Research using small tokamaks

Proceedings of a Technical Committee Meeting
held in Arlington, Virginia, USA, 27–28 September 1990
FOREWORD

The technical reports in this document on "Research Using Small Tokamaks" were presented at the International Atomic Energy Agency Technical Committee Meeting, with the same title, held in Washington, D.C., September 27-28, 1990. The meeting was attended by 47 participants from 17 countries.

One might well ask, "What is a Small Tokamak?" There may be no precise answer to this question, but one working definition is that a small tokamak is one with a plasma current less than 400 kA. Operational programmes of small tokamaks should contribute to research "targets of opportunity" aimed at

- improving the operation of standard (large) tokamaks;
- moving, through small steps, to more advanced magnetic confinement systems;
- investigating transport;
- developing diagnostics;
- training scientists, engineers, and technicians, particularly in developing countries.

Given this list of goals, one may also suppose that research using machines and facilities other than tokamaks might contribute as well. Indeed, the history of the series of TCM's on Research Using Small Tokamaks has shown that, while most of the research has indeed been conducted on tokamaks, some significant contributions to the meetings have come from non-tokamak-based research.

Researchers from programmes of various size in many countries are eager to participate in work directed toward one or more of the goals listed above. Some of the work is documented in this latest TECDOC on Research Using Small Tokamaks.

The programme of the Technical Committee Meeting was divided into three sessions: Plasma Modes, Control, and Internal Phenomena (nine papers), chaired by Dr. S. Luckhardt (US); Edge Phenomena (ten papers; one not submitted for publication), chaired by Dr. P. K. Kaw (India); and Advanced Configurations and New Facilities (four papers), chaired by Dr. J. Fujita (Japan). The chairmen of each of the sessions have prepared brief summaries that are printed at the beginning of the sections containing papers from their sessions.

This TECDOC is prepared from direct reproductions of the authors' copies. It is hoped that publication of the document in this way will provide timely information to fusion specialists in all member states that helps define the main directions of fusion research using small tokamaks and other small research facilities.
EDITORIAL NOTE

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PLASMA MODES, CONTROL, AND INTERNAL PHENOMENA

(Session 1)

Chairman

S. LUCKHARDT
United States of America
A number of groups reported experiments aimed at controlling low m and n modes. Feedback control using external coils was reported by AEA Fusion (Culham Laboratory) on the DITE and Compass-C tokamaks. Lower-hybrid current drive was successfully used to stabilize sawtooth oscillations and the m,n=1, 1 continuous mode on DITE, and both unmodulated and fast modulated ECRH mode stabilization experiments were reported. The characteristics of low m,n modes and magnetic island growth, as they relate to disruptions and transport, were discussed in the papers of A. Vannucci et al. (TBR-1 tokamak) and H.Y. Peng (IPP, Hefei, China). Static magnetic perturbations imposed by external helical windings were discussed by J. de Villiers et al. (Tokoloshe tokamak, AECSA) who find mode locking and sawtooth suppression. Detailed studies of the m=2 tearing mode were reported by J.Y. Chen et al. (TEXT) who measured island widths and mode frequencies and compared them to a neoclassical poloidal flow model. Pellet injection experiments on JIPP T-IIU were reported by Y. Ogawa, et al., in which an ultra fast cooling of the plasma center is observed during pellet injection. The JIPP T-IIU group attributes this cooling to an enhanced thermal conduction process.

Results of advanced diagnostics development were reported by F. Aumayr et al. and E. Unterreiter et al. using neutral lithium beam activated charge exchange spectroscopy on TEXTOR. Absolute impurity ion densities in the outer region of the plasma were successfully measured, and advances in atomic physics code modeling and cross section measurements were also reported. M.J. Ballico et al. (University of Sydney) reported studies of discrete Alfvén wave modes on TORTUS, in which highly localized resonant modes were excited and detected with laser scattering; q-profile information was inferred from measurements at three excitation frequencies.
THE USE OF ECRH, LHCD AND RESONANT HELICAL FIELDS TO INFLUENCE TOKAMAK INSTABILITIES

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Abstract

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 stabilisation and stimulation of disruptions by external field perturbations.

Introduction

The experiments described in this paper were performed on the DITE and COMPASS tokamaks at Culham over the last 18 months. DITE began operation in 1977, and its life was extended by 1 year in 1989 to perform a magnetic feedback experiment under contract to JET, and to conduct Lower Hybrid and Electron Cyclotron Current Drive experiments. Its main parameters for this series of experiments were: \(R_{\text{vessel}} = 1.17\) m, \(a_{\text{vessel}} = 0.3\) m, \(R_p = 1.19\) m, \(a_p = 0.23\) m, \(50\text{kA} \leq I_p \leq 200\text{kA}\), \(B_\phi \leq 2.8\text{T}\), \(P_{\text{ECRH}} \leq 600\text{kW}\), \(P_{\text{LH}} \leq 250\text{kW}\). ECRH is provided by three 200kW, 60GHz gyrotrons with high-field-side launch antennae in X-mode. Lower Hybrid power is provided by a 1.3GHz klystron from IPP Garching, with a 4-guide grill from the Petula-B tokamak, Grenoble.

COMPASS is a new device commissioned in 1989. There are two phases planned, first using a circular vacuum vessel (COMPASS-C) and then changing to a D-shaped vessel in 1991 (COMPASS-D). The parameters are given in Table I. Both devices have the same poloidal field coil set allowing many plasma shapes. Experiments on X-point plasmas have started on COMPASS-C with an inboard X-point [1], but these will not be discussed here. A major feature of the device is the complex array of external saddle coils, allowing many different resonant magnetic perturbations (RMPs) to be applied to well controlled shaped or circular tokamaks with strong auxiliary heating.
Table I: Parameters of COMPASS-C and COMPASS-D

<table>
<thead>
<tr>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>$I_p$</td>
<td>≤ 200kA</td>
<td>≤ 400kA</td>
</tr>
<tr>
<td>$B_\phi$</td>
<td>≤ 1.7T</td>
<td>≤ 2.1T</td>
</tr>
<tr>
<td>$R$</td>
<td>0.557m</td>
<td>0.557m</td>
</tr>
<tr>
<td>$a_{wall}$</td>
<td>0.22m</td>
<td>0.22m</td>
</tr>
<tr>
<td>$a_{limiter}$</td>
<td>0.196m</td>
<td>≤ 0.2m</td>
</tr>
<tr>
<td>Limiter</td>
<td>Poloidal ring</td>
<td>Inner belt + outer moveable</td>
</tr>
<tr>
<td>Material</td>
<td>Graphite</td>
<td>Graphite or composite</td>
</tr>
<tr>
<td>$b/a$</td>
<td>≤ 1</td>
<td>≤ 2</td>
</tr>
<tr>
<td>$A = R/a$</td>
<td>2.8</td>
<td>2.8</td>
</tr>
<tr>
<td>$t_{pulse}$</td>
<td>≤ 500ms</td>
<td>≤ 5s</td>
</tr>
<tr>
<td>$P_{ECHR}(60GHz)$</td>
<td>1MW</td>
<td>2MW</td>
</tr>
<tr>
<td>$P_{ECHR}(28GHz)$</td>
<td>200kW</td>
<td>200kW</td>
</tr>
<tr>
<td>$P_{LHCD}$</td>
<td>-</td>
<td>0.5MW</td>
</tr>
<tr>
<td>$f_{m=2,n=1}$</td>
<td>≤ 15kHz</td>
<td>-</td>
</tr>
<tr>
<td>$dI_p/dt$ (disruption)</td>
<td>≤ 2 x 10^8As⁻¹</td>
<td>-</td>
</tr>
<tr>
<td>$\tau_{sawtooth}$</td>
<td>~ 0.5 - 1ms</td>
<td>-</td>
</tr>
<tr>
<td>$I_{sawtooth}$</td>
<td>≤ 2kA</td>
<td>≤ 5kA</td>
</tr>
</tbody>
</table>

The Effect of Resonant Magnetic Perturbations on MHD activity

Fast magnetic feedback experiments have been performed on DITE [2, 3, 4] using driven saddle coils inside the vessel, with sensitive Mirnov coils to provide an $m = 2, n = 1$ feedback signal, and the the gain and phase of the feedback loop controlled by fast hybrid analogue-digital electronics. The saddle coils were driven by two 60V, 5kA 15kHz transistor amplifiers. This technique has allowed extension of the density limit by up to ~ 25% [4], and perhaps more importantly, disruption-free operation at or somewhat above the usual density limit (Fig. 1). The basic properties of the feedback system seem to be reasonably well modelled by a modification of the Rutherford theory of non-linear evolution of tearing modes wherein the boundary conditions are changed to include the applied $m = 2$ field. When a disruption does occur, it is caused by a sudden increase in the $m = 2$ amplitude, triggered at a sawtooth crash (i.e. not by the associated heatpulse) [5], which exceeds the control range of the feedback system. This sawtooth effect 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 stabilised with the fast $m = 2$ feedback system (Figure 2). A factor~ 8 reduction in $B_\phi$ and also a drop in $I_{sawtooth}$. 

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Figure 1: Prevention of density limit disruption on DITE with fast $m = 2$ magnetic feedback. $q(a) \sim 2.4$.

are achieved with the feedback, with an associated improvement in central confinement [6] ($I_{srr}(0)$ rises by $\sim 10\%$), indicating that such feedback systems have a wider application than just disruption avoidance. The $m = 1$ stabilisation is thought to be due to the global nature of this mode: the stability can be affected by changing $\Delta'$ at $q = 2$ [7]. In OH discharges, the $m = 1$ sawtooth precursor oscillation is at a lower frequency than the continuous $m = 2$ mode ($f_{m=1} \sim 4\text{kHz}, f_{m=2} \sim 10\text{kHz}$), and is not affected by the feedback. Experiments to measure dynamic mode-locking and the phase instability are described in [4].
Figure 2: Stabilisation by $m = 2, n = 1$ magnetic feedback of the $m = 1, m = 2$ mode produced by LHCD on DITE. The loop gain and phase are changed at 370ms.
On COMPASS-C \( (R = 0.557\text{m}, \ a_{\text{trim}} = 0.196\text{m}, \ I_p \leq 200\text{kA}, \ B_\phi \leq 1.75\text{T}) \) the emphasis has been on quasi-static perturbations using a complex set of external saddle coils [8]. In each quadrant of the tokamak there are 10 toroidal bars, 70\(^\circ\) long, with poloidal links. These are routed to linkboards allowing considerable flexibility in the choice of interconnections and hence the RMP spectrum. Feeder errors have been minimised, as have errors in the main poloidal field system - error fields have been reduced to below 0.5G \((n = 1)\) for each PF coil by \textit{in situ} magnetic alignment [9]. Most configurations have sideband levels below 10\% (calculated over a uniform \(\theta, \phi\) grid). Table II shows the range of experiments performed and summarises 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. RMPs tested so far also have little effect on the behaviour of runaway electrons.

Table II: Effect of RMP on COMPASS-C. \( I_p \simeq 100\text{kA} \) for all these OH discharges, with \( B_\phi \) being varied to change \( q(a) \). “stab” indicates reduction of the background \( m = 2 \) activity, “stim” indicates stimulated disruptions. When there is no observable effect on discharge evolution and the \( m = 1, m = 2 \) modes and sawtooth for the levels of \( I_{\text{saddle}} \) used, “nil” is entered. \( \tilde{B}_\theta \) is the rms value at \( r = 22\text{cm} \) for the unperturbed plasma. \( B_r(r = 20\text{cm}) \sim 10\text{Gauss}, \sim 1\%B_\theta(a) \) for the dominant applied perturbation: \( I_{\text{saddle}} = 500\text{A}-1\text{kA} \) (single conductor).

<table>
<thead>
<tr>
<th>Discharge type</th>
<th>RMP configuration</th>
<th>Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>( f_{m=2} ) kHz</td>
<td>( \tilde{B}_\theta(m = 2) ) Gauss</td>
</tr>
<tr>
<td>( q \simeq 2 )</td>
<td>1,1</td>
<td>stab,stim</td>
</tr>
<tr>
<td>( q \simeq 2 )</td>
<td>2,1</td>
<td>stab,stim</td>
</tr>
<tr>
<td>( 2 \leq q \leq 3 )</td>
<td>3,1</td>
<td>stab,stim</td>
</tr>
<tr>
<td>( 2 \leq q \leq 3 )</td>
<td>3,2</td>
<td>nil</td>
</tr>
<tr>
<td>( 3 \leq q \leq 4 )</td>
<td>3,1</td>
<td>weak stab</td>
</tr>
<tr>
<td>( 3 \leq q \leq 4 )</td>
<td>3,2</td>
<td>( \tau_{E,p} ) ( /n_e(max) )</td>
</tr>
</tbody>
</table>

Apart from the \( q(a) \sim 2 \) sequence, these experiments were mostly performed before boronisation 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 [10, 11]. Little driven reconnection is predicted when \( f_{\text{applied}} \) is more than a few percent 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 stabilising effect due to the applied RMP slowing the mode from $f_{\text{natural}}$, until $I_{\text{saddle}}$ exceeds a threshold $I_{s,\text{crit}}$ when $f_{\text{natural}} \rightarrow f_{\text{applied}}$ and full driven reconnection then occurs.

The “nil” entries are thus interpreted in terms of minimal driven tearing and accordingly little stabilisation 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, $m = 2$ stabilisation is seen at low density for $q(a) \leq 3$ (Fig. 3), usually accompanied by a reduction of $1 \sim 3\text{kHz}$ in the mode frequency ($\sim 12\text{kHz}$) 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 stabilising term is smaller. It is possible, however, to obtain a 15% enhancement in the density limit at intermediate $q(a)$ [8] - there is no strong stabilisation 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 due to the greater distance between the saddle coils and the resonant rational-$q$ surface.

---

![Graph](image)

**Figure 3:** Stabilisation of $m = 2$ instability by stationary 2,1 RMP on COMPASS-C. A variety of $I_{\text{saddle}}$ waveforms shows penetration to be fast. $I_p = 100\text{kA}$, $q_\alpha(a) \sim 3$.

At high values of $I_{\text{saddle}}$ disruptions are stimulated for 2,1 RMPs (Figure 4), as usual in such experiments, with $I_{\text{saddle}} \geq I_{s,\text{crit}}(\text{theory})$. A drop in central confinement is also seen just before disruption [11], consistent with substantial tearing. These disruptions do not have conventional rotating precursors, but the final explosive growth ($\tau \sim 50\mu\text{s}$) 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 ($I_{2,1} \approx 900A$, $I_{3,1} \approx 700A$) [8]. Stimulated and low $q(a)$ disruptions are generally faster than density limit disruptions on COMPASS-C ($dI_p/dt \lesssim 2 \times 10^8A/s$ for all types of disruptions).

Figure 4: Stimulated disruption by a stationary 2,1 RMP on COMPASS-C. Note the fall in central SXR emission shortly before disruption suggesting strong driven reconnection. $q_0(a) \simeq 3$. 

![Image of stimulated disruption](image-url)
Turning now to the other helicities used, 1,1 and 3,2, there are two main results from the experiments before boronisation. Firstly, there is no apparent effect on the sawtooth period, amplitude and precursor frequency, or on the \( m = 2 \) characteristics; and secondly some confinement degradation (reduction in \( d\bar{n}_e/dt \) for constant fuelling, 10 ~ 20\% drop in \( I_{\text{xtr}}(r \lesssim 10\text{cm}) \)) is seen with both RMPs. These two together, with observed changes in \( H_a \) signals, suggest that 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. After the vessel was boronised, initial results appear different with the (3,2) RMP increasing the disruptive density limit for \( 2 \leq q(a) \leq 3 \). This is accompanied by substantial changes in the profile evolution, when compared with discharges without the RMP. The 2,1 and sawtooth precursor frequencies are not much changed by the RMP, however, and there is no definite evidence of full driven reconnection occuring. Thus there are beneficial effects from these RMPs, as seen elsewhere, some of which appear to be related to wall conditions rather than MHD effects at the principle resonant \( q \)-surface.

![Figure 5: Suppression of sawteeth on COMPASS-C with ECRH for different resonance positions.](image)

\( I_p = 70\text{kA}, \bar{n}_e = 1.2 \times 10^{19}\text{m}^{-3}, q_{\text{cyt}}(a) \approx 3.5. \)
A final area where extensive use of RMPs was made was in an intensive campaign to attain \( q(a) < 2 \) in COMPASS-C which has a thin wall, \( (d = 0.7\text{mm st. steel}) \). A wide range of techniques have been used including: (i) programmed \( dI_p/dt \) (up to 10MA/s, 3MA/s in flat top); (ii) up to 600kW ECRH (variously localised); (iii) high \( dq/dt \) from position programming \((q^{-1}dq/dt \text{ up to } 100s^{-1})\); (iv) \( (2,1), (3,1), (3,2) \) RMPs during the \( q \) ramp; (v) connection of the saddle coils to form a passive lattice shell, or an \( m = 2, n = 1 \) shell; (vi) strong gas puffing during the \( q \) ramp and many combinations of these techniques. None of these procedures has so far reduced \( q(a) \) (determined from a filament-type reconstruction code) significantly below 2.0, but some techniques allow lower \( q \) than others. Reducing \( a_p \) by rapid outward motion produces \( q = 1.95 \) transiently. The various “conducting shells” do not influence the behaviour, the induced current in the passive conductors at \( r \sim 0.26\text{m} \) is \( \sim 100\text{A} \), which can be compared with \( I_{saddle} \sim 1\text{kA} \) (applied) necessary to stimulate disruptions. The oscillatory disruption precursor is removed by the RMP. As at higher \( q(a) \) the RMP has most effect at low density.

**Instability control by RF injection**

LHCD has been used on DITE [6] to suppress the sawtooth instability for \( P_{LHCD} \geq 80 \) – 100kW \((P_{OH} \sim 100 \text{ – } 200\text{kW before injection})\). 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 \( \sim 5\text{ms} \) 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 \( r_{inv} \) immediately after LHCD ends, due to strong Parail-Pogutse instabilities at that time. Increasing \( P_{LHCD} \) to above \( 160 \sim 200\text{kW} \) removes the \( m = 1 \) activity (as on other devices), as does applying \( m = 2 \) magnetic feedback. No effect on the sawtooth is seen with pure LH heating.

On COMPASS localised 60GHz ECRH [8] at the second harmonic, launched from low field side X-mode antennae \((85^\circ \text{ to } B)\), with several tangency radii has been used to remove the sawtooth for up to 60ms (the ECRH pulse length used). The results are summarised in Table III. Suppression occurs in \( \leq 1 \) sawtooth periods \((r_{ST} \sim 0.4\text{ms here})\) and may last for the whole ECRH pulse. After the ECRH pulse the sawteeth often appear with smaller \( r_{inv} \), but \( r_{inv} \rightarrow r_{inv, before \text{ ECRH}}(\sim 4\text{cm}) \) in a few periods. Suppression is obtained for a range of resonance positions \((-8\text{cm} \leq r_{res} \leq +5\text{cm})\) but is most reliable for \( r_{res} \sim +5\text{cm} \). A high frequency internal \( m = 1 \) mode at \( \sim 25 \text{ – } 30\text{kHz} \) can appear inside \( r \sim 6\text{cm} \) in addition to the sawtooth. Enhanced \( m = 2, n = 1 \) activity \((f = 15 \text{ – } 17\text{kHz})\) 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 destabilises the \( m = 2 \) mode may be consistent with this data, but will not explain the rapid sawtooth removal at the onset of ECRH.
Table III: MHD effects of $2\omega_{ce}$ ECRH on COMPASS-C (pre-boronisation). $I_p = 100\, \text{kA}$, $\bar{n}_e \sim 1.3 \times 10^{19}\, \text{m}^{-3}$, and $P_{ECRH} = 300 - 450\, \text{kW}$

<table>
<thead>
<tr>
<th>$B_\phi$</th>
<th>$r_{res}$ (cm)</th>
<th>Sawtooth removal</th>
<th>$m = 2$ mode $\sim 16, \text{kHz}$</th>
<th>$m = 1$ mode $25-30, \text{kHz}$</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>

A brief series of experiments have been performed on DITE using ECRH modulated at $f_{m=2}$ to attempt to control steady $m = 2$ oscillations by heating the magnetic islands [12, 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 localised 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 ($\sim 1 - 2\, \text{ms}$) stabilising effect of the mean ECRH power seen for $|r_{res} - r_{e=2}| \lesssim 2.5\, \text{cm}$ ($I_p \simeq 85\, \text{kA}$, $q_{oy}(a) \simeq 2.5$). The results (e.g. Figure 6) show that there is a definite sensitivity of the $m = 2$ amplitude to the phase in the feedback loop, when DC stabilisation is seen, and stabilisation/déstabilisation effects comparable to the DC effects are observed, with the optimum apparently being when the O-point is heated, as predicted by theory [12, 13].

**Conclusions**

A variety of MHD stabilisation methods have been attempted on the DITE and COMPASS tokamaks and sustainment of discharges above the usual density limit with magnetic feedback demonstrated. A range of effects from mode stabilisation to stimulated disruptions have 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 LHCD on DITE, and ECRH on COMPASS-C. 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 stabilisation 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.
Figure 6: Influence of fast feedback with modulated ECRH on the $m = 2$ mode level on DITE. $I_p = 80\text{kA}, r_{res} = -18\text{cm}, q_{eq}(a) = 2.5, n_e = 1.7 \times 10^{19}\text{m}^{-3}$. The separate trace shows the dependence of mode amplitude on the phase in the loop for the two phase sweeps in the shot. The phase is changed by $22.5^\circ$ every $2\text{ms}$. The optimum appears to be for heating near the O-point.
ACKNOWLEDGEMENTS

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EXTERNAL HELICAL COIL STUDIES
ON TOKOLOSHE TOKAMAK

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Abstract

Tokoloshe is a low aspect ratio tokamak \((R/a = 0.52m/0.24m)\) with a maximum possible toroidal field of 1.2 T. Due to the low aspect ratio, peak plasma currents of \(I_p \leq 150\) kA can be obtained at \(B^* \simeq 0.6\) T. The discharge duration is typically \(< 40\) ms when it is terminated by a major disruption due to strong gas puffing. Line average densities generally reach peak values of around \(6 \times 10^{13}\) \(cm^{-3}\). Central electron temperatures range from 550 eV early in the discharge to about 350 eV just before the disruption with a typical \(Z_{eff} \simeq 2\).

The machine is fitted with 4 sets of windings for detailed studies of the stability and confinement of tokamak plasmas. The work originally started with a pure poloidal winding \((\ell = m/n = 0/1)\) within the limiter [1]. Pulses yielding a maximum axial perturbation of the toroidal field of ±8.5% were possible. Subsequently, 3 sets of stellarator type helical windings \((\ell = 1, 2/1, 3)\) were fitted on the outside of the vacuum chamber. These conformed to a general winding law of \(\phi = m/n(\theta - \delta_{m/n}\sin \theta)\) (with \(\delta_{m/n} = 40, 48\) and \(0\) degrees respectively for \(\ell = m/n = 1/2, 1, 3/1)\) except for local deviations around ports and major protrusions [fig. 1]. The \(\delta\)-values were chosen initially [2] to minimize contributions to \(m/n = 3/1)\) islands. Subsequent comparison of the calculated mode spectra with those derived from vacuum field measurements confirmed the design expectations. In practice these 3 sets of windings can be driven individually or in pairs with peak current pulses of approximately 9, 6 and 5 kA respectively. The pulse rise times are about 1 ms and pulse widths are adjustable from a few ms to 10 ms or more with a falling amplitude due to resistive decay.

Results

\(\ell = 0\) Limiter coil: The observed interactions of the limiter coil with the plasma were largely confined to the edge region though under some conditions global effects were observed [1]. With the perturbing field direction parallel to \(B^*\), the edge density increased a few times while impurity radiation increased orders of magnitude with little or no other effect. In the antiparallel case, these effects were generally not seen but under slightly different initial conditions application of pulses invariably led to major disruptions.

\(\ell = 1\) Helical Coil: Effects of the \(\ell = 1\) coil have been explored for coil currents up to about 6% of \(I_p\) [3,4]. The direction of \(I_p\) was chosen for either resonant or non-resonant helicity with the field at the \(q = 1\) surface, with coil currents also in either direction. The most striking observed effect is the suppression of the relative sawtooth amplitude on a fast time scale \((\leq 1/5\tau_R)\) to the \(q = 1\) surface) with no change in the period at all. The amplitude suppression is linearly dependent on the
applied current and is up to 50% at the highest values irrespective of the direction of $I_p$ (fig. 2a). In contrast to this non-dependence on resonant helicity, investigations showed detailed differences in the sawtooth amplitude behaviour with the coil current direction reversed. Two diodes viewing the plasma tangentially at $r = 4.5$ cm above and below the plasma centre but well within the radius of inversion ($r_1 = 7.5$ cm) both show amplitude suppression but different detailed structure. Depending on the coil current direction, either the bottom or the top diode shows a marked hollow in the rising phase of the sawtooth lasting about 60% of the period followed by an $(m = 1)$-like spike just before the crash (fig. 2b).

A limited number of discharges where studied at a lower toroidal field to allow a higher relative coil current ($I_{c=1}/I_p \sim 8\%$). Here a decrease in the $m = 2$ mode amplitude was generally evident with an accompanying pulse lengthening (fig. 3) but detailed investigations are still to be done.

**$\ell = 2$ Helical Coil:** Sawtooth suppression with the $\ell = 2$ coil can be induced at much lower currents and faster time scales than with the $\ell = 1$ coil ($\sim 50\%$ reduction at $2$ kA). Effects on the $m = 2$ mode include the following. During that part of the discharge where a relatively high mode amplitude is present ($B_0 (2/1)/B_0 > 1\%$), application of a perturbing $\ell = 2$ field invariably caused mode slowing down and locking for varying periods depending on the current pulse amplitude and duration [5,6,7]. The poloidal phase of the locking is determined by the current direction and the locked phase ends with re-rotation having the same amplitude and velocity (fig. 4a). Higher currents produced mode locking followed by induced minor or major disruptions. During that part of the discharge where the $m = 2$ mode amplitude is an order of magnitude lower than mentioned above, the application of resonant helical perturbations lead to major disruptions when the current is larger than a critical threshold value ($\sim 2$ kA) [7]. The time to disruption decreased monotonically with higher coil currents (fig. 4b), where the dotted lines are theoretical fits based on the characteristic penetration times to the $q = 2$ surface. Detailed mode structure analyses of coil signals show $(m = 2)$-like modes locked in positions depending on the coil current direction (fig. 5). In the non-resonant case, no disruptions could be induced.
Figure 2: (a) Sawtooth suppression for resonant and non-resonant configuration and a positive coil current direction. (b) Details of top/bottom (T/B, see fig. 1) asymmetry in the soft X-ray emission for different $I_e=1$ directions.

Figure 3: Pulse lengthening observed at lower $B_\phi$ (.35 T).
Figure 4: (a) The locked mode duration is dependent on the coil current and the particular mode amplitude present. (b) Coil induced disruption characteristics for the $\ell = 2$ coil switched on at $t = 30$ ms where there is low Mirnov activity.

Figure 5: Polar plots of $\tilde{B}_p$ as a function of time until just before disruption for a resonant configuration and both coil current directions. The values in Gauss correspond to the radius of the circle in each case.
A series of experiments on mode locking were done with a high speed cine camera viewing the whole minor cross-section of the limiter. Pictures of a reference shot clearly show rather uniform poloidal $H_a$ intensity during high Mirnov activity and the expansion of the $H_a$ region during a thermal collapse which ultimately leads to a major disruption. When the mode is locked by the $\ell = 2$ coil, cine pictures show localised stationary bright regions with an ($m = 3$)-like mode structure which persist to the disruption and often become more intense just prior to the disruption. These regions show a change of position, intensity and poloidal extent when the helical coil current direction is reversed.

$\ell = 3$ Helical Coil: Sawtooth suppression occurs already at very low coil currents. The observed relative amplitude is almost zero at coil currents as low as 1 kA. Disruptions are also precipitated as is the case with the $\ell = 2$ coil with current threshold values about half as much but the time to disruption characteristic corresponds to the penetration time to the $q = 2$ surface. In the case of $\ell = 3$ however, the non-resonant configuration also precipitates disruptions but at 3 to 4 times higher current than in the resonant configuration [7].

Conclusions

The readily observable phenomena with helical coil fields are the reduction of sawtooth amplitudes, mode locking, precipitation of disruptions and to a limited degree delay of disruptions. To a greater or lesser extent, each helical coil can be used to induce these phenomena due to the low aspect ratio of Tokoloshe. The $\ell = 1$ has the only clearly resonant effect at the $q = 1$ surface, but a rather large coil current is required. The observations compared to the results of field line tracing calculations indicate that $q_0$ may be less than 0 during sawtothing on Tokoloshe. The ability to cause locked mode disruptions, particularly with the $\ell = 2$ and 3 coils, can be used to study the influence of mode orientation and wall stabilisation on the disruption mechanism.

References

Abstract

Observations are presented of discrete Alfven waves in the TORTUS tokamak. The results are interpreted in terms of the measured density profiles and an inferred central q.

1. INTRODUCTION

In principle, it is possible to combine observations of the discrete Alfven wave (DAW) spectrum with electron density profile measurements to obtain information on the q profile in a tokamak. Using very simple apparatus, several authors have already measured the central q by means of this technique [1-3]. In this paper, we present a number of results regarding DAW modes in TORTUS, obtained from measurements of the edge wave magnetic fields and laser scattering. These results were obtained at relatively low RF power levels, about 10kW, sufficient to obtain good scattering signals.

