a kinetic analysis, the physical reality of the preceding rational-$q$ structures, defining singular flux surfaces on which quasi-linear diffusion
FIG. 2. Adiabatic domains (dark) in \((v_l, q)\) space and corresponding D coefficients for \(q = 1.618\), for \(v_l/c = 0.5\) (a and c) and \(v_l/c = 0.7\) (b and d).

is prevented, must be questioned. In fact, extrinsic phenomena (e.g. collisions and radial diffusion) superimpose Markovian processes on the regular trajectories generating the ‘dark’ domains.

4.1. Effect of radial diffusion on rational-q structures

The anomalous radial diffusion, commonly observed during RF current drive experiments, induces random steps in q space and will decorrelate the relative wave-electron phase when

\[ \tau_q \ll \tau_{\text{wave}} \]

where \(\tau_q\) is the characteristic decorrelation time of an electron during its radial diffusion and \(\tau_{\text{wave}}\) the characteristic time for the same electron to complete its back and forth regular motion in the wave. We have
\[ \tau_q = \left( \frac{N_2}{q^2} \sqrt{D_r} \frac{\partial q}{\partial t} \frac{v_\perp}{R_0} \right)^{-2/3} \]

where \( D_r \) is the radial diffusion coefficient. On the other hand, the typical bounce time in the wave reads:

\[ \tau_{\text{wave}} = \frac{\pi v_{\text{fres}}}{2 \omega \sqrt{\frac{b_N}{B_0} |v_\perp^2 - v_{\text{crit}}^2|}} \]

where \( b_N \) is the N mode magnetic amplitude.

Quantitatively, for the parameters in Fig. 2(a), \( \tau_{\text{wave}} \approx 1 \mu s \) and, on the assumption of \( D_r \approx 2 \text{ m}^2/\text{s} \), \( \tau_q \) is smaller than \( \tau_{\text{wave}} \) for \( v_\parallel > 0.45 \text{ c} \) for \( q = 1 \) and \( v_\parallel > 0.8 \text{ c} \) for \( q = 2 \). The anomalous radial diffusion is not fast enough to wipe out the ‘diamond shaped’ structures in the whole velocity range but generates chaos at high parallel velocities, in the centre of the discharge.

### 4.2. Collisional regimes

Domains in velocity space where the collective motion can be described by a quasi-linear diffusion depend also on collisional considerations. Different regimes in velocity space exist, according to the Chirikov parameter \( S \), and to the ratio, \( D = D_{\text{QL}} \tau_{\text{slowing down}}/v_{\text{thermal}}^2 \) between the quasi-linear diffusion and the collisional diffusion coefficients; in particular:

- \( D \gg 1 \), over the region where \( S > 1 \): an effective plateau forms in the parallel velocity distribution, the highest parallel velocity being determined by the intrinsic stochasticity frontier.
- \( D < 1 \), where \( S > 1 \): the collisional drag is strong, the distribution function shows a small departure from the Maxwellian bulk, and damping is nearly linear.

While increasing their parallel velocity, electrons move from an \( S > 1 \) high collisionality region to an \( S < 1 \) low collisionality region. A high efficiency scheme requires an optimized spectrum in order to push the adiabatic barrier in the parallel velocity beyond the \( D = 1 \) frontier. However, because of the relative weakness of the electron–wave interaction, this is only possible for high perpendicular velocities. Typically, in the conditions of Fig. 2(a), \( D \) is of the order of 1 (Fig. 2(c)) and reaches 5 for \( v_\perp = 0.7 \text{ c} \) (Fig. 2(d)). \( D \) remains much lower than one under the critical perpendicular energy \( W_{\perp \text{crit}} \).
5. CONCLUSIONS AND PROSPECTS

Our theory shows that the extension of the quasi-linear diffusion is limited in parallel velocity because of the gaps between discrete resonances, especially for \( v_\perp \approx v_{\perp \text{crit}} \) and on integer-\( q \) flux surfaces. The absorption profile should, therefore, exhibit dips around the rational-\( q \) surfaces, and, interestingly enough, such a dip has indeed been observed around the sawtooth inversion radius in the JET TTMP heating experiment [2]. In addition, when quasi-linear diffusion holds, the formation of a plateau in the electron distribution function, which would further increase the efficiency, seems possible at high \( v_\perp \), where TTMP dominates the collisional drag.

To ensure the destruction of adiabaticity, the distribution of the resonant velocities is crucial. A natural enrichment of the poloidal spectrum is provided by toroidal coupling along the multiple passes of the wave through the plasma [5], but such an uncontrolled broadening of the spectrum may lead to a loss of directivity. This will be studied when our theory will be applied in the ALCYON full wave code.

In conclusion, we may state that the following prescriptions seem to be relevant on optimizing the FWCD scheme:

— Controlled enrichment of the resonance pattern, by launching different frequency waves from different sets of antennas, with asymmetric, broad spectra. A frequency ‘wobulation’, on a time-scale shorter than the collision time, can also be envisaged.

— Application of FWCD in synergy with ECRH, in order to take full advantage of TTMP [6].

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FAST-WAVE CURRENT DRIVE MODELLING FOR LARGE NON-CIRCULAR TOKAMAKS*

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FAST-WAVE CURRENT DRIVE MODELLING FOR LARGE NON-CIRCULAR TOKAMAKS.

ABSTRACT. The feasibility of using fast waves to drive substantial toroidal currents is investigated using ray tracing, two-dimensional (2-D) full-wave ICRF modeling and a 2-D code which self-consistently calculates current drive bootstrap current and MHD equilibrium. It is found that \( \eta_\parallel \) upshifts due to toroidal effects drastically decrease current drive efficiency \( \gamma \) for antennas mounted far off the midplane. With antennas on the midplane, \( \gamma \sim 0.3 \) to 0.4 should be obtainable. Self-consistent equilibria including bootstrap effect are found in which 18.2 MA of plasma current is driven by 120 MW of RF power: 65 MW of lower hybrid plus 55 MW of fast waves.

1. INTRODUCTION

It is widely recognized that a key element in the development of an attractive tokamak reactor, and in the successful achievement of the mission of ITER, is the development of an efficient steady-state current drive technique. Fast waves in the ion cyclotron range of frequencies hold the promise to drive steady-state currents with the required efficiency and to effectively heat the plasma to ignition. Advantages over other heating and current drive techniques include low cost per watt and the

ability to penetrate to the center of high-density plasmas. The following primary issues must be resolved:

1. Can an antenna array be designed to radiate the required spectrum of waves and have adequate coupling properties?
2. Will the RF power be efficiently absorbed by electrons in the desired velocity range without unacceptable parasitic damping by fuel ions or alpha particles?
3. What will the efficiency of current drive be when toroidal effects such as trapped particles are included?
4. Can a practical RF system be designed and integrated into the device?

We have addressed these issues by performing extensive theoretical calculations with ORION, a two-dimensional (2-D) code, and the ray tracing code RAYS, which calculate wave propagation, absorption, and current drive in tokamak geometry, and with RIP, a 2-D code that self-consistently calculates current drive and MHD equilibrium. An important figure of merit in this context is the integrated, normalized current drive efficiency, \( \gamma = \frac{\left( n_e \right)}{10^{20} \text{ m}^{-3}} I_{RF} (R/P_{RF}) \), where \( n_e \) is the average electron density, \( I_{RF} \) is the total RF-driven current, and \( P_{RF} \) is the RF power required. The calculations presented here emphasize the ITER device. We consider (1) a low-frequency scenario (\( f \approx 17 \text{ MHz} \)) such that no ion resonances appear in the machine and (2) a high-frequency scenario (\( f \approx 51-59 \text{ MHz} \)) such that the deuterium second harmonic resonance is just outside the plasma and the tritium second harmonic is in the plasma, midway between the magnetic axis and the inside edge. In both cases electron currents are driven by combined transit-time magnetic pumping (TTMP) and Landau damping of the fast waves. We find that fast waves are quite competitive with other current drive techniques.

2. NON-INDUCTIVE CURRENT DRIVE IN TOKAMAKS AT FINITE ASPECT RATIO

Electron trapping may severely reduce RF current drive (RFCD) efficiency for normalized wave phase speeds \( \omega / k || v_e < 1 \), while the reduction is more tolerable at \( \omega > 1 \). An analytic function for the normalized efficiency, \( \tilde{j} / \tilde{p} \equiv \tilde{\eta} \), has been provided in Ref. [1] by fitting a semi-analytic derivation to accurate numerical calculations. The expression for \( \tilde{\eta} \) is a function of inverse aspect ratio \( \epsilon \) on a flux surface, ion charge \( Z \), normalized phase speed \( \omega \), and poloidal angle at which the electrons receive heating power from the wave, \( \theta \). Two limiting forms of \( \tilde{\eta} \) were calculated, corresponding to Landau damping and Alfvén-wave damping (combined TTMP and Landau damping). Figure 1 illustrates the agreement between the functional form of \( \tilde{\eta} \) and several numerical calculations. The ray tracing studies reported here interpolate between the Landau and Alfvén-wave limits by monitoring
FIG. 1. Current drive efficiency versus parallel phase speed for Alfvén-wave current drive with $Z = 2.0$. Top curve is $\epsilon = 0$; dashed curve is $\epsilon = 0.03$ and $\theta = \pi$; and bottom curves (with open symbols) are $\epsilon = 0.10$ with wave damping at four different poloidal angles (triangles: $\pi$; circles: $3\pi/4$; squares: $\pi/2$; diamonds: 0). Curves are analytical formula and points are numerically derived.

the ratio $X = (\omega^2/\omega_p^2)(m_e^2/T_e)$. When $X \geq 1$, typical of lower hybrid current drive (LHCD), the Landau process dominates; $X \lesssim 1$, typical of low-frequency fast-wave current drive (FWCD), corresponds to the Alfvén-wave result.

