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I) Introduction:

Ohmic discharges in TORE-SUPRA are sufficiently long (≈ 6 s) for runaway electrons (R.E.) to reach a steady energy state: their energy limit is determined by the balance between parallel electric field acceleration (20 MeV/V.s in TORE-SUPRA) and radiation losses due to the curvature of the trajectories. When R.E. energy is supposed to be only parallel, this provides an estimate of order of 70 MeV (value usually called "synchrotron limit") reached in less than 2 seconds. Experimental observations on TORE-SUPRA of photoneutron emission together with residual induced radioactivity in the first wall components tend to prove that the actual value is much lower than 70 MeV (i.e., 15-35 MeV).

Earlier observations in ORMAK [1], PLT [2] and TFR [3] already showed R.E. energy slightly less than expected from standard loop voltage acceleration calculations. Explanations given for this lack of energy (as skin-effect lowering the electric field during the ramp-up phase or balance between continuous creation and losses) seem not to hold on TORE-SUPRA and therefore another mechanism must be considered to explain the R.E. energy limitation.

II) Experimental observations:

Ohmic discharges obtained with helium gas in TORE-SUPRA [4] exhibit a small amount of R.E. identified on both hard X-ray monitors and neutron detectors. The time behavior of a typical photoneutron signal is shown on figure 1 together with plasma current. Analysis of many shots shows that:

- R.E. are likely produced at the start-up phase of the discharge, when the break-down voltage reached values as high as 25 volts, for the following reasons: (i) when M.H.D. activity in the current ramp-up phase is strong, all R.E. are lost and no further production is observed afterwards; (ii) the photoneutron signal starts ≈ 30 ms after the plasma current which indicates an extreme early creation in addition to a high accelerating voltage during this period of time to reach the photonuclear reaction threshold (≈ 8 MeV) so rapidly (electric field must be at least 70% of the loop voltage).

- The photoneutron signal increases during the rising phase of the plasma current and sharp spikes (≤ 1 ms) are observed during it, corresponding to
extra R.E. losses when the edge safety factor $q_e$ is integer. This feature will not be discussed in this paper.

At the current plateau, the photoneutron signal exhibits an exponential decay with an e-folding time of order of 1 second. Assuming that R.E. have reached a steady energy state, this time can be interpreted in terms of a confinement time of the R.E. population.

On short discharges (as that shown on figure 1), there remain a sufficient number of R.E. during the decay phase of the plasma current so that a positive slope can be easily observed on the photoneutron signal for the lowest values of current. This enhanced losses interpreted as classical drift orbit losses suggest a R.E. energy of 10-20 MeV (for residual R.E. lost during this phase). Taken in account the negative loop-voltage during the current decay phase, this suggest that R.E. energy was between 20 and 35 MeV during the plateau phase.

A residual activation spot was identified on the stainless steel wall located in the bremsstrahlung beam behind the outboard carbon limiter. The ratio between radioisotopes corresponding to high threshold (Cr$^{50}$(y,n)$^{58}$: $T = 21.6$ MeV) to nuclei with lower ones (Mn$^{54}$(y,n)$^{54}$: $T = 10.2$ MeV) allows a rough estimate of the R.E. energy which is found to be 25±10 MeV. Note that C$^{14}$(y,n)$^{16}$Be reactions have been also identified in the carbon limiter, proving the existence of electrons at energies larger than 27 MeV.

III) Discussion:

The different observations which have been made on TORE-SUPRA:
a creation of R.E. only at the very beginning of the discharge ,

an efficient acceleration in the ramp-up phase ,

a relatively low mean energy (15-35 MeV) of those R.E. in comparison to classical synchrotron limit , suggests that a slowing down mechanism takes place counteracting the acceleration experienced by R.E. in the parallel electric field . It must be much more efficient than the synchrotron radiation process which becomes predominant only when R.E. energy is close to the limit (70 MeV).

An explanation for such a slowing down could be residual pitch angle scattering processes :

The radiation losses of a highly relativistic electron on a curved trajectory is given by: \( \frac{dy}{dt} = K \gamma^2/R \) where \( K = 5.632 \times 10^{-7} \text{ m}^2/\text{s} \), \( \gamma \) is the energy on mass ratio of the electron and \( R \) the curvature radius of the trajectory . In the classical limit calculation , \( R \) is taken equal to the field line curvature , i.e. roughly the large radius of the tokamak . But , with a small perpendicular velocity , the actual curvature of the trajectory can be strongly reduced , due to the small value of the larmor radius \( \lambda \) (on TORE-SUPRA , \( \lambda = \gamma \) when expressed in mm ) . The increase in radiation power is \( (1/R^2 \sin^2 \theta / \lambda^2) \) where \( R \) is the curvature of the field line (i.e. \( \approx 2.5 \text{ m} \)) and \( \theta \) the pitch-angle . Because of the large value of the \( R/A \) ratio , an \( 8^\circ \) pitch-angle is enough to reduce the radiation limit from 70 MeV to 25 MeV.