In the cold plasma, zero electron mass approximation, the dispersion relation for the shear Alfven wave is given very closely by

\[ \omega^2 = k_{||}^2 v_A^2 \left( 1 - \omega^2/\omega_{ci}^2 \right) \]  

where \( v_A \) is the Alfven speed and \( \omega_{ci} \) is the ion cyclotron frequency. Since very little wave energy in the shear wave propagates across B, this wave does not normally form a cavity mode in an inhomogenous plasma. Instead, shear wave disturbances tend to propagate along field lines at the local Alfven speed, thereby destroying any phase coherence on neighbouring field lines. However, if \( k_{||} v_A \) is essentially constant over a certain region of the plasma cross section, then it can be seen from Eq. (1) that a phase coherent or DAW mode will exist in this region. In cylindrical geometry, with periodic boundary conditions in the z direction \( (k_z = n/R) \), Eq. (1) has the form

\[ \omega^2 = \frac{1}{R^2} \left( n + \frac{m}{q(r)} \right)^2 \left( 1 - \frac{\omega^2}{\omega_{ci}^2} \right) - \frac{B^2}{\mu_0 n_i \mu_i(r)} \]  

where \( q = \frac{r B_z/(RB_0)}{n_i} \) is the safety factor and \( n_i \) is the ion density. This relation can be used to determine the location of an Alfven resonance layer \( (k_{||} = \infty) \) in a cylindrical, inhomogeneous plasma. Suppose that \( n_i = n_o(1 - r^2/a^2) \) where \( a \) is the wall radius. If all quantities in Eq. (2) except for \( n_o \) are held constant, then it is easy to show that (a) there is no resonance surface in the plasma at low \( n_o \), say \( n_o < n_A \) and (b) for \( n_o > n_A \), the resonance surface generally moves from the centre of the plasma towards the plasma edge as \( n_o \) increases. Of special interest is the form of Eq. (2) for \( n_o \) slightly less than \( n_A \), typically
1-5% smaller. Then there is no \( k_\perp = \infty \) layer in the plasma but Eq. (2) describes a shear wave of moderate to high \( k_\perp \). Provided that \( n \) has the same sign as \( m/q \), and provided that \( q(r) \) and \( v_A(r) \) increase with \( r \), then it is possible for \( k_\parallel u_A \) to remain relatively constant over a significant fraction of the plasma cross section. Numerical solutions of the cold plasma equations [4] show that significant wave fields at the plasma edge are obtained only for those profiles where \( n < 2m/q \) near the \( r = 0 \) axis. The wave fields extend further towards the plasma edge for \( m < 0 \) modes than for \( m > 0 \) modes, due to the favourable (left hand) wave field polarisation of \( m < 0 \) modes.

2. EXPERIMENTAL ARRANGEMENT

The results presented below were obtained in the TORTUS tokamak, which has a rectangular cross section, stainless steel vacuum vessel of dimensions \( 34 \text{cm} \times 26 \text{cm} \) and a major radius \( R = 44 \text{cm} \). All results were obtained with hydrogen as the filling gas, a toroidal magnetic field \( B_\phi = 1.2 \text{ Tesla} \) (at \( R = 0.44 \text{m} \)), and for a plasma current of amplitude \( 35 \text{kA} \) and duration \( 20 \text{ms} \). The electron density was ramped to an on-axis peak \( n_0 = 3.7 \times 10^{19} \text{ m}^{-3} \) and the density was then allowed to fall towards zero as shown in Fig. 1 in order to observe DAW modes under both rising and falling density conditions.

![Figure 1](image-url)

**Figure 1**: (a) Line average electron density \( <n_e> \) vs. time. (b)-(d) \( b_\theta \) waveforms at 3.2, 4.5 and 6.0 MHz for different antenna phasings, \( \Delta \phi \), as labelled on the vertical axis.

Six identical electrostatically shielded antennas were installed at the plasma edge. Two antennas were connected in series and mounted as a top–bottom pair at each of three toroidal locations, each pair being phased to launch odd \( m \) modes. The three antenna pairs were connected to three independent amplifiers via lumped element matching networks and were
all driven at the same frequency, 3.2 MHz, 4.5 MHz or 6.0 MHz. The three antenna pairs could be phased to couple to any desired toroidal \( n \) mode. In practice, the DAW modes observed were sufficiently pure and sparse that only one antenna pair was necessary to couple to a mode, but enhanced coupling was obtained using either two or all three pairs. Alternatively, the phasing could be adjusted to produce total extinction of a DAW mode in order to (a) identify the \( n \) number of the DAW mode and its sign and (b) to help identify other wave modes launched simultaneously with DAW modes.

Wave fields in the plasma edge were monitored with one toroidal and one poloidal array of magnetic probes. The toroidal array consisted of three identical probes inserted at three different toroidal locations, mounted vertically through the top wall of the vacuum vessel and inserted to the same \((r, \theta)\) coordinates. The poloidal array consisted of eight identical coils wound on a flexible tube and spaced at poloidal intervals of 10°. The array was housed in an 8mm OD quartz tube surrounding the plasma at minor radius \( r = 11 \text{ cm}, 0.5\text{ cm} \) outside the limiter radius. By relocating the array between successive discharges, it was possible to construct a complete 360° poloidal profile of the poloidal \((b_{\theta})\) component of the wave magnetic field at \( r = 11\text{ cm} \). All coil signals were fed to hybrid combiners, to eliminate the electrostatic component of the wave field, and to an array of mixers to monitor the phase and amplitude of the wave fields with respect to the antenna current.

The laser scattering system includes a formic acid laser operating at 433\( \mu \text{m} \), focused to a beam waist of 8mm at the centre of the plasma. The beam was scanned mechanically across the minor radius of the plasma on a shot to shot basis.

3. EXPERIMENTAL RESULTS

The most obvious indication of the presence of a DAW mode is provided from measurements of the wave fields at the plasma edge. Typical results, obtained at 3.2, 4.5 and 6.0 MHz, are shown in Fig. 1. At 3.2 and 4.5 MHz, two sharp peaks in the \( b_{\theta} \) waveforms are observed during the discharge. The first peak occurs during the density rise, and the second peak is simply a repeat of the first as the density falls. Both peaks occur at the same density and are observed simultaneously at all toroidal and poloidal locations.

The results shown in Fig. 1 were obtained by exciting antenna pairs 1 and 3 and by varying the phase difference, \( \Delta \phi \), between the currents in the two antennas. Antenna pairs 1 and 3 are separated toroidally by \( 3\pi/4 \) radians. An \( n = -1 \) mode is excited with maximum amplitude when \( \Delta \phi = -3\pi/4 \) or when \( \Delta \phi = 2\pi - 3\pi/4 = 5\pi/4 \). The \( n = -1 \) mode is extinguished if \( \Delta \phi = \pi/4 \). The results in Fig. 1 show a clear DAW mode extinction when \( \Delta \phi = \pi/4 \). Consequently, it is easy to identify the DAW mode at 3.2 MHz as an \( n = -1 \) mode (and not an \( n = +1 \) mode). Data from the poloidal probe shows clearly that \( m = -1 \) for the DAW modes at 3.2 MHz and 4.5 MHz.

The DAW mode observed at 4.5 MHz is extinguished when \( \Delta \phi = 3\pi/2 \), which identifies it as an \( n = -2 \) mode. We expected to observe an \( n = -3 \) mode at the highest operating frequency, 6.0 MHz. In fact, there is no evidence of any DAW mode at 6.0 MHz for any antenna phasing. The 6.0 MHz waveforms are different at different toroidal and poloidal
locations. We interpret this result as strong local or direct excitation of the shear wave at the plasma edge [5], rather than the excitation of a global mode.

The DAW modes in Fig. 1 are superimposed on a strong background $b_\theta$ wave field which does not disappear under any antenna phasing conditions. The background field cannot therefore be identified in terms of any particular $n$ mode, or even any dominant $n$ mode. The background field is due mainly to direct excitation of the shear wave in the low density edge plasma. This wave propagates as a narrow beam in both toroidal directions along helical field lines passing through or near each antenna, but the beam paths are intercepted by the limiters or by the antennas.

Laser scattering results are shown in Fig. 2. The scattered signal ($A_s$) rises to a maximum immediately after the appearance of a DAW mode, a result which suggests enhanced scattering from an Alfvén resonance layer in the plasma. However, there is no clear evidence of such a layer in the radial profiles of the scattered signal. Indeed, there is very little difference in the radial structure either before, during or after the appearance of a DAW mode, implying that the transverse wavelength of the shear wave remains comparable to the plasma minor radius at all times.

![Figure 2](image-url)

*Figure 2: (a) Waveforms of $n_e$, $b_\theta$ and scattering amplitude, $A_s$, vs. time, at 3.2 MHz. (b) Radial profile of $A_s$ at $t=7.0$ms. (c) Radial profile of $A_s$ at $t=9.0$ms.*
4. DISCUSSION

The absence of DAW modes at 6.0 MHz can be explained qualitatively in terms of equation (2). At 6.0 MHz, and over the observed density range, only the \( n = -3 \) mode was expected. However, if the density profile is approximately parabolic (as observed) then a relatively steep \( J_Z \) profile is needed to give a significant radial variation in \( n + m/q \) in order to maintain \( k_B V_A \) relatively constant. The \( J_Z \) profile need not be as steep when \( n = -1 \) or \( -2 \) since \( m/q \) is then significant compared to \( n \). The above data can be interpreted quantitatively in terms of the cold plasma equations [4], although a more appropriate analysis will require the inclusion of kinetic effects. Assuming a \( J_Z \) profile of the form \( J_Z = J_0(1 - r^2/a^2)^p \) where \( p \) is a variable parameter, we find that (a) for \( p \lesssim 1.5 \), there are no DAW modes with significant wavefields at the plasma edge for any of the conditions studied, (b) for \( p \gtrsim 2.5 \), strong DAW modes should exist for all of the conditions studied, and (c) for \( 1.8 \lesssim p \lesssim 2.2 \), low frequency DAW modes have significantly larger edge wave fields than high frequency DAW modes. Furthermore, the effective ion mass, when \( p = 2.0 \pm 0.2 \) is \( 1.1 \pm 0.1 \), when we compare the observed and predicted values of \( \langle n_e \rangle \) at which DAW modes appear. Our preliminary analysis, based on the cold plasma model, therefore indicates that \( q(0) \approx q(a)/3 \approx 1.5 \) for the above plasma conditions.

REFERENCES

MODE - COUPLING TRIGGERING MECHANISM IN TOKAMAK MAJOR DISRUPTIONS

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Abstract

The possible triggering mechanisms for the major disruptions in the small TBR-1 tokamak discharges and, latter on, in the JET high current pulses were studied. Evidences could be found, in both cases, that the coupling of the m/n=1/1 with other higher modes (but predominantly the m/n=2/1) would be the main responsible for triggering this instability.

I- Introduction

The major disruptions in tokamaks are characterized by an abrupt loss of confinement, leading to the completely destruction of the plasma column and limiting the tokamak operating conditions.

Although the intensive efforts already made the fundamental mechanisms that lead to a disruption is not yet well understood. Investigations carried out in different tokamaks have resulted in divergent explanations and, up to now, there is not a unique model which is able to describe all the experimental observations. Nevertheless it is generally accepted that radiative processes\(^1\)\(^2\) and MHD perturbations (mainly involving the m/n=2/1 mode either by directly contacting the limiter\(^3\) or by coupling with other MHD modes like m/n=1/1\(^4\) or m/n=3/2\(^5\), for example) may play important roles at some stage of the disruptive triggering process.

The major disruptions in the small TBR-1 tokamak discharges\(^6\) and, latter on, in the JET high current pulses (Ip\(^\text{max} \geq 5\text{MA})\(^7\) were investigated and the results will now be discussed.

II- Major disruptions in TBR-1 discharges

The TBR-1 tokamak is a small device located at the Physics Institute of the University of São Paulo-Brazil which started operation in 1980. Its main parameters are: R\(_0\)=0.30m, a=0.08m, B\(_\phi\)
0.4T, $T_e(o)\sim 200$ eV, 7KA $\leq I_p \leq 12$KA and discharge duration ranging up to 7ms.

The investigations of the disruptive phenomena was carried out using basically two diagnostic systems; one which detected the soft x-ray emitted by the plasma using six surface-barrier detectors, each one viewing a different chord of the plasma column section through a polypropylene covered slot, and the other which measured the perturbed poloidal magnetic field using a set of twenty magnetic pick-up coils, sixteen equally spaced in the poloidal direction and four in the toroidal direction.

The analysis of the experimental data showed that all the fluctuations identified in the magnetic coils signals exhibited the wave number component $n=1$. This implies that the explanation proposed in ref. 5 would not be adequate to describe the disruptive phenomena in TBR-1, although in TOSCA tokamak perturbations corresponding to the MHD modes $m/n=3/2$ and $m/n=5/3$ were effectively detected just before a major disruption event\(^{8}\). In TBR-1 it was very common to identify a dominant $m/n=2/1$ component increasing in amplitude just before the negative spike in the loop voltage signal. The perturbed poloidal field ratio would then reach $B_0/B_\theta \sim 4\%$. Estimations done, over the location of the $q=2$ magnetic surface and the size that a corresponding magnetic island would have, suggested no possible contact of island with the limiter.

On the other hand, the coupling between $m/n=2/1$ and $m/n=1/1$ MHD modes could be a realistic triggering mechanism for the major disruptions. This derives from the fact that these two components, identified in the magnetic and soft x-ray signals were observed to have their frequencies changed and finally assumed the same values right before the negative spike in the loop voltage (Fig.1),

![FIG.1 - Precursors $m/n=1/1$ (a) and $m/n=2/1$ (b) modes growing in amplitude and assuming the same frequency just before the negative spike in the loop voltage signal (c) is an indication of a mode-coupling in TBR-1.](image-url)
indicating an interaction between them. In others pulses, like the one shown in fig. 2, it can be verified that again the m/n=2/1 and m/n=1/1 fluctuations assume the same frequency before the confinement is lost, but now there is no amplitude growth of the m/n=1/1 mode. This indicates that the growing m/n=2/1 is the perturbation that, reaching certain amplitude, would couple with the m/n=1/1 mode, triggering a major disruptions.

![Diagram](image)

**FIG.2** - The m/n=1/1 (b) and m/n=2/1 (c) mode coupling is again identified in TBR-1 before the negative spike in loop voltage (a). However only the m/n=2/1 is now observed to grow in amplitude.

### III- Major Disruptions in JET High Current Pulses

More recently the major disruptions triggering mechanism in JET was also investigated for pulses that have achieved plasma current values above or approximatly equal to 5 MA, at some stage of the discharge.

In fig. 3, for example, a 7.1 MA pulse which disrupted at decay phase when I_p<5.2 MA is shown. Although there was some RF additional heating for 1.5 s it did not contribute to any substantial change in the radiated power which remained bellow 35% of the input power level.

The indication that a mode coupling would be the responsible triggering mechanism for this disruption comes from observing the sawteeth formation being strongly disturbed (fig. 3c) at the same time a mode-locking process initiates (fig. 3d). Since the m/n=1/1 perturbation is believed to be the basic cause for the sawteeth occurence then this mode may be considered to take part in the mechanism that triggers the disruption. The contribution of other MHD perturbations can be investigated by analysing the components of the
FIG. 3 - Sawteeth formation being disturbed (c) at the same time a mode-lock started (d) and an increase of the MHD activity is identified in the rectified magnetic signals before the negative spike in the loop voltage (g), is an indication that a major disruption was triggered in JET (at t ~ 12.15s, for this pulse).
perturbed magnetic field measured by the pick-up coils. In fig. 3e the corresponding rectified signals are presented and although they do indicate a stronger MHD activity mainly related to the n=2 perturbation, in this particular pulse, for the majority of the analysed cases the dominant component was the n=1.

A more convincing example of the interaction between MHD modes as the triggering mechanism for the major disruption is presented in fig. 4. The m/n=2/1 and m/n=1/1 perturbations detected in the magnetic coils (Fig. 4a) and in the soft x-ray (Fig. 4c) signals, respectively, were observed to grow in amplitude and reaching a maximum at the exact instant the mode-locking started (Fig. 4b), leading the plasma to a disruption approximately 35 ms latter.

![Figure 4](image)

**FIG.4** - The amplitude of the m/n=2/1 (a) and m/n=1/1 (c) modes growing until a mode-lock starts is a strong indication that a mode-coupling mechanism would trigger this disruption in JET.

Finally, it was interesting to note that all the disruptive pulses were invariably preceded by a mode-locking process, although the mode-lock itself did not cause the major disruption. Unlockings were sometimes observed and only later on, when a new locking took place, the plasma did disrupt.
Conclusion

Similar triggering mechanism for the major disruption in the small TBR-1 tokamak and in JET high current pulses could be found. The \( m/n=1/1 \) MHD mode would play an important role in the process by coupling with other higher modes but mainly the \( m/n=2/1 \). The coupling was identified, in TBR-1 discharges, when both the \( m/n=1/1 \) oscillations in the soft x-ray signals and the \( m/n=2/1 \) fluctuations in the magnetic coils signals changed their frequencies and finally assumed identical value just before the loop voltage spike. At JET the high current pulses were found to be preceded by a deformation in the sawteeth activity at the same time a mode-lock process started. In some cases the amplitude of the \( m/n=1/1 \) and \( m/n=2/1 \) precursors were observed to increase right before the plasma collapse and the same time a mode-locking started.

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THE m=2 TEARING MODE IN TEXT*

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Abstract

In TEXT an m=2 island is seen to grow linearly in time for a wide range of island widths before it saturates. This is consistent with the Rutherford regime and implies a weak dependence of \( \Delta' \) on the island width. Both experimental and numerical results show \( q(0) \) is an important parameter in the mode evolution.

The frequency of the mode changes significantly during the Rutherford growth. The scaling of frequency versus island width and machine size agree with a neoclassical poloidal rotation model. The enhanced particle transport observed is consistent with edge magnetic island transport produced by the Ergodic Magnetic Limiter but an order higher than that predicted by stochastic transport mechanisms.

Introduction

By carefully controlled operation of TEXT we can reproducibly produce magnetic fluctuations as high as 2% of the poloidal field without disruption. The tearing mode in these "high MHD" discharges is stable often persisting until the end of the shot. Almost all diagnostic signals are modulated at the mode rotation frequency.

The mode structure analysis from 2 toroidal and 24 poloidal Mirnov coils shows a dominant m=2, n=1 component. This is confirmed by soft x-ray and electron density tomography. The large-amplitude single tearing-mode structure is the key feature of the high MHD shots.

Expanding to order \( (W / a) \) in a cylindrical model, the island width \( W \) of the single mode will be\(^{(1)}\): \( W = 4 \sqrt{b_r r_s q / B_\theta m} \frac{dq}{dr} \bigg|_{r=r_s} \), where \( \hat{b}_r \bigg|_{q=2} \) can be determined from the measured \( \hat{b}_\theta \big|_{Edge} \)\(^{(2)}\). A model q profile is used which agrees with \( q(r) \) derived from Spitzer resistivity.

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Rutherford regime and tearing mode global features

It is pointed out by Rutherford\cite{3} that if the magnetic island is wide enough, the island growth will transition from exponential to linear. For \( m > 1 \) modes the island width \( W \) is governed by the equation:

\[
\frac{dW}{dt} = \eta \Delta'(W) + \alpha W
\]  

(1)

Fig. 1a. shows examples of island width evolution during mode growth in TEXT. Before the mode saturates there is a wide range (2-10 cm) in which the island width grows quite linearly. This agrees with Rutherford’s prediction. In this region \( \Delta' \) is weakly dependent on \( W(t) \), and the rates of the island growth are roughly proportional to the saturation level. The value of the rates typically is about 10 m/s.

Fig. 1b. presents numerical simulations of the island width evolution of a single \( m/n=2/1 \) mode using a fully nonlinear 2-D reduced MHD code. The curves a, b and c correspond to initial \( q(0) = 1.2, 1.1, \) and 0.93. The linearity and saturation is similar to Fig. 1a and the rates are of the same order as experimentally observed.

The similarity between Figs. 1a and 1b emphasizes the importance of \( q(0) \) to the overall dynamics of the tearing modes. Fig. 1b shows that MHD (tearing mode) activity is favored by \( q(0) > 1.0 \). This is consistent with the experimental observation in TEXT that sawtooth activity, favored by \( q(0) < 1 \), is never observed simultaneously with large Mirnov signals.

Fig. 1a. Linear growth and saturation of \( m=2 \) magnetic island in Rutherford regime.

Fig. 1b. Simulation of \( m=2, n=1 \) island evolution in high MHD discharge with \( q(a)=3.3 \) and various \( q(0) \).

This strong anti-correlation even persists in auxiliary heating experiments. Fig. 2 shows ECRH applied near the plasma center on a high MHD shot (no sawteeth). The ECRH power is turned on at 300ms. As the central temperature (measured by ECE) rises, the Mirnov signal falls, disappearing concurrently with the appearance of strong sawteeth on the SBD and ECE signals.
Scaling law of mode frequency

Fig. 3 shows mode frequency versus island width in the Rutherford regime from 79 shots. The frequency is inversely proportional to the island width. Near the high frequency limit, MHD activity is very low indicating a narrow island. The low frequency limit corresponds to the nonlinear mode saturation with a large island or rotation slowing due to externally applied stochastic field by Ergodic Magnetic Limiter.

A neoclassical poloidal rotation model is developed to explain this scaling. Consistent with \( E_x \) data from the HIBP[4] we account for the Doppler effect of plasma rotation with an additional term in the tearing mode eigenfrequency[5]:

\[
    f = \alpha f_{\text{Tearing}} + f_D = \alpha f_e + \frac{E_r k_\theta}{B_{\phi 0}},
\]

where \( \alpha \) is a model dependent factor of order unity.

In steady state the radial ion momentum balance equation[6] is

\[
    e_i n_i (E_r + V_\theta B_\phi) - \frac{\partial P_i}{\partial r} = 0.
\]

For ohmic heating usually \( V_\phi << V_\theta \). Using the neoclassical poloidal flow[7] and combining (1) and (2) we have

\[
    f = \left( \alpha + \frac{T_i}{T_e} [1 + (1 + g)] n_i \right) f_e^* \tag{3}
\]

with \( f_e^* = \frac{k_B k_\theta T_e}{2\pi e B_{\phi 0} L_n} \), \( g \) is a collisionality-related factor.

From (3) and Fig. 3, we expect a proportional relation between \( L_n \) near \( r_{q=2} \) and \( W \). This can be identified in Fig. 5a. Also, the normalized frequency \( (B_{\phi 0}/T_e) f \) would be proportional to the factor \( (k_\theta/L_n) \), i.e. roughly \( (1/a^2) \) machine to machine. This is consistent with Fig. 4 in which a linear dependence of \( (B_{\phi 0}/T_e) f \) on \( 1/a^2 \) is shown with the maximum frequencies obtainable for each machine.
Enhanced transport by perturbed magnetic fields

In the presence of the magnetic islands the $n_e$ and $T_e$ profiles are broadened by a central value decrease and $T_i(0)$ is reduced by 150eV (Fig.5b). If the island is wider than $\sim 5$ cm the local density profile near $q=2$ surface is flattened around the island O-point (Fig.5a).

Using source data from H$\alpha$ detectors and a 1D transport code, we deduce that the particle transport coefficient increases by a factor of two to three in the island region. This is consistent with the effect of islands produced externally by the EML but is an order of magnitude higher than that expected from stochastic transport\cite{8}.

Conclusions

1) Before mode saturation a Rutherford regime is observed.
2) $q(0)$ has a strong effect on the mode growth.
3) Mode frequency scaling is consistent with neoclassical model for poloidal rotation velocity.
4) Local particle transport is enhanced by a factor of 2-3 for large island width.
REFERENCES

ENERGY CONFINEMENT AND COUPLING OF TRANSPORT PROCESS WITH MHD MODES

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Abstract

Starting from Tearing Mode equation, the growth rate equations of magnetic islands, the energy and particle balance equations with large effective thermocondutivity and diffusion in magnetic islands, it has been shown that the width of magnetic islands of different tearing modes and the temperature profile all oscillate with time. In some cases the oscillation frequency is the same as of the sawteeth, but in other cases it is in the range of observed Mirnov oscillations. From these equations we have calculated the energy confinement time of the plasma, and compared with the confinement time purely due to the anomalous thermoconductivity in $q > 1$ region. The scaling laws of these two kinds of confinement time are quite different. This result shows clearly the inpermission to compare the experimental results directly with anomalous transport analysis, as usually done before.

Energy confinement of the Tokamak Plasma has been widely studied since Sixties on small as well as big machines. In Seventies, even there were still lot of problems, the physical picture of energy loss in Tokamak Plasma seemed to be clear. It has been supposed that, the energy loss in central core ($q < 1$) of the plasma is mainly controlled by Sawteeth Process; in low temperature edge region, it is determined by impurity radiation and anomalous thermo-conductivity due to some kind of instability; and in outer region of the plasma ($1 < q < 2$ or $3$), the energy loss process seems to be anomalous transport due to some drift instability for electron, and to be neo-classical with some anomalies for ion. From this picture, a lot of numerical codes have been developed, and used to determine the transport coefficients by comparing with the results of experimental measurements. People generally thought that, according to
the behavior of these obtained transport coefficients, the mechanism of anomalous transport could be determined.

But since the end of the Seventies, the more experimental observations on different kinds of devices were got, the more confusive the energy confinement picture was. At the present time, there are already more than ten scaling laws for energy confinement time, which give quite different extrapolated behavior. The significant fact is that, more and more observations have shown the very close relation between energy confinement and macroscopic plasma motion. Such discorvaris, including: Profile Consistency, L \to H mode transition, Helical field experiments, \ldots, have made many scientists to suggest that, the energy loss in plasma outer region could be due to the activities of tearing modes or chaotic magnetic configuration. But most of the scientists have still insisted in the local anomalous transport processes in the plasma outer region due to the experiments of heat pulse propagation, even the prevailed picture could not take into account the forementioned phenomena connected with the macro-processes.

In Institute of Plasma Physics, Academia Sinica (IPPAS), we have believed that it is most important to have a picture which could uniformly take into account all the main experimental results from different devices, as different aspects of the plasma behavior or at different conditions. In this picture the anomalous transport should be coupled with MHD processes to determined the energy loss, especially in the outer region of the plasma. Such problem has been studied experimentally and theoretically in our Lab.
On our small Tokamak HT-6B (a = 12.5 cm, R = 45 cm, B_t = 1 T), there are L=2/n=1 and L=3/n=1 helical winding. These helical fields have good configurations and all the other harmonic components are less than 15% of the main component. In the course to study the mechanism of suppression of m=2 or m=3 modes by weak helical fields, we have found very strong coupling between different m modes. In some cases L=2 and L=3 fields all could suppress not only m=2 and m=3 but also m=1 modes. An unexpected fact has been observed that, the weak helical fields, which could only change the magnetic configurations near q=2 or q=3 surfaces respectively, can amplify the Sawteeth significantly, but change the sawteeth period a little. These experiments have shown the improvement of energy confinement by helical field.

The successive measurements of the thermo-conductivity by propagation of Sawteeth wave form and impurity radiation also show the ability of weak helical fields to improve the energy confinement (Fig. 1 and Tab. 1). The temperature profile was broadened by WHF.

FIG. 1. From top to bottom: line emission signals of O(III) (at 1.5 and 8.7 cm), C(III) (at 1.5 and 8.7 cm) and Hα (at 1.5 and 8.7 cm), and RHF current (inverted).
TABLE I. REDUCTION OF $x_c$ ($m^2 \cdot s^{-1}$) BY RHF

<table>
<thead>
<tr>
<th>Shot 11 176</th>
<th>$r = 0$ cm</th>
<th>$r = 7.6$ cm</th>
</tr>
</thead>
<tbody>
<tr>
<td>Without RHF</td>
<td>3.38 ± 0.49</td>
<td>7.76 ± 0.15</td>
</tr>
<tr>
<td>With RHF</td>
<td>2.37 ± 0.38</td>
<td>4.30 ± 0.30</td>
</tr>
</tbody>
</table>

It should be mentioned again that, the only way, by which the WHF could directly influence the plasma, is to change the magnetic configuration near the corresponding resonant surface. Such changes of local magnetic structure could only influence the MHD processes and it has really influenced (Fig. 2). Thus, these experiments very clearly show the close relation between MHD process and energy loss process.

FIG. 2.
It is well known that the instability and the growth rate of tearing modes are mainly determined by the plasma current gradient at the corresponding resonant surfaces. Inside the magnetic islands, energy or particle transport processes are very fast due to the convection. In the discussions related with the profile consistency, people have found some temperature profiles or current profiles, which could stabilize all the tearing modes. Because there is no magnetic island, these profiles are unstable due to the transport processes. In decaying process of these profiles, some MHD modes become more and more unstable and the corresponding magnetic islands will grow. These growing modes will take part to change temperature profile cooperating with transport process, and gradually to make themself stable. In this process, some other modes could become unstable and new magnetic islands appear, which will change the profile further. The formentioned consideration makes the picture of strong coupling between MHD modes and anomalous transport process clear and natural. It should be considered if anyone want to analyze the Tokamak Plasma behavior.

To see what results we could get from the coupling of MHD modes with transport process, let us neglect the toroidal effect starting from cylindrical geometry. The following model has been used:
1. Transport Processes

\[ \frac{3 N_e}{\partial r} = -\frac{1}{r} \frac{\partial}{\partial r} (r \Gamma_e) \] .... Diffusion: \( \Gamma_e = -D \frac{\partial N_e}{\partial r} \)

\[ + S \] .... Source: \( S = N_e N_o <\sigma_{\text{ion}} N_e> \)

\[ \frac{3}{2} \frac{\partial (N_e T_e)}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} r (K_e \frac{\partial T_e}{\partial r} - \frac{3}{2} T_e \Gamma_e) \] .... Thermoconductivity and Diffusion

\[ + \frac{nc^2}{(4\pi r)^2} \left( \frac{\partial (r B_0)}{\partial r} \right)^2 \] .... Ohmic heating

\[ - \frac{3m_e N_e}{m_i \epsilon_i} (T_e - T_i) \] .... e + i Thermoexchange

\[ - W_R \] ... Radiation Loss: \( W_R = W_{BR} \cdot A \)

\[ - W_I \] ... Ionization Loss: \( W_I = S \epsilon_i \)

\[ \frac{3}{2} \frac{\partial (N_i T_i)}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} r (K_i \frac{\partial T_i}{\partial r} - \frac{3}{2} T_i \Gamma_i) \] .... Thermoconductivity and Diffusion

\[ + \frac{3m_e N_e}{m_i \epsilon_i} (T_e - T_i) \] .... e + i Thermoexchange

\[ - W_{ch} \] ... Charge Exchange Loss

\[ W_{ch} = <\sigma_{ex} V_i>, \frac{3}{2} N_i T_i N_o \]

\[ \frac{\partial}{\partial t} B_0 = \frac{c^2}{4\pi} \frac{\partial}{\partial r} r \left( \frac{n}{r} \cdot \frac{\partial (r B_0)}{\partial r} \right) \] .... Magnetic Diffusion

Where \( N_o \): Neutral Hydrogen Density; \( <\sigma_{\text{ion}} N_e> \): Ionization Rate;

\( W_{BR} \): Bremsstrahlung Power Density; \( \epsilon_i \): Ionization Potential;

\( A = (1 + \frac{37.9 Z^2}{T_e} \cdot \frac{m_e}{A} + \frac{850 Z^4}{T_e^2} \cdot \tau_{ei} \) \( \tau_{ei} \): e - i Heat Exchange Time,

and \( N_i = N_e, \Gamma_e = \Gamma_i \) (Neglecting Impurity)
2. Transport Coefficients

\( \eta: \) Spitzer Resistivity with impurity effect in \( Z_{\text{eff}} \).