3. FWCD EFFICIENCY IN ITER AT 17 MHz

In these calculations we have taken ITER parameters as given for the steady-state technology phase: major radius $R_0 = 6$ m, $R_{\text{min}} = 2.15$ m, $B_0 = 4.8$ T, $\kappa = 2.0$, $\delta = 0.4$. There is an additional 10 cm scrape-off region between the plasma and the antenna. We have assumed a peak temperature of $T_e(0) = 33$ keV, $\langle T_e \rangle \sim 20$ keV, and peak density $n_e(0) \sim 1.1 \times 10^{20}$, $\langle n_e \rangle = 0.7 \times 10^{20}$, where $\alpha_T = 1.0$, $\alpha_n = 0.5$. The antenna arrays under consideration consist of a great many (24–60) individually phased straps. Therefore, the power spectrum excited is very narrow in toroidal mode number $N_T$. The designs considered include an array of in-port and wall-mounted antennas at the equatorial plane of the tokamak and an array of antennas integral with the blanket modules.
located above the equatorial plane at θ ≈ 40°. With the antennas located at the equatorial plane we obtain power deposition and RF-driven current profiles that are peaked on axis. From the RAYS code we obtain γ in the range 0.25 to 0.39 and one-pass absorption of 0.27 to 0.59 for toroidal mode numbers 4 ≤ N_T ≤ 7. These results are in substantial agreement with calculations using the ORION full-wave code. We have also considered a poloidal array of antennas (four individually phased straps) filling the entire port and focused poloidally toward the magnetic axis. For example, rays launched from 20° above the equatorial plane have γ in the range 0.26 to 0.33 with one-pass power absorption in the range 0.25 to 0.54. These rays also produce a current profile that is peaked on axis.

When the launch point is not at the equatorial plane, n_∥ varies widely along the ray trajectory owing to the large projection of k in the poloidal direction. For example, when the initial launch point is at θ = 40°, rays launched with n_∥ ~ 1 rapidly upshift to n_∥ ~ 5. Since the current drive efficiency depends sensitively upon v_{phase}/v_{thermal}, the effect is to drastically reduce γ. Also, rays launched at this location do not pass through the magnetic axis, resulting in hollow profiles for power deposition and driven current. As a result, the maximum γ that we were able to attain with the launch point at 40° was 0.24 and more typically was of order 0.1.

Calculations with ORION show strong evidence of toroidal and poloidal eigenmode structure. By tracking these modes it is possible to obtain very peaked power deposition profiles and therefore large values of γ (∼0.45) and high antenna loading resistance. Unfortunately, because of an approximation for k_∥ used in the code (k_∥ = k_{toroidal}), which neglects the k_∥ upshift, the values obtained for γ in ITER are not reliable for waves propagating very far above or below the equatorial plane. Also, the ray tracing calculations do not include the coherent interference structure of the multiple rays. Thus, any ability to control the deposition by tracking eigenmodes does not appear in the ray tracing. We conclude that γ in the range 0.3 to 0.4 should be obtainable with an antenna array located at the equatorial plane and that improvements may be possible by tracking modes.

4. SELF-CONSISTENT ITER MHD EQUILIBRIA WITH FWCD AT 59 MHz

Self-consistency of RFCD and the plasma equilibrium is found for a fixed pressure, p(Ψ), as a function of the normalized poloidal flux ˜Ψ. The flux-surface-averaged bootstrap current density, H = ⟨j_∥^B⟩/⟨B^2⟩, and RF current density, G = ⟨j_∥^{RF}B⟩/⟨B^2⟩, determine the equilibrium [2], and these functions are calculated on the equilibrium flux surfaces. Near the edge (˜Ψ = 0) LHCD contributes to G, while near the magnetic axis
FIG. 2. Self-consistent RFCD/MHD equilibrium with $I_0 = 18.2$ MA, $\bar{n}_e = 0.82 \times 10^{20}$ m$^{-3}$, Troyon coefficient of 3.4, and RFCD power of 120 MW. (a) Contributions to $\langle j_B \rangle$ from LHCD and FWCD (dashed), bootstrap (chain-dotted), and total (dotted). (b) Toroidal component of current density in midplane. (c) Ray trajectories. (d) Safety factor.

(\bar{\psi} = 1) FWCD provides the 'seed current' for the bootstrap effect. The current drive efficiency, $\gamma \equiv \bar{n}_e R_0 I_0 / P_{RF}^F$, computed from the resulting equilibrium current $I_0$ and the total launched RF power, is a composite figure of merit that accounts for all RF contributions. A superscript $B$ ($\gamma^B$) indicates that the bootstrap contribution to $I_0$ is also included.

In contrast to wave-driven currents, the bootstrap current increases with smaller aspect ratio. The present results are based on the finite-aspect-ratio formulation of Hirshman [3], which is valid in the banana regime for a single ion species. Whereas the RFCD theory is easily generalized for multiple ion species, the accuracy of the bootstrap results remains uncertain for arbitrary mixtures of ions, especially when similar masses are involved (D, T, He). Nonthermal ions (alphas) are also treated in an approximate fashion in the calculation.

In these ray tracing calculations we use the ITER specifications: $A = 2.8$, $\kappa = 2.0$, $\delta = 0.55$, $R_0 = 6.0$ m, $B_0 = 4.85$ T, and peak pressure $p_0 = 1.58$ MPa. The temperature profile is $T(\Psi) = T_0 \bar{\Psi}^{0.9}$ and the density
profile is \( n(\Psi) = n_0 \Psi^{0.5} \). The temperatures \( T_i = T_e = 38 \text{ keV} \) and densities \( n_e = 1.3 \times 10^{20} \text{ m}^{-3} \) were adjusted to have \( \beta_T = 6\% \) and \( Z \approx 2 \). Based on the divertor location and preliminary antenna studies, we limited the fast-wave antenna location to the outboard first wall within the bounds of \( \pm 2.8 \text{ m} \) from the midplane. For frequencies \( \sim 51-59 \text{ MHz} \) and \( n_0 \approx 1.2 \), the fast wave is strongly damped on electrons in a single pass through the plasma, and ion damping (mainly second harmonic tritium) is small—typically only \( \sim 5\% \) of the launched power. The lower hybrid waves (8.0 GHz) achieve maximum penetration at \( n_\parallel = 1.9 \), but at high \( T_e \) they cannot reach very deeply into the plasma. The wave spectra versus \( n_\parallel \) are adjusted to achieve equilibria with monotonic safety factors \( q(\Psi) \) and minimal RF power.

With no bootstrap contribution, additional high-frequency (\( \sim 800 \text{ MHz} \)) fast-wave power is needed to fill in the current density profile around the mid-minor radius (\( \Psi \sim 0.5 \)). In a typical calculation without the bootstrap effect, \( I_0 = 16.7 \text{ MA} \) is generated by a total RF power \( P_{RF} = P_{LH} + P_{HHF} + P_{HFF} = 338 \text{ MW} \), with \( \gamma = 0.26 \).

Of course \( P_{CD} > 200 \text{ MW} \) is unacceptable for ITER, but the bootstrap effect provides large savings in external power. Figure 2 illustrates an equilibrium solution including the bootstrap current, and the result is much more reasonable, \( I_0 = 18.2 \text{ MA} \) and \( P_{RF} = 120 \text{ MW} \), provided by \( P_{LH} = 65 \text{ MW} \) and \( P_{FW} = 55 \text{ MW} \). In this case \( \gamma^B = 0.75 \). Thus, roughly two-thirds of the total current is supplied by the bootstrap effect. Further reductions in the fast-wave power and complete elimination of the lower hybrid system allow the generation of equilibria with almost all the equilibrium current provided by the bootstrap effect. In the extreme case, equilibrium is achieved with \( P_{RF} = 7.5 \text{ MW} \), provided by only fast waves at 51 MHz, with \( \gamma^B = 8.8 \); the drawback in this limit is that the total current is low \( (< 14 \text{ MA}) \). However, MHD stability of similar low-current equilibria was tested and found acceptable in the context of the ARIES tokamak reactor study [4].

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EFFICIENT CURRENT DRIVE WITH HIGH FREQUENCY FREE ELECTRON LASER RADIATION

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Abstract

EFFICIENT CURRENT DRIVE WITH HIGH FREQUENCY FREE ELECTRON LASER RADIATION.

Current drive in tokamak plasmas with high frequency free electron laser (FEL) radiation is investigated. High power pulsed FELs at millimetre wavelengths are capable of exciting non-linear wave-wave interactions that produce high phase velocity electrostatic waves. These accelerate electrons giving rise to a slowly decaying current. The best current drive efficiencies are obtained with stimulated Raman forward scattering and left hand, circularly polarized pump wave. Ray tracing studies show that FEL beams can penetrate into the plasma centre, which makes the method ideal for bootstrap current seeding. Bootstrap effect enhances the overall current drive efficiency considerably.

1. INTRODUCTION

Promising experimental results have been achieved with non-inductive current drive (CD) techniques, such as neutral beams and lower hybrid waves. In reactor relevant plasma densities, these methods, however, suffer from serious problems. The progress in free electron laser (FEL) technology has opened new interesting possibilities to current drive with intense microwave pulses ($\lambda_0 < 2$ mm) [1, 2]. High frequency microwaves can easily penetrate into the plasma centre, which makes a control of the current profile possible. This feature is particularly attractive in bootstrap current seeding. The bootstrap amplification of the FEL driven current may enhance the overall current drive efficiency substantially.

Free electron lasers have a potential for GW peak powers, large pulse repetition rates, high efficiencies and low cost per watt as compared with high power gyrotrons [3]. Large peak powers are appropriate to the excitation of high phase velocity ($v_{ph} \gg v_e$), electrostatic modes by various wave–wave interactions [1, 2]. Fast
electrostatic waves accelerate electrons to large parallel velocities \( v_\parallel = v_{ph} \), generating a slowly decaying current. Electrons with \( v_\parallel > v_\perp \) are not trapped into banana orbits, which is a clear advantage over the electron cyclotron current drive.

2. NON-LINEAR PLASMA WAVE GENERATION

Large amplitude electrostatic plasma waves can be excited by beat wave (BW) interaction or by parametric decay of intensive FEL radiation. In stimulated Raman scattering (SRS), an intense electromagnetic pump wave (T) decays into a scattered wave (T') and a longitudinal plasma wave (L): \( T \rightarrow T' + L \). Energy and momentum conservation relations for the waves are: \( \omega_0 + \omega_s = \omega + \omega_p \) and \( \mathbf{k}_0 = \mathbf{k}_s + \mathbf{k}_p \), where '0' and 's' refer to the pump wave and to the scattered wave, respectively. In SRS the lower frequency electromagnetic wave grows up from the noise level or from a weak broadband seed beam. The advantage of the SRS current drive over the BW method is that only one high power FEL beam is needed. Furthermore, the Raman process is less sensitive to density variations.

Both Raman forward (SRS-F) and backward (SRS-B) scattering can be used to current drive. In SRS-B the gain and the momentum of the plasma wave are larger than in SRS-F. The phase velocity of the SRS-B plasmon is, however, fairly small in a hot, well underdense plasma so that Landau damping severely limits the growth of the SRS-B instability. In SRS-F the plasmon phase velocity \( v_{ph} = \omega/k \) is much higher and its linear damping is negligible. Thus SRS-F can dominate over SRS-B, even though its gain is smaller. The high phase velocity, \( v_{ph} > 0.6 c \), means that SRS-F generates nearly collisionless electrons. This leads to a slow decay of the resulting current.