The question is then raised of the origin for such a process , and a quantitative analysis has been carried out to compare different candidates for pitch-angle scattering : magnetic fluctuation (\( 8B \approx 10^{-3} \text{ B} \)) and electric fluctuation (\( 8E \approx 1 \text{ V/mm} \)) effects have been calculated to be many orders of magnitude under the residual coulomb collisions . The corresponding diffusion coefficient is \( D = 8mc(Z+1)n_0^2 \ln \Lambda/\gamma \approx 3.35 \times 10^{20} n^2 / \gamma^2 \) (rd^2/s) , where \( n \) is the electron density (m^-3) . The mean pitch-angle reached by R.E. at equilibrium between diffusion and electric field (which reduces it ) is around \( 60^\circ/\gamma \) , a value much lower than the \( 8^\circ \) needed.

The R.E. energy is shown on figure 2 as a function of time for different assumptions illustrating the cases discussed before:

(1) Pure acceleration by the electric field (the experimental loop-voltage signal from shot 485 has been used , as shown on the bottom curve ).

(2) With addition of synchrotron radiation , at zero pitch-angle .

(3) The same as (2) , but with a 30% reduction in loop-voltage (to evaluate the role of a skin-effect ) : increase in energy is delayed but final values are not very different.

(4) With full loop-voltage , and synchrotron radiation increase due to the mean pitch-angle obtained with coulomb scattering (i.e. \( \approx 60^\circ/\gamma \) ).

(5) Full loop-voltage and synchrotron radiation obtained with a permanent pitch-angle around \( 7^\circ \) . This last value as been adjusted to insure a final energy of 15 MeV , coherent with experimental value (see § II). Note that this value cannot be obtained in any other cases (1 to 4).
IV) Conclusion:

Two different observations have shown an important lack of energy for runaway electrons confined in TORE-SUPRA. This has been assumed due to a small pitch-angle scattering (a few degrees), and many candidates for this have been compared: the strongest known one, collisions, seems not to be enough by an order of magnitude.

Density and magnetic scans on TORE-SUPRA will be needed to discriminate between enhanced collisional scattering processes ($D_c \propto n$) and purely magnetic phenomena.

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A POSSIBLE EXPLANATION FOR THE RUNAWAY ENERGY LIMIT: RESONANT INTERACTION WITH THE RIPPLE FIELD

L. Laurent, J.M. Rax

1- Resonance condition

In the rest frame of the guiding center motion the ripple of the magnetic configuration and its $k^{th}$ toroidal harmonics behave as intense electromagnetic waves. They are characterized by $\omega = \frac{kNc}{R}$ (R is the tokamak major radius, c the velocity of light, N the number of toroidal coils). For a given electron energy this frequency match the electron cyclotron one $\omega_c = \frac{eB}{m\gamma}$ ($\gamma mc^2$ is the usual relativistic energy). For the TORE-SUPRA case the resonance condition is

$$mc^2\gamma = \frac{eBrc}{kN} = \frac{40.B[T]}{k} [\text{MeV}]$$ (1)

Taking present TORE-SUPRA discharges ($B = 1.8$ T) and the $k = 2$ resonance one obtains critical energy about 36. MeV which agrees with the observed energy limitation. Thus we are naturally led to investigate the resonant interaction between the runaway population and the ripple. This interaction which can be viewed as an intense cyclotron heating in the moving frame appears in the laboratory frame as an angular diffusion ultimately responsible of the radiative slowing down. This picture is meaningfull only if a phase randomizing mechanism allows an irreversible transfer of energy between the wave and the particle. This is investigated in section 2.

2- Stochasticity threshold and quasi-linear diffusion

Because of the rotational transform, and the poloidal variation of the ripple field, the magnetic perturbation experienced by the fast electrons can be expanded as a set of waves with different $k_{||}$.