The effective resistivity (due to \( m = 1 \) internal mode);

\[
\eta = c_0 (1 - q_0) \left( 1 - \frac{3}{4} \right) \left( \frac{r}{r_1} \right)^2 \quad \text{for} \quad r^2 < \frac{3}{4} r_1^2
\]

where \( r_1 \) is the radius of \( q = 1 \) surface.

\( K_e: \) Pseudo-classical Electron Thermoconductivity, it has been phenomenologically enhanced by

\[
k = k_0 \left[ 1 + c \left( 1 - \left( \frac{x_s}{h} \right)^2 \right) \right] \quad \text{for} \quad x_s \leq h
\]

in magnetic islands with width \( h \), \( c \approx 100 \);

in \( q < 1 \) region \( k = k_0 \left[ 1 + c \left( 1 - 0.75 \left( \frac{r}{r_1} \right)^2 \right) \right] \left( r^2 \leq \frac{4}{3} r_1^2 \right) \)

\( K_i: \) Neo-classical Ion Thermoconductivity, it has been anomalous enhanced in magnetic islands and \( q \leq 1 \) region.

\( D: \) Diffusion Coefficient, Alcator scaling and has been enhanced in magnetic islands and \( q \leq 1 \) region.

To use these transport coefficients (Neo-classical Ion Thermoconductivity, pseudo-classical Electron Thermoconductivity etc.), it is only to begin more smoothly the concrete calculation.

3. Tearing Mode Equation:

The equation of perturbed radium magnetic field \( B_r \) of \( m/n \) mode:

\[
\frac{d}{dr} \left( r \frac{d}{dr} \left( r B_r \right) \right) - \left( m^2 + \frac{d j / dr}{(B_0/r^2)} (1 - nq/m) \right) B_r = 0
\]

and

\[
j = \frac{1}{\mu_0} \frac{1}{r^2} \frac{\partial}{\partial r} \left( r B_0 \right);
\]

The evolution equation of the width of magnetic islands:

\[
\frac{d}{dt} w = 1.66 \frac{\psi(r_s) \Delta(w)}{\mu_0 \psi_s \frac{r_s}{w}},
\]

Where

\[
\Delta(w) = \frac{\psi'(r_s + \frac{w}{2}) - \psi'(r_s - \frac{w}{2})}{\psi_s / r_s},
\]

\( r_s \): the radius of resonant surfaces, \( W \) the width of magnetic island,

\( \Delta \): stability factor.
The following plasma parameters were used for concrete numerical calculation: \( R = 45 \text{ cm}, \ a = 12 \text{ cm}, \ B_z = 7 \text{ kG}, \ q_\phi = 2.8 - 5.0, \)
\( Ne = (1-3) \times 10^{18} \text{ cm}^{-3} \). Only three MHD modes (2/1, 3/2 and 4/3 or 3/1) have been taken into account. Because all the equations could be scaled by the length \( L \), the following results also could be used qualitatively to discuss the processes in big machine.

(IV)

The plasma behavior of different discharge parameters have been calculated.

For high density, low \( q \) discharge (\( Ne = 3 \times 10^{18} \text{ cm}^{-3}, \ q_\phi = 3.2 \)), Fig 3.1 shows the time-evolution of the positions and amplitudes of the 2/1, 3/2 and 4/3 modes. After 1 ms from the beginning, the 2/1 magnetic islands keep the width (\( w \)) near 20% of the minor radius, but oscillate within the amplitude of 0.15\( w \). It is the real oscillation of the magnetic island (i.e., expansion and shrinkage) due to the strong coupling of transport process with MHD modes, and its frequency is the same as of sawteeth. Such oscillation was observed in \( m = 2 \) and \( m = 3 \) Mirnov's signals but has been thought as due to the propagation of sawteeth. The behaviors of 3/2 and 4/3 tearing modes are different from that of 2/1 mode: the magnetic islands appear only in part of the sawteeth period, in other words the magnetic islands appear and annihilate alternately. When they are maximum, all islands of three modes are overlapped, and make energy very quickly to transfer through a wide region. Fig. 3.2 shows the time variation of the energy confinement time, the deep valley in each sawteeth period indicates such overlapping of magnetic islands. The movement of \( q = 1 \) surface in Fig. 3.1 does not represent the sawteeth oscillation in our
FIG. 3.1. The time evolution of 2/1, 3/2, 4/3 modes.

FIG. 3.2. The time variation of energy confinement time $\tau_{ge}$.

FIG. 3.3. The time variation of temperature at different positions ($q_a = 3.2$, $N_e = 3 \times 10^{13}$ cm$^{-3}$).

model. In Fig 3.3 the time variation of the temperature shows the typical sawteeth behavior, and the phase shifts between sawteeth waves at different places could determine a formal thermo-conductivity.

Fig 4 has been obtained for low density and low q discharge (Ne = $1.8 \times 10^{15}$ cm$^{-3}$, qa = 3.4). Only 2/1 and 3/2 modes have been considered. The sawteeth type oscillation of the magnetic islands is very obvious. The only difference with Fig 3 is that the 2/1 magnetic islands appear and
annihilate alternately within each sawteeth period. The lack of lasting 2/1 magnetic islands, of course, should make energy confinement better than the case of Fig. 3.

For high density high $q$ case ($N_e = 3 \times 10^{13} \text{ cm}^{-3}$, $q_a = 4.1$), the 2/1 mode grows up and then saturates to stable magnetic islands with width of near 0.1 $a$ (Fig. 5). The behavior of 3/2 mode is quite different. Firstly it is stable for near 1.5 ms, and a high frequency oscillation appears lasting for two ms, then it turns to be stable. The oscillation frequency of 3/2 mode is in the frequency range of Mirnov's oscillation observed on small Tokamak, but the amplitude is very small. There is no oscillation in time variation of temperature. It shows the possibility that, the Mirnov oscillation observed in magnetic measurement could be due to the coupling of tearing modes with transport process.

All the forementioned results show a general trend, which has been confirmed by many experiments on small devices, that the sawteeth oscillation will disappear and $m=1$ mode will arise gradually when the $q$ value or electron density increases.
It is valuable try to analyze the energy confinement in the major part of the plasma (q < 2 or 3) starting from the coupling of MHD modes with transport process.

At the present time, it is main task to show the necessity of introducing this new picture, or the significant differences between these two pictures, thus our forementioned simple model could be used.

The energy confinement time has been defined as:

\[ \tau_{Je} = \frac{W}{(P_{in} - dW/dt)} \]

Fig. 6.1 shows the energy confinement times, obtained by considering the coupling of tearing modes with pseudo-classical transport processes, for different \( N_e \) and \( q_a \). Fig. 6.2 is the usual pseudo-classical results, i.e. considering only the pseudo-classical transport in outer region, for reference. For fixed \( q_a \) value the pseudo-classical confinement time
increases with electron density as usually obtained but the MHD effect totally changes the scaling, in Fig 6.1 the confinement time decreases with electron density increased For fixed $N_e$ these two dependences of $t_{ge}$ on $q_a$ are also totally different Fig 7 shows the sensitive dependence of $t_{ge}$ on the temperature of the plasma boundary in our new picture Such sensibility is due to taking part of MHD modes The MHD processes coupled with anomalous transport not only change the value of global confinement time but, perhaps more important also change the scaling law

It should be noticed that, the popular way to analyze transport process has been combining the usual transport calculation with experimental measurements to get the local thermo-conductivity or diffusion coefficients, which will be further compared with the results of
turbulence theory of different micro-instabilities. But due to the strong effects of MHD processes, as mentioned before, these transport coefficients only have formal meaning and could be quite different from the real ones. They could not directly determine the mechanism of the micro-turbulence. Fig. 8 shows the propagation of a heat pulse in Tokamak plasma in the new picture, it is easily to be simulated by a formal thermoconductivity.

![Graph showing the sensitive dependence of $T_{ge}$ on the boundary temperature $T_a$ with the tearing modes.](image)

**FIG. 7.** The sensitive dependence of $T_{ge}$ on the boundary temperature $T_a$ with the tearing modes.

![Graph showing the propagation of the heat pulse with the tearing modes.](image)

**FIG. 8.** The propagation of the heat pulse with the tearing modes.
We have not made any attempt to simulate the real plasma process, but only emphasized the necessity to consider the coupling of transport processes with MHD modes, and to renew our point of view to analyze the problem of energy confinement.
Numerical simulation of current instability development and saturation in a plasma with two ion species under the relative motion of ion components transverse to the strong magnetic field are discussed. It is shown that saturation of ion-ion current instability is of strong turbulent nature. This fact is based on the generation in configuration space of non-stationary density holes (cavitons) and regular structures (vortexies) in space phase. Resonant interaction of ions with plasma waves trapped into the cavity leads to the destruction of regular structures and to the stochastic (turbulent) plasma heating.

Particle simulation of two-ion hybrid cyclotron resonance heating of a magnetized hydrogen plasma with deuteron minority of FMSW launched from the low magnetic field side is presented. Depending on the minority concentration, partial transmission and partial reflection of the incoming waves off the two-ion hybrid resonance layer occur, in contrast to the mode conversion mainly taking place during incidence from the high field side. Preferential minority heating is observed, as the minority cyclotron resonance is close to the two-ion hybrid resonance layer.

I. Turbulent heating of plasma ions under the development of ion-ion instability (computer modelling)

1.1 Introduction. As is known [1-5], relative motion of ion components under the effect of low-frequency (\( \omega \approx \omega_{\text{ei}} \)) electromagnetic waves across static magnetic field leads to the buildup of various ion-ion instabilities. The lowest threshold in the frequency region \( \omega \sim \omega_{\text{ei}} \) has a small-scale instability (\( k \rho_L >> 1 \), \( k \)-wavenumber of unstable oscillations, \( \rho_L \)-ion Larmor radius) for which \( \omega >> \kappa_{\text{ei}}, \kappa_{\text{e}}, \kappa_{\text{i}} \), where

\[
\kappa_{\text{e,i}} = \left( \frac{m_i}{m_e} \right)^{\frac{1}{2}} \kappa_0
\]
represents a component of wavevector in the direction of external magnetic field. By the order of magnitude the growth rate of this instability is determined by relation \( \gamma \sim Re \omega \sim \omega_{pl} \) (\( \omega_{pl} \)-ion Langmuir frequency) with wavenumber \( k \sim k_0 \sim \omega_{pl}/u_{ti} \). For strong magnetic field an electron exerts negligible effects on the instability development, so it is conceivable that electrons could be treated as an immovable background necessary to sustain plasma quasineutrality.

In this section the results of numerical experiment modelling physical processes connected with development and saturation of ion-ion instability in a magnetized plasma containing two ion species plasma is reported.

1.2 Numerical model and parameters. Let us consider the uniform unbounded plasma with two ion species of different temperature \( T_2 > T_1 \) and masses \( m_2 > m_1 \). More "light" ion component \( m_1, T_1 \) has initial velocity \( v_1 \) across the external magnetic field. The numerical scheme of the code has been described in [6]. The kinetic equations for ion components and Poisson equation are solved with periodical boundary conditions on the interval \( L = 500 \, r_{D1} \), where \( r_{D1} \)-Debye length for "light" ions. Basic parameters of the model are follows: \( m_2/m_1 = 2, T_0/T_0 = 1, n_2/n_1 = n_e/2 \) (\( n_e \)-density of electrons), "half-width" of model particle \( \Delta = r_{D1} \), time step \( \Delta t = 0.04 \, \omega_{pe}^{-1} \).

1.3 Linear theory of ion-ion instability in a magnetized plasma. In paper [1] the following dispersion equation describing electrostatic plasma oscillations in a magnetized plasma has been obtained:

\[
\frac{1}{1 + \sum_{j=1}^{\infty} \frac{\omega_{pl}^2}{\kappa^2 \sqrt{\gamma_j^2}} \left[ \frac{1}{1 + i \sqrt{\pi} \, \Psi(\gamma_j)} \right]} = 0, \tag{1}
\]

where \( \gamma_j = (\omega - K \, \bar{u}_j)/\sqrt{\kappa} \, \kappa \right)^{1/2} \), \( \Psi(\gamma_j) \)-Kramp's function, \( \bar{u}_j \)-directed velocity of the \( j \)-th ion's sort. It was supposed that \( \omega V \ll \omega_{pe} \), \( \lambda \ll \lambda_j \), where \( \omega \) and \( \lambda \)-frequency and wavelength of the unstable oscillations.

The results of the numerical solution of Eq (1) are reproduced on Fig. 1.1. These results support the theoretical predictions of Ref. [1]: for critical value of relative velocity \( |\bar{u}_{rel}| \approx 0.5 \, \sqrt{T_1} \) the build-up of instability takes place. Unstable oscillations occupies the region in the
Fig. 1.1 Dependence of frequency ($\operatorname{Re} \omega$) and growth rate ($\gamma = \operatorname{Im} \omega$) from wavenumber $k$ for following set of parameters: $n_{o1} = n_{o2}$, $m_2 = 2m_1$, $|\tilde{u}| = 5|\tilde{u}_c| = 4.5 \nu_c$. Solid line corresponds to $T_{o1} = T_{o2}$; the dashed line $- T_{o2} = 1.5 T_{o1}$.

wavenumber space determined by relation $0 < k < k_c \approx 0.3 \nu_c^{-1}$ with $(\operatorname{Re} \omega)_{\text{max}} \approx \omega_{pA}$ and growth rate $(\operatorname{Im} \omega) = \gamma_{\text{max}} = 0.34 \omega_{pA}$ with corresponding value of wavenumber $k = k_{\text{max}} = 0.1 k_0$.

1.4. Results of the numerical experiment. The overall dynamics of ion-ion instability at the strongly supercritical regime ($\nu = 5 \nu_c$) can be described by the analysis of plots representing time dependence of the energy density $\gamma$ and temperature $T_\alpha$ of ion components (Fig. 1.2).

The distribution functions of ions at the initial state ($\omega = 40$) corresponds to Maxwellian ones shifted on the value equals to the directional velocity (Fig. 1.4). From Fig. 1.3a it follows that at $\omega = 40$ build-up of narrowband ($0.1 < k_r < 0.16$) unstable oscillations takes place with the growth rate of the most unstable mode ($k_r = 0.11$) which coincides with the high degree of accuracy with the results of linear theory (see Fig. 1.5).
Fig. 1.2  Time evolution of electrostatic field energy $W_F$ and ion temperatures $T_{1,2}$ ($n_{01} = n_{02}$, $T_{01} = T_{02}$; sign × marks temperature of "high" ions, sign • temperature of "heavy" ions).

(a) $E(\omega) \times 10^{-4}$

(b) $E(\kappa)$

Fig. 1.3  Spatial spectrum of the Fourier-harmonics of the electrical field $E_\kappa$ at different moment of time: a) $\omega_p t = 40$ and b) $\omega_p t = 60$. 
In accordance with Fig. 1.2 it is possible to distinguish some characteristic stages in the instability development process. The first one ($40 < \omega_{pl} t < 55$) is associated with exponential growth of electrostatic energy $W_F$ and monotonic increase of the ion's temperature. Simultaneously, the relative velocity of ion components decreases essentially (see Fig. 1.6). The phase space picture at the moment $t = 55 \omega_{pl}^{-1}$ (Fig. 1.7) shows the formation of the phase "holes" for the phase space of the "light" ions. Inside the holes the motion of the ions is characterized by closed trajectories (see Fig. 1.8). The formation of the phase space holes are accompanied by the development of the localized distribution of the electric field amplitude.
Fig. 1.6 Time evolution of current velocity of ion components \((x - u_1, *, - u_2)\) and their relative velocity \(|\vec{u}| = |\vec{u}_1 - \vec{u}_2|\).

Fig. 1.7 Plot of "light" ion phase space \((v, x)\) at time \(t = 56 \omega p_1^4\).

Fig. 1.8 Phase trajectory of single "light" ion from the region of phase space "hole".
(Fig. 1.9). From Figs. 1.8 and 1.9 it is easy to calculate the intensity of the electric field $E$ necessary to trap the "light" ions:

$$\gamma_{\text{theor}} = \frac{E}{(\pi n_0 T_0)^{1/2}} \approx \left(\frac{2\pi}{T_{tr} \omega_{pl}}\right)^2,$$

where in accordance with Fig. 1.7 the characteristic period of the trapped ion oscillation $T_{tr} \approx 10 \omega_{pl}^{-1}$. By this means at $\omega T_{tr} = 0.2$ one finds $\gamma_{\text{theor}} = 1.8$ which is in accordance with the result of the numerical experiment ($\gamma_{\text{exp}} = 1.85$). At the time the interval $55 < \omega T_{tr} < 60$ preferential heating of "heavy" ions is observed because of the "light" ions in view of their trapping by waves falls essentially. At $\omega T_{tr} = 60$ the difference in the temperature $\Delta T = T_2 - T_1 / T_0$ has attained the value $\Delta T = 10$ which is high enough for the development of the "non-isothermal" ion-ion instability [1,5,6]. A distinguishing feature of this instability is its low threshold of excitation: the instability develops even under the condition $|\Delta| < \delta T_3 (\gamma)$ which as can be seen from Fig. 1.6 easily holds for $\omega T_{tr} = 60$. From Fig. 1.1 it follows that for $T_2 = 1.5 T_1$ (dashed lines) the growth rate of the instability decreases slightly and the width of the instability domain increases as compared to the case of the isothermal plasma. The build-up of the "non-isothermal" instability at the interval $60 < \omega T_{tr} < 65$ leads to the substantial growth of the electrostatic energy $W_P$ and heating of the "light" ions which by its nature is very similar to the turbulent ion heating in
non-isothermal plasma \((T_i \gg T_e)\) under the development of the electron-sound instability \([7]\). This instability is characterized by considerable extension in the Fourier-spectrum of the electrical field (Fig. 1.3b) and the formation of the phase space "holes" for "heavy" ions. The intensity of the electrostatic field reaches its peak value at \(\omega_{pl}t = 74\) when the relative velocity of the ion components becomes below the critical velocity \(u\) (see Fig. 1.6). At this time the instability saturates definitively through the trapping of both the ion species by the fields of unstable waves as determined by the phase space plots (Figs. 1.10) At this time the interval \(74 < \omega_{pl}t < 105\) gradual destruction of the localized structures is observed which accompanied by considerable reduction of electrostatic field energy and slow heating of the balk ions. The correlation analysis of the time-dependence of the electrical field amplitude shows that on the time interval \(83 < \omega_{pl}t < 100\) along with the regular oscillations of the low frequency the high-frequency oscillations with random phases occur. In accordance with Fig. 1.11 the time of the fields phase decorrelation \(\tau_{DC} = 10 \omega_{pl}^{-4}\). Hence, slow heating of ions at this stage can be associated with stochastic mechanism which is based on the scattering of ions on the ununiformities of the turbulent pulsations. It should be particularly emphasized that stochastic heating of the ions as a result of the modified two-stream instability has been described before \([8-10]\).

Fig. 1.10 Plots of ion phase space at time \(t = 74 \omega_{pl}^{-4}\).
Fig. 1.11 a) frequency spectrum of the Fourier-harmonics $E_n$ of the electrical field, b) time evolution of autocorrelation function $A_c$, c) time dependence of the electric field under analysis.

II. Simulation of ICRH through resonant absorption in two-ion species plasma

2.1 Introduction. We study a numerical simulation of two-ion hybrid resonance heating that plays an important practical role in many experiments of ion cyclotron resonance heating in tokamaks and stellarators [2, 13]. In a paper [12], which discussed the results of numerical study of two-ion hybrid resonance heating, Riyopoulos and Tajima observed mode conversion of magnetosonic wave that is excited by a current sheet and is propagating across the magnetic field into the two-ion electrostatic modes. The interaction of plasma ions with strong electrostatic oscillation leads to efficient heating of plasma ions. The results of Ref. [12] were obtained with hydrogen ion (H)
as minority and deuteron (D) as majority. In the case of fast magnetosonic wave excitation from the high magnetic field side, the electromagnetic wave attains the ion-ion hybrid resonance region before the minority cyclotron resonance, being mode-converted to the ion-ion short-wavelength electrostatic mode. In this case efficient heating of both majority (D) and minority (H) ions is observed. The rates of heating depend mainly on relative concentrations and the ion plasma to cyclotron frequency ratios, but in fact are not very different for both ion species. The plasma fluctuations excited through the complete mode conversion of the incoming magnetosonic mode to the ion-ion electrostatic mode cause strong velocity space diffusion for both ion species. The wave incidence from the low field side, on the other hand, shows preferential minority heating and much smaller heating of majority ions. This is because for the case of low field side excitation the wave encounters the minority cyclotron resonance layer before the ion-ion resonance layer. In this case, consequently, effective minority heating connected with usual ion cyclotron resonance absorption takes place. A reflection of the incoming magnetosonic mode takes place at the cut-off layer, located between the minority cyclotron resonance and the ion-ion hybrid resonance, instead of a mode conversion to the ion-ion hybrid mode.

2.2 Low-field side excitation of ICRH. In the present section, we consider an experimentally created scenario of low field side excitation during recent ion-cyclotron resonance heating (ICRH) experiments in the T-10 tokamak [11]; The wave propagates perpendicular ($\kappa_n = 0$) to the magnetic field in a two-ion species plasma with hydrogen, as a majority and deuteron as a minority. This case differs from the low field side launching of Ref. [12] in that the heavy ion (D), rather than the light (H), is now the minority and that, therefore, the sequence between ion-ion hybrid and minority cyclotron layers is reversed. The profiles of the majority and minority cyclotron frequencies across the varying magnetic field are shown in Fig. 2.1. The ion-ion hybrid frequency curve is given by

$$\omega_{\text{hyb}} = \frac{2}{\kappa_n} \left( \frac{\omega_p^2}{\omega_p^2 + \Omega_\alpha^2} + \frac{\omega_p^2}{\omega_p^2 + \Omega_\beta^2} \right) \left( 1 + \frac{\omega_p^2}{\sqrt{\mu}} \right)^{-1/2},$$

(2)

where $\omega_\alpha$, $\omega_\beta$ are the ion plasma frequencies and $\Omega_\alpha$, $\Omega_\beta$ are the ion cyclotron frequencies of the majority ($\alpha$) and minority ($\beta$) ion species, respectively. According to Eq. (2) the ion-ion-hybrid frequency lies between the cyclotron frequencies for each ion species and falls much closer
Fig. 2.1 The one-dimensional plasma model in our simulations. Shown are the ion cyclotron frequency profiles and the two-ion hybrid frequency $\omega_{1h}$ profile as a function of the distance $x/L$. The line $\omega = \omega_{pump}$ matches the local $\omega_{1h}$ at $x = 0$. The deuteron cyclotron resonance point is also within the simulation system ($x = x_D$).

to the minority cyclotron frequency $\Omega_\beta$ ($\omega_{\gamma}, \Omega_\beta, \Omega_{1h} \rightarrow \omega_{1h}$). The ion hybrid layer $x_{1h}$ and the cyclotron resonance $x_{\beta}$ ($x_D$ for the present case) are defined by $\omega_{1h}(x_{1h}) = \omega_p$ and $\Omega_\beta(x_{\beta}) = \omega_p$ respectively, with $\omega_p$ the pump frequency. Consequently during propagation from the low field side (on the right in Fig. 2.1) encounters a cut-off layer $x_C$ close to but before the ion-ion hybrid layer $x_{1h}$. A complete treatment [13,14] through the differential equation for propagation including warm plasma and inhomogeneity effects reveals that an almost total reflection of the incoming magnetosonic wave off the resonant regime and back towards the exciting antenna will occur in case $k \parallel \Omega_{1h} \leq 0$. Transmission and mode conversion into the ion-ion hybrid mode can be very small depending on the effective width of the non-propagation region between the cut-off and the resonance. This region gets narrower with decreasing minority content and a considerable

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tunneling through may occur for small minority percentage. This aspect of the low field side incidence differs from the high field side incidence where mode conversion of the incoming radiation into the ion-ion hybrid mode rather than reflection or tunneling is the dominant outcome. It also implies that the wave does not reach the exact minority cyclotron resonance in the case with relatively high percentage of deuteron as minority.

2.3 Electromagnetic particle simulation. The present investigation is based on the fully self-consistent electromagnetic particle code with electrons treated as guiding centers in order to eradicate the excessive noise that would otherwise arise due to electron thermal motions [12]. The full ion dynamics is followed in our simulation. The system is 1-1/2 dimensional (one spatial and two velocity space coordinates) with absorbing boundaries in the χ-direction [15]. We investigate two cases of relatively large and small minority concentration.

Figure 2.2 (a) represents time dependence of Poynting vector for the following set of parameters:

\[ B_0, \quad m_{maj} = 0.5, \quad \frac{T_{maj}}{T_{min}} = 1, \quad \frac{h_{maj}}{h_{min}} = 9, \]

corresponding to \( \omega_p = 1.076 \Omega_d \) at the ion hybrid layer \( \chi_{ih} \). From this picture an almost complete reflection of the incoming electromagnetic flux near the ion-ion resonance-cut-off region is observed. No significant conversion to the ion-ion hybrid mode is observed as the electrostatic field profiles indicate no activity significantly above noise levels. The increase of the electromagnetic energy tapers off at time \( t \approx 180 \omega_{ph} \), which is the arrival time of the wave near the hybrid layer, as shown in Figs. 2.3(a)-2.3(b). Beyond this time the stored energy changes a little bit slower, since the wave cannot penetrate further to the left and the net incoming flux is balanced by particle heating. The electrostatic energy keeps increasing in accordance with the increasing level of thermal fluctuations caused by the increasing temperature. There is only a small amount of electromagnetic transmission through the middle as shown in Fig. 2.2(a). The phase space plots \( v_x, v_\chi \) vs. \( \chi \) at the end of the run (Figs. 2.3(a)-2.3(f)) show the ion response to the pump wave. There is little activity on the left of the resonance-cut-off layer as expected. The heating rate for the ions is shown in the energy versus time plots in Figs. 2.3(b)-2.3(d) in comparison with the field energy density history. In spite of the near total reflection...
ion heating takes place. The relative increase in deuteron temperature is 50% over a time span of 10 cyclotron periods. We observe a relative increase in temperature \((\Delta e_D/e_D \sim 1.5)\) rate three times faster for deuteron than for hydrogen, although more absolute amount of energy goes into majority ions. A physical interpretation of this phenomenon is the following: Nonlinear particle trapping combined [12] with quasilinear diffusion due to enhanced turbulence heats the minority (D) more effectively than majority as the wave frequency near the hybrid resonance is close to the minority cyclotron frequency.
Fig. 2.3 Time history of field energies and ion temperatures during ICRH ($t = 0+400\omega_{ph}^{-1}$) for $n_D/n_H = 1/9$; (a) The electrostatic field energy $E_L/4\pi$, (b) the transverse electric field energy $E_{tr}/4\pi$, (c) majority (hydrogen) ion kinetic energy $e_H$, (d) minority (deuteron) ion kinetic energy $e_D$, (e) $(p_x,x)$ phase space plot hydrogen at $t = 418\omega_{ph}^{-1}$, (f) $(p_x,x)$ space plot of deuteron at $t = 418\omega_{ph}^{-1}$.
We show the results of a similar case of lower minority concentration $\eta_{\text{maj}}/\eta_{\text{min}} = 39$ in Figs. 2.2(b) and 2.4(a)-2.4(d). It is obvious from the electromagnetic Poynting flux plot of Fig. 2.2(b) that a significant wave penetration to the left of the cut-off region takes place with the wave propagating all the way to the left boundary. The stored electromagnetic energy does not saturate until $t \sim 300\omega_{\text{ph}}^{-1}$, roughly the arrival time at the left boundary. The majority and minority phase plots (2.4(c) and 2.4(d)) at the end of the run $t = 418\omega_{\text{ph}}^{-1}$ verify propagation to the left half, as opposed to Figs. 2.3(e) and 2.3(f). As the wave now reaches the deuteron cyclotron layer, it causes intense selective heating of the minority. The relative increase in deuteron temperature is $\Delta e_D/e_D = 2.5$ over ten cyclotron periods, much larger than in the previous case. The total (combined) ion energy absorption however is not significantly different in

![Diagram](https://via.placeholder.com/150)

**Fig. 2.4** Time history of ion temperatures during ICRH ($t = 0$-$400\omega_{\text{ph}}^{-1}$) for $\eta_d/\eta_H = 1/39$. (a) majority (hydrogen) ion kinetic energy $e_H$, (b) minority (deuteron) ion kinetic energy $e_D$, (c) $(p_x', x)$ space plot of hydrogen at $t = 418\omega_{\text{ph}}^{-1}$, (d) $(p_x', x)$ space plot of deuteron at $t = 418\omega_{\text{ph}}^{-1}$. 75
both cases. In contrast, in case of hydrogen minority in a deuteron background when the minority cyclotron resonance always falls to the low field side of the two ion hybrid layer, similar, but even stronger, preferential minority heating has been observed numerically.

2.4 Discussion. It has been suggested that buildup of temperature difference, combined with the local counter-drifting between the two ion species under the influence of the wave, may trigger velocity space instabilities [16,17] when $T_D/T_H$ becomes large. This will result in an effective redistribution of the deuteron energy into the hydrogen ions on a time scale much faster than the Coulomb collision time scale.

The anomalous heating of plasmas containing two ion species has been observed to date in stellarators [2,3] and tokamaks [11] during generation of intense ICRH and fast magnetoacoustic waves; in a rotating plasma under the conditions of crossing electric and magnetic fields [18]; in mirrors during the buildup of ion-ion instabilities [19,20]. Considerable enhancement in neutron production, attributed to the heating of the minority ions of the beam, was recently reported in combined beam - ICRH heating experiments in JET [21].