An apparent problem is the weak coupling of the fast plasma wave to the electron distribution. In the strong pump limit the plasmon amplitude can be large enough to trap electrons from the tail of the distribution. Also at high electron temperatures, \( T_e > 20 \) keV, the wave–particle interactions are stronger, because \( k\lambda_D > 0.2 \). An interesting possibility is simultaneous operation of SRS-B and SRS-F, which may substantially enhance the fast electron production \[4, 5\]. The SRS-B generates a large number of medium fast electrons, which can further be accelerated by the SRS-F plasma wave.

We assume a simple geometry where all waves propagate parallel to the magnetic field. This gives the best current drive efficiency according to 2-D calculations \[6\]. The pump wave can be either right (RCP) or left (LCP) hand circularly polarized. The magnetic field effects are particularly important at low frequencies where the current drive is expected to be most efficient.

The basic gain length for the Raman process is

\[
L_g = \frac{2 \sqrt{\xi}}{\omega B} \left( \frac{c}{v_0} \right) \left( \frac{\omega_p}{\omega} \right) \left( \frac{k_s}{k} \right)^{1/2} \lambda_D
\]
where \( v_0 = e |E_0|/m_e \omega_0 \) and \( E_0 \) is the input pump amplitude. The magnetic field dependence is included in the factor \( \alpha_B \) [2] and in the wavenumbers. The gain length gives a scale for the required interaction length. In most cases the SRS-F plasmon is weakly damped and the pump is completely depleted (i.e. full conversion \( R_c \sim 1 \)) at a distance

\[
L_{\text{depl}} = \frac{L_g}{2} \ln \left( \frac{4}{\epsilon_N} \right) \tag{2}
\]

where \( \epsilon_N \) is the electromagnetic noise level. When the linear damping is heavy, the depletion length is given by \( L_{\text{depl}} \approx \left( L_g^2 / 2L_e \right) \ln(4/\epsilon_N) \), where \( L_e \) is the linear damping length of the Raman plasmon. In SRS current drive \( \epsilon_N \) can be increased with the aid of a weak, external, broadband seed beam at \( \omega_s \). Recall that in the beat wave case \( \epsilon_N \sim 1 \).

If the FEL peak intensity is \( I_0 \lambda_0^2 = 10^9 \text{ W \cdot cm}^{-2} \cdot \text{mm}^2 \), the gain lengths are below 10 cm for SRS-F and around 1 cm for SRS-B. The ratio \( L_g / L_e \) is, however, very large in the SRS-B case and thus the heavy damping suppresses the growth of the process. On the other hand, in SRS-F, full conversion \( R_c \sim 1 \) occurs fairly easily at a distance of 0.5–2 m.

### 3. CURRENT DRIVE EFFICIENCY

In the steady state, the momentum transferred from the pump to the plasma wave equals the momentum transferred from the plasma wave to the resonant electrons. The efficiency of the momentum transfer from the pump to the plasma wave is described by the relative action transfer \( R_c \) [1]. We shall assume that the plasma wave deposits all its momentum to the fast electrons that have \( v = v_{ph} \) and \( v = v_f \).

The current drive efficiency can be determined by balancing the momentum obtained by the resonant electrons with the momentum lost by them in collisions. Including relativistic corrections and both like and unlike particle collisions, we write the resulting CD efficiency as

\[
\frac{n_e R_c}{P} = \frac{2\epsilon N c^2}{e^3 \ln \Lambda} \frac{\gamma(v/c)^2}{1 + Z + \gamma} \left( \frac{\omega}{\omega_b} \right) R_c \tag{3}
\]

where \( \ln \Lambda = 17 \) is the Coulomb logarithm and \( \gamma = [1 - (v/c)^2]^{-1/2} \). In the relativistic limit (\( \gamma \to \infty \)), the current drive efficiency (3) is given by \( 1.84 \times 10^{20} \times (\omega/\omega_b)R_c \left( A \cdot W^{-1} \cdot \text{m}^{-2} \right) \), which agrees with the result of Ref. [7] if the quantum efficiency and \( R_c \) are set equal to unity. In the non-relativistic limit (\( \gamma \to 1 \)), the efficiency (3) is somewhat lower than predicted by a 2-D analysis [7].
Current drive efficiency (3) versus pump wavelength for LCP and RCP pump waves when \( R_c = 1.0 \). Solid curves: SRS-F; dashed curves: SRS-B, at temperatures of 3, 10 and 30 keV (\( n_e = 10^{20} \text{ m}^{-3} \), \( B = 4.8 \text{ T} \), \( Z = 1.3 \)).

We emphasize that \( P \) in the efficiency (3) is the average pump power fed into the plasma. A high peak power \( \hat{P} \) (or peak intensity \( \hat{I} \)) is required to achieve \( R_c \sim 1.0 \). The relation between the peak and the average power is \( P = \nu_{\text{rep}} t_p \hat{P} \), where \( t_p \) is the pulse length and \( \nu_{\text{rep}} \) is the repetition rate.

Figure 1 illustrates current drive efficiencies (3) for various FEL wavelengths \( \lambda_0 \). Notice that the quantum efficiency \( \omega/\omega_0 \) and \( \nu \approx \nu_{\text{ph}} \) in Eq. (3) depend on the pump wavelength. Figure 1 predicts much better current drive efficiencies for the SRS-F process, because it produces faster (less collisional) electrons than SRS-B. At short wavelengths the RCP and LCP cases are similar; at long wavelengths the LCP pump gives better efficiencies.
Ray tracing calculations reveal that at long wavelengths the beam propagation becomes complicated. The beams diverge, which can be seen by comparing the diameters of the pump and seed beams at the input plane ($d_{in}$) and one metre after the resonance point ($d_{out}$). Figure 2 illustrates the beam divergences, $d_{out}/d_{in}$, for SRS-F. The LCP seed beam suffers severely from the beam curvature above $\lambda_0 = 1.5$ mm, the RCP beams propagate almost straight at $\lambda_0 = 0.6$ mm, where the maximum CD efficiency occurs.

4. BOOTSTRAP AMPLIFICATION

The FEL driven current can be localized around the magnetic axis, which allows an amplification of the current by the neoclassical bootstrap effect [8]. Assuming a form $1 - (r/a)^\nu$ for the tokamak density and temperature profiles, we obtain the total bootstrap amplified current in the limit of a small seed current ($J_s < 0.5 J_B$) [9]:

$$J_B = 40.3 f(\nu) \left( \frac{1 + \kappa^2}{2} \right)^{1/2} RA^{-5/4} \left( \frac{n_0 k_B T_0}{\mu_0} \right)^{1/2}$$

(4)

where $f(\nu) = \nu (2\nu + 5)^{-1/2} (4\nu + 5)^{-1/2}$, $A$ is the aspect ratio $R/a$ and $\kappa$ is the elongation. $T_0$ and $n_0$ refer to the peak values on the magnetic axis. It is straightforward to generalize Eq. (4) for unequal temperatures, $T_e0 \neq T_i0$, and profiles, $\nu_e \neq \nu_i$ [9].

The total current $J_{tot}$ and safety factor profiles for various seed currents $J_s$ are shown in Fig. 3 for an ITER/NET size tokamak. The seed current radius is chosen such that $q(0) = 1$. The $q$-profiles in Fig. 3 reveal that too low seed currents ($J_s < 2$ MA) lead to MHD unstable situations, i.e. to hollow current profiles. Bootstrap amplification by a factor of five to seven is, however, possible without MHD stability being lost. Thus the CD-efficiency, $nRJ/P$, in Fig. 1, may be further improved by the same factor. Figures 1 and 3 predict a rather high value $J_{tot}/P \approx 0.5$ A/W for an ITER/NET size tokamak.

The beta limit $\langle \beta \rangle = 0.028 \mu_0 J/aB_T$ (a Troyon factor of 3.5 assumed), with the bootstrap current (4) reduces to a very simple form [9]:

$$\langle \beta \rangle = 0.033 \left( \frac{a}{R} \right)^{1/2} \frac{1 + \kappa^2}{2}$$

(5)

which depends only on the plasma geometry. The pressure profile dependence disappears, because the factor $f(\nu)^2 \rho/\rho (\nu) = 0.1$ is nearly constant for all values of $\nu$. The relation (5) calls for a tight aspect ratio and high elongation for $\langle \beta \rangle > 5\%$. 
5. SUMMARY

In a reactor grade tokamak plasma \((n = 10^{20} \text{ m}^{-3}, T > 10 \text{ keV}, B = 5 \text{ T})\), the best current drive efficiency is obtained with the Raman forward process. If an LCP pump wave is used, a CD efficiency of \(nRJ/P = 4 \times 10^{19} \text{ A} \cdot \text{W}^{-1} \cdot \text{m}^{-2}\) can be achieved. When the propagation limitations of the FEL beam are taken into account, the optimal pump wavelengths are 1–1.5 mm.

When an RCP pump wave and the Raman forward process are used, the efficiency is somewhat lower than in the LCP case: \(nRJ/P \approx 2 \times 10^{19} \text{ A} \cdot \text{W}^{-1} \cdot \text{m}^{-2}\). Ray tracing calculations indicate that the RCP beams propagate almost straight, when the pump wavelength has its optimum value of \(\lambda_0 = 0.4–0.8 \text{ mm}\). Both in the LCP and RCP cases the CD efficiency can be improved by a factor of 5–7 with the bootstrap amplification. This makes the overall current drive efficiency excellent.

REFERENCES

CURRENT DRIVE BY ALFVÉN WAVES (HELCITY INJECTION)

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Abstract

CURRENT DRIVE BY ALFVÉN WAVES (HELCITY INJECTION).

The results of theoretical and experimental research on current generation by Alfvén waves are presented. Averaged RF forces generating steady state currents additional to drag currents are theoretically analysed. Additional currents are due to RF helicity injection and gradient forces. DC current drive is demonstrated experimentally using standing helical Alfvén waves. The experimental data indicate the possible relation of this current to RF helicity injection.

1. INTRODUCTION

Recently Ohkawa [1] and other authors [2] proposed a method of current drive by Alfvén wave (AW) helicity injection. The DC current may be maintained by transferring the momentum from the wave to the plasma electrons, but neoclassical effects due to trapped electrons may essentially reduce the drag current generation efficiency by the AW [3]. Under such conditions, current drive using RF helicity injection may be more efficient.