4 bis
where $s$ is the curvilinear abscissa of the guiding center, the integer $p$ represents the $p^{th}$ poloidal Fourier component of the field and $q$ is the safety factor averaged over the trajectory. This field induces neighbouring resonances which may overlap. Using the usual Hamiltonian formalism, the associated Chirikov parameter is derived:

$$S = \sqrt{\frac{a_{pk}V_\perp}{c}qkN}.$$  

(3)

Under the previous experimental condition with $k = 2$ this parameter is larger than 1. When this overlapping criterion is fullfilled an irreversible transfer of parallel energy into perpendicular one occurs. An angular diffusion coefficient describes this process:

$$D_{\theta\theta}[S^{-2}S^{-1}] = \frac{2\pi a_{pk}q^2N^2c}{R}.$$  

(3)

For typical TORE-SUPRA conditions the value of this coefficient reaches values larger than $10^6$ $S^{-1}$.

3- Conclusion

The resonance (1) in the ripple field (fundamental or harmonic $k$) plays a major role in the runaway dynamics. Because of its strength and reasonable corresponding energy, the $k = 2$ component is the leading term. Around this resonance the electron behaviour is driven by three effects, the loop voltage acceleration, the previous diffusion and the cyclotron radiative cooling (approximately at constant pitch angle). When the diffusion coefficient is large enough the electrons cannot reach energy larger than the resonant one. This model explains the energy limitation in TORE-SUPRA. This limitation is expected to disappear in the high field operating regime because the usual synchrotron limit will appear before the resonance.
1. **Introduction.**

Tore Supra (1) is a large tokamak ($R=238$ cm, $a=75$ cm) with superconducting coils. This tokamak is designed for quasi-steady-state operation and therefore the properties of the edge plasma and its interactions with edge components are of primary interest. All surfaces in contact with the plasma are of graphite construction. In its first year of operation only a small number of reproducible shots have been obtained. However, these shots are representative of plasmas in the ohmic phase of the machine and deserve to be studied before the next phase of operation with additional heating. In the initial experiments the plasma was leaning on the inner axisymmetric graphite wall, which covers about $+45^\circ$ poloidally. To simplify particle control, the working gas in these discharges was helium. Past studies of particle confinement times in tokamaks have involved hydrogen, and the collisional-radiative model of Johnson and Hinnov (2) has been employed to relate the measured intensity of H-lines to the ground state population $n(1)$. The electron source $n(1)\lambda e_{\text{eff}}$ is then calculated using an effective rate coefficient $\lambda e_{\text{eff}}$. In this study the model is extended to helium by including the work of Drawin et al (3,4). For the first Tore-Supra discharges only one line of sight was available. The spatial distributions are modeled by a 3D neutral transport code which is coupled to a 2D equilibrium code. In a previous study this code was benchmarked to experimental results by modeling complete spatial ionization profiles from TEXT (5).

2. **Experimental setup.**

A 0.64 m Czerny-Turner spectrometer equipped with an Optical Multichannel Analyzer (OMA) and a grating with 2400 gr/mm (dispersion=3.14 Å/pixel) is used. Due to the fused silica lenses of the telescope ($f/2$, $f=50$ cm) and the silica fibers (PCS-1000, $d=1$ mm) the spectral range available extends from 300 to 800 nm. Each fiber views a chord through the plasma of about 5 cm in width. Nine radial views are achieved. They scan the bottom half of the poloidal cross-section.
The first is slightly above the horizontal midplane, and the last is just outside the last closed magnetic surface. This system is absolutely calibrated by means of a tungsten filament lamp and a Lambertian diffuser plate, both calibrated against standard sources. For the calibration, the diffuser plate was placed inside the tokamak at the actual center of the plasma chord observed.

3. Experimental results:

In this paper we study some typical discharges during the early phase of TS operation. Figures 1 and 2 show plasma current (\(I_p\)), line-averaged density (\(n_e\)), loop-voltage (\(V_p\)) and a typical Halpha signal. During the current plateau (\(I_p=600\) KA) a gas injection is used to raise the density from 1.5 to 2.1x10^{13} (cm^{-3}). The radius of the plasma is maintained constant at 70cm with the plasma limited by the inside carbon wall. The continuum signal given by the OMA in the range of 520nm is used to calculate an averaged Zeff. The initial value of Zeff=3 (+50 -102) before gas puffing drops to 2.5 and remains at this value until the end of the current plateau. This calculation uses only the Te-profile obtained with the ECE system, and the absolute value of Te(o) was derived from the transport code (MAKOKOT). This leads to the large error bar quoted above. We have verified spectroscopically that no heavy impurities contributed to the Zeff measurement presented here. The dominant contribution to the Zeff are carbon, oxygen, and hydrogen.