The experiments on Uragan [2,3] were designed to study the heating of H + D plasmas with Alfvén waves of moderate power coupling (~ 300 kW). The plasma density $(4 \times 10^{12}) \text{cm}^{-3}$ was achieved by resonant wave excitation across a broad range of the longitudinal magnetic field strength $B_0 = (5 \times 20) \text{kG}$. Provided that the conditions are met for IC-resonance of deuterium ions ($\omega = \omega_{\text{CD}}$, $B_0 = 13\text{kG}$), the ions of hydrogen ($T_H < 1 \text{keV}$ at $n \sim 10^{12} \text{cm}^{-3}$) were found to heat up anomalously fast (less than 1 ms). The measured times of proton heating were less than those of energy exchange between the ions of H and D by far on account of Coulomb collisions ($\tau_0 \sim 4 \text{ ms at } n \sim 4 \times 10^{12} \text{cm}^{-3}$, $T_D \geq 300 \text{ keV}$). The fluctuational magnetic field strength in the IC-resonance region being 15G, the deuterons in the wave field picked up the directional speed $v_{\perp,\text{B}} = \sqrt{\omega_B} B_0 / 4\pi n_e c \sim 4 \times 10^7 \text{cm/sec in excess of the thermal one}$ $v_{\perp,\text{B}} \sim 10^7 \text{cm/s})$. Such conditions are conducive to growth of a microscale ion-ion instability whose phase of saturation is peculiar for high level turbulent pulsations. Scatter of both the ion species upon pulsations leads to an effective collisionless heating of the ions.
The paper [19] gives the results of experiment on growth and saturation of the ion-ion instability in a mirror \((\mathcal{R} = 60 \text{ cm}, L = 120 \text{ cm})\). The plasma volume is separated with a double negative potential mesh (tandem mirror). The experiments were run in a cold \((T_0 \sim 7 \text{ eV})\) low density \((n \sim 10^8 \text{ cm}^{-3})\) plasma. Changing the potential sign of the guiding mesh brought the beam of ions into the "target" plasma with a velocity close to the ion acoustic one. Such procedure discovered density wells of the space scale \(L \sim (5 + 10) \text{ cm} \sim 100 \lambda_{pi}\) and associated long time-of-life structures (vortices) \((D \sim 50 \omega_{pi}^{-1})\) localized in phase space which is indicative of trapping of the bulk of ions by the unstable wave electric field. The measurements of the ion velocity distribution function is a direct experimental evidence of ion trapping, bringing to light a rise in the velocity range corresponding to the central part of the vortex.

Worth specific reference are the results of experimental programme on rotating plasma [18] and a triple mirror [20] which also attested to ion-ion instability growth and anomalous heating of ions.

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Abstract

Transport phenomena associated with pellet injection have been intensively studied on JIPP T-IIU tokamak. When a large pellet has been injected, marked peaking of the density profile has been observed just after the hollow density profile produced by the pellet injection. This marked inward pinch of the density might suggest some instability associated with an $\eta_i$ mode, since density fluctuations on the ion mode has simultaneously observed in FIR scattering. By injecting a small pellet as a source of a density perturbation, the particle diffusion has been analyzed, and diffusion coefficients and inward pinch velocity have been derived with 1-D transport code. Next, we have studied the ultra fast cooling phenomenon (so-called pre-cooling) at the pellet ablation phase, where the propagation speed of the cooling front is extremely faster than the pellet injection speed. This pre-cooling phenomenon has been confirmed with two independent diagnostics (ECE and Soft-X ray). While, the density around the plasma center has never increased during this pre-cooling, contrary to the clear decrease of the electron temperature.

I. Transport analysis associated with pellet injection

On JIPP T-IIU tokamak (R=0.91m/a=0.23m and Bt=3T), particle and thermal transport associated with pellet injection has been intensively studied. Diagnostics particularly prepared for this study are ECE grating polychromator ( spatially 10 channel with $\Delta t=2\mu s$) and HCN interferometer ( spatically 6 channel ).

Various sizes of pellets are injected into plasmas, to examine the particle transport characteristics. Figure 1 shows the time evolutions of the density profile factor, $k(t)$, defined by the fitting of the density profile as $n(r,t)=n(0,t)(1-(r/a)^2)^k(t)$. When the small pellet was injected, the density profile did not change so much except the transient phase just after the pellet injection ($\Delta t < 10$ ms), and the confinement improvement is not observed.
Fig. 1 density profile for (a) large and (b) small pellets, where the profile factor, $k(t)$, is defined by $n(r,t)=n(0,t)(1-(r/a)^2)k(t)$.

While, when a large pellet has been injected, marked peaking of the density profile has been observed just after the hollow density profile produced by the pellet injection. Since density fluctuations on the ion mode has simultaneously observed in FIR scattering[1], this marked inward pinch of the density might suggest some instability associated with an $\eta_1$ mode.

We have analyzed particle transport coefficients with a small deuterium pellet injection, which is introduced as a source of an electron density perturbation. Figure 2 shows the time evolution of the density profile and the profile factor, $k(t)$. The integral form of the particle diffusion equation is written as

$$\frac{\partial}{\partial t} \int_0^r n(r) dV_p = 4\pi^2 R_0 (rD \frac{\partial n}{\partial r} + vn)$$

where $D$ and $v$ are a particle diffusion coefficient and an inward pinch velocity. The source term is neglected in eq.(1), because at the density relaxation phase (after the pellet ablation) neutrals around $r/a=0.5$ is negligibly small. The values $D$ and $v$ are determined by evaluating the above equation at two different times $t_1$ and $t_2$ in Fig. 2, resulting in $D=0.4$ m$^2$/s and $v = 4$ m/s at $r/a=0.5$.

Using a one-dimensional transport code, the time evolution of the density profile has been analyzed. The experimental data just after the pellet injection (i.e., $t=145$ms) has been employed into the computer code as the initial density profile. The diffusion coefficient $D(r)$ and the inward velocity $v(r)$ are assumed as

$$D(r) = \frac{D(0)}{(n(r)/n(0))} \quad (2a)$$

$$v(r) = v(a)(r/a)^n, \quad n=2 \quad (2b)$$
Figure 3 shows the density profiles at the steady-state (t=153 ms) for various $D(0)$ and $v(a)$ values. A good agreement with experimental data has been obtained in the case of $D(0) = 0.4 \text{ m}^2/\text{s}$ and $v(a)= 20 \text{ m/s}$, which gives the values of $D=0.5 \text{ m}^2/\text{s}$ and $v=5 \text{ m/s}$ at $r/a=0.5$. It is found that the transport coefficients simply calculated with two density profiles at different times, as described above, are in good agreement with those evaluated by 1-D transport code. We should remark that for $n=1$ case in eq.(2b) we could not realized the experimental data.

**Fig. 2** Time evolution of the density profile $n(r,t)$ with a small pellet injection.

**Fig. 3** Density profiles calculated by a 1-D transport code for various particle diffusion coefficients $D(0)$ and inward velocity $v(a)$. 

![Diagram](image-url)
II. Pre-cooling phenomena at the pellet ablation phase

At the ablation phase of injected pellet it had been reported in TFR[2] and JET[3] that the propagation speed of the cooling front due to the pellet penetration into core plasmas is much faster than the injected speed of pellet. This so-called pre-cooling phenomena has been also observed in JIPP T-IIU plasmas[4]. To study this phenomena more in detail, time- and space-resolved diagnostics are prepared, and simultaneous measurements of the density and the electron temperature profiles have been conducted.

Figure 4 shows time evolution of the electron temperature measured by ECE system at different major radii just at the pellet injection. Although the pellet has been injected from large major-radius side, the signals of small major-radius side begin to decrease almost simultaneously. This symmetric decay of the temperature profile indicates that ECE signals are free from the cut-off problem due to the quick density rise.

Cooling front (start of temperature drop at each radius) has been indicated by arrows in Fig.4, and replotted as a function of arrival time of the cooling front in Fig.5. In the outer region of the plasma column (R> 96cm and R<80cm chords), the cooling front penetrates into the plasma with a moderate and constant speed, which is comparable with the injection speed of the pellet. While, in the central region the propagation speed is significantly faster than the pellet velocity. Soft X-ray signals, as plotted in Fig.5, have also represented the similar characteristics.

![Fig.4](image1.png)  ![Fig.5](image2.png)

Fig.4 Electron temperature measured with ECE grating polychromator at different major radius.

Fig. 5 Properties of cooling front measured with ECE (shown by arrows in Fig.4) and Soft X-ray.
To examine a relation between the q=1 surface and the pre-cooling, pellets have been injected into various kinds of plasmas, by changing plasma current and sawtooth phase. When the pellet is injected into no sawtooth plasmas, the cooling front propagates to the plasma center with constant speed and the pre-cooling has never occurred. When the pellet has never reached the sawtooth inversion radius, the pre-cooling has never occurred, too. These experiments, therefore, tell us that for the pre-cooling it is necessary for the pellet to reach the sawtooth inversion radius.

Next, the correlation between the pre-cooling and the sawtooth oscillation has been examined. The dependence of the pre-cooling characteristics on the sawtooth phase (0% and 100% mean just after and before the sawtooth crash, respectively) are shown in Fig. 6. If the pre-cooling is assumed to have a strong correlation with the q=1 surface, the standard sawtooth model suggests that in 0% phase (just after the sawtooth crash) the pre-cooling should start at the plasma center, because the q-value at the plasma center is relaxed to unity. This is, however, contradicting to the results given in Fig. 6, where in the case of 0% phase the pre-cooling start around the sawtooth inversion radius, not at the plasma center. Another possibility to explain these data is that q(0) value is kept sufficiently below unity even just after the sawtooth crash, as observed experimentally in TEXTOR[5] and TEXT[6] devices. According to this model, it may be explained that the pre-cooling starts at this q=1 surface far from the plasma center.

Simultaneously the density profile has been measured with 6 ch. HCN interferometer system. The density profiles just before and after the pellet injection are shown in Fig. 7, where the pellet is ablated mainly at r/a=0.5 and the profile has become hollow. The electron temperature profiles corresponding to this and the resultant electron pressure, n(r)Te(r), are plotted, too. It is found that just after the injection the central temperature has already decreased, although the central density has not yet changed.
As for the mechanism of the pre-cooling phenomena, it seems difficult to predict theoretically. Since we have identified that the pre-cooling has a strong correlation with the q=1 surface, we could say that some MHD instability might play an important role on the pre-cooling phenomena.

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EVALUATION OF ABSOLUTE IMPURITY ION DENSITIES IN TOKAMAK EDGE PLASMAS BY MEANS OF LITHIUM BEAM-ACTIVATED CHARGE EXCHANGE SPECTROSCOPY (Li-CXS)

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Abstract

Impurity ion concentrations in the edge plasma of tokamak discharges can be determined with excellent spatial and temporal resolution by means of lithium beam-activated charge exchange spectroscopy (Li-CXS), i.e. by utilizing electron capture from injected fast Li atoms into excited states of the impurity ions of interest, and spectroscopic observation of the resulting characteristic impurity ion line emission. First demonstration measurements of the new diagnostic method are being presented. Furthermore, we discuss important aspects regarding the practical use of Li-CXS to achieve precise information on impurity ion concentrations in tokamak edge plasmas.

1) Attenuation of the injected lithium beam has precisely to be taken into account. For this purpose we have developed an attenuation code which includes the excitation, deexcitation and ionization of injected Li atoms via collisions with electrons and protons in the edge plasma environment. The attenuation can be calculated for different numbers of excited LiI states.

2) After their population by electron capture from Li, some of the resulting excited impurity ion states before radiating may become mixed due to collisions with plasma particles. Consequently, the relation between impurity ion density and characteristic line radiation intensity can deviate considerably from the one derived in experiments involving single collisions between the impurity ions of interest and Li ground state atoms.

Attenuation calculations have been carried out to evaluate first measurements of C\textsuperscript{5+} and C\textsuperscript{6+} concentrations in the edge plasma of TEXTOR ohmic discharges (KFA Jülich, Germany), as will be demonstrated in the present report.

Regarding the actual importance of excited level mixing, specific experiments are needed to compare our calculations with measured distributions of charge exchange-excited impurity ion states produced within typical tokamak edge plasmas.

1) Introduction

The Li-CXS method for measuring impurity ion distributions in tokamak edge plasmas has been presented in a number of publications /1/ - /3/. For a proof-of-principle experiment for this novel diagnostic technique a 20 keV/20 mA neutral Li beam injector /4/ has been installed at the tokamak experiment TEXTOR of KFA Jülich, Germany. TEXTOR is very well suited for demonstration of Li-CXS because of its special dedication to plasma-wall interaction phenomena.
and edge plasma physics and its extensive edge plasma diagnostic equipment /5/. Fig. 1 shows the present experimental setup for Li-CXS measurements.

![Fig. 1 - Schematic view of the Li-CXS experiment at TEXTOR.](image)

The neutral Li beam is injected radially in the midplane of the TEXTOR discharge, with its path observed via a movable scanning mirror by filter-photomultiplier detectors operating at preselectable wavelengths in the visible spectral region.

Fig. 2 shows typical signals for line radiation from injected excited Li atoms and excited C\(^{4+}\) ions which are produced by charge exchange (CX) from C\(^{5+}\) impurity ions in the observed edge plasma region. Determination of absolute impurity ion concentrations calls for evaluation of the local density and composition of the injected Li atoms, for which the attenuation of the Li beam along the observation region has to be calculated (cf. next section).

2) Attenuation of the injected Li beam

The injected Li atoms interact with electrons and protons in the edge plasma (in first approximation we may neglect the presence of plasma impurities), and consequently become excited as well as ionized via collisional processes as impact excitation and ionisation, and charge exchange. Furthermore, the excited Li\(^{l}\) states undergo radiative as well as collisional deexcitation. Fig. 3 shows the Li\(^{l}\) states which have been considered in our attenuation calculations. For these calculations we have assembled a data base containing all collisional processes of interest. We used
Fig 2 - Signals from the pulsed (250 Hz/50% duty cycle) neutral Li current (top), the Li I 670.8 nm resonance line radiation (bottom) and the CX-produced CV 494.5 nm radiation linked to the C⁵⁺ impurity ions in the TEXTOR edge plasma.

Fig. 3 - Li states considered for calculation of the Li atom beam attenuation in the edge plasma region.

Experimental informations whenever available, but for most of the processes the needed cross sections had to be calculated /6/. By considering different numbers of excited LiI states in our calculations we checked the importance of including also higher excited states and the convergence of the developed attenuation code.

Fig. 4 shows a comparison between a measured LiI 670.8 nm resonance line profile and a calculated one. We took into account the relevant TEXTOR edge plasma conditions (profiles for electron- and ion density and temperature as measured by the TEXTOR standard diagnostics) to derive the resulting attenuation of the injected neutral Li beam /6/, /7/. The good agreement seems
to demonstrate that our calculations probably included all important mechanisms responsible for Li beam attenuation. However, we plan to carry out more critical tests involving observation of other LiI lines originating from excited LiI(n=3) states. In this way our present attenuation code might be further developed for eventually being capable to derive from observation of different LiI line profiles on a shot by shot basis the corresponding Li atom beam composition along the injection line together with the respective plasma density distribution.

3) Evaluation of impurity ion profiles

In Fig. 5 measured characteristic impurity line emission profiles for C$^+$ (529 nm resulting from CVI excited state) and C$^+$ (494.5 nm resulting from CV excited state) are shown. To derive from these profiles the corresponding impurity ion concentrations, the earlier described attenuation calculations for the injected Li atom beam have been carried out and delivered the local densities of the various excited LiI states. By using appropriate scaling we estimated the cross sections for electron capture from different excited Li atoms into the relevant excited impurity ion states /7/. Finally, we derived the local impurity ion densities as they are shown in fig. 6. Different numbers of excited LiI states have been used in these calculations to check the convergence of the method.

4) Collisional mixing of Li-CX-induced excited impurity ion states

In fig. 7 we show an example for mixing of CX-excited impurity ion states. By means of electron capture from injected Li atoms into C$^+$ impurity ions various excited CIV states are produced /8/, which in the edge plasma environment may become mixed due to collisions with electrons and ions before decaying via emission of characteristic line radiation. We have selected a
Fig. 5 - Spectroscopically observed Li-CXS-induced characteristic line profiles related to carbon impurity ions in the edge plasma of TEXTOR

Fig. 6 - Carbon impurity ion concentrations derived from the measured data shown in fig. 5, after calculation of Li atom beam attenuation and state composition (cf. text)

Fig. 7 - Mixing of CX-produced excited impurity ion states in the edge plasma (selected case for Li-CXS of C\textsuperscript{4+} ions)
particular case, i.e. mixing among the CIV 6g and 6h states. By considering the relevant radiative lifetimes in comparison with the mean collision times for impact excitation and -deexcitation processes, the modification of the initial CX-induced population pattern can be evaluated. For these calculations we take into account the local distributions for plasma density and -temperature in the considered edge region. Fig. 8 shows results of such a calculation /9/ for different local plasma density. There is also a relatively weak dependence on the electron temperature. By measuring ratios of selected emission lines as observed at different edge plasma locations, we intend to check the reliability of these calculations in the near future.

![Fig. 8 - Calculation for mixing of CX-produced CIV states in the edge plasma vs. local edge plasma density (local electron temperature 100 eV)](image)

**Summary**

We have presented first experimental data for carbon impurity ion concentrations in the TEXTOR edge plasma, as determined by means of lithium atom beam-activated charge exchange spectroscopy (Li-CXS). Evaluation of impurity ion concentrations from the measured raw data involves calculation of the attenuation and varying LiI excited state fractions along the Li injection line, for which purpose an attenuation code has been developed. In addition, collisional mixing of excited impurity ion states between their CX-production in the edge plasma and their radiative decay can provide considerable experimental uncertainty, but might be removed after further studies involving Li-CXS of selected impurity ion species under different edge plasma conditions.
ACKNOWLEDGEMENTS

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EDGE PHENOMENA

(Session 2)

Chairman
P.K. KAW
India
SUMMARY OF THE SESSION
P.K. KAW

The papers presented in this session may be conveniently subdivided into those dealing with experimental results on existing and newly commissioned tokamaks, diagnostic development experiments and theoretical papers on edge physics.

The session started with a general introduction on the importance of edge phenomena in tokamaks (P. Kaw, India). Edge plasma physics is strongly influenced by the presence of impurities and neutrals which lead to radiative and ionization phenomena. In the past the edge was only supposed to perform a few useful functions such as fueling, heat and ash removal, shielding from impurities, etc. However, experiments with moderate size tokamaks in the past decade or so have shown that the edge often controls crucial properties of the core plasma, sometimes in a totally unexpected manner. New phenomena such as the H-mode, MARFES and detached plasmas, improved core confinements with biased probes and limiters, second stable regime access using current pedestals at the edge, etc. all point to the increasingly important role played by the edge control in optimizing the performance of the core plasma in the tokamak. This realization coupled to the fact that the edge is more directly accessible to manipulation using novel limiters, divertors, gas puffing, pellet and low Z impurity injection, RF heating, RF current drive, RF plasma charging, feedback stabilization methods, etc. is going to make the study of edge physics a very fertile area of investigation in the coming years. It was asserted that experiments on devices similar to the ones to be discussed at the meeting can play an important role in these developments.

Experimental results from existing machines were presented by Takamura, et al (HYBTOK II, Japan), Stöckel (CASTOR, Czechoslovakia) and Xie, et al (HT-6B, China). Limiter biasing experiments have been done in HYBTOK II. Positive biasing was found to enhance the edge radial electric field and to increase particle confinement (as inferred from decreased $N_q$ and increased $\eta_e$). Lower hybrid current drive experiments were reported from HT-6B (2.4 GHz, 110 kilowatt, 60 msec.) and CASTOR (1.25 GHz, 40 kwatts). Castor experiments showed considerable suppression of edge turbulence, sharpening of edge density gradients and improvement of global particle confinement time. The HT-6B experiment showed good current drive efficiency up to $\omega_p(\omega)/\omega_{ce}(\omega) = 4$ and also demonstrated stabilization of sawteeth and Mirnov oscillations; start-up using lower hybrid waves was also demonstrated.

Two recently commissioned tokamaks also presented new results. The ADITYA team (India) gave an account of the status of the machine which is now routinely producing 50 kA., 20 msec discharges. Experimental results on breakdown studies, self-consistent electric field generation in the initial phase and current termination in low q discharges were reported. An account of the experiments planned on the control of tokamak edge fluctuations using feedback and biased double-probe systems was also presented. Breun, et al (USA) presented results from the Phaedrus-T experiment in which 100 kA 60 msec discharges have been already produced. Experiments using a 100 kilowatt, 20 msec ICRF heating source were also reported; one sees an increase of line density, impurity radiation, edge $T_e$ and edge $\eta_e$.

The diagnostic development experiments were presented by Kardon, et al (Hungary), Winter, et al (Austria) and Breun, et al (USA). Kardon presented results of multichord time resolved (2 usec) plasma radiation measurements (soft and ultraviolet x-ray region) with a combination of absorber foils and microchannel plates on the MT-1 tokamak. The technique gives useful
information on $T_e$, magnetic structures and impurity concentrations in a tokamak and is especially useful for small tokamaks. Winter described experiments on the measurement of absolute electron capture cross-sections in multicharged ion-atom collisions. This data is important for modeling of the edge plasma as well as for the Lithium beam diagnostic. Breun, et al discussed the development of a diagnostic neutral beam (4.5 amps, 21 keV) on Phaedrus-T which is going to be used to measure $\bar{n}_e (r)$, $T_e (r)$, $q (r)$ and $E_{RF}$.

Theoretical papers on edge physics were presented by Kaw, et al, (India), Spineanu and Vlad (Romania) and Jankowicz (Poland). Kaw, et al investigated the influence of sheared poloidal flows and/or an active feedback source in the continuity equation on the linearized theory of dissipative drift instability, rippling instability and the radiative condensation instability. Conditions for stabilization were derived. Spineanu presented theoretical calculations on the setting up of large radial electric fields at the tokamak edge. In their model, particle percolation in quasipotential periodic structures associated with drift modes, was shown to aid in the development of potentials higher than those coming from neo-classical theory. Jankowicz presented a summary of his two-fluid modelling of non-ambipolar scrape-off layers in realistic toroidal geometry. His model is presently being implemented in the JET 2D fluid modelling codes for the edge.
This paper deals with two attractive methods which have been proposed for the control of edge fluctuations in a tokamak plasma, viz., (i) active feedback stabilization; and (ii) introduction of sheared poloidal flow. Eigen mode equations for the rippling mode, the drift-dissipative mode and the radiative condensation instability in the presence of a sheared poloidal flow and a phase sensitive feedback source have been set up and analytically solved in some limiting cases. Conditions for stabilization of the above mentioned instabilities are discussed.

Edge fluctuations are widely believed to influence the overall particle and energy confinement in a tokamak discharge\(^1\). Methods to control the level of these fluctuations are therefore of great interest in tokamak physics. Recent experiments\(^2\) with biased probes have demonstrated that shear in poloidal velocity flow can introduce a significant stabilizing influence on fluctuations in the tokamak edge. In fact, these experiments have lent credibility to the suggestion\(^3\) that the improved confinement performance of the H-mode in tokamaks is due to introduction of sheared poloidal flow in the edge region. Similarly, the use of phase sensitive feedback methods to stabilize low-frequency fluctuations in a laboratory plasma, has received wide experimental verification\(^4\). Such methods have not yet been attempted in tokamaks. However, feedback stabilization methods seem to be particularly attractive for the tokamak edge region because (a) it appears that the particle confinement time is largely governed by a relatively narrow band of low-frequency long wavelength fluctuations in the edge region, and (b) the edge region is accessible to probes, limiter bias potentials, etc., and so feedback methods could be readily implemented. Proposals to carry out such experiments therefore already exist\(^5\).

The important mechanisms which may be responsible for the observed edge fluctuations in a tokamak are: (i) rippling instabilities which are driven by resistivity gradients associated with temperature and/or impurity density gradients in the plasma; (ii) radiative-condensation instabilities which arise due to dependence of radiative capacity \(L\) on the density and temperature \((L \propto n^2\) and \(\partial L/\partial T < 0\); and (iii) drift-dissipative instability driven by pressure gradients and depending on parallel resistivity to give a phase difference between \(\vec{n}\) and \(\vec{\phi}\) (which is necessary for wave growth). In this paper we have examined the influence of poloidal velocity shear and a phase sensitive feedback source in the electron continuity equation (and hence in the charge neutrality condition) on the linear growth of the above mentioned three instabilities. Taking account of magnetic shear, a simple model of velocity shear and dependence of \(\omega\) on \(z\) (the radial distance \(r_0 - r\) from a mode rational surface), the slab-model eigenmode equations have been set up in

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\(^1\) Columbia University, New York, USA.
each case. Limiting cases have been solved by analytic methods using perturbation theory, WKB approach, matched asymptotic schemes, etc. Expressions for modified growth rates have been obtained and conditions for stabilization discussed. A detailed analysis of these equations, together with a numerical integration using shooting methods, will be published elsewhere.

We first give a general treatment for rippling, radiative condensation and drift-dissipative instabilities, using the Braginskii two-fluid equations. We take linear perturbations to vary as

\[ \phi \sim \exp(-i\omega t + i m \theta + i l \psi) \phi(x) \]

and use the slab approximation, where the parallel wavevector is written \( \hat{n} \cdot \nabla = k || x \), \( k' = k_0 / L_s \), \( L_s \simeq q R \), \( x = x - r_s \), \( k_0 = m / r_s \) and \( r_s \) is the radial position of the mode-rational surface defined by \( m - l q(r_s) = 0 \).

Using the continuity equation, the equations of motion (neglecting electron inertia), the energy equations, an equation for the impurity charge number dynamics, and the quasi-neutrality condition, we get the following equations:

\[ A \nabla_x^2 \phi - \frac{\omega_e}{\omega} \left( \frac{\nu_d}{c_s} \right) \phi - i \frac{S_e}{\omega} \hat{n} = - \left( - \frac{\xi^2}{\xi R} \left( \hat{\phi} - \hat{n} - 1.71 \hat{t}_e \right) + a_1 \xi \hat{t}_e - \frac{g_1(\xi)}{A} a_2 \xi \right) \]  

(1)

\[ A \hat{n} = A \nabla_x^2 \phi - \frac{\omega_e}{\omega} \left( \frac{\nu_d}{c_s} \right) \phi + \frac{\omega_s}{\omega} \phi + \frac{\xi^2}{A \xi_s^2} [\hat{n}(1 + \tau) + \tau \hat{t}_e + \hat{t}_e] \]  

(2)

\[ A \hat{t}_e = - \frac{\eta_0}{\omega} \frac{\omega_e}{\omega} \phi - \frac{2}{3} \frac{\xi^2}{A \xi_s^2} [\hat{n}(1 + \tau) + \tau \hat{t}_e] \]  

\[ = -i \frac{\xi^2}{\xi_{k||}^2} \hat{t}_e + \frac{1}{\xi_{k||}^2} \nabla_x^2 \hat{t}_e - i \gamma \hat{n} \hat{t}_e + i \gamma_T \hat{t}_e \]  

\[ - i \gamma_{el} (\hat{t}_e - \hat{t}_i) \left( 1 - 1.14 \left[ - \frac{\xi^2}{\xi_{k||}^2} \left( \hat{\phi} - \hat{n} - 1.71 \hat{t}_e \right) + a_1 \xi \hat{t}_e - \frac{g_1(\xi)}{A} a_2 \xi \right] \right) \]  

(3)

and

\[ A \hat{t}_i = - \frac{\eta_0}{\omega} \frac{\omega_e}{\omega} \phi - \frac{2}{3} \frac{\xi^2}{A \xi_s^2} [\hat{n}(1 + \tau) + \tau \hat{t}_i + \hat{t}_e] \]  

\[ = -i \frac{\xi^2}{\xi_{k||}^2} \hat{t}_i + \frac{1}{\xi_{k||}^2} \nabla_x^2 \hat{t}_i + i (\gamma_{el}/\tau) (\hat{t}_e - \hat{t}_i) \]  

(4)

where \( S_e = |S_e(\omega)| \exp(i\phi) \) is a phase sensitive feedback source in the continuity equation, \( \xi = \frac{x}{a_s} \), \( \nabla_x^2 = (\partial^2 / \partial x^2 + k^2 a_s^2) \), \( \xi_{k||} = \omega v_{el} / k^2 a_s^2 \), \( \xi_{k||}^2 = T_s / m_s \), \( g_1 = -i \xi_{k||} (A^{-1} \xi^2 - i \xi_{k||}^{-1}) \) takes account of impurity diffusion along field lines, \( \xi_{k||} = Z v_{el} / k^2 a_s^2 \), \( \omega = \omega - k_0 v_{el} \), \( A = 1 - \beta \xi_{k||} = k_0 v_{el} / \omega \), \( a_1 = 3 J_{||} k^2 a_s / 2 e N \omega \), \( a_2 = (3 N \omega / L_s)(\omega / L_s) \) (takes account of impurity gradient) \( L_s^{-1} = \left( \frac{a_s}{L_s} \ln Z \right) \), \( \omega_s = k_0 a_s / L_s \), \( \xi_{k||} = \omega / k^2 a_s^2 \), \( \xi_{k||} = 0.95 v_{el} / k^2 a_s^2 \), \( \xi_{k||} = 0.95 (3 / 2) N \omega a_s / \chi_{L||} \), \( \gamma_n = (2 / 3)(L / N T_e \omega) \), \( \gamma_T = -1 / \left( 3 A \right) \), \( \gamma_{el} = 2 \left( \frac{m_e e^2}{m_e v^2} \right) \), \( \xi_{k||} = \xi_{k||}^2 = 0.5 (N \omega / L_s)(a_s / k^2 a_s^2) \) and the definitions of \( \chi_{L||} \) and \( \chi_{L||} \) follow from Braginskii. Eqs. (1)-(4) may now be utilized to derive the eigenmode equations for the various instabilities in the plasma.
A. Rippling Modes

We consider the simplest derivation of rippling modes driven by impurity gradients alone where \( \tilde{t} \to 0 \) because of large parallel thermal conductivity and \( \xi_l^2/\xi_e^2 \ll 1 \) because impurity parallel diffusion is ignorable\(^7\). Taking \( |\omega| \gg \omega_0 \) so that pressure gradient effects may be neglected, eqn. (1) leads to the eigenmode equation

\[
\begin{bmatrix}
\frac{d^2}{d\xi^2} + \frac{\dot{b}}{1 - \beta \xi} - \frac{ic \xi^2}{4(1 - \beta \xi)} - \frac{a_2 \xi}{(1 - \beta \xi)^2}
\end{bmatrix} \phi = 0
\]

(5)

where we have neglected FLR terms, \( \dot{b} = -\frac{\omega_0}{\omega_c} \frac{\partial a_s}{\partial s} - iS_e \omega_0 / \omega^2 \), \( c = 4/\xi_R^2 \) and as before, we have used a linear expansion of the velocity around the mode-rational surface: \( v_0 = v_o(r_s) + v'_o(r - r_s) \).