The experiments [4–6] on AW current drive by travelling and standing AW in the R-0 stellarator have shown that, besides the drag current, additional currents appear.

2. THEORETICAL STUDIES

We give here the derivation of the stationary current from the analysis of averaged RF forces (see, for example, Refs [2, 3]) using the kinetic equation with the collision integral in Landau form. We shall use a model of a cylindrical, radially inhomogeneous plasma column with identified ends and constant parameters along the homogeneous magnetic field, $\mathbf{B}_0$, which is directed along the z axis.

† Deceased.
After averaging kinetic equations for electrons and ions over the time and space periods of the RF field, multiplying them by \( m_e, m_i \), combining them and integrating over the velocity space, we obtain an equation for the \( z \) component of the averaged current \( \langle j_z \rangle \):

\[
\langle j_z \rangle = m_e \frac{\nu_{ei}^{(0)}}{e^2} = \left\langle E_z \left( \frac{n_e - m_e}{m_i} \bar{n}_i \right) \right\rangle - \frac{1}{c |e|} \left[ \left( \frac{\bar{\sigma}_e - m_e}{m_i} \bar{\sigma}_i \right) \times \vec{B} \right]_z.
\]

(1)

Note that this equation follows from the generalized Ohms's law; the corrections connected with ion oscillations may be neglected.

Further, we use the continuity (\( i \varepsilon q \) \( \bar{n} = \text{div} \bar{j} \)) and induction (\( i \Omega \) \( \vec{B} = c \text{rot} \vec{E} \)) equations for calculating the power density absorbed by electrons \( P_e = (1/2) x \Re(\bar{E}_e \bar{E}_* \bar{E}_e) \), where \( \bar{E}_e = (-i \Omega/4\pi) e^{(e)} \bar{E}_0 \) (s = 1, 2, 3) is the oscillating electron current and \( \epsilon_{0e} \) is the electron part of the plasma dielectric tensor operator.

Let us consider the RF wave field in the geometric optics approximation:

\[
\delta E_{x,z} = \bar{E}_{x,z}(r) \exp \left[ i \left( \int_0^r k_z \text{d}r + m_0 + k_x z - \Omega t \right) \right]
\]

From the Maxwell equations for the fast wave we obtain:

\[
\epsilon_{33} E_z + \epsilon_{32} E_\phi = 0
\]

(2)

The components of the tensor \( \epsilon_{ij} \) are taken from Ref. [7]:

\[
\epsilon_1 = \omega_{pi}^2 / (\omega_{ci}^2 - \Omega^2); \quad \epsilon_2 = i \epsilon_{12} = i \epsilon_1 (\Omega / \omega_{ci})
\]

\[
\bar{\epsilon}_{22}^{(e)} = \epsilon_1 \left[ 1 - \frac{2 \partial}{n_e r \partial r} \left[ r \eta \bar{\nu} \bar{\nu} \right] \right]; \quad \bar{\epsilon}_{32}^{(e)} = (1 - \eta) \frac{\omega_{pi}^2}{k_x \omega_{ci}} \left[ \frac{\partial}{r \partial r} - \frac{r}{r} \right]
\]

\[
\bar{\epsilon}_{23}^{(e)} = - \frac{1 - \eta}{\Omega} \frac{\partial}{\partial r} \left[ \frac{\omega_{pi}^2}{k_x \omega_{ci}} \right]
\]

\[
\eta = - i \sqrt{\frac{\pi}{2}} \frac{\bar{\nu}}{\nu_{Te}} W \left( \sqrt{2 \nu_{Te}} \right)
\]

is the Kramp function, where \( \nu_\phi = \Omega / |k_x| \) and \( \nu_{Te} = \sqrt{T_e / m_e} \). Then,

\[
\langle j_z \rangle = - \left( \frac{|e| k_z}{m_e \nu_{ei} \Omega} \right) \left[ P_e + \frac{1}{2 k_z} \frac{\partial}{\partial r} \text{Im} (\bar{E}_z \bar{j}_z^{(e)}) \right]
\]

(3)
Here, the first term in brackets corresponds in homogeneous plasma to drag current \([3]\), and the second consists of the current corresponding to helicity injection \([1, 2]\) and a gradient current:

\[
j_{\text{add}} = - \left( \frac{|e|}{8\pi e_i m_e} \right) \text{Im} \left\{ \varepsilon_{12}^{(e)} \left[ \frac{\partial}{\partial r} (\vec{E}_x^* \vec{E}_y) + \frac{1}{n} (\vec{E}_x^* \vec{E}_y) \frac{\partial n_e}{\partial r} \right] \right\}
\]  (4)

Note that a part of the helicity current in (4) is compensated with the part of the drag current in homogeneous plasma, owing to divergence- and gradient-like terms in the expression for \(\text{Im}(\vec{E}_x^* \varepsilon_{22}^e \vec{E}_y + \vec{E}_x^* \varepsilon_{23}^e \vec{E}_y)\) in \(P_e\), which do not enter in the real absorbed power \([8]\).

The value of generated current is reduced owing to the factor \(\eta\), which is small at \(v_\phi \ll v_{Te}\), and \(\eta = 1\) at \(v_\phi = v_{Te}\). Let us compare now the efficiency of current drive due to helicity injection with that of the drag current, \(j_{dr}\), for the fast wave in plane geometry:

\[
\frac{j}{P_e} = \frac{j_{dr}}{P_e} \left[ 1 + \left( \frac{\text{Im} k_r^2 - \kappa_p \text{Re} k_r}{\text{Im} \eta \text{Re} k_r^2} \right) + \frac{\kappa_p \text{Im} k_r}{\text{Re} k_r^2} \right]
\]  (5)

where \(\kappa_p = \partial (\ln nT) / \partial r\). The efficiency of additional current drive by the kinetic AW is lower than by fast modes, by the factor \(\beta = 8\pi nT / B_0^2\).

Numerical calculations have been carried out using the 1-D cylindrical code EPSI for ITER parameters \(m = -1\), \(n = 1\), \(q_0 = 0.85\), \(n_0 = 2 \times 10^{20} \text{ m}^{-3}\), \(B_0 = 4.85 \text{ T}\). They show that radial distribution of the current density due to helicity injection for the global AW changes its sign at the absorbed power density maximum. The total current achieves a magnitude of 18 MA with 1.2 kA/m surface current density in the RF helical antenna.

Thus the current (4), and in particular the helicity injection current, is a part of the total current (3) resulting from the averaged RF forces acting on the plasma. Owing to the helicity injection current and the gradient current, it is possible to sustain a current which may be comparable to drag currents in a direct magnetic field; such a current would be sufficient for effective sustainment of the DC current in tokamak reactors. Note that there is no need for circularly polarized waves \([1, 2]\) in order to sustain this current; wave damping can take place on both electrons and ions. A drawback of the helicity injection method consists in the changing sign of the current density over the radius.

3. EXPERIMENTAL STUDIES OF ALFVÉN WAVE CURRENT DRIVE

The experiments were carried out in the \(\ell = 3\) stellarator R-0 \([4–6]\) (\(R = 0.5 \text{ m}, b = 0.05 \text{ m}, a = 0.035 \text{ m}, B_i \leq 0.8 \text{ T}, \eta_0 \leq 0.8, n_{19} = 0.5–0.8, T \leq 100 \text{ eV}, \bar{P} \leq 400 \text{ kW}, I_{OH} = 0, H_2\)). The helical RF antenna makes it possible
to excite helical travelling and standing modes with m/n = ±1/±1, ±2/±2 or a standing mode with m/n = ±4/±4 (k_e = m/r, k_z = −n/R), f = 1.2 MHz. A special generator feeds the same antenna with m/n = 0/0 currents at f = 0.4–0.7 MHz in order to produce an initial plasma.

3.1. Current drive by travelling Alfven waves

A travelling AW generates a drag current, I_{TW}, owing to its momentum absorption by electrons [4], and also an additional DC current, I_{H}, which increases the rotational transform angle (\bar{I}_{TW} \bar{B}_t ≅ 0, \bar{I}_{H} \bar{B}_t > 0). As a rule, the plasma current direction coincides with that of the drag current at low plasma densities. For high plasma densities, in the case where the drag current is directed against the toroidal magnetic field (\bar{I}_{TW} \bar{B}_t < 0; n < 0), this additional current may exceed the drag current, so the total plasma current may be positive. Figure 1 shows the plasma current variations when density changes take place during a single discharge.

3.2. Current drive by standing Alfven waves

The experiments have been performed for standing helical AW modes: m/n = ±1/±1; ±2/±2; ±4/±4; f = 1.2 MHz (I_0H = 0). In all cases plasma densities were higher than the threshold value (n_0 > n_A), i.e. the local Alfven resonance (LAR) condition was met in the plasma: \Omega = k_1(r_A)c_A(r_A). For all of the standing helical m/n modes, within the broad range of variation of the parameters (\bar{n}, \bar{T}, B_t, \nu_0), a positive DC current was generated which increased the rotational transform angle (\bar{I}_{pl} \bar{B}_t > 0).

In the experiments on the m/n = ±1/±1 standing mode, the DC current achieved a value of 0.4 kA at P_{max} = 350 kW. For this mode, the regime of the ‘cold’ LAR with \Omega/k_1 > \nu_T (threshold T_e > 80 eV), \nu_e/\Omega < 1 or \nu_e/\Omega ≥ 1 has been realized, depending upon \bar{n}_e and \bar{T}_e.

For the case of m/n = ±2/±2, the maximum DC current was 0.6 kA at \bar{P}_{max} ≅ 400 kW. Here, both ‘cold’ and ‘hot’ LARs might take place with \Omega/k_1 \equiv \nu_T (threshold T_e ≥ 20 eV), \nu_e/\Omega < 1 and \nu_e/\Omega ≥ 1.

For the m/n = ±4/±4 mode, the hot LAR condition was met with \Omega/k_1 < \nu_T (threshold T_e ≥ 5 eV). The absorption regime was collisional with \nu_e/\Omega ≥ 1. The threshold density for this mode is high enough: n_{19}(r_A) ≤ 0.5 \times B_t^2 (kG). The antenna–plasma couplings were smaller than for the m/n = ±1/±1 and ±2/±2 modes, because of the sharper drop of the RF field over the radius from the antenna towards the plasma, \bar{B}(r) \propto r^{-m-1}. The plasma current was 100 A at \bar{P}_{max} ≅ 100 kW.