4. Method of analysis:

Neutral impurity influxes are calculated from measurements of absolute intensities of neutral or weakly ionized ions. For hydrogen the model of Johnson and Hinnov (2) has been used. The carbon influx was derived from a simple model using the ratio of ionization to excitation coefficient JET (6). Drawin extended the collisional-radiative model to HeI and HeII. We used the effective ionization rate coefficient \(\text{Seff}\) compiled by Bell and al. (7). We find (fig.3) that the ratio of ionization events per photon, unlike the case of hydrogen, is strongly dependent on the electron temperature. Since we did not dispose of an electron temperature profile we have assumed an electron temperature corresponding to half the ionization potential of the considered ion i.e. 7eV for H, 12eV for HeI and CII, and 25eV for HeII.

Neutral transport modeling is used to relate the locally measured flux to the poloidally averaged flux. In these calculations,
an axisymmetric equilibrium is coupled with a 3D Monte-Carlo code (5) to estimate the relative poloidal distributions of HeI and HeII. HeI is assumed incident from the inner limiter, emitted uniformly in the toroidal direction. An incident energy (typical of reflected particles) of 30 eV is assumed. The calculation uses ionization rates published by Janev and Langer (8) and neglects Helium charge exchange. Ionized Helium is assumed to move on a flux surface, subject to random cross-field diffusion (D=1M2/s). The HeI distribution is localised near the limiter, while HeII exists in a poloidal shell. Temperature and density profiles will affect the quantitative results. The predicted local brightness in the model is integrated along the same chord as used experimentally.

5. Results and discussion.

The particle confinement time has been computed for the typical early Tore-Supra discharge. The ionization source used in the computation includes contribution of not only the working gas He, but also of H which is always present due to desorption from the graphite surfaces. The measured local sources and their ratios to the modeled poloidally weighted sources are shown in the table (shot 570 fig.1,2 and shot 511):

<table>
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<tr>
<th>SPECIES</th>
<th>ne(10^{13})</th>
<th>LAMBDA(nm)</th>
<th>LOCAL SOURCE</th>
<th>WEIGHTING FACTOR</th>
</tr>
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<tr>
<td>H</td>
<td>2.</td>
<td>656.3</td>
<td>4.4 \times 10^{15} cm^{-3}s^{-1}</td>
<td>6.3</td>
</tr>
<tr>
<td>HeI</td>
<td>1.</td>
<td>388.9</td>
<td>5.0 \times 10^{15} -</td>
<td>2.4</td>
</tr>
<tr>
<td>HeII</td>
<td>2.</td>
<td>656.0</td>
<td>4.0 \times 10^{14} -</td>
<td>&gt;1.0</td>
</tr>
<tr>
<td>CII</td>
<td>2.</td>
<td>657.8</td>
<td>8.0 \times 10^{14} -</td>
<td>-</td>
</tr>
</tbody>
</table>

The viewing chord used for this measurement had a normal radius of 44 cm. For the neutrals, the source comes primarily from the inner carbon limiter. For the Tp calculation, the HeII was not modeled, but it can be assumed that its emission is poloidally symmetric due to the high ionization potential (54 eV), and hence the weighting factor is assumed to be close to 1.0.

Analysis of the partial pressures of the recombined gas, right after the end of the discharge has been achieved with an absolutely calibrated quadrupole mass spectrometer. This analysis shows a 91 ratio of H to He atoms. This should be compared with 181 ratio measured from recycling neutrals.

In the computation of Tp the absence of the HeII contribution to the source is overcome with the approximation that 2 electrons are contributing for each He atom. Using only the neutrals, a total
ionization rate of $1.6 \times 10^{21}$ particles/s is found and a particle confinement time of 127 ms is computed. An uncertainty of $\pm 100\%$ is estimated primarily due to the sensitivity of these calculations to the edge electron temperature. This value for $T_p$ is reasonable for a low density, low current He plasma.

Ergodic Divertor Experiments on Tore Supra

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1. INTRODUCTION AND DESCRIPTION OF THE EXPERIMENTAL HARDWARE

The ergodic divertor (ED) configuration is expected to provide substantial benefits in terms of controlling edge parameters and impurity concentrations in high temperature fusion plasmas. Although previous ED experiments [1,2] have provided stimulating results the Tore Supra ED program is uniquely positioned to yield significantly new physics with which to advance our understanding of plasma surface interactions and their relationship to the overall confinement properties of tokamaks plasmas. Tore Supra is a relatively large machine (i.e., \( R_0 = 2.37 \) m, \( a = 0.8 \) m, \( B_T = 4.5 \) T) with a long plasma pulse and high power auxiliary heating. The ED coils are mounted inside the vacuum vessel (6 coils equally spaced toroidally) and are equipped with neutralizer plates attached to pumping channels with getter pumps for exhausting particles [3] and heat. By matching the angle of the field lines on the resonant surface they create a magnetic island structure centered at higher poloidal mode numbers than in previous experiments. An ED layer, consisting of several overlapping island chains, is produced within a well localized radial interval \( \approx 10 \) cm at the edge. These elements provide flexibility for modifying the boundary conditions and interactions with a variety of plasma facing components (pump limiters, RF couplers, ...). In addition to edge studies, the ergodic layer provides an opportunity for MHD and equilibrium experiments (due to modifications in the current profile and boundary conditions) and for runaway electron confinement experiments.