In the absence of velocity shear \( (v'_o = 0) \) we take \( \beta = 0 \) and the eigen mode problem may be exactly solved to give

\[
(\frac{\gamma_{\text{Rip}}}{\gamma})^{5/2} - (\frac{\gamma_{\text{Rip}}}{\gamma})^{3/2} \left( \frac{\nu_{ei}}{\gamma_{\text{Rip}}} \frac{m_e \xi_e^2}{m_i \xi_i^2} \right)^{1/2} \frac{|S_e|}{\gamma_{\text{Rip}}} e^{-i(\pi/2 + \phi)} = 1
\]

(6)

where \( \gamma_{\text{Rip}} \) denotes the standard rippling mode growth rate in the absence of feedback sources. For \( \phi = \pi/2 \), the rippling mode growth is reduced, as may be seen by a perturbative calculation of \( \gamma \) as well as a limiting calculation in which \( |S_e| \) is so large that the unity of RHS may be neglected. When \( \beta \neq 0 \), we may solve the eigenvalue problem perturbatively to get the dispersion relation:

\[
(\frac{\gamma_{\text{Rip}}}{\gamma})^{5/2} = 1 - \theta^{1/4} - \frac{1}{\beta} \theta^{1/2} [2 + \delta^* + (\delta^* - 1)(\delta^* + 3)Z(\delta^*)]
\]

(7)

where \( \delta^* = 1 + \theta^{1/4}/\beta \), \( \theta = ic/4 \), and \( Z(\delta^*) \) is the plasma dispersion function with argument \( \delta^* \). The weak shear limit corresponds to \( \delta^* \to \infty \) and we can use the asymptotic form of the \( Z \) function to show that mode stabilization occurs irrespective of the sign of the shear.

B. Radiative Condensation Instability

For \( \nabla N = \nabla T = 0 \), \( k_0^2 C_s^2 / \omega^2 > 1 \), \( m_e \nu_{ei} / m_i \omega > 1 \), we find by adding the electron and ion energy equations (because electrons and ions are strongly coupled) and using the condition that \( \tilde{p} = \tilde{p}_e + \tilde{p}_i \equiv 0 \) the eigenmode equation

\[
\left[ \frac{\partial^2}{\partial \xi^2} - k_0^2 a_s^2 + \xi_{Li}^2 \left( \gamma_n + \gamma_T - 1.14 iS_e - \frac{\xi_e^2}{\xi_{Li}^2} \right) + 2i(1 - \beta \xi) \xi_{Li}^2 \right] \tilde{t}_e = 0
\]

(8)

This is a standard Hermite equation which yields the eigenvalue condition

\[
\omega_r = k_0 v_o(r_s)
\]

\[
\gamma = \frac{1}{3} \left[ \frac{2L}{NT} - \frac{1}{N} \frac{\partial L}{\partial T} \right] - \frac{1}{3} \left( \frac{k_0^2 L_s^2}{L_e^2} \chi_{Li} \chi_{le} \right)^{1/2}
\]

\[
- \frac{1}{3} k_0^2 \chi_{Li} - \frac{3}{4} \left( \frac{v'_o L_e^2}{\chi_{le}} \right) - 0.57 |S_e| \exp[i(\pi + \phi)]
\]

(9)
where $\chi_{\perp \perp}$ and $\chi_{\parallel e}$ are the heat diffusion coefficients from Braginskii. It may be noted that the velocity shear produces a stabilizing effect on the growth irrespective of the sign. The feedback term is found to be stabilizing for $\phi = -\pi$.

C. Drift-Dissipative Instability

It is well known that in the slab approximation the drift mode is stable in a resistive plasma in a sheared magnetic field. However, if the drift frequency $\omega_*$ has a parabolic profile such that $\omega_*(x) = \omega_*(1 - \xi^2 / L_*^2)$ where $L_* \ll \xi_*$ (the ion sound turning point), the instability is recovered. The basic physical reason for this effect is that in the former case, the asymptotic solution is a propagating sound wave which introduces shear damping, whereas in the latter case, the eigenmode is confined. We first write the eigenmode equation for drift waves with feedback terms but $\beta = 0$:

$$\frac{d^2 \phi}{d\xi^2} - \left[ 1 + k^2 - \frac{\omega_*}{\omega} \left( 1 - \frac{\xi^2}{L_*^2} \right) \right] \phi - \frac{i e_R^2}{\xi^2 - i e_R(1 - \tilde{S}_a)} \left[ 1 - \frac{\omega_*}{\omega} \right] \phi = 0$$

where $\tilde{S}_a = (i |S_\parallel| / \omega) \exp(i \phi)$ and we have assumed $\xi_R \ll \xi_*$. This equation may be solved by the method of matched asymptotic expansions and shows that the growth rate is multiplied by a factor $(1 - \tilde{S}_a)^{1/2}$. Thus for a $\phi$ such that $\tilde{S}_a > 0$ and real, there is a significant reduction of the growth rate using feedback methods. We now consider the case where $\tilde{S}_a = 0$ and $v'_0 \neq 0$ so that $\beta$ dependent terms are important. The new eigenmode equation is

$$\frac{d^2 \phi}{d\xi^2} - \left[ 1 + k^2 - \frac{\omega_*}{\omega} \left( 1 - \frac{\omega/2}{\omega_0 C_s} \right) \left( 1 - \frac{\xi^2}{L_*^2} \right) \right] \phi - \frac{i e_R^2(1 - \beta \xi)}{\xi^2 - i e_R(1 - \beta \xi)} \left[ 1 - \frac{\omega_*}{\omega(1 - \beta \xi)} \right] \phi = 0$$

For $\xi_* = \theta$ we get the eigenmode equation for the drift wave without resistive effects. If $\beta$ is small i.e. weak shear case such that $\beta^{-1} > L_*$, the mode is essentially confined by $\omega_*$ variation and the eigenmode has similar properties as the zero shear case. When $0 < \beta^{-1} < L_*$, we get the strong velocity shear case. In this case we find that the asymptotic behavior is like $\exp \left( \frac{2 i}{3 t} \xi^{3/2} \sqrt{\beta} L_*^{-1} \right)$, i.e., it is like a propagating wave. Again, in this case we get a shear damped mode which cannot then be destabilized by resistivity.

Conclusion

We have investigated the linear stability of rippling modes, radiative condensation instabilities, and the drift-dissipative mode in the presence of velocity shear. A feedback source term in the electron continuity equation, and conditions for stabilization of the modes, have been obtained. In weak shear, the stabilization is typically independent on the sign of $v'_0$. 100
References


5. ADITYA team, *ADITYA: Initial Results and Status of Edge Control Experiments* (enclosed paper).


ADITYA: INITIAL RESULTS AND STATUS OF EDGE CONTROL EXPERIMENTS

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Abstract

The initial results for ADITYA and status of edge control experiments are presented. Breakdown at low loop voltages ($\approx 15V$), peak currents up to 50 kA and discharge duration of 22 ms have been achieved at nominal operating parameters. The studies on the convective losses preceding the formation of rotational transform indicate that convection plays a more significant role than error magnetic fields in ADITYA. The current termination shows a sequence of positive spikes and decrease in plasma current until the external applied vertical magnetic field overcompensates the hoop force and causes a crashing of current channel towards inboard side of vacuum vessel.

1. ADITYA TOKAMAK

Tokamak ADITYA was commissioned in September 1989. The peak design parameters of ADITYA are: Toroidal magnetic field 15 kG; major radius 75 cm; plasma radius 25 cm; plasma current $\approx 250$ kA and the plasma duration $\approx 300$ ms [1]. Figure 1 is a schematic of the device. A Glow discharge cleaning system and various diagnostics - Rogowski Coil, Loop, Visible Spectrometer, Microwave interferometer, Soft X-ray detector array, Langmuir Probes, Normal and Grazing incidence spectrometers have been deployed. Charge Exchange diagnostics is expected to be operational soon. Bolometer, Thomson Scattering and Li fluorescence diagnostics are getting ready. The device is operated with a graphite limiter. The base vacuum after baking the vessel to 100°C and 50 hours of glow discharge cleaning was $10^{-8}$ Torr. Partial pressures of the impurities ($Z > 2$) were reduced to a value below $10^{-9}$ Torr in the background and during glow discharge cleaning below $10^{-7}$ Torr. For the initial experiments in ADITYA, the ohmic transformer was powered using a capacitor bank system which consists of a fast bank (408 $\mu$F, 10 kV) power crowbarred with a slow bank (21 mF, 5 kV). The vertical magnetic field was produced by a condenser bank system (0.7-5.6 mF, 5 kV) and the effect of applying vertical field on the current evolution was studied.

2. INITIAL RESULTS

The plasma breakdown was found to occur at a peak loop voltage of 15 volts. After optimisation of various parameters like the gas pressure, synchronisation of vertical magnetic field with current rise etc. a plasma duration of ≈ 21 ms with a peak current of ≈ 50 kA was obtained (Figure 2). The peak plasma current rise achieved was ≈ 20 MA/s. The rate of rise of vertical magnetic field determines both the duration and peak plasma attained during a shot. Over compensation at peak current leads to minor disruptions and ultimately uncontrolled sharp decay of plasma current. A chord averaged density of 5 × 10^{11} to 8 × 10^{12} cm^{-3} was measured by Microwave Interferometer. An edge plasma temperature of 10 eV was measured by Langmuir probe and the core temperature was estimated to be 25 - 50 eV using spectroscopy techniques.

3. PLASMA BREAKDOWN

The breakdown studies have been made in the pressure range of 2 × 10^{-5} to 3 × 10^{-4} Torr. The experiments show that successful discharges can be obtained for loop voltage of 10 to 27 V with initial gas fill pressure of 2.5 × 10^{-5} to 3 × 10^{-4} Torr. The minimum loop voltage required for a successful breakdown increases with pressure for the fill pressure exceeding 5 × 10^{-5} Torr (Figure 3). The delay between the applied ohmic field and the breakdown is also minimum at minimum breakdown voltage. The maximum plasma current is obtained at a fill pressure of approximately 5 × 10^{-5} Torr for all loop voltages; the peak current increases with increasing charging voltage. The volt-seconds consumed for breakdown is minimum for fill gas pressure in the range of 5 × 10^{-5} Torr and it does not vary significantly upto 2 × 10^{-4} Torr. For fill pressures below 5 × 10^{-5} Torr, the volt-seconds consumption
in breakdown increases with increasing loop voltage. The peak plasma current obtained without $B_v$ is 6.5 kA. The breakdown studies have been made with externally applied vertical magnetic fields also. The observed results in ADITYA present the following scenario for breakdown. The build-up of the discharge is decided by production by electron-neutral collisions and loss predominantly due to magnetic field errors [2]. An optimum regime for pressure exists where the avalanche is favoured. At very low pressures, the electrons get accelerated without undergoing sufficient ionising collisions which does not lead to an avalanche. At higher pressures, the electrons do not get much energy due to small collision mean free path which is detrimental to avalanche.
Figure 3: Plots of the breakdown voltage and time delay versus the neutral pressure. In top left plot X and Y indicate unsuccessful and successful breakdowns respectively. For other plots the curves correspond to peak loop voltages: + : 27V; : 23V; Y : 20V; X : 18V; * : 16V

4. CONVECTIVE LOSSES DURING INITIAL PHASE

The loss of electrons during gas breakdown and current initiation in a tokamak plays an important role in determining the success of discharge, formation of rotational transform and ultimately in achieving a maximum of plasma current for a given loop voltage [3]. It is known that during the formation of rotational transform, the losses of electrons take place by convection even though the exact mechanism is not well understood. A set of four carbon-tip single Langmuir probes has been placed at four locations on the poloidal cross-section protruding within the limiter by 4 mm. The probes were operated for measuring floating potential and ion saturation current.

The vertical and horizontal electric fields are measured by recording the floating potentials of top-, bottom-, in- and out- probes and taking differences appropriately (Figure 4 & Figure 5). The vertical electrical field is found to reverse with the direction of the toroidal magnetic field. The value of applied $B_v$ at the time of peak plasma current of 7.2 kA and 11 kA are 20 G and 58 G respectively. Within this range the vertical electric field decreases for higher values of $B_v$. When the direction of $B_v$ is reversed, the discharge fails even for $B_v = 8$ G. It is possible to estimate the value of equivalent error field ($B_z$) from the observed loss rate and attributing it to $B_z$. The estimated value for ADITYA is 5 to 10 G and is from top to bottom for the case when $E_z$ is from bottom to top and both $B_T$ and $I_p$ are in clockwise direction. The consequent convective loss is radially outward, i.e. in the same direction as the loss due to hoop force. If $B_h$ is the equivalent magnetic (error) field corresponding to hoop force, then $B_z + B_h$ determines the fate of the discharge in the absence of externally applied vertical magnetic field. If the poloidal magnetic field $B_p$ associated with the
plasma current exceeds $B_x + B_h$, the rotational transform is established. If $B_x + B_h$ is larger than $B_p$, the discharge may terminate or continue without rotational transform being established depending on whether the difference between the two is greater or smaller than the critical error field necessary to terminate the discharge.
5. CURRENT TERMINATION IN LOW-Q DISCHARGES

A typical signature of current termination is a sharp positive spike in $I_p$ followed by a negative spike in $V_l$. From the analysis of about 30 shots in ADITYA, with peak plasma current in the range of 30-50 kA, it is seen that a sequence of positive spike in current and negative spike in loop voltage is repeated 2 to 4 times each separated by about 1 ms (Figure 6). The delay between the positive current spike and negative voltage spike is 100-300 μs. Hard X-ray burst coincident with current spikes are also observed. The appearance of a positive spike in current is an indication of the flattening of the current profile [4]. The subsequent appearance of the negative spike in loop voltage is an indication of ejection of poloidal flux outside the plasma. ADITYA current termination shows a sequence of flattening of the profile. The electric field measurement indicates the movement of plasma column towards the outboard side of the vessel during this phase. An increase in optical emission is also observed at this stage. The plasma resistivity increases as indicated by the rise in the mean level of the loop voltage. The plasma current decreases until the externally applied vertical magnetic field overcompensates the hoop force and causes crashing of current channel towards the inboard side of the vacuum vessel. This is inferred from the observation of reversal and increase in the vertical electric field. The ejection of the poloidal flux in a time scale of 100-300 μs requires the plasma temperature to be 2-5 eV. Since the plasma current remains high prior to the crash, the core temperature might not have fallen to such low values.

A possible alternative scenario is as follows: At plasma currents of 30-50 kA, the value of $q$ at limiter is 3-2. Hence in these discharges, the island width at $q = 2$ surface grows and touches the limiter, leading to loss of current and consequent redistribution of current profile. The $q=2$ surface, then, moves in and the growth of the island restarts. This sequence of events

![Figure 6: Typical profiles of plasma current, loop voltage, Hard X-Ray, and vertical electric fields as a function of time during a low-q discharge.](image)
is repeated till the plasma current channel crashes due to overcompensation of hoop force by vertical magnetic field. Further experiments are in progress to understand this current termination in ADITYA.

6. EDGE MODIFICATION EXPERIMENTS

Experiments in various tokamaks measuring the edge density, temperature and their fluctuations have clearly shown that these fluctuations play an important role in determining the particle and in some cases heat transport. The impurity generation and impurity transport is also closely linked to this. The influence of the edge on the core plasma parameters have been well brought out by the experiments which correlated the high confinement discharges with the low edge transport. The edge control and modification experiments thus achieve an importance in the tokamak plasma optimisation.

Detailed edge fluctuation studies are being planned using arrays of Langmuir probes. Using the fluctuation measurements, a feedback experiment to modify the edge plasma (and consequently the core plasma) has been planned. ADITYA poloidal limiter has fourteen discrete carbon tiles part of which will be used for detecting the signals and derive the dominant modes in the fluctuation spectrum and the signal after band-pass filtering will be applied using the remaining tiles with controlled phase shift and amplification for feedback control. Use of such a method has been found to reduce the core fluctuations in a mirror [5].

Radial currents driven by negatively biased electrodes placed within the tokamak plasma column have been found to induce plasma rotation and subsequently induce a transition from L to H mode [6]. An experiment is planned in ADITYA where the role of the radial current could be studied in a controlled way. The current injection scheme involves a two-electrode structure immersed in the tokamak plasma leading to an ‘ion extraction’ like region. The diagnostics for this experiment include a pair of rake probes, energy analysers for poloidal ion velocity measurement, Doppler shift measurement, edge fluctuations besides the routine diagnostics. In Divertor tokamaks [7], the establishment of radial electric field is attributed to the loss of ions at the edge as evidenced by the electric field layer width equal to the ion Larmor radius. In ADITYA, the radial electric field generation in the plasma is also being planned to be achieved by producing large Larmor radius impurity ions using radio-frequency heating at the edge near the ion cyclotron range of frequencies and expelling them from the plasma.

REFERENCES

EDGE STUDIES, HELICITY CURRENT DRIVE AND Diagnostic Development on the Phaedrus-T Tokamak


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Abstract

Phaedrus-T is a moderate sized tokamak with low toroidal field and substantial radio frequency power capabilities. This machine has only recently become operational and full RF power is expected in the next year. The initial studies include edge modifications due to variable phased antenna structures and RF helicity current drive experiments. The tokamak is also being used as a test site for diagnostic development.

MACHINE DESCRIPTION

The basic design parameters of the Phaedrus-T device are as follows: \( R = .93 \) m, \( a = .26 \) m, \( B_t = .8-1.2 \) T, \( I_p = 60-140 \) kA, \( <n_e> = .5 - 3 \times 10^{13} \) cm\(^{-3}\), \( T_e = .4-.9 \) keV. At the time of this report, we have achieved \( .9 \) T, \( 100 \) kA, \( 1-2 \times 10^{13} \) cm\(^{-3}\) operation with discharge current flattops of 70 msec out of a 150 msec total duration. The three independent poloidal field coil-power supply systems (ohmic, vertical, and horizontal) are fully programmable. Variable wound Rogowski loops, sine and cosine coils, are used to measure the centroid position of the current channel and to provide feedback signals for the poloidal field power supplies. The vacuum chamber, originally a part of the ISX-A device [1], was modified so that large access ports would be available for multiple RF antennas and diagnostics. The plasma is circular in poloidal cross section with a single poloidal limiter composed of graphite blocks. The limiter can be independently baked to 300°C. The RF transmitter will be capable of 2 MW, 80 msec pulse length operation. At present, we are not using the final stage of the amplifier chain and hence have 300 kW available. The present installed antenna is a two strap low field side antenna system enclosed in a single stainless steel Faraday shield. The two straps are fed independently and a tuning element is present to minimize the mutual inductance between the two antennas and allow

1 On leave from Southwestern Institute of Physics, Leshan, Sichuan, China.
operation at continuously variable phasing. The antenna and Faraday shield system has been designed to have excellent access for probes and optical viewing. The plasma diagnostic system includes a 30-40 keV, 2-5 amp, 10 ms pulse length neutral beam - the FIDE [2] beam system. Excellent viewing access is available for viewing emission lines generated by the beam - plasma interaction. Figure 1 is a top view of the tokamak and shows the location of the RF antenna, the diagnostic beam and other major diagnostics.

![Figure 1. Top View of Phaedrus-T. The RF Beam Viewing Antenna is at Box #3; Diagnostic Neutral Beam, #5; Thomson Scattering, #8; VUV-Spred System, #1; Probes at Port Box #2, #3, #4, and #9; the Poloidal Limiter; #7. Microwave Interferometers, #7; Charge Exchange Analyzer, #9; Poloidal Mirnov Coil Array, #6; Soft-X-Ray Array, #1; Z_{eff} Array, #4.](image)

**OHMIC RESULTS AND RUNNING CONDITIONS**

Present ohmically generated plasmas are summarized in Figure 2. The energy confinement as measured from the equilibrium fields and ohmic input power is about 10 ms and is consistent with the ISX-A data [1]. Our best running conditions are achieved after glow discharge cleaning in Helium: \( I_d = 1.2 \) amps, \( V_d = 300 \) V, and \( p(He) = 1-2 \) mTorr, for 1/2 to 1 hour and concurrent baking of the graphite limiters to 300 C. The glow bombards the walls and graphite limiters with 100-300 eV helium ions. The ensuing tokamak discharges have reduced oxygen emission lines by a factor of 3-10, reduced loop voltages by .2-.5 volts, and reduced hydrogen or deuterium wall recycling by a factor of 2-10.
Figure 2. Time history of plasma parameters during a typical discharge; I, the plasma current; U, Loop voltage; $\bar{n}_e$, the line average density; $Q_{gp}$, the equivalent atomic amp of deuterium gas puffed into the device.

The water levels in the background neutral gas measured two minutes after a discharge are reduced by a factor of 10. After 20-60 discharges, the glow must be repeated to retain good discharges. The number of discharges between glows depends on the amount of gas puffing and the ultimate plasma density achieved. This operation is consistent with results from large tokamaks that have extensive graphite surfaces exposed to the plasma. For the discharge shown in Figure 2, 20-30 successive, repeatable shots without a disruption have been obtained, although the gas puffing is reduced by a factor of three towards the end of the run. For line averaged densities of $1.5 - 2 \times 10^{13}$ cm$^{-3}$, only 2-5 successive discharges can be obtained without heavy Mirnov activity appearing and disruptions occurring. When this occurs, 1/2 hour of helium glow discharge cleaning is necessary to set up optimum conditions again. The first two shots after helium GDC show the presence of helium in line emission diagnostics but after these shots the He signals disappear into the noise.

**RF RESULTS**

At the present time, we have reached 100 kW, 20 ms pulse length operation with the RF systems and are in the process of conditioning the antennas with plasma to reach 200-300 kW. At the 100 kW level, we have observed results similar to results obtained on TEXTOR [3] before that tokamak's inner surfaces were carbonized. For 100 kW, 20 ms pulses, density increases of $.5 \times 10^{13}$ cm$^{-3}$ have been observed along with loop voltage increases of .2 to .4 volts, and increased emission lines of both metals and light impurities. Probe measurements show that the edge plasma temperature and density is increasing with RF power. Floating potential measurements using a capacitive probe for the same poloidal location as the center of the antenna are always more positive when the RF is turned on. This indicates electrons are being driven out of the plasma either by direct
heating or an RF sheath effect. Edge effects can dominate RF experiments and so it will be necessary to understand the edge phenomena consistent with good RF helicity current drive. Since the Faraday shield is constructed out of stainless steel and the limiter, out of graphite, we intend to use the VUV emission line diagnostic, SPRED [4], to help us unravel the role each plays in injecting impurities into the plasma at various antenna phasings. So far the results are ambiguous; both types of emission increase. We are considering boronization/carbonization of the Faraday shield and inner vacuum surfaces of the tokamak to reduce impurity levels in a way similar to TEXTOR [3].

HELICITY CURRENT DRIVE PLANS

RF Helicity current drive is an idea born from the constant helicity arguments used to explain the RFP dynamo effect. The argument for RF helicity current drive goes as follows [5]: Circularly polarized RF waves contain helicity and as the waves are dissipated in the plasma or as the wave vector changes (e.g. at an Alfven resonance), they impart this helicity to the plasma in the form of a steady-state current. The drive is then not limited by density except in the generation and the dissipation of the wave fields. Both circularly polarized Alfven waves and ICRF fast waves contain helicity and theoretically do not have a density limit. There is a difference, however, the AW helicity current drive is in the same direction as the momentum drive for that wave while it is opposite for the FW. We have used the antenna code FASTWA [6] to show that the present two strap antenna can excite a polarized FW with helicity if the phasing between the straps is 90 degrees. Our first experiments on RF helicity current drive will be in this configuration. We will then lower the frequency and excite AW with the same antenna structure. The theory for helicity current drive shows that the efficiency increases with reduced wave number because the wave helicity density is proportional to $l/k_w$. For that reason we intend to install a second antenna system with a quasi-helical structure.

ADVANCED DIAGNOSTICS WITH NEUTRAL BEAMS

In order to measure helicity current drive we intend to measure the polarization of the hydrogen neutral emission lines due to the interaction of a neutral beam with the plasma in the tokamak's magnetic field. The emission lines are polarized due to the Motional Stark Effect [6]. Along with this diagnostic, several variations of a Beam Emission Spectroscopy (BES) measurement of density fluctuations will be tested. In particular, the Phaedrus-T BES diagnostic allows very high spatial resolution and measurements using either an $H^0$ or $He^0$ beam. In addition vacuum ultraviolet observations at Lyman-$\alpha$ are possible. A more sensitive BES system is necessary to reduce the beam requirements and make the diagnostic viable on the next generation tokamaks. The neutral beam provides the capability of measuring many other parameters in the plasma. Calculations suggest that RF
field induced Stark shifts measured from the beam emission lines can be used to measure the RF wave fields. If a helium neutral beam is injected, the ratio of the line emission can be used to measure the electron temperature.

ACKNOWLEDGEMENT

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POTENTIAL FORMATION IN A TOKAMAK PLASMA
BY LIMITER BIASING AND ITS EFFECT ON
CONFINEMENT PROPERTIES

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Abstract

It is observed in HYBTK-II tokamak that a negative change of the electric potential by negative limiter biasing is dependent sensitively on the magnetic configuration in the scrape off region while the positive change is not. In the positive biasing the plasma potential increases in proportion to the bias voltage and enhances the radial electric field at the edge. On the other hand the plasma potential does not decrease deeply in the negative biasing. The positive radial electric field enhanced by the positive biasing seems to improve the particle confinement associated with a stabilization of the magnetic and electrostatic fluctuations.

1. Introduction

It has been found from the experimental[1],[2] and theoretical studies[3],[4] that the edge electric field has an important role in the plasma confinement properties. It is considered that the onset of the transition from the L mode to the H mode in tokamaks may be triggered by a sudden change of the radial electric field $E_T$, which subsequently suppresses the turbulent fluctuations and improves the plasma confinement in the H mode.

Active control of the plasma potential, such as limiter biasing and electron beam injection, has a potential to improve the plasma confinement by modifying $E_T$. However, the physical mechanism for the potential formation in the toroidal configuration with closed magnetic surface has not been understood satisfactory. In the present study the potential formation in the edge region by limiter biasing is investigated experimentally and the effects of the limiter biasing on the tokamak plasma properties are also presented.
2. Experimental Setup

Limiter bias experiment was carried out on a small research tokamak HYBTOK-II. The bias voltage from a charged capacitor is introduced to a poloidal circular limiter made of molybdenum. A movable Langmuir probe on the equatorial plane and a poloidal array of Langmuir probe are employed to obtain radial and poloidal profiles of the edge plasma parameters. Typical discharge parameters of HYBTOK-II are summarized in Table-I.

<table>
<thead>
<tr>
<th>Table-I Discharge parameter of HYBTOK-II tokamak</th>
</tr>
</thead>
<tbody>
<tr>
<td>Major Radius ( R = 40 \text{ cm} )</td>
</tr>
<tr>
<td>Minor Radius ( a = 11 \text{ cm} )</td>
</tr>
<tr>
<td>Plasma Current ( I_p \leq 15 \text{ kA} )</td>
</tr>
<tr>
<td>Density ( \bar{n}_e \leq 1.5 \times 10^{19} \text{ m}^{-3} )</td>
</tr>
<tr>
<td>Toroidal Field ( B_t \leq 0.5 \text{ T} )</td>
</tr>
<tr>
<td>Discharge Duration ( \tau \leq 12 \text{ ms} )</td>
</tr>
<tr>
<td>Bias Voltage (</td>
</tr>
</tbody>
</table>

3. Experimental Results

3.1 Potential Formation

Typical time traces of discharges parameters of the limiter biasing are shown in Figs. 1 for positive(a) and negative(b) bias voltage with respect to the vessel wall. The bias voltage is turned on at \( t=5 \text{ ms} \). The floating potential \( V_f \) of the Langmuir probe at the edge shows a prompt increase for positive biasing, while \( V_f \) changes in time for negative biasing. This time scale is much longer than the transit time for protons with \( T_j=10 \text{ eV} \) at the edge to circulate the torus(\( t \sim 60 \mu \text{s} \)). Here, it should be noted that the change of \( V_f \) corresponds well to that of the plasma potential \( V_s \) since the electron temperature at the edge is not affected significantly by limiter biasing and \( V_f = V_s - 3T_e \). The temporal change of \( V_f \) in the negative biasing is strongly correlated to the plasma position. In Fig. 2 \( V_f \) measured with the poloidal Langmuir probe array in the negative biasing is compared with that without biasing as a function of the horizontal plasma shift. The top probe shows that the negative change is maximum when the plasma column is located at the center. The outside probe shows a maximum change when the plasma shift is outward while in the inside probe the negative change becomes large with inward shift. The ion current flowing to the wall is much larger...
Fig. 1 Typical time traces of the plasma current, limiter voltage, limiter current, floating potential, average electron density, and $H_\alpha$ intensities from the wall and the limiter for the positive(a) and negative(b) biasing.

Fig. 2 Floating potential as a function of the plasma position with and without biasing. $V_b=-250\,V$
than the theoretically estimated value based on the friction induced ion current. From these observations it is predicted that a direct connection by magnetic field lines between the limiter and the wall plays an important role for the formation of the electric potential by limiter biasing. The physical mechanism of the potential formation in the scrape off layer has been discussed in Ref.[5]. The observed results on the plasma potential agree with the theoretical calculations qualitatively.

Radial profiles of $V_s$ and $V_f$ at the outside edge are shown in Figs. 3(a),(b). The positive biasing brings up the plasma potential efficiently, but the negative biasing slightly decreases $V_s$. The radial electric field calculated from the measured $V_s$ is shown in Fig. 4. The positive biasing enhances $E_r$ drastically and the negative biasing does not change the sign of $E_r$ in the present bias experiment.

![Graph showing radial profiles of floating potential and space potential](image1)

**Fig. 3** Radial profiles of the floating potential and space potential, taking the bias voltage as a parameter.

![Graph showing radial profile of radial electric field](image2)

**Fig. 4** Radial profile of the radial electric field obtained from the space potential profile shown in Fig. 3.
3.2 Effect of Limiter Biasing on Tokamak Discharge

For both biasing polarities $n_e$ increases and $H_\alpha$ intensity looking at the wall decreases. The drop of $H_\alpha$ in the positive biasing is much larger than that in the negative biasing. The fact that $n_e$ increases with a decrease of $H_\alpha$ intensity might be related to an improvement of the particle confinement. In addition to $n_e$ and $H_\alpha$ the impurity influx from the limiter and the wall may affect the particle confinement properties. As mentioned before, the theoretical studies indicate that the improvement of the confinement is caused by the stabilization of turbulent fluctuations. The positive biasing effectively suppresses both the magnetic and electrostatic fluctuations. In Fig.5 the effective stabilization of the magnetic fluctuation is shown. The radial profile of the decrease of the electrostatic fluctuation intensity normalized by that without biasing is shown in Fig.6. The electrostatic fluctuations are stabilized effectively beyond the limiter face where the $E_r$ is enhanced drastically by positive biasing. In the negative biasing the suppression of $\tilde{B}_p$ and $\tilde{V}_f$ is also observed but the decrease of their intensities is much smaller that that in the positive biasing. This difference should be related to the difference of the modified $E_r$ that the modification of $E_r$ in the positive biasing is much larger than that in the negative one.