Experiments have been carried out to generate a steady state current by the m/n = 0/0 mode at f = 0.4–0.7 MHz. This mode produces a plasma with \bar{n}_e < 3 \times 10^{19} m^{-3}, \bar{T} < 10 eV, giving rise to plasma DC current, both positively and negatively directed: \bar{I}_{pl} \bar{B}_t ≅ 0.
4. DISCUSSION OF THE EXPERIMENTAL RESULTS

In the experiments without Ohmic currents, DC currents are generated by standing AWs. These currents are considerably greater in magnitude than those resulting from the toroidal plasma column equilibrium. The experimental scaling [4], $I_L \propto \frac{P}{T^{3/2}} B_t$, obtained for these currents is in fairly good agreement with the estimations from Ref. [2] for current drive by RF helicity injection for the case of a cold plasma, $\nu_e/\Omega \gg 1$. This fact allows us to suppose that the current drive in our experiments is due to RF helicity injection into the plasma.

The theory [2] predicts that the current drive efficiency with RF helicity injection will be much lower for $\nu_e/\Omega \ll 1$. The above theoretical analysis within
the frame of linear theory shows that in operation within the Alfvén continuum, at \( \Omega/k_\parallel < v_{Te} \), the \( I_H \) current drive efficiency (owing to AW helicity absorption) is lower than that of drag current generation, \( I_{TW} \). At the same time, the experimental data are nearly kept within the broad range of values \( \Omega/k_\parallel \approx v_{Te}, \nu_e/\Omega \approx 1 \), in the regimes with \( I_H > I_{TW} \).

Even in the experiments with relatively low levels of absorbed RF power, \( \tilde{P}/V \approx 5 \text{ MW/m}^3 \), strong non-linear phenomena are observed: modulations of the RF field in the plasma interior, threshold phenomena [4] and birth of main frequency harmonics. These facts indicate that non-linearities make essential contributions to AW–plasma interaction. Signals of an electric probe measuring the electric RF field component in the plasma record fast spikes (see Fig. 2) which are probably due to

**FIG. 2.** RF antenna current and RF field \( E_z \) component for standing AW, \( m/n = \pm 2/\pm 2; B_0 = 3 \text{ kG, } \nu_e = 0.4. \)
the relaxation processes in the plasma arising when the stellarator magnetic configuration is perturbed by the RF field of the wave. Magnetic field line reconnections may give rise to steady state currents [9].

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RESULTS OF dc HELICITY INJECTION EXPERIMENTS FOR TOKAMAK CURRENT DRIVE

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Abstract

RESULTS OF dc HELICITY INJECTION EXPERIMENTS FOR TOKAMAK CURRENT DRIVE.

Helicity injection as a means of non-inductive tokamak current drive has been studied on both the Current Drive Experiment (CDX) and the Continuous Current Tokamak (CCT) devices. Currents up to 1 kA and 6 kA, respectively, have been obtained, corresponding to q's in a range useful to tokamaks. At these levels, the method has an efficiency of up to 23% of ohmic current drive. Some correlations have been observed between current penetration into the plasma and fluctuations near \( f_{ij} \). The Current Drive Experiment Upgrade (CDX-U) device has recently replaced CDX, and initial operation has yielded 3.4 kA of plasma current from dc helicity injection.

1. Introduction

The desire for a steady-state tokamak plasma has motivated research into many possibilities for non-inductive current drive, dc helicity injection among them. Helicity, denoted \( \kappa \), is defined by \( \kappa = \int_V \vec{A} \cdot \vec{B} dV \), where \( \vec{A} \) is the magnetic vector potential, \( \vec{B} \) is the magnetic field, and \( V \) is the plasma volume. The helicity of a plasma is a measure of the intertwining of field lines, and for a tokamak plasma, this corresponds to the amount of plasma current. The equation governing the evolution of helicity in a plasma is derived from the resistive MHD equations[1]:

\[
\frac{\partial \kappa}{\partial t} + \int_S \vec{Q} \cdot \vec{n} dS = -2 \int_V \eta \vec{J} \cdot \vec{B} dV.
\]
where $\vec{Q} = 2\phi \vec{B} + A \times \frac{\partial \vec{A}}{\partial t}$ (for plasmas with fixed boundaries; see Ref. 2 for a form appropriate to plasmas with moving boundaries). Here $\vec{Q}$ is the helicity flux into or out of the volume. The first term in $\vec{Q}$ is the dc helicity injection term, and the second is the ac helicity injection term. If solely dc helicity injection is utilized, and the plasma is in steady state, then the equation reduces to

$$\int_S \phi \vec{B} \cdot \hat{n} dS = - \int_V \eta \vec{J} \cdot \vec{B} dV.$$

This equation demonstrates that dc helicity injection depends upon biased objects at the plasma edge to supply the plasma with helicity and thus maintain it against resistive decay. Furthermore, it predicts no reduction in effectiveness at high density. In practice, the method uses simple apparatus, and can be quite efficient.

2. Experimental Arrangement

In experiments on the Current Drive Experiment (CDX) and Current Drive Experiment Upgrade (CDX-U) at Princeton, and on the Continuous Current Tokamak (CCT) at UCLA, a biased, emissive cathode located at the plasma's edge has been utilized for dc helicity injection. For non-ohmic plasmas, an emissive cathode is an obvious necessity, in order to provide electrons to break down the fill gas at reasonable voltages. Once the plasma is formed, the presence of an emissive cathode in the helicity injection arrangement is still thought to be advantageous as more current, and hence more power, may be delivered to the plasma for a given voltage. The cathode is located to one side of the X-point of a divertor, at one toroidal location. Current emitted from the cathode flows along field lines in the edge, interacting with the plasma and, eventually, returning to the anode at the other side of the X-point. Often, the vessel is also allowed to serve as an anode.

By this method, plasmas which have the features of tokamak plasmas[3] have been formed and maintained in CDX with currents up to 1 kA ($q_a = 5$). The characteristic (current vs. cathode voltage) of these discharges and of discharges in CCT is shown in Fig. 1.

More recently, experiments have been conducted on the Continuous Current Tokamak (CCT) at UCLA. CCT plasmas have typical parameters of $R = 1.5$ m, $a = 0.15$ m ($a_{\text{vessel}} = 0.4$ m), $B_T = 0.3$ T, $n_e = 5 \times 10^{18}$ m$^{-3}$, and $T_e = 50$ eV. In these plasmas, the cathode has driven as much as 6 kA of current, corresponding to $q_a = 3.5$. 
FIG. 1. Current versus voltage for CDX and CCT discharges. Solid circles are data from CDX, open circles are data from cathode-only discharges in CCT, and solid triangles are data from ohmic startup discharges in CCT. Solid curves are the algebraic helicity model, assuming $\eta \propto V_{\text{cath}}^2$.

3. Results from CDX and CCT

Figure 2 depicts several plasma parameters as functions of time for a discharge in CCT with 6 kA of cathode-driven current. The discharge is started by the ohmic heating transformer, and driven to 25 kA, after which the transformer is shut off, and the plasma current is allowed to decay until the cathode begins to drive current. At the point where the cathode starts
FIG. 2. Plasma evolution during a 6 kA, cathode-sustained plasma in CCT. Plasma startup was performed with the ohmic heating system, after which current was driven by the cathode. (a) Loop voltage. Note reversal of the loop voltage when the cathode begins to drive current. (b) Plasma current. Current is driven first by the ohmic heating system, then by the cathode. (c) Cathode voltage. The cathode voltage is turned on early in the pulse to draw ion current so as to heat the cathode surface and improve emission. (d) Injected current. The injected current increases substantially in the later portion of the discharge when the plasma moves close to the cathode.
to function, the loop voltage becomes negative, indicating that any (small) ohmic drive thereafter is opposite the direction of current flow. Thus, the effect seen is non-ohmic current drive.

It is possible to simplify further the helicity equation given previously, assuming that the vessel and anode define the zero of potential, that $B_T$ is sufficiently uniform to be removed from the dissipation integral, that the plasma may be characterized by a volume averaged resistivity, $\bar{\eta}$, and that helicity dissipation due to fluctuations is much smaller than that due to mean-field effects. Then,

$$I_p = V_{cath} A_{cath} B_{cath} / 2\pi R\bar{\eta} B_T.$$ 

Here, $V_{cath}$ is the potential applied to the cathode, $A_{cath}$ is the cathode area, and $B_{cath}$ is the normal magnetic field at the cathode surface. For the experimental arrangements in CDX and CCT, $B_{cath} = B_T$. If all of the other parameters are known, then it is possible to compute $\bar{\eta}$ and, from it, the resistance of the plasma column. Knowing this, the power that would be necessary to sustain the plasma with an ohmic heating system may be computed. The ratio of this (estimated ohmic) power to the actual injected power, $I_{inj} V_{cath}$, is termed the power efficiency relative to ohmic current drive, and is given by

$$\epsilon_E = \frac{I_p^2 R_p}{I_{inj} V_{cath}} \simeq \frac{I_p A_{cath}}{I_{inj} A_{xsect}},$$

where $A_{xsect}$ is the area of the cross-section of the plasma.

The 6 kA discharge described above was sustained by 365 A of injected current at an applied cathode voltage of 400 V, meaning that the injected power was 146 kW. An ohmic discharge of the same resistivity would have required 34 kW to sustain, so helicity injection drives current at 23% of the efficiency of ohmic in this case. (By way of comparison, a typical discharge of about 580 A in CDX had an efficiency relative to ohmic of 9.3%[4].)

In CDX, the helicity balance of a 585 A discharge was studied by making Langmuir probe measurements of the temperature to infer the resistivity and current density profiles (assuming Spitzer resistivity and that $j \propto 1/\eta$). From the calculated resistivity and current profiles, the dissipation rate of helicity was computed to be 0.020 Wb$^2$/sec. In addition, since the cathode voltage and area were known, the helicity input rate was calculable and was found to be 0.0239 Wb$^2$/sec. The uncertainty in each quantity is about a factor of 2, so the rates balance, in agreement with the helicity equation.
The helicity equation also provides a model for the current-voltage characteristic of the discharge. The helicity equation, reduced by the assumptions enumerated above (including $B_{\text{cath}} = B_T$) is:

$$I_p = \frac{V_{\text{cath}} A_{\text{cath}}}{2\pi R \bar{\eta}}.$$  

The data from CDX fits this model, assuming $\bar{\eta} \propto \eta_{\text{cath}}^{-1}$. This model and the data points from CDX are shown in Fig. 1, and are in reasonable agreement. Data from the CCT experiments is also shown on the same plot, as open circles and triangles. The solid curve in the vicinity of the CCT points is the same model used for CDX, with the assumption $\bar{\eta}_{\text{CCT}} = \bar{\eta}_{\text{CDX}}$. In this case, the curve predicts the trend of the data, but there is greater scatter. This may be due to a failure to account for a wider range of operating conditions (the triangles, which denote plasmas started by the ohmic heating system, and are therefore likely to be hotter, are significantly farther off the curve than the others), or by a shortcoming of the model.

Previous measurements of the poloidal field profile\[3\] indicated that the plasma current density was peaked on axis, in spite of the fact that the cathode supplies current at the edge. In order to understand the physics of this apparent current transport, a survey of fluctuations in CDX discharges was conducted with a shielded double-loop radiofrequency probe,\[5\] which was inserted into the plasma, providing measurements of $E_r$ and $B_T$. Attention was most strongly focused on fluctuations near $\Omega_i$ (5.3 MHz for hydrogen in a 0.39 T field), where the probe measurements showed a strong enhancement of amplitude when $\frac{dI_p}{dt}$ was large. In particular, the power in fluctuations in this frequency range was larger by a factor of 50 during the startup phase of the discharge and during the interval when $I_p$ was rapidly increasing than it was during steady state conditions. Furthermore, the integral of the enhanced power over the duration of the current ramp-up was found to be proportional to the total change in plasma current. These observations suggest that modes in the ion cyclotron frequency range may be driven by (or may drive) rearrangements of the current profile; more detailed investigations of this phenomenon are underway.

4. The CDX-U Device

The CDX device has recently been replaced with the CDX-Upgrade. CDX-U is a low-aspect-ratio, completely dc helicity-driven device. Figure 3 shows a diagram of it. The machine has $R = 0.33$ m, $a = 0.24$ m, $\kappa \leq 1.5$ (vertical elongation), and $B_T = 0.15$ T on axis, and it will provide a facility
for studies of dc helicity injection at the 10-to-30 kA level. Preliminary operation of the machine has yielded 3.4 kA of plasma current with $B_T = 0.09$ T, $a \sim 0.15$ m, $V_{cath} = 500$ V, and $I_{inj} = 120$ A. This gives an approximate safety factor of $q \sim 9$, at an aspect ratio of 2.2.

Planned experiments include aspect ratio scans to determine the effect of aspect ratio upon dc helicity injection efficiency and the construction of an axisymmetric cathode. Since CDX-U will be able to accommodate plasmas with aspect ratios as low as 1.3, the predicted properties of low aspect ratio tokamaks[6,7] can be tested. These include plasma paramagnetism, large bootstrap currents, and less robust microinstabilities.

5. Summary

Helicity injection as a possible means of non-inductive current drive for tokamaks is being studied on the CDX, CCT, and, now, CDX-U devices.
Currents as large as 1 kA ($q_a = 5$) in CDX, 3.4 kA ($q_a \sim 9$) in CDX-U, and 6 kA ($q_a = 3.5$) in CCT have been sustained. A measurement of the helicity influx and consumption has been made in the CDX plasma, and balance was found, in keeping with the predictions of helicity theory. The current-voltage characteristic is also in keeping with the predictions of theory. Fluctuations at frequencies near the ion cyclotron frequency have been seen to become markedly stronger during times when the plasma current is changing rapidly, suggesting that such fluctuations may be related to current transport processes. Experiments are planned on the CDX-U device to develop the physics and technology of helicity injection at higher currents.

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HIGH FREQUENCY SURFACE HELICITY INJECTION IN TOROIDS WITH MOVING BOUNDARIES

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Abstract

HIGH FREQUENCY SURFACE HELICITY INJECTION IN TOROIDS WITH MOVING BOUNDARIES.

Recently it has been argued that the negative results of F-θ pumping experiments on helicity injection may be due to the effects of a moving plasma boundary. The paper re-examines the possibility of helicity injection at the plasma edge using oscillating fields around the ion cyclotron frequency. It is found that a combination of ion inertial effects and ion dissipation can lead to a significant helicity injection at the plasma surface. Estimates of the surface helicity injection indicate that steady state currents in a tokamak may be driven efficiently using this method.

1. INTRODUCTION

An attractive scheme for steady state current drive in tokamaks and reverse field pinches (RFP) is that associated with injection of magnetic helicity at the plasma surface [1]. In DC helicity injection schemes, one may use DC biased electrodes [2] or pressure gradients [3] on open lines at the plasma edge to drive steady plasma currents; the feasibility of this scheme has been clearly demonstrated by recent experiments [4]. On the other hand, in AC helicity injection schemes, one uses low frequency (ω < ωi, the ion cyclotron frequency) modulation of toroidal and poloidal fluxes in a plasma with closed magnetic surfaces to drive toroidal plasma currents. Experimentally, AC helicity injection schemes have been attempted on the ZT-40 RFP at Los Alamos National Laboratory [5] and on the ENCORE tokamak at the California Institute of Technology [6], but without much success. No significant currents could be maintained in either of these experiments using AC helicity injection. Recently, Bellan et al. [7] and Bellan [8] have proposed an interpretation of this negative result. Bellan [8] has demonstrated analytically that if one takes account of moving plasma boundaries which naturally result when low frequency compressional modes are excited in the plasma by the oscillating magnetic fields, then additional terms arise in the analysis which cancel the usual surface helicity injection terms, leaving only a small residual effect associated with resistive dissipation.

In this paper we examine the injection of AC helicity at the plasma surface using relatively higher frequency modes (ω ~ ωi) in which ion inertial effects (from the Hall terms in Ohm's law) become important. We demonstrate that if non-resistive
damping of the modes is included (which is done in our analysis using a model viscosity term), then significant current drive results. Our analysis demonstrates that the ion inertial effects lead to a crucial separation of the plasma fluid motion and that of the magnetic field, and this results in significant current drive. We also find that the important current drive terms may be interpreted as cross-helicity injection terms proportional to $\vec{V} \cdot \vec{A}$. Finally, we demonstrate the importance of incorporating surface wave mode effects in the analysis, since the boundary conditions at the plasma-vacuum boundary play a critical role in determining the magnitude of current drive. We should emphasize here that the present investigation assumes that the modes are periodic in the toroidal direction so that there is no contribution from terms proportional to the toroidal gradient of radiation pressure; this is in contrast to our earlier report [9]. In the final discussion, we compare our conclusions with those of the recently proposed bulk helicity injection scheme of Chan et al. [10].

2. AC HELICITY INJECTION WITH MOVING PLASMA BOUNDARIES

Following Moffat [11] and Bellan [8] we evaluate the time rate of change of the magnetic helicity function  $K = \int \vec{A} \cdot \vec{B} \, dr$ to obtain:

$$\frac{dK}{dt} = -2 \int \eta \, \vec{n} \cdot \vec{B} \, d\tau - \int \vec{n} \cdot (\nabla \times \vec{A}) \, dS$$

$$+ \int \vec{n} \cdot \vec{V} \, (\vec{A} \cdot \vec{B}) \, dS$$

(1)

where $\vec{n}$ is the unit normal to the plasma vacuum interface, $\vec{V}$ is the plasma velocity, $\nabla \times \vec{A}$ is the vector potential and $\eta$ and $\vec{n}$ are the current density and resistivity. Using Ohm's law, $\vec{E} + \vec{V} \times \vec{B} = \eta \, \vec{J}$, Eq. (1) can be rewritten as

$$\frac{dK}{dt} = -2 \int \eta \, \vec{n} \cdot \vec{B} \, d\tau - \int \vec{n} \cdot (\vec{A} \cdot \nabla \times \vec{A}) \, dS$$

$$+ \int \eta \, \vec{n} \cdot (\vec{A} \times \vec{J}) \, dS + \int \vec{n} \cdot [(\vec{A} \times (\vec{E} + \vec{V} \times \vec{B})) \, dS$$

(2)

In Eq. (2), the first term on the right side gives the resistive decay of helicity, the second term denotes the helicity injection on open lines and the last two terms give the helicity injection on closed magnetic surfaces. For low frequency AC helicity injection, note that inclusion of moving plasma boundary effects (i.e. $\vec{V} \neq 0$ at the plasma surface $S$) makes a significant modification of the helicity injection physics. For $\vec{V} = 0$, the last term gives a major contribution to AC helicity injection.
However, if the plasma and the magnetic field move together, as in an ideal fluid at low frequencies, the contribution of the last term vanishes; for $\eta \neq 0$ a weak effect proportional to $\varphi$ survives owing to resistive slippage of the plasma with respect to field lines. This, as Bellan [8] argues, explains the negative results of the AC helicity injection experiments.

We now consider the use of high frequency fields ($\omega \sim \Omega_i$) for helicity injection at the plasma surface. The basic idea is that in this case the plasma and the magnetic field do not move together and hence there may be a net injection of helicity. It turns out that the magnitude of net helicity injection crucially depends on the wave dissipation in the plasma. To take account of non-resistive damping of waves in the ion cyclotron range of frequencies by collisionless mechanisms (cyclotron damping, minority ion damping, etc.) we shall later incorporate a model ion viscosity term in the fluid equations. For fields with $\omega \sim \Omega_i$, the Ohm's law is modified to include the ion inertial effects, viz. $\vec{E} + \vec{V} \times \vec{B} = \eta \vec{J} + (\vec{J} \times \vec{B})/qn$, where $n$ is the number density. Note that since the Hall term is also orthogonal to the magnetic field $\vec{B}$, the magnetic helicity conservation theorem is not modified nor does it lead to any breaking of lines of force, i.e. $\vec{E} \cdot \vec{B}$ is still proportional to plasma resistivity. Using the modified Ohm's law, Eq. (2) can be written as:

\[
\frac{dK}{dt} = -2 \int \eta \vec{J} \cdot \vec{B} \, d\tau + \int \eta \cdot \left( \vec{A} \times \vec{J} \right) \, dS
- \int \hat{n} \cdot \vec{B} \left[ \varphi - \vec{A} \cdot \left( \vec{V} - \frac{\vec{J}}{qn} \right) \right] \, dS
+ \frac{1}{n q} \int \hat{n} \cdot \vec{J} (\vec{A} \cdot \vec{B}) \, dS
\]  

(3)

Ignoring the helicity injection due to $\hat{n} \cdot \vec{B}$ ('open' field line) terms, we may write the steady state helicity balance equation:

\[
2 \int \eta \vec{J} \cdot \vec{B} \, d\tau = \frac{1}{q n_0} \int \langle \hat{n} \cdot \vec{J}_1 \rangle \left( \vec{A}_0 \cdot \vec{B}_1 + \vec{A}_1 \cdot \vec{B}_0 \right) \, dS
\]

(4)

where subscript 1 denotes an oscillating term due to applied wave fields, $\langle \cdots \rangle$ denotes the time average and we have retained only the non-resistive helicity injection terms on the right side because of their larger magnitude in hot collisionless plasmas. Note also that we have dropped the term proportional to $\vec{A}_0 \cdot \vec{B}_0$ on the right side since it can be made to vanish by an appropriate gauge choice.