Fig. 5 Time traces of the magnetic fluctuation picked up by a magnetic probe without(a) and with biasing of 50 V(b).

Fig. 6 Radial profile of the decrease of the electrostatic fluctuation normalized by that without biasing. The bias voltage is 50 V.
4. Conclusion

The radial electric field strongly enhanced by positive biasing might be related to the improvement of the particle confinement associated with the stabilization of the magnetic and electric fluctuations. The effect of the negative biasing on the plasma confinement properties is not clear since we could not make the negative plasma potential by the present cold limiter.

Acknowledgement

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References

LHCD EXPERIMENT ON THE HT-6B TOKAMAK

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Abstract

LHCD experiment was performed on the HT-6B tokamak recently. The rf system consisted of a 100kW, 2.45GHz power source and a multijunction antenna with relatively narrow power spectrum. At relatively high density, \( \eta \sim 0.2 \left( 10^{19} \text{M}^{-2} \text{kA/kW} \right) \) was achieved. The MHD modes and sawteeth were strongly affected by LHW.

1. INTRODUCTION

Low Hybrid Current Drive (LHCD) has been proved an effective way for driving current in tokamak plasma. Combining LHCD with RF heating seems one of the promising options for a steady state fusion power plant. We are carrying out a project for exploration of the physics in that regime. A 100kW rf system is recently assembled on the HT-6B tokamak for LHCD experiment. The data from that system will be introduced to a 500 kW rf system on the HT-6M tokamak. The leading subjects are the different optimising operational conditions for LHCD and ICRF and some new physics concerning the wave distorted velocity distributions of electrons and ions.

2. EXPERIMENTAL ARRANGEMENT

The rf system on HT-6B consisted of two main parts, a microwave source and a set of transmitter including a multi-junction grill type antenna. Figure 1 gives a brief view of the rf alignment.

The rf source was made of a NG 4883 magnetron and its power supply. At present this kind of magnetron can operate at the following parameters:

- anode voltage: 20 kV
- anode current: 20 A
- magnetic field: 2 kG
- total efficiency: 50%
- power output: 120 kW
- pulse duration: 60 ms
The parameters of power supply for megnetron are

- output voltage: 25 kV
- output current: 23 A
- pulse length: 100 ms
- repetition period: 2-4 minutes

A multi-junction grill was used as the slow wave antenna. According to the parameter of HT-6B tokamak plasma we took 1*8 subwaveguide structure. The dimension of each subwaveguide was 109.2*8.0 mm, with 1.5 mm thickness of the wall. Thus gave the period of grill 9.5 mm. The phase shift $\Delta \Phi$ between two adjacent subwaveguide was $\pi/2$. The calculated $N_/$ spectrum is shown in Fig. 2. It shows that the maximum energy peak at $<N_/>=3.22$ and the width of $N_/$ spectrum is 2.0.
The antenna was connected to the magnetron by standard waveguide of rectangular cross section and high power microwave components, such as isolater, directional coupler, DC breaker, flexible waveguide and ceramic sealed window.

The output rf power and pulse length were adjustable, but the phase shift was fixed.

3. EXPERIMENTAL RESULTS

The HT-6B Tokamak in which the LHCD experiment was conducted has the following parameters:

- $R = 45 \text{ cm}$
- $a = 12.5 \text{ cm}$
- $I_p = 20-40 \text{ kA}$
- $B_t = 6-9 \text{ kG}$
- $N_e = 0.4-2.0 \times 10^{13} \text{ cm}^{-3}$
- $T_e = 150-250 \text{ eV}$

The Diagnostics used in LHCD experiment is listed below:

- Measurement of $I_p$, $V_l$, Displacement, $N_e$
- Soft X-ray Diode Arrays
- Mirnov Probes
- Langmuir Probes
- ECE (Multi-channel)
- Hard X-ray Spectrum
- Spectroscopy (UV, VUV, Doppler Broadening)
- Bolometer
- etc.
In tokamak discharges with LHCD and inductive electrical field, the plasma properties were investigated in some aspects. Figure 3 gives a typical discharge with LHW.

When LHW launched in the existing ohmic plasma the loop voltage dropped. As $V_l$ dropped to zero, the equivalent $I_{\text{rf}}=23\text{kA}$ was obtained. The driven efficiency reached $\eta_{\text{cd}}=0.2 \left(10^{19} \text{m}^{-2}\text{kA/kW}\right)$ with the dielectric parameter $\varepsilon = \omega^2_{\text{pe}(0)}/\omega^2_{\text{ce}(0)}$ in the range of 1.6-2.1. $\eta_{\text{cd}}$ decreased with the decrease of $B_t$ and increase of $N_e$. The current drive effect could be found clearly until $\varepsilon$ up to 4. [Fig. 4]

![Fig. 3 A typical LHCD discharge](image-url)
Dependence of the electron velocity distributions on the direction of launching traveling waves were observed. In parallel launch discharges (Vf anti-parallel to Eohmic), the velocity of the driven electrons was limited lower than 40 keV which was equivalent to the phase velocity of LHW. But in counter launch case (by changing the phase shifts between wave guides), the electron velocities could be accelerated to much higher value. The energetic electrons with energy up to 1 MeV made the velocity distribution rather flat and often caused runaway discharges.[Fig.5]

The profile of driven current depended on the plasma density. When Ne<1x10^{13} cm^{-3} the driven current distributed mainly within 6 cm [Fig.6a]. When Ne>1.5x10^{13} cm^{-3} rf power deposited in outer region.[Fig.6b] The phenomena meant that the highest driven efficiency existed in a proper density window, neither in high density (cut off), nor in low density.
Heating effect could be observed from the rise of $T_i$. Typically, $T_i$ increment was more than 20 eV which is shown in the table below.

<table>
<thead>
<tr>
<th>shot no.</th>
<th>$P_{\text{LHCD}}$ (kW)</th>
<th>$T_i$ before LHCD (eV)</th>
<th>during LHCD (eV)</th>
<th>after LHCD (eV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>2D75</td>
<td>40</td>
<td>31</td>
<td>66</td>
<td>24</td>
</tr>
<tr>
<td>2D45</td>
<td>60</td>
<td>35</td>
<td>57</td>
<td>40</td>
</tr>
<tr>
<td>4D252</td>
<td>70</td>
<td>41</td>
<td>71</td>
<td>38</td>
</tr>
</tbody>
</table>

The MHD instability property varied by LHCD and also depended on plasma density. Fig. 7 showed about 15 kHz Mirnov-like oscillation occurred on SX signals (detected by surface barrier detector arrays) when $N_e < 1 \times 10^{13} \text{cm}^{-3}$. 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 $N_e > 1.5 \times 10^{13} \text{cm}^{-3}$ Mirnov-like oscillations were suppressed, and sawteeth changed in two ways. Figure 8a showed the sawtooth repetition period increased from 250 $\mu$s to 400 $\mu$s during LHCD. the amplitudes increased too. But it didn't propagated to outer region ($r > 6 \text{ cm}$), quite different from the large sawteeth in Ohmic plasma and RHF experiment. Fig.8b showed another effect by which the sawteeth were suppressed during LHCD and recovered or even amplified right after LHW stopped. Those two effects occurred occasionally, in other words, we can not identify the difference of the two discharge conditions. According to the theory, the energetic particles from LHCD and change of $J(r)$ profiles are two possible reasons for those phenomena. Further experiment is undergoing.

Fig. 7 15kHz Mirnov oscillations occur on SK signals during LHCD

Fig. 8a Sawteeth amplified by LHCD
Fig. 8b Sawteeth suppressed by LHCD
RADIAL PARTICLE TRANSPORT DUE TO THE TURBULENT AND QUASI-STATIONARY CROSS FIELD DRIFT ON THE CASTOR TOKAMAK

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Abstract

The relative importance of the turbulent and quasi-stationary cross field drift induced particle transport and their poloidal asymmetries are experimentally investigated at low density discharges on the CASTOR tokamak.

INTRODUCTION

Recent experiments on TEXT tokamak /1/ identified the electrostatic turbulence at the plasma edge as a candidate for inducing anomalous particle transport at standard ohmic heating (OH) regimes of tokamak operation. The link between the global particle confinement and the level of density and poloidal electric field fluctuations has been established experimentally also at combined inductive and lower hybrid current drive (OH/LHCD) regimes on CASTOR /2/ and ASDEX /3/ tokamaks. It was found for $P_{RF} \ll P_{OH}$ that the global particle confinement improves while the cross-field drift turbulent flux $\Gamma_{turb}$ at separatrix

$$\Gamma_{turb} (a) = \langle \delta n_e \cdot \delta v_r \rangle = \frac{1}{B_{TOR}} \langle \delta E_p \cdot \delta n_e \rangle$$

(1)

decreases noticeably. Here $n_e$, $v_r$, $E_p$ are the electron density, radial cross-field drift velocity and poloidal electric field at the separatrix, symbol "\(\delta\)" denotes the fluctuating part of the corresponding quantity and the bracket $\langle \rangle$ means an averaging over a time interval sufficiently longer than the characteristic period of fluctuations.
To make possible a quantitative comparison of the turbulent data with those deduced from the global particle balance, however, two conditions should be fulfilled:
1) the value of the turbulent flux $f_{\text{turb}}$, generally measured by a multiple Langmuir probe at a single point in the region of the separatrix, should be representative for the whole surface of the plasma torus;
2) the quasistationary value of the poloidal electric field $E_p$ must be so low that the quasistationary cross-field drift-induced flux $\langle \gamma \rangle$

$$\langle \gamma \rangle = \frac{1}{B_{\text{TOR}}} \langle n_e \rangle \langle E_p \rangle$$

(2)
can be neglected with respect to the turbulent flux $\int f_{\text{turb}}$

This report is devoted to the experimental test of this two conditions fulfilment on the CASTOR tokamak.

EXPERIMENT

The experiment was performed on the CASTOR tokamak (the major radius $R=40$ cm, the radius of the liner $b=10$ cm, of one aperture limiter $a=8.5$ cm and of two top/bottom movable rail limiters at $r_l=7.9$ cm). The low density ($n_e \leq 10^{13}$ cm$^{-3}$) OH regimes with the feed-back stabilized position of the plasma column inside the vacuum vessel were operated.

The parameters of the edge plasma are measured by means of an array of triple Langmuir probes, uniformly distributed in one poloidal plane of the torus at the radius $r=7.9$ cm (see Fig. 1).

Fig. 1. Arrangement of the array of the eight triple probes (denoted as A - H) in the vacuum vessel of the CASTOR tokamak.
The tips of each probe form an equilateral triangle with sides $d = 3$ mm. The tips are from molybdenum wire 0.5 mm in diameter, 3 mm being the active length of each tip. The two tips (denoted as 2 and 3), separated in poloidal direction, are on the floating potential with the respect to the liner. The third tip is negatively biased and serves for determination of the ion saturation current. Simultaneously, the density and poloidal electric field fluctuations (assuming $\mathcal{J}_e \approx 0$) are monitored by means of an analog correlator. This correlator allows to determine directly the turbulent flux $\mathcal{F}_{turb}$ as well. The details of the edge electrostatic turbulence measurement have already been described elsewhere /4/.

**EXPERIMENTAL RESULTS**

Fig. 2 shows the temporal evolution of plasma parameters measured at two different poloidal angles.

![Temporal evolutions of two shots with the identical global plasma parameters, but with the different triple probes in operation.](image)
The global characteristics (the loop voltage $U_{\text{loop}}$, the plasma current $I_p$, the line-average density $n_e$ and the intensity of $H_\alpha$-spectral line $I_{H_\alpha}$) are nearly identical, while the quasistationary edge parameters (the ion saturation current $\langle I_s \rangle$, the floating potential $\langle U_f \rangle$ and the difference $\langle \Delta U_f \rangle = \langle U_{f2} - U_{f1} \rangle$) differ from each other noticeably. Note namely that the signals of the probe at position D differ from those of the probe G in the polarity of the quantity $\langle U_f \rangle$ and in the response on the anomalous Doppler instability at $t=17$ ms. This kinetic instability, manifested by the positive spike on the loop voltage, is a general feature of low density discharges in the CASTOR tokamak.

Fig. 3 presents the poloidal dependence of the edge parameters at $t=15$ ms. The dependences of RMS-value of the ion saturation current fluctuations $\langle I_s \rangle$ and the RMS-value of the poloidal electric field fluctuations $\epsilon_p$ are given as well. The turbulent flux $\Gamma_{\text{turb}}$ is not included as its poloidal variation is similar to the variation of $\langle \Delta U_f \rangle$.

Fig. 3. Poloidal variation of plasma parameters at the plasma edge.
DUSCUSSION AND CONCLUSION

1) It follows at first from the above given results that the experimental data measured at only one point at the plasma edge can be hardly taken as representative for the whole surface of the plasma. The Fig. 3 indicates that a strong poloidal asymmetry of the plasma edge exists in CASTOR tokamak.

2) A surprisingly high value of the difference of the potentials $\Delta U_f$ between two floating tips is observed at some poloidal angles. The $\Delta U_f$ can even change the sign as documented in Figs. 2 and 3. The measured difference in $U_f$ can be interpreted in two ways:

a) If the floating tips are on such near-by magnetic surfaces that

$$T_{e1} - T_{e2} \ll (1/3)(U_{f1} - U_{f2})$$

(for maxwellian hydrogen plasma), the high measured value of the $\Delta U_f$ will imply the high value of poloidal electric field. Such field can induce a poloidally localized quasistationary radial flux of particles. If this flux is directed outward ($\Delta U_f > 0$ in our case), the region with an enhanced recycling would be created in this way.

b) If relation opposite to (3) is valid, the measured value of $\Delta U_f$ indicates predominantly a noncircularity of magnetic surfaces at the plasma edge.

At this phase of experiment we can not distinguish which possibility is more realistic in our case without an independent determination of the form of magnetic surfaces. High values of $\Delta U_f$ has been measured on the TF-1 /5/ and during ICRH on JET /6/ tokamaks. An existence of the quasistationary poloidal electric field has been concluded in these cases.

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PHENOMENA CONNECTED WITH THE RADIAL ELECTRIC FIELD IN TOKAMAK PLASMAS

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Abstract

The magnitude of the radial electric field and its role in the improved confinement regimes are examined in situations mainly relevant to small tokamaks.

Recent experimental results obtained on small tokamaks have revealed the importance of a quasistationary radial electric field \( E_r \) and of the corresponding plasma rotation in the transition to the H-mode and also the fact that the additional heating and the edge magnetic configuration play a secondary role in this process /1/. The experimental observation of a better confinement regime with an imposed radial electric field confirms the importance of this field in the transition /2/. In these conditions, researches connected with improved confinement regimes analog to the H-mode could be tackled also on small tokamaks. To this aim, we present three problems concerning the radial electric field in small tokamaks. First, the selfproduction of such fields in ohmic discharges is analysed and an alternative model which explains the experimental observations on CASTOR tokamak \((R=0.4m, a=0.085m, B=1T)\) is presented (Cap.1). The possibility of obtaining improved confinement regimes is demonstrated in Cap.2 by studying the evolution equation for the poloidal rotation velocity. Finally, the partial suppression of the edge turbulence in the presence of the LH wave (observed in the CASTOR tokamak) is presented as an improved confinement regime produced by an additional electric field (Cap.3).

1. The motion of the plasma in the direction of the magnetic field associated with its viscosity determines a radial flux. In the frame of the neoclassical theory, due to the fact that the parallel ion viscosity is much higher than the electron viscosity, the radial ion flux is the dominant loss. This flux produces a radial electric field whose asymptotic value results from the condition of ambipolarity. In the neoclassical model, this condition reduces to the vanishing of the ion flux \( J_{ir} = 0 \) and the resulting electric field (for the collisional regime) is:

\[
\frac{\varepsilon}{T_i} \frac{d\phi}{dr} = (\epsilon_R)^2 \frac{d}{dr}(\ln T_i)
\]  

(1)
This is a negative electric field as in experiments but its value is much smaller than the experimental one (e.g., for the CASTOR tokamak, $E_r^{nc}=0.2\text{V/cm}$ while $E_r^{exp}=2.\text{V/cm}$). It is also inconsistent with the experiment to neglect the electron flux which is anomalously large. Imposing the exact ambipolarity condition ($\Gamma_r^e + \Gamma_r^i = 0$) with $\Gamma_r^e$ containing the anomalous diffusion coefficient for electrons and $\Gamma_r^i$ determined by the loss cone, a different expression for the radial electric field is obtained. Moreover, $E_r$ shows a threshold in its dependence on the values of the plasma edge gradients, which is followed by a sudden reduction of the radial flux and by a large change in the radial electric field /3/. For the parameters of the CASTOR tokamak, this model gives too small values for the selfconsistent radial electric field in the L-mode, showing that the actual processes are more complicated.

An alternative mechanism for the selfproduction of the radial electric field results from the model we have developed for the plasma evolution in the CASTOR tokamak /4/. At the beginning, the discharge is localised in a hot channel and there are large gradients which drive the drift (collisional) instability. Due to the large linear sinks (magnetic shear and ion damping), the mode is kept for about 1 ms in the linear stage, the nonlinear mode coupling being negligible. As a result, a quasiperiodic potential structure appears over a limited zone of the plasma cross section. The small nonlinear mode coupling determines a stochasticization of the separatrices of the potential cells. A fraction of the particles which perform $E_\times B$ drift motion in this potential can percolate through the potential structure and determines the fast density decay observed in this tokamak. But, due to the large differences in the parallel (thermal) velocities of the electrons and ions, the particle percolation flux is dominantly ionic leading to a negative radial electric field. This is determined by the fact that the electrons traverse the toroidal correlation length in a time much shorter than the ions ($\tau_e=1.2 \times 10^{-7}\text{s}$, $\tau_i=1.1 \times 10^{-5}\text{s}$) which determines much shorter excursions in the $E_\times B$ drift motion ($d_e=d_i/90$). The resulting radial electric field perturbs the cell structure of the drift wave potential determining a displacement of the local maxima to the cell borders. When the radial electric field equals the maximum poloidal drift wave field, the potential structure is essentially modified passing from the cell shape to a family of waved circles around the magnetic axis. The $E_\times B$ drift is now a poloidal rotation and the large particle cross flux ends. $E_r$ cannot grow anymore and begins to decay due to viscosity and electron diffusion. As a result, the electric field produced by this mechanism has a maximum at the moment of maximum decay rate of the average density.
which equals the field of the drift wave (2 V/cm). This is in good agreement with the experiments on the CASTOR tokamak.

Both the development of the radial electric field and the evolution of the drift mode to the turbulent stage contribute to a natural inhibition of the large particle loss flux characteristic for this tokamak. The effectiveness of each of the two processes depends on the specific time scales. The radial electric field reaches a maximum value when the nonlinear evolution of the drift wave is slow enough. The large particle loss at the beginning of the CASTOR discharge could be inhibited by applying an external radial electric field of the order of the self-produced field. A practical method which have to be considered would consist in the onset of the LH wave from the first stage of the discharge (see Cap. 3).

2. As suggested by the above considerations, the rotation of the plasma and the radial electric field must be considered as evolution processes. In fact, the competing mechanisms in reaching a steady state, parallel and perpendicular viscosities, have very different time scales. Consequently, an arbitrary initial rotation has a phase of rapid readjustment consisting in the evolution of the parallel velocity \( v_u \) under the parallel ion viscosity:

\[
\frac{\partial v_u}{\partial t} = \frac{\lambda}{3} v_{th} (v_u v_u)
\]

where \( v_u \) is the gradient along \( B \). In this phase (of duration given by \( \tau = r B^2 / (\lambda v_{th} B^2) \)), the velocity becomes consistent with the amount of conserved angular velocity on the magnetic surface. The longer time scale arises from the slow transport of angular momentum between magnetic surfaces and consists in a slow readjustment of the electric potential and of a relaxation toward solid rotation. The poloidal rotation \( v_\phi \) decays through poloidal viscosity damping:

\[
\frac{m \frac{\partial v_\phi}{\partial t}}{\partial r} = e (\Gamma_{ir} - \Gamma_{er}) B_t - \frac{2}{\partial r} F_{\phi}
\]

where \( F_{\phi} \) is the component of the stress tensor and \( B_t \) is the toroidal magnetic field. The fluxes \( \Gamma_{ir}, \Gamma_{er} \) have neoclassical and anomalous components, the latter being nonlinearly connected with the poloidal velocity and its gradient. It was demonstrated that the velocity shear modifies the turbulence (the spectrum and the saturation level) /5/. The viscosity (i.e. the diffusion of poloidal velocity between magnetic surfaces) is also velocity dependent /6/, having a maximum which permits run-away states. Introducing model expressions for these terms, an equation which describes the evolution of the poloidal rotation velocity as it starts
from an arbitrary initial value (determined in the initial phase of the
discharge) is obtained. The simplified asymptotic form of this equation
is the ambipolarity condition, \( f_{1r} - f_{er} = 0 \).

A qualitative analysis of this equation shows that it has two
branches of solutions. On one branch, the initial rotation velocity
evolves to a small stationary value (corresponding to a small radial
electric field) and the particle fluxes remain rather high. On the other
branch, the (higher) initial rotation velocity saturates at a large va-
lue and the particle fluxes are significantly smaller than in the first
case. In fact, this behaviour reflects also in the ambipolarity condition
which, as shown in Ref.3, has two solutions corresponding to the L and H
modes. Quantitative results could be obtained by a numerical study of
the system formed by coupling Eq.(3) with the equations for density and
- temperature profile evolution in order to have a realistic selfconsistent
description of the fluxes (which are functions of both poloidal velocity
and plasma gradients). However, this model equation shows that there is
a critical initial poloidal velocity (or radial electric field) above
which plasma evolves to an improved confinement regime.

3. In order to understand the experimental observation that in the
presence of the LH wave the edge plasma turbulence is partially suppres-
sed in the CASTOR tokamak /4/, we have examined various possibilities.
It resulted that the direct influence of the LH wave on the linear drift
instability /7/ cannot explain this effect since it can compensate the
growth rate only for a too high value of the wave electric field.

The observed effect should be connected to the LH waves interac-
tion with the turbulence. As resulting from all experiments, the basic
characteristics of the turbulence are: high level of density fluctuati-
ons, strong nonadiabatic component \( \nabla n \neq e \times B / T \) and broad wave num-
ber and frequency spectra.

The mode coupling nonlinearity in the equation for the perturbed
density gives an incoherent contribution which broadens the frequency
spectra. The renormalised two-point equation reveals the scale dependence
of the correlation time. The two-point density correlation is coherent
at large scale, i.e. the correlation time equals the decay time of the
mode calculated from the renormalised one point (coherent) equation. At
small scale, the correlation time exceeds the coherent value due to the
space correlation in the turbulent field at small scale. The turbulent
mixing arising from the ExB convective nonlinearity destroys the corre-
lation (with the exception of those of small scale) on the time scale of
the correlation time. The dissipative parallel viscosity is scale independent and reduces the small scale correlations as well.

Since the presence of the small scale correlations enhances the rate of extraction of the free energy (the relaxation of the mean density), any process acting to reduce these small scale correlations may determine a lower saturation level of the turbulent fluctuations. The LH waves represent an additional source of scattering which enhances the effect of collisional viscosity and produces a diffusion which is not vanishing in the small scale limit. Thus, the LH wave acts to lower the fluctuation level but as a high order mechanism and cannot explain the observed degree of fluctuation suppression.

The experimental determination of the edge density turbulence spectrum shows that only the amplitudes of the small frequency (\(\leq 100\)kHz) fluctuations are reduced (up to a factor of 0.2) in the presence of the LH wave. But, the low and high frequency parts of the spectrum are determined by different mechanisms: by the drift wave turbulence and respectively by the sheared poloidal rotation of the plasma (i.e. by the edge radial electric field). These facts and the previous analysis support the idea that, at the onset of the LH power injection, an additional radial (quasistationary) electric field appears in the plasma and produces a transition to an improved confinement regime.

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Abstract
Absorber foils have been used with combination MCP (microchannel plate) detector for the measurements of soft and ultrasoft X-ray radiation of tokamak plasma with 2 μsec time resolution along 12 chords covering 2/3 of the plasma minor radius. Usefulness of such diagnostics has been demonstrated.

Introduction
Results of the measurements of soft X-ray and ultrasoft X-ray radiation of tokamak plasmas using imaging arrays are some of the main sources of our knowledge on the MHD instabilities on most tokamaks. Besides the information on the internal structure of magnetically confined plasma it can provide also valuable experience on electron temperature and on impurity concentration.

The standard method for getting such results is based mainly on silicon diodes (1). Some special devices have been elaborated for lower energies, too (2,3). In small scale tokamaks where electron temperature is not higher than some hundred eV and the time scale of phenomena is in the microsecond region. This circumstances led us to design a special detector imaging system to get the similar results as are available at greater machines with higher temperature and longer time scale.

The imaging array system consists of foils and a microchannel plate detector (4). The energy sensitivity of the MCP detector and the foils used allow us to select special energy bands. In fig.1. one can see the energy sensitivity of the detector system determined by the transmission of absorber foil and the response of the detector material. A position sensitive measurement is achieved by using more then one anod (see Fig.2.).
Measurement set up. The schematic view of the measurement set up is seen in fig. 2. The detector used is a Russian made VEU-7 chevron type MCP with 27 mm active diameter. The resistivity of the detector is couple gigaohm. The typical anod current is in the order of ten nanoamps. For the measurements of plasma radiation the imaging array detector was used in d.c. mode. Because of the noise reduction the anod current was amplified in two step and then digitized. One amplifier is next to the anods in vacuum chamber the second remote one is outside near to the digitizer. For the digitization the sampling time was typical 1-10μsec.

Measurement. For the demonstration of the capability of the detector system periodic density limit disruption were measured with different plasma current and limiter radius with fixed toroidal magnetic field. (15 kA < I_p < 30 kA, a_L_L = 6.5 - 9 cm, B_T = 1 T, discharge time = 8 ms, R = 40 cm). In fig. 3. the plasma main parameters are shown during a disruptive discharge.

For the measurement three energy bands were used

- 2 μm Al 30-80 eV and 400 eV-10 keV
- 0.2 μm C 80-300 eV and
- 10 μm Be 400-10 keV.
In our case the electron temperature of the plasma was about 200 eV which was corroborated by the beryllium and carbon foils measurements. This fact indicated, that the line radiation in these two energy intervals was negligible small. But intensive line radiation could be concluded because of the radiation measured in the 30-80 eV energy internal was about 300 times higher.

To evaluate the radial emissivity profile from the line integrated signals, these were averaged over many (10-20) disruption using the loop voltage spikes as a time base then a tomographic inversion method was used.

In such mode of operation of MT-1 tokamak three phases of the disruption could be distinguished. The first phase is a rotating MHD mode with increasing amplitude and with constant rotating frequency. The second phase lasts for 30-50 μsec. During this interval the plasma lost its energy content considerably. The last phase was the time of recreation. The plasma reached approximately the same stage as before. This recovery was occurred only if \( q_{\text{lim}} > 3 \).

Before disruption precursor oscillations could be seen both in radiation and in the Mirnov coil signals, which indicated the appearance of a rotating \( m = 2, n = 1 \) magnetic island. The highest modulation on the emissivity could be seen in the \( 80 < \hbar \nu < 300 \) eV energy range the smallest one in the lowest energy interval. The time evolution of emissivity profile of a disruptive discharge

Fig. 3. Time dependence of plasma parameters during disruptive discharge (Ip- plasma current, Ul- loop voltage, IX- X-ray intensity with Al foil, Pr- total radiated power).
measured in the three energy internal can be seen on fig.4. The main character at such a disruption is the hollow profile before it in the low energy internal.

Conclusion  Soft and ultrasoft X-ray emissivity measurements have been made for a number of chords using different filters and MCP detector. The time resolved radial profile measurements played an important in the investigation of disruption on MT-1 tokamak. The potential time resolution was 1-10 μsec and the spatial resolution is about 2 cm. The fast time resolution and the multichord possibility in different energy ranges may permit us to study fast events of plasma radiation.

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MULTICHARGED ION-ATOM COLLISIONS: APPLICATION FOR TOKAMAK EDGE PLASMA PHYSICS

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Abstract

Multicharged ion-atom collisions are of interest for tokamak edge plasma physics because of several reasons. Electron capture reactions taking place very probably in such collisions give rise to production of highly excited ion states from which characteristic line radiation is emitted, which produces a localized efficient energy drain from the plasma. On the other hand, by controlled application of such electron capture processes the temperature profile of the plasma may probably be modified to achieve optimization of plasma confinement conditions. In addition, line radiation from such excited impurity ions can be utilized for plasma diagnostics, using both passive and neutral beam-activated photon spectroscopies.

In recent years we have studied electron capture in ion-atom collisions toward the following goals.

1) Electron capture by C\textsuperscript{q+} and O\textsuperscript{q+} (q = 1, 2) from H\textsubscript{2} and He is of interest in the outermost region of tokamak discharge plasmas. Unfortunately, available cross section data are scarce and not very reliable, which is primarily caused by unknown long-lived excited (metastable) ion beam fractions in the corresponding cross section measurements.

2) State-selective cross sections for electron capture by multicharged ions (in particular C\textsuperscript{q+}) from Li(2s) are needed for the application of Li beam-activated charge exchange spectroscopy (Li-CXS), by which means the impurity ion concentrations in tokamak edge plasmas can be measured with excellent spatial and temporal resolution.

In this paper we present recently obtained cross sections for processes listed under (1). The experimental methods for measuring electron capture cross sections separately for ground state and metastable ions are described. We demonstrate that for low primary ion charge states these electron capture cross sections can be much larger for the metastable than for the ground state ions, which stresses the importance for careful consideration of possible effects as caused by the participation of metastable ions in such collisional processes.

1) Experimental methods for consideration of metastable primary ion states in measurements of electron capture cross sections for slow collisions of C\textsuperscript{q+} and O\textsuperscript{q+} (q = 1, 2) with H\textsubscript{2} and He

Figs. 1a,b and 2a,b show the binding energy schemes for single electron capture in the collision systems of our present interest. For both C\textsuperscript{+} and C\textsuperscript{2+} only one metastable primary ion state has to be considered, whereas for O\textsuperscript{+} two and for O\textsuperscript{2+} even three metastable states might be admixed in the primary ion beam /1/.
Fig. 1 - Binding energy level diagrams for electron capture in collisions of C$^+$ (a) and C$^{2+}$ (b) with H$_2$ and He, considering the respective ground state and metastable ions.
Fig. 2 - Binding energy level diagrams for electron capture in collisions of O$^+$ (a) and O$^{2+}$ (b) with H$_2$ and He, considering the respective ground state ion and the metastable ions.
The actual metastable fractions depend crucially on the ion production mechanisms in the ion source. We have applied a Nier type electron impact source, where the electron impact energy $E_e$ could be precisely controlled. By keeping $E_e$ below the respective threshold impact energy for metastable ion production, ion beams containing exclusively ground state ions can be produced. When elevating $E_e$ gradually above these threshold values, increasing fractions of metastable admixtures will be produced. These fractions have been measured by applying ion beam attenuation- as well as excitation-deexcitation techniques /2/, /3/.

Figs. 3a,b show our experimental setups which permit application of these techniques in conjunction with both state-selective measurements via translational energy spectroscopy /4/ and total cross section measurements /5/.

![Diagram](image)

Fig. 3 - Experimental arrangements for
(a - top) attenuation measurements, and
(b - bottom) translational energy spectroscopy alternatively with measurement of total cross sections for single electron capture.

As an example, in fig. 4 the attenuation of a $C^{2+}$ ion beam in He is demonstrated. When producing $C^{2+}$ by electron impact on CH$_4$, for $E_e \geq 150$ eV the respective metastable fraction was found to assume a constant value of about 10% (cf. fig. 4). For $C^+$ ions similar measurements delivered at high electron impact energy in the ion source a metastable fraction of about 12%. In a similar way also for oxygen ion beams the metastable admixtures can be controlled and determined. The above described attenuation- and excitation-deexcitation measurements are generally carried out with different gas targets to assure the evaluation of reliable results.
Furthermore, we have applied different ion source feeding gases for production of the desired primary ions. Different molecular feeding gases can give rise to considerably different kinetic energy spread of the produced ion beams, which is related to the energy sharing among atomic components of the dissociatively ionized feeding gas molecules in the ion source. As a general rule, it has been found that the relatively narrowest energy spread (which also dominates the available energy resolution for translational energy spectroscopy) is achieved with feeding gas molecules containing only hydrogen together with the desired species, as CH₄ for carbon ion production and H₂O for oxygen ion production.

Measurements for the various collision systems have been carried out by first using pure ground state ion beams, from which the corresponding cross sections for total electron capture could be deduced. Absolute data were obtained by comparison of these measurements with well established data for other collision systems involving the same target gases /6/. Subsequently these measurements were continued with ion beams with defined metastable fractions (see above), delivering so-called "apparent" electron capture cross sections, from which the cross sections for the particular metastable species could be derived.

2) Presentation of total single electron capture cross sections for C⁺ - H₂ - and C²⁺ - H₂, He collisions

Figs. 5, 6 and 7 show total single electron capture cross sections for collision of singly /7/ and doubly charged carbon ions /6/ with given target species, in comparison with results from other groups. We have not carried out such measurements for C⁺ - He collisions, since the related cross sections are probably rather small (≤ 10⁻¹⁸ cm²) because of the comparably large endothermic reaction energy defects, cf. fig. 1a.
Fig. 5 - Absolute total cross sections for single electron capture by ground state (full squares) and metastable (full circles) C$^+$ ions from H$_2$/$//$ in comparison with data from other groups.

Fig. 6 - Absolute total cross sections for single electron capture by ground state (full squares) and metastable (full circles) C$^{2+}$ ions from H$_2$/$//$ in comparison with data from other groups.
Only for C\(^+\) - H\(_2\) other primary ion state-selective measurements are available /8/,/9/. They happen to agree with ours within the combined experimental errors. For C\(^{2+}\) primary ions the present measurements are the first ones involving separation between ground state- and metastable primary ions. In all three cases the cross sections for the metastable ions are much larger than those for the ground state ions, with the difference surpassing a factor of ten at impact energies of \(\leq 500\) eV.

Similar work is in progress for O\(^+\) and O\(^{2+}\) primary ions, for which the measurements are considerably more difficult because of the larger number of metastable primary ion states involved.

Summary

In the present report we have described measurements on total single electron capture from He and H\(_2\) by singly and doubly charged carbon and oxygen ions at impact energies between 0.4 and 5 keV. Most notably, these measurements have been made separately for the ground state and the metastable primary ions, involving determination of the respective metastable primary ion beam fractions. It is shown that for the collision systems investigated especially at lower impact energies (which are relevant for such processes in tokamak edge plasmas) the cross sections for the metastable primary ions become much larger than those for the ground state ions. This underlines the importance for a careful consideration of possible influences due to the participation of metastable ion beam admixtures in such electron capture studies.
Acknowledgments

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SUMMARY OF THE SESSION

J. FUJITA

This session included the following four papers, which range from a very small aspect ratio tokamak to a large one, and even from a spheromak to a reversed field pinch. The work has been carried out in several countries: Brazil, Turkey, Japan and the United States.

The first paper, "Physics issues in Proto-Eta", presented by G.O. Ludwig, is a proposal for constructing a spherical tokamak called Proto-Eta to investigate the dependence of confinement, stability and operational limits of tokamaks on aspect ratios as small as 1.5.

The second paper, "Eccentric Operation Modes in Small Tori and Their Physics", presented by Z. Yoshida, deals with detailed experimental results and physical analysis in ultra-low-q discharges on TORIUT-6, REPUTE-1 and REPUTE-1Q, and in runaway discharges on TORIUT-6 and NOVA-II. What they found was: 1) self stabilizing nonlinear dynamics of m=1 kink modes, 2) resistance anomaly, and 3) anomalous ion heating of ions in ultra-low-q discharge, 4) intense runaway discharge with large Shafranov's A, and 5) acceleration of runaways by inductive coupling of the plasma and beam in "runaway discharges".

The third paper, "Studies of Turbulence and Wave Coupling Based on Self Organization Process at SK/CG-1 Machine", presented by A. Sinman, is a detailed analysis of experimental results on a spheromak named SK/CG-1. The main results are concerned with a sudden occurrence of self organization resulting from the transfer of external circuit current to the plasma, production of turbulent plasma column, associated instabilities and plasma-wave coupling in a shock-heated plasma.

The last paper, "Design and Fabrication of the HBT-EP Tokamak: a Large Aspect Ratio, High ß, 10 msec Pulsed Tokamak", presented by M.E. Mauel, is on the design and construction of a versatile tokamak of large aspect ratio 5 ~ 6. The scientific programme is to study 1) the role of a conducting shell on MHD instability, 2) instabilities at density limit, 3) role of plasma cross-sectional shape and of a large aspect ratio on MHD instability, and 4) tokamak operation at high εβp ~ 1 to assess the second stability regime. The machine is under construction, and the first plasma is expected to be produced in June 1991.

The work presented in this session and related discussions is quite valuable and useful when people think about the construction of the next device in developing countries as well as in a small scale laboratory in developed countries, for the purpose of promoting contributions of small tokamak research to the international fusion effort.
THE PROTO-ETA SMALL-ASPECT-RATIO EXPERIMENT

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Abstract

Proto-ETA is a small-aspect-ratio tokamak (or “spherical torus”) currently under design at the “Instituto Nacional de Pesquisas Espaciais” (National Institute for Space Research - INPE) in collaboration with the Oak Ridge National Laboratory. The main objective of the device is to characterize the properties and the performance of low-aspect-ratio tokamak plasmas at relevant current and temperatures. This paper reports the current status of the conceptual design and discusses the models used to predict the operational regime and to design the main components of the device.

I. INTRODUCTION

Small-aspect-ratio tokamaks have been proposed to enhance the value of the parameter beta from a maximum of around 5% on current standard devices to 10% or more needed for economical fusion reactors \(^1,2\). Here beta is defined as \(\beta = 2\mu_0\langle p \rangle/B_0^2\), where \(\langle p \rangle\) is the average value of the kinetic pressure and \(B_0\) is the value of the vacuum toroidal magnetic field at the geometric center of the cross-section of the plasma column. Recently, high values of \(\beta\) have been obtained in highly elongated \((k > 2)\) discharges in D-IIID\(^3\). These discharges require careful feedback control of the shape of the cross-section and vertical position of the plasma column. Equilibrium calculations indicate that small-aspect-ratio axisymmetric configurations have substantial elongation in a straight vertical field. This “natural elongation” leads to enhanced vertical stability and the tight aspect ratio also enhances the stability to ballooning and global ideal MHD modes\(^4,5\). Consequently, achievement of high values of \(\beta\) is facilitated in these configurations. The saturated width of magnetic islands due to the quasilinear evolution of single tearing modes is also predicted to decrease with aspect ratio\(^6\). This can lead to enhanced energy confinement if anomalous transport is indeed related to the magnetic topology of the equilibrium configuration\(^7,8\).
One important issue with regard to the performance of compact configurations is anomalous transport. The trapped particle population increases as the aspect ratio decreases whereas the width of their orbits is expected to decrease\textsuperscript{9}. The net effect on drift modes and related anomalous transport remains to be thoroughly investigated. Neoclassical effects upon the parallel plasma resistivity and bootstrap current, among others, are also expected to be greatly enhanced in small-aspect-ratio tokamaks. Thus, the experimental and theoretical investigation of the performance of spherical tori can substantially broaden the knowledge of the physics of tokamaks and have influence on the design of fusion reactors. In particular, current scaling laws for the energy confinement time in tokamaks are based upon data with insufficient discrimination of geometry related parameters\textsuperscript{10}. The shape-size index $f_s = 0.32A(ak^2)^{1/4}$, where $A = R_0/a$ is the aspect ratio, $R_0$ and $a$ are the major radius and halfwidth of the plasma column, and $k$ is the elongation of the plasma cross-section, is approximately $1.0 \pm 0.1$ in almost all devices currently in operation. This factor can be substantially reduced in spherical tori offering the possibility to extend the database for scaling laws in the least varied direction of the parameter space\textsuperscript{10}.

The main objective of the Proto-ETA project is to design and construct an inexpensive device to investigate the performance of small-aspect-ratio axisymmetric configurations and study relevant physical processes such as stability and transport. A schematic drawing of the device is shown in Fig. 1 and the main parameters are

FIG. 1. Schematic diagram of Proto-ETA.
listed in Table I. The crucial component in the design of a spherical torus is the main induction solenoid due to the small space available between the vacuum chamber and the center post of the toroidal field coils. Furthermore, currents induced in the center post can reduce the flux available to drive the plasma current. The design of the induction solenoid of Proto-ETA follows the concept developed in Oak Ridge for the solenoid of STX. A detailed analysis of the conceptual design is presented in the sequel. A maximum plasma current $I_p = 240$ kA can be induced in the plasma for at least 100 ms. Thus, the values of characteristic plasma quantities and dimensionless parameters to be obtained in Proto-ETA discharges are expected to be relevant for comparison with discharges in standard tokamaks.

<table>
<thead>
<tr>
<th>TABLE I. Major parameters of Proto-ETA conceptual design</th>
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<tbody>
<tr>
<td><strong>design values</strong></td>
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<tr>
<td>major radius $R_0 = 0.36$ m</td>
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<tr>
<td>half-width of plasma column $a = 0.24$ m</td>
</tr>
<tr>
<td>aspect ratio $A = 1.5$</td>
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<tr>
<td>elongation $k = 1.8$</td>
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<tr>
<td>triangularity $\delta = 0.35$</td>
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<tr>
<td>toroidal magnetic field (vacuum) $B_0 = 0.36$ T</td>
</tr>
<tr>
<td>plasma current $I_p = 0.24$ MA</td>
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<tr>
<td>cylindrical safety factor $q_c = 2.5$</td>
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<tr>
<td><strong>expected values</strong></td>
</tr>
<tr>
<td>maximum average density (Hugill-Murakami) $n = 10^{20}$ m$^{-3}$</td>
</tr>
<tr>
<td>density averaged temperature (for 1.5 V loop voltage) $T = 390$ eV</td>
</tr>
<tr>
<td>average poloidal beta (Shafranov) $\beta_I = 0.35$</td>
</tr>
<tr>
<td>maximum total $\beta$ (MHD limit) $\beta = 0.12$</td>
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<tr>
<td>effective charge number $Z_{eff} = 1.5 \sim 2.5$</td>
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</table>

II. PLASMA PERFORMANCE

One of the physics issues of major relevance for Proto-ETA is plasma performance, i.e., energy confinement, stability and operational limits. Full inductive operation is predicted for the experiment in its basic role; auxiliary heating and current drive will be considered in an extended phase of the research program.

Recent experimental results indicate that, at least for large tokamaks, the density limit is an edge density limit. For insufficient heating power, densities at the plasma edge above a critical value lead to rapid cooling and equilibrium collapse. Empirical scalings, such as the Hugill-Murakami, do not give the correct dependence on the parameters and the connection between central and edge densities. Even though, the maximum density achievable in Proto-ETA can be estimated using the empirical expression

$$n_{20} = \nu \frac{B_0}{R_0 q_c},$$  \hspace{1cm} (1)
where \( n_{20} \) is the average plasma density in units of \( 10^{20} \text{ m}^{-3} \), \( \nu \) is a constant and 
\[ q_e = \pi a^2 B_0 (1 + k^2)/(\mu_0 \rho a L_p) \]
is the cylindrical safety factor. Taking \( \nu \approx 2.5 \), an
average density limit \( n \approx 10^{20} \text{ m}^{-3} \) is obtained.

To estimate the density averaged temperature, \( T = (n_e T_e + n_i T_i)/n \), where \( n_e \), \( T_e \)
and \( n_i \), \( T_i \) are the densities and temperatures of electrons and ions, respectively, one
uses a simple model for the global power balance which is described by the equation

\[
\frac{\partial Q}{\partial t} = -\frac{Q}{\tau_E} + W_Q. \tag{2}
\]

In this equation \( Q = 3nTV_p \) is the plasma energy content, \( V_p \) is the plasma volume,
\( \tau_E \) is the energy confinement time, and \( W_Q \) is the net heating power given by

\[
W_Q = W_\Omega - W_R + W_A, \tag{3}
\]

where \( W_\Omega \) is the ohmic heating power, \( W_R \) is the radiated power due to brems-
strahlung, and \( W_A \) is the auxiliary heating power, taken equal to zero for ohmic
operation. The model takes into account profile, impurity, and neoclassical effects on
\( W_\Omega \) and \( W_R \) but neglects the strong paramagnetism and bootstrap current expected
in small-aspect-ratio configurations. The energy confinement time \( \tau_E \) is calculated
using different L-Mode scaling laws. Assuming the Neo-Alcator scaling\(^{12}\), one obtains
\( \tau_E = 2.2 \text{ ms} \) and a loop voltage \( V_l = 4.0 \text{ V} \) for \( Z_{\text{eff}} = 2.0 \). This value of the loop
voltage is considered excessive as compared to experimental results which indicate
that a value of \( V_l \) within the range \( 1.0 \sim 2.0 \text{ V} \) should be expected\(^{13}\). Accordingly, a
model was developed for ohmic operation in which the value of \( V_l \) is fixed a priori;
neglecting the radiated power in steady-state one obtains

\[
\tau_E \sim \frac{n R_o^{2} k^{1/3} q_e^{1/3}}{V_l^{5/3} B_0^{1/3} (1 + k^2)^{1/3}}. \tag{4}
\]

The plasma operation contours in the ohmic regime are shown in Fig. 2 for different
values of the loop voltage. An average plasma temperature in the range \( 390 < T < 510 \text{ eV} \)
can be expected for \( 1.5 > V_l > 1.0 \text{ V} \), which corresponds to \( 11 < \tau_E < 22 \text{ ms} \).
Using the standard expression for \( \tau_E \), one obtains \( T = 200 \text{ eV} \) and \( V_l = 4.0 \text{ V} \), as
mentioned above. Provision is being made for additional power for auxiliary heating
and current drive after the initial ohmic heating phase of Proto-ETA. With auxiliary
heating, the energy confinement time is taken to scale as \( \tau_E \sim W_Q^\alpha \), where \( \alpha = -0.5 \)
for most empirical laws. The operation contours for the Kaye-big scaling law\(^{14}\) are
shown in Fig. 3 for ohmic equilibrium and for additional power levels \( W_A = 1.0 \) and
2.0 MW.

It must be pointed out that large discrepancies result from the application of different
scaling laws. The reason for these discrepancies is two-fold. For basically similar
empirical scaling laws, a geometrical factor of about two results in the values predicted
for the energy confinement time due to the value of the shape-size index \( f_s \)
mentioned in the introduction. For scaling laws with different dependence on the
plasma parameters (different physics, as Rebut-Lallia and Lackner-Gottardi, for in-
stance), one order of magnitude discrepancy in the value of \( \tau_E \) can result. Of course,
FIG. 2. Plasma operation contours for Proto-ETA in ohmic regime and for different values of the loop voltage.

FIG. 3. Plasma operation contours for Proto-ETA with auxiliary heating assuming Kaye-big scaling\textsuperscript{14}.
one of the main objectives of a small-aspect-ratio experiment such as Proto-ETA is to help clarify the uncertainties in the scaling of $\tau_E$.

Considering the MHD operational limits for $\beta$ and for the safety factor $q$ and assuming the Troyon-Sykes scaling, $\beta = 10^{-8} g I_p / (a B_0)$, a two-fold increase in the value of $\beta$ can be immediately obtained if the aspect ratio is reduced from the standard value $R_0/a = 3.0$ to 1.5, for the same value of the cylindrical safety factor. Actually, $\beta$ values over 10% are expected for stable discharges in Proto-ETA with careful control of the current-density profile. Figure 4 shows a stable equilibrium obtained for Proto-ETA using a fixed boundary equilibrium and stability code. It can be noted that the pressure profile is rather broad and the current-density profile is almost hollow, producing a very flat $q$ profile. For this equilibrium $q_0 = 1.3$, $q_a = 4.8$ and $\epsilon \beta_I = 0.23$ giving $\beta = 0.12$ with a reasonable paramagnetic effect ($B_t/B_0 \simeq 1.2$ at the magnetic axis). Here $q_0$ and $q_a$ are the values of the safety factor at the magnetic

![Flux contours](a)
![Pressure profile](b)
![Toroidal current density](c)
![Normalized q profile](d)
![Toroidal magnetic field](e)
![Poloidal magnetic field](f)

**FIG. 4.** Stable MHD equilibrium profiles for Proto-ETA: (a) flux surfaces; (b) pressure; (c) toroidal current density; (d) normalized $q (q/q_a)$; (e) toroidal magnetic field; (f) poloidal magnetic field. In (e) the dashed line corresponds to the field in vacuum.
axis and at the plasma boundary, respectively, and the poloidal beta is defined as
\[ \beta_T = 8\pi S(p) S/(\mu_0 I_p^2), \]
where \( S \) is the area of the poloidal cross-section of the plasma column. To maintain the value of \( q_0 \) above one, non-inductive current drive or strong bootstrap current will certainly be required.

III. START-UP MODEL

A simple zero-dimensional model is used to calculate the power supply requirements for the start-up and maintenance of the full plasma current in Proto-ETA. The plasma column is modelled as a single current loop with a temperature-dependent resistance coupled to the main induction solenoid. The circuit equations are advanced in time together with the power balance equation, Eq. 2, and the particle balance equations

\[ \frac{\partial n_0}{\partial t} = -n n_0 S(T) + P_0 \]

and

\[ \frac{\partial n}{\partial t} = -\frac{n}{\tau_p} + n n_0 S(T), \]

where \( n_0 \) is the density of neutral particles, \( S(T) \) is the ionization rate, \( P_0 \) is the particle refuelling rate, and \( \tau_p \) is the particle confinement time; one takes \( \tau_p \approx 2\tau_E \). The equivalent circuit and the current waveform in the primary circuit of the induction transformer are shown in Fig. 5. The lumped circuit parameters \( R_{Q} \) and \( L_{Q} \) refer to the main induction solenoid and \( R_{m} \) and \( L_{m} \) to the correction coils of the primary circuit (Section IV). One of the secondary circuits, with parameters \( R_{t} \) and \( L_{t} \), models the induced currents in the center post of the toroidal field coils which is closely coupled with the main induction solenoid, as indicated in Fig. 6. The other secondary circuit, with parameters \( R_{p} \) and \( L_{p} \), models the plasma loop. The plasma inductance \( L_{p} \) is calculated using the model of Hirschman\textsuperscript{15} for the external self-inductance valid for small-aspect-ratio plasma loops.

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FIG. 5. Equivalent circuit (a) and schematic variation of the current in the primary circuit (b) of the induction transformer.
To estimate the parameters of the induction current loops in the central legs of the toroidal field coils, a simplified model is used. Each one of the $N$ cylindrical wedged sectors and cooling holes which constitute the center post, as shown in Fig. 6, is approximated by confocal elliptic cylinders. The magnetic diffusion equation is solved in elliptic-cylinder coordinates in the low magnetic Reynolds number approximation. This model is thus valid only if the time scale of variation of the external flux is larger than the characteristic decay time of the induced currents. The resistance and self-inductance associated with each induction current loop is calculated from the eddy current distribution. The mutual inductance of the sectors and the main induction solenoid is calculated assuming that the main solenoid links the total magnetic flux generated by the eddy currents in each sector. With this model one finds that the decay time of the induced currents is less than 2 ms for $N = 16$ coils in Proto-ETA. In Table II the values of all the parameters in the equivalent circuit of the induction transformer are given.

| TABLE II. Lumped circuit parameters of the circuit model for the induction transformer |
|---------------------------------|---------------------------------|-----------------|
| main solenoid                   | $R_\Omega = 9.55\,m\Omega \ (\sim 65 \,\text{C})$ | $L_\Omega = 315\,\mu\text{H}$ |
| magnetizing (correction) coils  | $R_m = 0.678\,m\Omega \ (\sim 35 \,\text{C})$       | $L_m = 76.4\,\mu\text{H}$    |
| induced current loops           | $R_t = 4.14\,\mu\Omega$                        | $L_t = 5.76\,\text{nH}$      |
| plasma loop                     | $R_p = 1.50/T_{keV}^{3/2} \,\mu\Omega$           | $L_p = 0.229\,\mu\text{H}$  |
| mutual inductances              | $M_{tp} = 2.82\,\mu\text{H}$                     | $M_{th} = 0.930\,\mu\text{H}$|
|                                | $M_{tp} < 7\,\text{nH}$                         |                 |
The temporal profile of the current in the primary circuit is shown in Fig. 5(b). Initially, the induction coils are energized by switching the power supply $V$. At the end of the charging period, the induction coils are discharged on the load resistance $R_s$ to provide fast flux variation for plasma breakdown (ECRH assisted) and current ramp-up. When the plasma current reaches the desired flat-top value a feedback controlled power supply $e$ is switched on to keep the current constant against resistive decay. Finally, at the end of the flat-top period the discharge is quenched by letting the primary current decay on the load resistance $R_r$. The duration of the plasma current is limited by imposing a maximum temperature of 100°C or a maximum current density of 36.6 kA/cm$^2$ in the main solenoid, which corresponds to a peak current of 54 kA in the power supply. In Fig. 7 the temporal profiles of the primary current, plasma current, plasma temperature, neutral particle density, and plasma density are shown for a typical simulation. The particle refuelling rate $P_0$ is adjusted in the simulation to give a desired plasma density during flat-top. For this case, the power supply has to provide a peak power of 36 MW and an average power of $\sim 10$ MW.

FIG. 7. Temporal profiles of the currents in the primary circuit and in the plasma (a), the plasma temperature (b), the neutral particle density (c), and the plasma density (d) for a typical discharge in Proto-ETA. The densities are in units of $10^{19}$ m$^{-3}$.
IV. POLOIDAL COIL SYSTEM

The poloidal coil system includes the main induction solenoid, the magnetizing or correction coils to minimize the strong field in the plasma region, and the equilibrium coils. The positions of the different coils have been chosen such as to have a reasonable space between coils for diagnostic purposes.

The main solenoid will be made of GlidCop AL-15 conductor and has two layers of 66 turns each. The inner radius is 6.6 cm, the external radius is 8.8 cm, the total length is 1.11 m, and the packing factor is 0.78. For a maximum current density of 36.6 kA/cm$^2$, the tangential stress in the inner surface of the solenoid is 97 MPa, well below the yield strength of 310 MPa for GlidCop.

The magnetizing and the equilibrium coils will be made of OFHC copper in half hard condition with a central hole for water cooling. They are designed to carry a maximum current density of 8 kA/cm$^2$. The positions of the magnetizing coils have been chosen to yield a high null at the expected position of the magnetic axis, as shown in Fig. 8. A sixth-order null is produced with a field smaller than 1.2 mT in a region of approximately 10 cm radius around the magnetic axis. The inner coils, closer to the main solenoid, have 8 turns and the outer ones, packed together with the shaping and radial field coils, have just one turn.

![Field pattern produced by the magnetizing coils (black) in Proto-ETA.](image)
The main equilibrium field is produced by the pair of coils at \( R = 0.96 \) m and \( Z = \pm 0.43 \) m. They produce a vertical field with an almost zero decay index at the plasma region, as indicated in Fig. 9(a). A small control of elongation is provided by the shaping coils at \( R = 0.585 \) m and \( Z = \pm 0.685 \) m, as shown in Fig. 9(b). The vertical field coils have 36 turns and the shaping coils 9 turns. Control of the vertical position of the plasma column is provided by the radial field coil just above the shaping field coil.

V. TOROIDAL FIELD COIL

A simple design based on a picture frame geometry is adopted for the toroidal field coils, with a butt joint connecting the external legs to the central column, and lap joints connecting each of the 16 external legs to the bus bar of the power supply. The thickness of each coil is 2.24 cm and the other major dimensions are shown in Fig. 10. The inherent ripple produced with this configuration is about 0.01% at the outer edge of the plasma column \( (R = 0.6 \) m). Taking into account the small space available, the central column (Fig. 6) has been designed for a maximum temperature rise of 60 C during the 1.5 s effective pulse duration of the toroidal field. Cooling channels are introduced to restore the coil temperature to the initial value in a short time interval. The cooling channels have the added advantage of increasing the electrical resistance of the central column to eddy currents. The central legs are to be fabricated by extrusion of OFCH copper and bonded together as one assembly for strength and dimensional control. The power dissipated in the central column due to eddy currents
is estimated to be a small fraction of the total ohmic power. Including the cooling channels and the insulation gap between legs (1 mm), the current density at maximum field is about 8.2 kA/cm$^2$. A simple structural analysis of the coil system indicates that the coils will operate at stress levels about three times less than the maximum admissible stress for copper, even at the critical points near the external joints. The stress at the butt joint is kept at a much lower level. Presently one is considering the possibility of adopting C shaped instead of rectangular coils to minimize the bending moment along the coil, and at the butt joint in particular.

**VI. POWER SUPPLIES**

The power supplies needed are basically for the toroidal field coils, the induction transformer and the equilibrium field coils. The high-voltage power grid has been selected as power source for the three types of supplies. For the toroidal field supply a 12 pulses bridge rectifier using a 2.5 MVA transformer with a 90 V RMS no-load voltage is perfectly adequate. A 2.5 MVA autotransformer with adjusting taps will be introduced to cover the $B_0 = 0.1$ to 0.4 T range.

The power supply for the magnetizing coils is more demanding. Two dual secondary 5 MVA transformers are needed to attain a peak power level of 45 MW and peak voltage level of 800 V. A 24 pulses transformer/rectifier system will be adopted to reach a rise time of about 3 ms for the current control loop.

FIG. 10. Poloidal cross-section of Proto-ETA showing a toroidal field coil, and cross-sections of the vacuum chamber, the magnetizing coils M1, M2, and M3, the equilibrium coils V and S, and the radial field coil R for feedback control of the vertical position of the plasma column.
VII. VACUUM SYSTEM

A cross-section of the vacuum chamber is shown in Fig. 10. The vessel will be continuous and thus it has to be sufficiently resistive to avoid large induced currents. The maximum allowed thickness of the inner wall is estimated from the maximum value of the loop voltage during start-up. The current flowing in the wall is given by \( I = \frac{Vh}{(2\pi\rho)}\Delta r \), where \( h \), \( r \), \( \Delta r \), and \( \rho \) are respectively the height, radius, thickness and resistivity of the wall. To estimate the maximum allowed current, the resulting error field at the geometric axis of the plasma column is imposed to be smaller than 0.1 mT. Considering the wall as a short solenoid, the field produced at the geometric axis \( (R_0 = 0.36\,\text{m}) \) is \( B(T) = 6.7 \times 10^{-9}I(A) \), for \( h = 1.32\,\text{m} \) and \( r = 0.09\,\text{m} \). Thus the maximum allowed current is \( I = 14.8\,\text{kA} \). Assuming that the maximum loop voltage during start-up is \( V \approx 7\,\text{V} \), one obtains \( \Delta r = 0.9\,\text{mm} \). Considering however that the stress during disruptions can be of the order of 1 atm, the design value is taken to be \( \Delta r = 1\,\text{mm} \).

The vacuum chamber has large diagnostic ports which allow good access into the chamber. The large vertical ports also provide good access for neutral beam injection. The inner wall will be made from a 1 mm thick inconel tube; the other walls will be made of 5 mm thick inconel plates. Presently a modification in the design of the chamber is being considered in order to use BRIGHTON 80-10 ASME heads for the top and bottom of the chamber. These heads come in standard sizes and are made of stainless steel or inconel. Spacers will be used to provide a constant clearance between the inner wall of the vacuum chamber and the induction solenoid. Conflat seals will be used at the diagnostic ports so that baking and discharge cleaning at high wall temperatures will be possible.

Graphite limiters covered with a layer of titanium carbide will be used covering the inner wall of the chamber. However, similar stainless stell limiters will be used for commissioning and initial operation of the device.

A total pumping speed of 4000 l/s will be made available by two turbomolecular pumps, a 500 m³/h Roots pump, and a 80 m³/h two-stage rotary pump.

VIII. AUXILIARY HEATING

Wave heating and current drive in Proto-ETA will probably be inefficient because lower hybrid accessibility and electron cyclotron wave absorption are made difficult by the low field and high density. Alfvén wave heating and helicity injection are interesting schemes to be investigated in small-aspect-ratio tokamaks. However, because they are not yet fully demonstrated schemes, neutral beam injection is being considered as the first alternative for auxiliary heating in the extended phase of Proto-ETA (high-beta operation). Two neutral beam injectors of 20 keV particle energy are envisaged for Proto-ETA. The beamlines are at right angles, as shown in Fig. 11. The beams can penetrate into a large portion of the plasma column without hitting the inner wall of the vacuum chamber. A simple pencil model and a Solev'ev equilibrium model were used to calculate the neutral beam shinethrough. The results are shown in Fig. 12 for a central plasma density \( n(0) = 0.5 \times 10^{20}\,\text{m}^{-3} \) and for different im-
FIG. 11. Top view of Proto-ETA showing access of neutral beam injectors.