Equation (4) demonstrates that the current drive in steady state is maintained by oscillating terms proportional to $\hat{n} \cdot \vec{J}_1$ at the plasma surface. These are the terms which generate oscillating surface charge at the plasma boundary and therefore couple to surface waves. It is therefore of importance to consider the excitation of plasma
current by surface wave modes near the ion cyclotron frequency. Finally, it is worth pointing out that the $\hat{n} \cdot \vec{J}$ terms are related to ion polarization current across the $\vec{B}$ field and hence also to the ion motion $\hat{n} \cdot \vec{V}$. By writing an overall conservation theorem for kinetic helicity defined as $K_i = \int \vec{P} \cdot \nabla \times \vec{P} \, d\tau$, where $\vec{P} = (m_i \vec{V}/q + \vec{A})$ we have been able to demonstrate that the helicity injection terms in Eq. (4) can be interpreted in terms of injection of cross-helicity terms (proportional to $\vec{V} \cdot \nabla \times \vec{A}$) at the plasma surface; this will be described in a separate publication [12].

3. CURRENT DRIVE BY SURFACE WAVE MODES

We now wish to calculate the magnitude of the helicity injection effect given by Eq. (4) when surface wave modes near the ion cyclotron frequency are launched in the plasma. Let $x = 0$ be the plasma-vacuum interface with plasma occupying the $x > 0$ semi-infinite region. We take the magnetic field $\vec{B}$ to be in the interface (say along the $z$ direction) and assume that the propagation direction of a surface wave in the $y-z$ plane makes an arbitrary angle with $\vec{B}$. We now follow the standard derivation for surface 'shear' Alfvén modes [13] with $\omega \ll \Omega_i$ and $k_y \ll k_z$. Writing the wave equations in vacuum and the plasma, demonstrating that in the $\eta \to 0$ limit $E_{z1}$ plays a negligible role, solving for $E_{y1}$ and matching across the plasma vacuum interface we obtain the dispersion relation

$$\omega^2 = \frac{2k_y^2v_A^2}{1 - \frac{\omega}{\Omega_i}} \left(1 - \frac{\Omega_i}{\omega}ight)$$

and the ratio of electric fields determining wave polarization as

$$\frac{E_{y1}}{E_{z1}} = i \frac{(\omega A/\Omega_i) + \alpha k_z}{k^2 - A}$$

where $\alpha^2 = k^2 - a - \omega^2 A^2/\Omega_i^2 (k_z^2 - A)$, $A = -\omega_0^2 k_z^2 c^2 (\omega^2 - \Omega_i^2)$ and $\omega = \omega + i\mu k^2$. Note that $\mu k^2$ gives the 'viscous' damping of the mode and $\alpha$ is the inverse decay length in the plasma. Equation (6) demonstrates that the mode is generally elliptically polarized and thus has a 'mixed' nature (i.e has significant right and left handed circularly polarized components in it).

We now estimate the magnitude of the toroidal current driven by the surface helicity injection scheme described in this paper. We consider the excitation of surface wave modes near the ion cyclotron frequency. In the steady state problem, $\omega$ and $k_z$ are to be treated as real and determined by the antenna configuration. The
dispersion relation, Eq. (5), then determines real and imaginary parts of $k_y$ (and hence $\alpha$ and $\Delta$). Noting that $\mathbf{h}$ is in the $x$ direction, ignoring the weak effects due to the parallel electric field of the wave, we may write Eq. (4) as:

$$2 \int \eta (\mathbf{J} \cdot \mathbf{B}) \, d\tau \approx \frac{1}{q_n} \int \langle (J_{x1}B_{y1}^* + J_{y1}B_{x1}^*) \rangle \, dS \quad (7)$$

We now imagine that the parallel current driven at the plasma surface is also transported to the plasma core, so that the volume loss on the left side is over the whole plasma volume. Using expressions of $J_{x1}, B_{y1}$ in terms of $E_{x1}, E_{y1}$, etc., and approximating the various integrals by simple geometrical factors, we obtain:

$$J_t \approx 4 \left( \frac{R}{a} \right) \frac{q_k |E|^2}{\eta n_0 (\Omega_f^2 - \omega^2)} \left( \frac{\mu k^2}{\omega} \right) \quad (8)$$

Note that the driven current is proportional to the wave dissipation factor ($\mu k^2/\omega$) and to $k_z$. The power absorbed per unit volume may be determined from the expression $p = J_x^* E_x^* + J_y^* E_y^*$ + complex conjugate, which may similarly be approximated as

$$p \approx \frac{q^2 n_0 |E|^2 \omega}{M(\Omega_f^2 - \omega^2)} \left( \frac{\mu k^2}{\omega} \right) \quad (9)$$

We may now calculate the usual figure of merit for current drive schemes by using the appropriate geometrical factors ($I = I \pi a^2$ and $P = p(2\pi a)\Delta a(2\pi R)$, where $\Delta a$ is the width of the surface region where the wave is becoming dissipated), viz.

$$\frac{IR}{P} \approx \frac{R}{\pi \Delta a} \frac{k_z}{\omega \eta q n_0} \approx \left( \frac{c}{2 \eta \omega_{pi}} \right) \frac{1}{\eta B_0} \left( \frac{c}{\omega_{pi}} \right) \quad (10)$$

where we have approximately substituted for $\omega$ from the dispersion relation, Eq. (5). In dimensional units this gives

$$- \frac{I}{P} \approx \frac{T_e^{9/2}}{B_0 n_{13}^{1/2} \ln \Lambda} \left( \frac{1}{\Delta a_m} \right) (A/W) \quad (11)$$

where $T_e$ is the temperature measured in $eV$, $B_0$ is the toroidal field in gauss, $n_{13}$ is the density in units of $10^{13}$ cm$^{-3}$, $\ln \Lambda$ is the Coulomb logarithm and $\Delta a_m$ is the width of the surface wave excitation region in metres. For $T_e \sim 10^4$, $B_0 \sim 5 \times 10^4$, $n_{12} \sim 20$, $\Delta a_m \sim 0.1$, we obtain $I/P \sim 2$ A/W. The AC surface helicity injection scheme thus provides us with an efficient mechanism for toroidal current drive.
4. DISCUSSION

We have investigated the AC helicity injection at the surface of a plasma volume using wave fields near the ion cyclotron frequency. We demonstrate that finite ion inertia and non-resistive wave dissipation lead to a significant current drive which should be experimentally observable in contrast to the low frequency F-θ pumping experiments.

We may now compare the present scheme of helicity injection with that recently proposed by Chan et al. [10] and others. Investigating bulk helicity injection by pure modes, Chan et al. [10] have concluded that for a plasma periodic in the toroidal direction, a finite toroidal current can be steadily driven if and only if the plasma is warm and has anisotropic pressure. Thus the overall efficiency of this current drive scheme is determined by the temperature and is somewhat low. In a separate calculation, they have considered the current drive by mixed modes in a cold plasma. Higher efficiencies are reported in this case. This calculation has many common features with the calculation presented above and it will be useful to highlight the similarities and differences. Both calculations concern waves with mixed polarization, but whereas Chan et al. [10] mix the modes in an arbitrary manner, in our case mode mixing naturally occurs because we deal with surface waves in the plasma (see discussion after Eq. (6)). Both calculations demonstrate that for waves periodic in the toroidal direction (k_z real), only perpendicular damping can result in a net helicity injection; this is in contrast with our earlier report [9], where k_z was taken as complex. The most important difference between our calculation and that of Chan et al. [10] is that they obtain a finite helicity injection even without ion inertia, whereas we have demonstrated that the ion inertia effect is crucial. We believe that the result in Ref. [10] is in error in this respect and that the error arises because the authors have not properly accounted for the moving boundary terms at the plasma surface (although for all practical purposes they are working with surface modes which evanesce away from the surface). For comparison, we write here the approximate expressions obtained by Chan et al. for J_z and IR/P (in place of our equations (8) and (10)). They find J_z ~ (kv_i^2/2ηB_0cωc) |B|^2 and IR/P ~ 4π/ηB_0k. We note from these expressions that the current drive efficiency I/P for the calculation of Ref. [10] is independent of the density, whereas it decreases as n_0^{1/2} in our case.

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NON-RESONANT CURRENT DRIVE BY RF HELICITY INJECTION

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Abstract

NON-RESONANT CURRENT DRIVE BY RF HELICITY INJECTION.

Current drive via non-resonant interaction between RF waves and plasma is studied by using kinetic wave analysis and computation in a stationary state. The non-resonant interaction generates a force, in addition to the usual ponderomotive force. The force under consideration mainly acts as an internal force among plasma species; the net momentum input from the wave to the plasma is small in this process. The current driven by this process is not included in the conventional current drive scheme (i.e. resonant absorption of the parallel wave momentum or wave energy) and appears to be associated with the change in the RF wave helicity ($\vec{A} \cdot \vec{B}$). The analysis is applied to ICRF waves in large tokamaks. A one-dimensional wave analysis code (TASK/W1) shows that the conversion ratio of the RF wave helicity to the DC helicity becomes close to one in the case of hot plasmas. The current drive efficiency of RF helicity injection is not necessarily bounded by the conventional RF current drive efficiency. A preliminary study suggests that it scales strongly with electron temperature and toroidal magnetic field but only weakly with plasma density.

1. INTRODUCTION

One of the key issues in tokamak reactor design is current drive efficiency. A consistency analysis on ITER grade plasma has shown that the current drive efficiency is critical in determining circulating power, Q-value and divertor heat load [1]. A new scheme of current drive by RF wave helicity injection has been proposed [2]. If RF helicity is totally converted into DC helicity, i.e. the conversion ratio of $\lambda = \Delta \vec{E} \cdot \vec{B}_0 / (\vec{E} \cdot \vec{B})$ is unity, a high current drive efficiency is expected. ($\Delta \vec{E}$ is the reduction of DC electric field due to the current drive). Extensions have been made in MHD theory [3–5], the mechanism of wave helicity current drive [6] and the ICRF wave application [7]. However, the relation with the conventional momentum input schemes [8] had not been clear, so far.