PROTO-ETA
Solovev Equilibrium Model
$E_b = 20.0 \text{ keV}$
$N_e = 0.50 \times 10^{20} \text{ m}^{-3}$

FIG. 12. Neutral beam shinethrough for a 20 keV beam, $0.5 \times 10^{20} \text{ m}^{-3}$ central plasma density, and different impact parameters.
pact parameters. It can be verified that for the large densities expected in high beta operation there is good beam absorption.

IX. DIAGNOSTIC SYSTEMS

The different diagnostic systems for Proto-ETA are divided in two groups: the first includes diagnostic systems that are considered essential for the ohmic operation phase of the device in which maximizing the value of $\beta$ is not considered a major goal; the second is for the extended phase of operation in which auxiliary heating is envisaged. In the first group are included magnetic loops, electrostatic probes, Thomson scattering system, 2mm microwave interferometer, spectrometer (ion temperature), thermocouples, bolometers and TV/VCR. In the second group are included upgraded magnetic loops, 1mm microwave interferometer, heavy-ion beam probe, multi-channel Thomson scattering system, soft X-ray diodes, neutral beam probe, upgraded bolometer system, upgraded spectrometer and polarimeter.

X. CONTROL AND DATA ACQUISITION SYSTEM

The Proto-ETA Control and Data Acquisiton System (CODAS) is comprised of three main sub-systems: the plasma data acquisition system (DAC), the machine control and monitoring system (MCMS), and the interlocking system. The CODAS primary objective is to integrate the basic functions of plasma data acquisition and storage, machine control and sequencing of operations, and safety and security interlocking for personnel and sub-systems integrity.

The plasma data acquisition system (DAC) is based on a micro-Vax computer interfaced with a serial highway CAMAC link. The connection between the computer and the serial highway driver is a standard asynchronous line. The CAMAC crates are controlled with standard L-2 crate controllers. The link between the crate controllers is done with optical cables through U-type adapters. A master timing generator and time delay modules provide the triggering and synchronization signals to the CAMAC crates.

The connection between the main DAC computer and the machine control and monitoring system (MCMS) is done through a PC based supervisory work station. The MCMS is a distributed control and data acquisition system with three hierarchical levels. The basic functions of machine control and subsystems operational sequencing is done through a multi-drop network of programmable logic controllers (PLC). The dedicated PLC's (level 3) are coordinated by a master PLC which is connected to the supervisory workstation (level 2). The main DAC computer (level 1) is linked to a host computer in a Vax cluster.

XI. CONCLUSIONS

An overview of the conceptual design of Proto-ETA has been presented. Cost estimates are not yet completed but preliminary figures indicate that the basic device shall cost around five millions dollars. Thus, the device will allow relevant tokamak physics to be investigated at low aspect-ratios with a very good cost/benefit ratio. The
conceptual design has recently been reviewed by a team of physicists and engineers from INPE and Oak Ridge National Laboratory; the final report on the design will be completed by the end of this year. Preliminary arrangements are being made in Brazil between various plasma groups and financing agencies to assemble technical and financial support for the project.

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ECCENTRIC OPERATION MODES IN SMALL TORI
AND THEIR PHYSICS

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Abstract

Some different eccentric operation modes have been tried on small tori. These experiments are aimed at understanding strongly nonlinear processes in tokamak-type discharges with enhancing and extending their anomaly. In the ultra low q (q < 1) discharges, anomalous ion heating has been observed. The discharge is self-stabilized against kink instabilities through MHD relaxation process. Intense runaway discharges have been also studied to highlight the dynamics of high energy beam component in a tokamak type equilibrium.

1. INTRODUCTION

This paper describes some trials of eccentric operation modes in small tori, and their physical implications for anomalous behaviors of discharges. While detailed measurements are limited in small experiments, deviations from standard parameters extends general understandings of the plasma physics.

We used four different devices; REPUTE-1 [1], REPUTE-1Q [2], TORIUT-6 [3] and NOVA-II [4]. The REPUTE-1 device is basically an RFP, which also operates in the ultra low q (safety factor q < 1; ULQ ) regime. REPUTE-1Q and TORIUT-6 have been also used to study the ULQ plasmas. The ULQ regime is the intermediate of the tokamak and RFP regimes of the toroidal discharges, and ULQ plasma exhibits characteristics of both the tokamak and the RFP in its relaxation phenomena [5,6]. In Sec. 2, we shall describe experimental observations of ULQ plasmas. TORIUT-6

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and NOVA-II have been operated in the strongly runaway regime. When the runaway component of the toroidal current becomes comparable to the plasma component, the toroidal equilibrium deviates from the normal plasma equilibria, and the significance of plasma-beam systems appears. The strongly runaway discharge is, therefore, in the intermediate of the tokamak and betatron [4]. We shall discuss runaway discharges in Sec. 3.

2. ULTRA LOW $q$ DISCHARGES

2-1. Resistance Anomaly

The low $q$ regime of toroidal discharges has been studied on TORIUT-6, REPUTE-1, and REPUTE-1Q. In the ULQ regime, the MHD relaxation process associated with the $m=1$ kink fluctuations dominates the dynamics of the plasma [6]. A paramagnetic configuration is self-organized through the self-stabilizing behavior of the kink instabilities, which has similarity with the RFP. There is significant resistance anomaly in ULQ discharge, which is also similar to the RFP. Figure 1 shows typical waveforms of a ULQ discharge in REPUTE-1. In the stabilized period, the surface safety factor $q_a$ is around 0.3. Typical electron temperature (Thomson scattering) is around 100 eV. The effective resistivity is in the order of 10 times the classical resistivity evaluated by the kinematic electron temperature. The anomalous resistance is scaled by the toroidal magnetic field $B_t$. In Fig. 2, we plot the effective resistivity

$$\eta^* = \frac{\text{loop voltage}}{\text{toroidal current}} \times \frac{(2R/a^2)}{B_t}$$

versus $B_t$ for data from REPUTE-1, REPUTE-1Q, OHTE [7], STP-3M [8], HBTX-1C [9], and TORIUT-6 [3]. We obtain a resistance scaling

$$\log \eta^* (\mu \Omega \text{m}) = 0.60 - 1.0 \log B_t (T).$$

The anomalous resistance of pinch-type discharge is considered to be related to the MHD relaxation process, and which is also related to the anomaly in the ion heating mechanism.

![Fig. 1 Typical waveforms of an ULQ discharge.](image1)

![Fig. 2 Experimental scaling of anomalous resistance in ULQ discharges.](image2)
2-2. Anomalous Ion Heating

Remarkable anomaly is observed in the ion temperature $T_i$. The ions are rapidly heated up to 1 keV level (REPUTE-1), and the direct ion heating is correlated with the low frequency MHD fluctuations. Figure 3 shows typical waveforms of a high-$T_i$ ULQ discharge in REPUTE-1. We observe a significant differences in temperatures of different species. A quasilinear model of random accelerations of ions by instability electric fields accounts well for the anomalous ion heating [10].

![Fig. 3 Typical waveforms of a high $T_i$ ULQ discharge.](image)

3. RUNAWAY DISCHARGES

3-1. Plasma-Beam Equilibrium

A runaway discharge with intense relativistic electron current is a typical plasma-beam system, which is considered to be a bifurcation from the normal tokamak equilibrium to a betatron type configuration [11,12]. Figure 4 shows typical plasma-beam equilibrium solution with strong beam pressure and self-field effects. The beam pressure effect is experimentally detectable in Shafranov's $\Lambda$. In Fig. 4, $\Lambda=2$. Using TORIUT-6 and NOVA-II, strong runaway discharges have been studied. Self-consistent plasma-beam equilibrium is achieved by applying stronger vertical field. In NOVA-II runaway equilibria with $\Lambda$ up to 3 have been obtained.

![Fig. 4 Typical plasma-beam equilibrium. $\psi$ is the flux function, and $\phi$ is the streaming function of beam velocities with 10 to 164 MeV (five components). Plasma current is 1 MA.](image)

3-2. Acceleration by Inductive Coupling with Plasma Component

The beam component of the toroidal current is enhanced by the decay of the plasma-current component through the in-
Intense runaway discharge is studied as a typical plasma-beam system. Present small experiment simulates the disruption produced runaways, and the acceleration mechanism has been studied. There is also a possibility of applications the runaway beams in the beam-core concepts. The self-field effect of intense electron beam has strong influence to the structure of the toroidal equilibrium.

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STUDIES OF TURBULENCE AND WAVE COUPLING 
BASED ON SELF ORGANIZATION PROCESS 
AT SK/CG-1 MACHINE

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Abstract

In the SK/CG-1 Machine, using the He gas pressures between 30 and 50 mTorr, a complex turbulence regime has been determined. During the on-set phase of 0.35 us, besides the anomalous plasma resistivity, the ion acoustic instability is also observed. At the C-gun the plasma belt which contains the hot electrons at the energy range of 60-80 eV, breaking from the gun electrodes, the translation into plasmoid case occurs and than it is produced the high electric fields about 5 kV/cm in the poloidal frame of the gun. Thus, the electron-ion two stream instability and just after of 0.1 us, the ion acoustic instability are started. On the other hand, in the second phase, the beam plasma instability and the ion acoustic instability are parallely sustained. The beam plasma instability includes also the W3 ion cyclotron instability. In addition, due to the toroidal magnetic field of the poloidal frame and the electric field perpendicular to it, the density and temperature gradients are come into existence. At the end of the resonance of the drift wave and the ion acoustic wave probably, the drift wave turbulence commences. The time constant of the octagonal flux conserver is about 28-30 ms. After the turbulence, the electric field on the poloidal frame drops down to 900 V/cm and it is kept for 300 ms (E/e). Thus a wave coupling mechanism depending on the self organization process is sustained. Nevertheless, it has not yet been understood exactly if the propagation at the flux conserver is realized due to the reflections or the helicity transfer from the wave to the plasma. The frequency of the closed wave packet changes in the range of 25 - 5 kHz depending on the electron temperatures in the range of 30 - 5 eV. The average electron density is about 10^14 cm^-3 and the global energy confinement times are found in the range of 65 - 250 us.

1. INTRODUCTION

In magnetic confinement systems, the self organization process can generally be defined by the important mechanism which develops the variation of controlled and programmable system parameters impressed by the principal plasma parameters.
The role of the self organization process may be seen in the examples such as the change of q profiles at high beta systems /1/; the sensitivity of magnetic Reynolds number S at RFP schemes /2/; the influence of F-O field reversal ratio and pinch parameter /3/ and the tokamak formation in circular cross section from a cylindrical transverse plasma sheet in the DC helicity driven tokamak /4/.

As it has been emphasized in the previous studies /5,6/, in order to generate different types of compact toroids such as spheromak, spherical torus and spherical pinch tokamak in the hybrid SK/CG-1 machine (World Survey of Activities in Controlled Fusion Research, 1986 Edition, IAEA Vienna, p.197) a magnetically driven C-gun having a high helicity injection efficiency is used.

In the content of second section; after the system description of SK/CG-1 machine and the diagnostic techniques used, the conserver as an active device having its poloidal frame and self generated multipole cusped toroidal field are submitted. The phases in operation are given in third section. Here, the on-set, the turbulence and sustainment phases are discussed. In the fourth section the heating and confinement mechanisms are evaluated depending on the self organization principle. Lastly, the conclusions are presented in section five.

2. SK/CG-1

2.1. System and Diagnostics

The C-gun is a noval and alternative version of conventional magnetically driven plasma gun. The vertical two electrodes at the toroidal plane of flux conserver are the main structure of the C-gun.

The electrical characteristics of the gun conforms with the Critical Damping and Under Damping Modes (COM, UDM) of operation depending upon the back ground gas pressure ranges of 40-70 mTorr and 75-250 mTorr respectively. The operating period of the C-gun at UDM is about 10-12 us, but at the COM this is only 2-3 us.

Without using crowbaring technique, the operational regime of this gun can be transformed from UDM scheme into COM. For this purpose, it is sufficient only to be changed the back ground gas pressure. The working voltage range of the C-gun is 10-15 kV and maximum operating loop current is 150 kA. The logarithmic decrement is about 0.58 at the UDM.

The main diagnostic techniques used are: the Langmuir and the floating surface electrical probes; the magnetic probes and the paramagnetic loops including the L-R fast operational integrators, the charge exchange cell and lastly the visible light and high frequency spectrum analyzers. In addition, for data recording and processing the fast analog and digital storage oscilloscopes are used.

2.2. Poloidal Frame

The vertical electrodes of the C-gun are in the plasma slab where it extends from the center of octagonally shaped flux conserver up to the vertical surface of conserver which is inductively coupled to the back strap. Cross section of this plasma slab is the poloidal frame of the C-gun. The boundaries of the poloidal frame are defined by the surfaces of flux conserver and the passive rod located across the top and bottom walls of conserver at the centre. Thus the poloidal frame is a closed loop.
On the construction of flux conserver there exist two symmetrical poloidal frames at the angular distance of 180 degrees.

2.3. Self Generated Magnetic Field

Due to the inductive coupling between the back strap and the vertical surface of conserver and the eddy currents on it, during the operational phase of the C-gun, the conserver becomes active. By the calculations using data measured, the following typical results are obtained. The minimum energy storage time is 288 ns; the characteristic impedance is 106 ohms; the LC time constant is 2.12 ns (471 MHz); the storage energy is 23.4 Joules for a gun current of 150 kA; the L/R time constant of conserver with decaying flux density is 28-30 ms for I/e.

By the experiments made using a practical network model of flux conserver, it is understood that the eddy currents on flux conserver can produce a multi pole (1.16 kG/pole) toroidal cusped type of magnetic field which affects the plasma directly.

3. PHASES IN OPERATION

At the beginning in 0.35 us, from the C-gun electrodes breaking the plasma belt which contains of shock heated electrons having 60-80 eV energies created by the gun itself and therefore producing the electromotive force up to 70 kV (Fig.1), 5 kV/cm typically, the electron-ion two stream instability is initiated. Just after this, the ion sound instability is formed by the electrons with thermal velocities. In the same duration, due to the conventional beam plasma instability and under the influence of toroidal magnetic field generated by the active poloidal frame of the C-gun, the ion sound wave instability is sustained in a configuration of two phase. Thus after the on-set phase the turbulence regime starts. In this phase, the anomalous plasma impedance is approximately 20-30 ohms.

Fig. 1. Variation of plasma belt potential in time when C-gun is open circuited. $p = 40 - 45$ mTorr.
X axis = 200 ns/div., Y axis = 20 kV/div.
In the content of beam plasma instability, W3 ion cyclotron wave packet in the frequency range of 10-15 MHz has a pulse power of 30 MW and it damps in 1-2 us as is being in a magnetic beach model and seen in Fig. 2. Because of its special construction, the flux conserver becomes active during the on-set phase. By means of the toroidal magnetic field in poloidal frame and the perpendicular electric field of C-gun with a duration of 200-300 ms, the pressure gradients towards the azimuthal direction are produced.

![Fig. 2. W3 Ion cyclotron wave damping.](image)

Cyclotron frequency = 10 -15 MHz.
X axis = 200 ns/div., Y axis = 50 mV/div.

Consequently in the sustainment phase, the period of resonance produced by the wave coupling between drift wave and ion acoustic wave depends on the electron temperature and the electrical characteristics of the poloidal frame. The drift wave turbulence regime lasts as long as the Bohm diffusion. So this complex and hybrid phenomena is based on the self organization process.

4. HEATING AND CONFINEMENT

It should be noted that, in SK/CG-1 machine a fast acting valve is not used. All of the operating modes are realized by the back ground gas pressures. In other words, at the beginning the plasma belt in the poloidal frame is surrounded by a gas mantle. Thus all the plasma systems in the flux conserver such as spheromak, spherical torus and drift mode are protected from the conserver by the gas mantle and therefore MHD instabilities are reduced.

On the other hand, before the turbulence and micro instabilities in drift mode, the neutral-electron and electron-ion collision frequencies are increased at the end of the interaction between the hot plasma belt and the gas mantle. As a consequence, the threshold plasma resistivity is found much higher than the Spitzer's resistivity.

The magnetic loop and magnetic probes do not give any response except the on-set phase. Figure 3 represents a probe signal as a typical example. Thus, at the on-set phase the shock heating may...
come into existence together with the ohmic heating. It is possible to calculate \( T_e \) the electron temperature of the shock heated plasma belt using the experimental data obtained. For a constant ionization probability, as a function of \( E/p \) and the aspect ratio of the belt h/w, it is found 60-80 eV where \( E \) is the electric field of the C gun at the onset phase, \( p \) is the operating gas pressure and lastly \( h \) and \( w \) are the height and the width of the belt respectively.

The hybrid instabilities and the signals in sustainment phase can only be detected by the electrical surface probe. As a matter of fact, the axial and radial voltage gradients of the wave frame at a fixed position of electrical surface probe can be characterized by the expression in the forms

\[
dV/dr = E(r) = \left( \frac{C(r)}{T_e^{3/2}} \right) Q_e N_e \left( k \frac{T_e}{M_i} \right)^{1/2}
\]

\[
dV/dz = E(z) = \left( \frac{C(z)}{T_e^{3/2}} \right) Q_e N_e \left( k \frac{T_e}{M_i} \right)^{1/2}
\]

where \( Q_e \), \( N_e \), \( k \), \( T_e \) and \( M_i \) are the electronic charge, the electron density, the Boltzmann's constant, the electron temperature and the ion mass respectively. As a typical result, the average electron temperatures between 17 and 25 eV are calculated. The detailed traces from the electrical surface and the Langmuir probes are given in Fig.4.

The frequency of drift wave turbulence changes in the range of 25 - 5 kHz depending on the electron temperatures between 25 and 5 eV. This result is indicated clearly in Fig.5a. Figure 5b shows the onset phase of this signal.

Angular variations of the surface probe affect the signals obtained in only the case of phase shifts and it is not observed any difference on the amplitude of the signals. Figures 6 (a,b) show these characteristic results. Consequently, it is understood that the surface probe takes place in the electric field of the drift wave turbulence frame.

The confinement mechanism has not yet been exactly clear. Nevertheless, this mechanism may be realized in the sustainment phase due to the flux conserver can become active and the energy can be stored in the belt at the on-set phase. The plasma energy
Fig. 4. A typical result from electrical surface probe. X axis = 2 ms/div., Y axis = relative units. Lower trace indicates the variation of electron density in time taken by Langmuir probe at space potential. Here the maximum electron density is approximately $5 \times 10^{14}$ cm$^{-3}$ at $p = 35 - 40$ mTorr.

Fig. 5. Signals obtained from charge exchange cell at $p = 40-45$ mTorr. (a) X axis = 1 ms/div., Y axis = relative units. (b) X axis = 20 us/div., Y axis = relative units.
Fig. 6. Typical electrica surface probe signals at the beginning of sustainment phase. $p = 40$ mTorr.
X axis = 10 us/div., Y axis = relative units.
(a) Reference position,
(b) after 180 degrees rotation with respect to (a) position.

balance on the SK/CG-1 shows that the global energy confinement times can be increased in the range of 60-250 us.

5. CONCLUSIONS

The main results achieved on SK/CG-1 have been: i) During the on-set phase (CDM, 0.35 us), the current on external circuit of the C-gun transforms onto closed loop current and than the current flowing on the gun fed by the capacitor bank circuit is suddenly broken and the self organization process starts. Due to the electromotive force between gun electrodes, a potential difference up to 70 kV (5 kV/cm) and thus a turbulent plasma column are produced. In this phase, the anomalous plasma impedance is about 20-30 ohms; ii) Parallelly, the beam plasma instabilities and consequently the ion cyclotron, ion acoustic waves and an electron pressure gradient along the azimuthal direction are come into existence; iii) The time and frequency domain analysis of drift wave turbulence oscillograms show that the electron temperatures of 17-25 eV at that of the beginning are about 2-5 eV when it draws near the damping case. This characteristic result do not vary in the sustainment period of 5-25 ms; iv) It is observed that the sustainment period depends
on the on-set gas pressure and the electric field strength in poloidal frame; v) As an average electron density, $10^{14}$ cm$^{-3}$ is approximately found. The maximum diffusion rate is limited by the initial electron density and the critical value of toroidal magnetic flux density in time at poloidal frame; vi) The sustainment phenomenon is independent from the recombination of three body collision scheme. The recombination time corresponding to the pressure range of 30-50 mTorr is almost 607 us. Thus it is exactly understood that this mechanism is not an afterglow. vii) Due to the sustainment period can be controlled, for the optimum conditions it is found about 100-150 ms. The equilibrium condition depends upon the quasi steady state self organization wave coupling; viii) It is determined that the input turbulent power on the global energy confinement time is an important parameter. So it is possible to increase the global energy confinement time up to 250 us at the experimental conditions of SK/CG-1; ix) In the sustainment phase, the wave impedance of the gun increases up to kohms orders due to the wave coupling.

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DESIGN AND FABRICATION OF THE HBT-EP TOKAMAK: A LARGE-ASPECT RATIO, HIGH BETA, 10 msec PULSED TOKAMAK

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Abstract
An upgrade of the Columbia HBT tokamak is now under construction. Called HBT-EP (for extended pulse), it uses the toroidal field field coil set from the former CLEO tokamak/stellerator at Culham which, together with the HBT power supplies, permits significantly extended performance over the HBT device. HBT-EP will have a major radius of \( R = 0.92 \) m, an aspect ratio between 6 and 7, a toroidal field capability to 1 T, and an ohmic pulse length of \( \geq 10 \) msec. Other key features of the device include (1) an adjustable internal conducting shell for external mode stability studies, (2) a variable plasma cross-section, and, ultimately, (3) 2 MW of pulsed ICRF heating for pressure profile control and adjustment. Since HBT-EP is presently under construction, this report emphasizes our design procedure and calculations. Since HBT-EP is pulsed and designed to operate at high density, a rapid high-voltage start-up technique and a very low impurity environment are needed. A novel segmented vacuum chamber has been designed to be compatible with these requirements as well as the need for easy assembly.

1. Introduction

The HBT and TPE-2 tokamaks have demonstrated the ability to sustain high-beta operation near and above the Troyon beta limit using only ohmic heating\(^1,2\). In both cases, high-beta occurred during pulsed operation (100-500 \( \mu \)sec) at densities higher than that which would be normally expected by applying conventional Murakami and Hugil limits\(^3\). Although the precise reasons permitting this ultra-high-density operation are not known, we believe discharge lifetimes short compared with resistive time scales may be responsible. This unusually high density has allowed both devices to study ideal MHD instabilities, and, in the case of the smaller (and cooler) HBT device, internal magnetic probes have been used to measure the internal structure of fast-growing, external kink instabilities for the first time\(^4\).

HBT-EP has been conceived to extend this operating regime to longer discharge lifetimes (the acronym "EP" refers to "extended pulse") and to a larger plasma size (permitting improved diagnostic access). The updated experiment will make use of toroidal field coils from the CLEO tokamak/stellerator built at
Culham laboratories\(^5\), a fast ohmic heating system capable of rapid discharge startup, a flexible poloidal field coil system permitting control of plasma position and shape, and a modular vacuum vessel. A major goal of the first phase of the HBT-EP operation is to reach the Troyon beta limit with peak electron temperature near 100 eV using only ohmic heating. As the impurity level increases beyond about 0.2% of oxygen, high beta operation will still be possible at either higher density and lower temperature or by using auxiliary ion cyclotron resonance heating.

In the following sections, we describe both the mechanical features of HBT-EP and the calculations performed in order to determine attainable plasma parameters.

2. Description of HBT-EP

Figure 1 shows a top view of the HBT-EP tokamak illustrating the toroidal field coils and the modular vacuum chamber. The twenty TF coils are grouped into ten coil pairs, and each coil pair is linked by one stainless-steel vacuum chamber segment. The vacuum segments are made from standard 20 inch diameter 90° elbows significantly reducing the cost of fabrication. This modular arrangement allows easy assembly of the tokamak, since each coil pair can be handled separately, and it allows the use of large diagnostic ports. Each vacuum segment is electrically-insulated and supported at both ends by two rigid supports, one connected horizontally from the inside of the vacuum vessel to the large centering ring and the other connected vertically from the bottom of the vacuum vessel to the cement experimental basepad. These supports also prevent relative motion between segments.

After assembly and alignment of the TF coil pairs, the ten vacuum segments will be connected by five straight stainless-steel spool pieces and five straight sections of quartz tubing. The quartz segments reduce the peak electric field strength along the surface of an insulator to less than 45 V/cm at the fastest toroidal current ramp-up (0.4 MA/msec). Double, pre-baked Viton O-rings are used to seal the quartz tubes to stainless-steel sealing rings, creating ultra-high vacuum conditions while retaining the many advantages of a modular design. Since the seals for the quartz tubes allow for axial motion, a separate bellows is not needed to compensate for manufacturing tolerances.

The poloidal field system is shown in Figure 2. The PF coil sets consist of a six-turn ohmic heating coil, a five-turn vertical field coil, and a multi-turn shaping coil system located within the bore of the TF coils in order to strongly couple to the plasma. The ohmic heating coil has been designed to eliminate the poloidal field
with the plasma, and the vertical field coil system has been designed to eliminate the mutual inductance between the OH and VF coils allowing fast plasma start-up. The shaping coils can be used to produce circular, dee-shaped, and diverted discharges. The coil locations and the magnitude of the currents were first determined with static equilibrium calculations. The start-up and control of the equilibria were then verified using TSC.

The two other major vacuum components are the conducting shell segments and the discharge limiters. Figure 3 shows a cross-sectional view of the adjustable conducting shell segments. Twenty segment pieces will be cut from spun, 3/8 inch thick, aluminum, and the aluminum will be nickel plated. The discharge limiters will be located in the straight stainless-steel spool pieces, and they will be made of stainless steel due to the expected low edge electron temperature and the need for precise density control.

3. Vacuum Procedures

Two 1000 l/sec turbomolecular pumps and several titanium getters will be available to maintain the base pressure of HBT-EP. Only the surfaces between the conducting shell and the outboard side of the vacuum vessel will be gettered in order to keep the surfaces facing the plasma (as well as the quartz tubes and the limiters) free from titanium. Using well-known measurements of outgassing rates, the base-pressure of HBT-EP should be a few times $10^{-8}$ T within five to seven days after a major opening—even without the use of gettering or discharge-cleaning. With titanium gettering, the base-pressure will likely fall by an additional order of magnitude.

After a major vacuum break, HBT-EP will reach a pressure of $1 \times 10^{-7}$ T within 24 hours, and glow discharge cleaning will proceed immediately according to the procedures described by Dylla\(^6\). The chamber walls will be warmed to 100 °C, since this removes oxygen and water impurities from surfaces up to ten times faster than room temperature walls. After a few days, titanium gettering will be applied and several discharge cleaning runs will be performed. We also plan to investigate schemes used to deposit boron onto the ungettered surfaces facing the plasma.
4. Plasma Performance

The operating plasma parameters and the role impurities play in reducing the temperature and beta of the HBT-EP discharges has been investigated using (1) a simple, zero-dimensional, global power balance, and (2) time-dependent, 1-D transport calculations. The global power balance relation readily permits the inclusion of empirical confinement laws and the ability to benchmark these relationships with previously obtained experimental results. The conduction losses are modeled by a combination of the neo-Alcator and L-mode scaling laws (proposed by Goldston to model energy confinement both of auxiliary heated discharges and of high-density, saturated ohmic discharges\(^7\)) and neo-classical ion conduction. The convection losses are assumed to result from a particle confinement time which is 10% of the global energy confinement time for densities below the Murakami limit and 100% for larger densities. This approximation is based on our high-density experience with the HBT tokamak. The radiation losses are modeled in a standard manner. The model has been shown to reproduce the parameters of the HBT and TPE-2 high-density tokamaks.

Figure 4 shows the results of applying this same model to the future HBT-EP discharges. For this figure, we use the initial 0.6 T toroidal field strength and examined the temperature and normalized beta for impurity fractions of 0.1%, 0.2%, and 0.5%. The results illustrate a significant reduction of plasma temperature at large impurity fraction and high density and when the plasma temperature is at the peak of the “radiation barrier” for oxygen \((T \sim 15 \text{ eV})\). However, for lower impurity fractions and lower densities, the model supports our design goal of reaching high-beta ohmic operation provided that the impurity fraction is less than 0.5%.

The time-dependent, 1-D transport calculations were performed using the HERMES code. This code was developed by Hughes in order to model the effect of impurity radiation in the conceptual test reactor INTOR\(^8\). When simulating the ohmic startup of an HBT-EP discharge, the HERMES code held constant the plasma current and density and calculated the time evolution of the electron and ion temperatures and the impurity concentration. The impurity influx was pre-set to be a constant fraction of the escaping plasma effluent, and the impurity specie was chosen to be oxygen. When the impurity influx was small, the impurity fraction within the plasma would also remain small, and the plasma temperature would increase from an initial 5 eV to a value comparable to that predicted by global power balance. On the other hand, when the impurity influx increased so would the impurity density within the plasma. If the radiation losses approached the central ohmic heating power, the electron temperature would pre-maturely saturate relative to the lower impurity concentrations. Since the temperature profiles between “radiation-dominated” and non-radiation dominated discharges are easy to distinguish, the code was used to determine
those conditions enabling both high-beta (near the Troyon limit) and high-temperature (~ 100 eV) ohmic startup. These calculations indicate that high-beta and high-temperature ohmic operation will be possible in HBT-EP provided that the impurity fraction remains below 0.2%. Details of the calculations using this code for the HBT-EP tokamak have been reported separately. Although this low impurity fraction is sometimes difficult to achieve in conventional low-density tokamaks, the higher-density planned for HBT-EP should be able to achieve these low impurity influx rates provided that the ultra-high vacuum procedures described previously are followed.

5. Conclusions

The HBT-EP tokamak has been designed to explore MHD instabilities near the Troyon beta limit and to reduce the impurity influx to the lowest possible level. The low impurity influx rates will be achieved by operating at ultra-high vacuum levels using double O-rings on all large seal surfaces, titanium gettering, and glow-discharge cleaning. The modular design was chosen because of its low cost, easy handling, and good diagnostic capability. Using empirical confinement models, the average equilibrium discharge parameters were estimated for HBT-EP. In addition, by using a time-dependent one-dimensional transport simulation, we were able to determine the maximum impurity influx compatible with high-beta and high-temperature operation. Impurity fractions of the order or less than 0.2% are required in order to obtain these conditions (with ohmic heating alone) in HBT-EP.

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