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We investigate an internal force acting on a plasma, induced by RF waves. It is generated by non-resonant wave interaction in addition to the resonant force and the conservative ponderomotive force. This force is associated with a source or a sink of the RF wave helicity, \( \langle \hat{A} \cdot \hat{B} \rangle \). The case of minority heating by ICRF waves is studied numerically. The conversion ratio approaches unity in hot plasmas. It has become clear that the current drive by RF helicity injection is an independent scheme of the usual RF current drive and that the efficiency is not be bounded by the conventional limit.

2. NON-RESONANT INTERNAL FORCE

Wave momentum conservation in an inhomogeneous and dispersive medium has been discussed in the literature [9–11]. We write the force exerted \( \vec{F} \) by the wave, according to Ref. [11], as

\[
\vec{F} = n q \left[ \frac{k}{\omega} (\vec{E}^* \cdot \vec{M}_H \cdot \vec{E}) - i \frac{k}{2\omega} \frac{\partial}{\partial \vec{r}} \left( \vec{E}^* \cdot \frac{\partial \vec{M}_A}{\partial \vec{k}} \cdot \vec{E} \right) + i \frac{1}{2\omega} \frac{\partial}{\partial \vec{r}} (\vec{E}^* \cdot \vec{M}_A \cdot \vec{E}) - i \frac{1}{\omega} \frac{\partial}{\partial \vec{r}} \cdot (\vec{M}_A \cdot \vec{E} \vec{E}^*) \right]
\]

(1)

where \( n \) is the density, \( q \) the charge, \( \vec{k} \) the wavenumber, \( \omega/2\pi \) the wave frequency, \( \vec{M} \) the mobility tensor and subscripts \( H \) and \( A \) stand for the Hermitian and anti-Hermitian parts, respectively. The wave is assumed to be stationary, i.e. \( \omega \) is real. The first term is the force due to wave dissipation (in the direction of \( \vec{k} \)), the second is due to the wave dispersion (also in the direction of \( \vec{k} \)), the third is the ponderomotive force (in the direction of the gradient) and the last is the force in which we are interested here (in the direction of \( \vec{E} \)). This force is due to the non-resonant interaction. The ponderomotive force is written as a potential gradient and disappears by flux surface averaging, but this non-resonant force does not. We call this force internal (polarizing) force. In the case where a single wave of \( k_y = 0 \) is propagating in the \( x \)-direction (the \( z \)-axis is in the direction of the static magnetic field), the force in the \( z \)-direction, which generates a parallel current, can be written as

\[
F_z = F_z^R + F_z^{NR}
\]

\[
F_z^R = \frac{n q \vec{k}_z \vec{E}^*}{\omega} \cdot \vec{E}, \quad F_z^{NR} = - \frac{n q}{\omega} \left[ (k_x - k_x^*) \vec{v}_x^* + (k_z - k_z^*) \vec{v}_z^* \right] \vec{E}_z
\]

(2)

where \( R \) and \( NR \) denote resonant and non-resonant, respectively. We note that the sum of \( F_z^{NR} \) over the particle species corresponds to the change of the momentum of
FIG. 1. Waveform in the equatorial midplane. The magnetic axis corresponds to \( x = 0 \); the \( x \)-axis is in the direction of the major radius. The fast wave antenna is located on the low field side. Parabolic density and temperature profiles are used. The plasma parameters are chosen according to the JET plasma: \( n_e(0) = 10^{20} \text{ m}^{-3} \), \( n_e(a) = 10^{19} \text{ m}^{-3} \), \( n_H/n_e = 5\% \), \( T_e(0) = 10 \text{ keV} \), \( T_e(a) = 0.1 \text{ keV} \), \( T_e = T_D = T_H \), \( B(0) = 3.5 \text{ T} \), \( \omega/2\pi = 52.5 \text{ MHz} \), \( k_z = 8 \text{ m}^{-1} \) (\( \omega/k_z v_{Te0} = 0.96 \)), and total input power is 1 MW. The \( E_x \) component (a), the \( E_y \) component (b), the \( E_z \) component (c), and the power absorption profile with the Poynting flux (d) are shown.
FIG. 2. Forces and driven current for the case of Fig. 1. Resonant and non-resonant forces on electrons are compared in (a). Resonant forces and non-resonant forces on electrons, deuterons and protons are shown in (b) and (c), respectively. The driven current is shown in (d); the majority is driven by $F_R^e$. 
FIG. 3. Comparison with RF wave helicity for the case of Fig. 1. Helicity flux (a), and helicity source/sink term (b). The change of the DC helicity, $\Delta E \cdot B_0$, is compared to the RF helicity sink/source term in (c). The conversion ratio, $\gamma = \Delta E \cdot B_0 / E \cdot B_0$, is shown in (d).
the vacuum electromagnetic field and is usually small. This is the reason why we call this force as an internal force. The RF helicity $H = \langle \vec{A} \cdot \vec{B} \rangle$ satisfies the conservation relation, $\dot{H} = -\nabla \cdot \dot{Q}_{H} - 2(\vec{E} \cdot \vec{B})$, where the helicity flux $Q_{H}$ is defined as $Q_{H} = \langle \vec{A} \times \delta \vec{A}/\delta t + 2\phi \vec{B} \rangle$. The helicity source/sink term is given by $\langle \vec{E} \cdot \vec{B} \rangle$.

3. APPLICATION TO ICRF WAVES

We calculate the non-resonant force, the helicity conversion rate and the driven current. The one-dimensional slab model of a tokamak plasma is used, and the ICRF wave propagation is studied by using the TASK/W1 code [7, 12]. The $x$-axis is taken in the direction of the major radius, and $x = 0$, $x = a$, and $x = b$ correspond to the magnetic axis, the plasma surface and the chamber radius, respectively. The fast wave antenna is placed at $x = d$ on the low field side of the torus, carrying the oscillating current in the y-direction.

Figure 1 shows the wave structure in the case of the JET plasma. The parameters are: $B_{T} = 3.5$ T, $\omega/2\pi = 52.5$ MHz, $D(H)$ plasma, $n_e(0) = 10^{20}$ m$^{-3}$, $T(0) = 10$ keV. The wave is absorbed by electron Landau damping and ion cyclotron damping. Figure 2 illustrates the spatial profile of the forces: $F_{Rz}^R$ and $F_{Rz}^{NR}$ are shown in (a), $F_{z}^{R}$ in (b) and $F_{z}^{NR}$ in (c). The driven current is represented in (d). The helicity flux [7] $Q_{H}$ and the change rate of the helicity, $\langle \vec{E} \cdot \vec{B} \rangle$, are also shown in Figs 3(a) and (b). On balancing the $F_{Rz}^{NR}$ term with the parallel ion drag, the current induced by the non-resonant force, $f^{NR}$, and $AE$ are calculated as

$$J^{NR} = \frac{e}{m_{e}V_{ei}} F_{z}^{NR}, \quad AE = \frac{1}{ne} F_{z}^{NR}$$

Figure 3 also shows $\Delta \vec{E} \cdot \vec{B}_{0}$ and $\langle \vec{E} \cdot \vec{B} \rangle$ in (c) and the conversion ratio $\lambda$ in (d). The self-consistent calculation of wave structure and forces confirms that the non-resonant force appears to be associated with the helicity change. The helicity conversion coefficient is close to unity near the magnetic axis. The resonant interaction of electrons with fast waves also generates a current [13] as is shown in Fig. 2(d). The driven current is given by the sum of these two mechanisms.

4. SCALING STUDY

The scaling of the helicity current drive efficiency was derived previously [6] for an ICRF fast wave launched from the low field side as

$$\frac{1^{NR}}{P_{abs}} \approx 0.004 \frac{N_{e} T_{e}^{5/2}}{B_{T}^{2}}$$
FIG. 4. Helicity current drive efficiency versus (a) density $n_D$, (b) electron temperature $T_e$, (c) toroidal magnetic field $B_T$, and (d) toroidal wavenumber $k_z$. For all cases, $B_T = 5T$, $\omega = \omega_{ce0}$. $T_{e0} = T_{D0} = 10$ keV, $T_{H0} = 80$ keV, $n_H/n_D = 0.1$, and $k_z = 25$ m$^{-1}$, unless specified otherwise.

which is valid near the ion-ion hybrid layer. Using a 1-D wave propagation code [14] which treats Landau and cyclotron dampings in WKB approximations, we study the global helicity current drive efficiency by integrating over the plasma slab. The current is calculated from helicity balance, and typical density and temperature profiles are used. As an example, we choose a deuterium plasma with a hot proton minority species to achieve strong single absorption. Figure 4(a) shows $I_{NR}/P_{abs}$ versus $n_D$. As is expected, the efficiency is insensitive to the density within this range. We have also plotted $\omega_{ce0}^2/c^2k_z^2$ for the same densities. At lower densities, its value drops below unity bringing the magnetosonic cut-off close to the plasma centre, and the wave becomes largely evanescent. At higher densities, the mode conversion layer widens and a full wave treatment is needed to calculate absorption and current...
accurately. Figure 4(b) shows the efficiency as a function of $T_e$. It scales as $T_e^{5/2}$ over a broad range of temperatures, in good agreement with Eq. (4). Note that $\omega_{pe}^2/c^2k_z^2$ is constant for this scan; hence, the propagation condition is held fixed. In Fig. 4(c), the efficiency is computed with $B_T$ being varied. Since $N_z = c k_z/\omega$ and $\omega$ is varied accordingly to keep the cyclotron resonance location fixed, Eq. (4) predicts a scaling of $B_T^3$, which fits the computed result reasonably well, with the latter showing a slightly weaker dependence. In Fig. 4(d), we vary $k_z$, with $\omega_{pe}^2/c^2k_z^2$ held constant to avoid a change of the evanescent layer. The efficiency indicates only a weak dependence on $k_z$, which disagrees with Eq. (4). Further study is needed to understand this result.

5. SUMMARY AND DISCUSSION

In this paper, we have identified an RF driven current by non-resonant interaction. The total current is the sum of two contributions: (1) the current driven by absorption of RF wave momentum through resonant interaction, and (2) a new current driven by a non-resonant internal (polarizing) force. The latter current appears to be associated with the source/sink of RF wave helicity. The case of ICRF waves is studied. The conversion ratio $\lambda$ is calculated and found to be close to unity in large and hot plasmas. The analytic prediction of the efficiency [2, 6] based on the assumption that $\lambda = 1$ has been compared with a 1-D calculation. The strong scaling with $T_e$ and $B_T$ and insensitivity to $n_D$ are in agreement, whereas linear scaling with $k_z$ has not been obtained.

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