The technical specifications for the ED Coils have been described by Lipa et al., [4]. The key physics points are that the coils produce a relatively broad toroidal (\( n \)) and poloidal (\( m \)) Fourier mode spectrum (e.g., a full poloidal width at half maximum of \( \pm 0.15 m_0 \), where \( m_0 \) is the fundamental poloidal mode number) and can be phased such that \( m/n \) resonant surfaces occur at several \( q \) values between 1.33 and 8. The width of the ergodic layer as well as the degree of stochasticity in the layer can be varied by changing either the ED current or the radial position of the \( m_0/n_0 \) resonant surface with respect to the coils. The overlap or Chirikov (\( S_C \)) parameter may be varied from 0 to approximately 1.7–2.0 with a maximum ED coil current of 45 kA.
2. THEORETICAL HARDWARE AND BACKGROUND MODELS  Plasma effects in the ergodic layer are calculated in terms of the usual magnetic diffusion coefficient \(D_T = L_0 \sum_{m,n} \delta(m-n/q)(b_{m,n}/B_T)^2\), where \(b_{m,n}\) is the resonant part of the imposed perturbation field and \(L_0\) is the correlation length) with a quasilinear model. Such calculations have been validated via a 3-D variational method [5]. Additional analytic and numerical \(E\Phi\) models, including several types of field line tracing/mapping codes, and a modified Grad-Shafranov (G-S) equation are used to guide the experiments. These are briefly discussed here with respect to their implications on various aspects of the experimental program. In particular, the numerical codes are used to evaluate the vacuum magnetic field structure in the ergodic layer as a function of the coil current, direction, and magnitude for various plasma \(q(r)\) profiles while the modified G-S equation is used for MHD stability and equilibrium studies with an ergodic layer. The line tracing and mapping codes give us a surface image of the magnetic patterns on the walls and plasma facing components. This provides an important link between data obtained with visual and infrared imaging diagnostics and the 3 dimensional field structure in the ergodic layer.

The field line tracing codes are also used to evaluate the spectral properties of the \(E\Phi\) coils and to study the size, orientation, and relative coherence of island chains on each resonant surface as well as for the computation of exact field line trajectories to any boundary point. For instance, in one case, a tracing code, has been used to follow field lines started at the \(r=74.45\) cm surface for a distance of 200 meters. The result show that, given 400 initial points, 36% hit the midplane pump limiter. Of this, about 4% enter the limiter throat, 31% hit the leading edge of the limiter blade, and 1% hit the face of the limiter positioned at \(r=74.5\) cm [6].

Global MHD stability and equilibrium properties are also an important consideration for the planning of the \(E\Phi\) experiments, especially with respect to ideal kink modes and resistive tearing modes as the current profile is reconfigured during the formation of the ergodic layer. A modified Grad-Shafranov model, in which an exponential term is used to include changes in the boundary conditions, is employed to study discharge equilibrium properties with \(E\Phi\) layer effects included. This equation is known to exhibit bifurcated solutions [7] indicating that the introduction of the \(E\Phi\) layer may trigger a switch between two possible current profiles. If these solutions prove to be stable, we may find that one profile is highly peaked (giving a \(q\) significantly less than 1 on axis) while the other is relatively flat similar to those observed during L and H-mode discharges. In addition, the model may be used to determine a parametric condition for the loss of toroidal equilibrium. A stability parameter applicable to the \(E\Phi\) geometry includes changes in the effective plasma radius with the introduction of the ergodic layer.
3. ERGODIC DIVERTOR PARTICLE CONFINEMENT TIME MODEL. An essential part of the \( \text{ed} \) experiments is to determine the effects of stochastic fields on the behavior of heat and particle transport (especially impurities) to the machine walls and plasma facing components. Thus an experimental figure of merit \( \Sigma = \frac{\delta C}{T_p \text{ed}} \), where \( T_p \text{ed} \) the \( \text{ed} \) layer particle confinement time:

\[
T_p \text{ed} = \frac{N \text{ed}}{\frac{2S_a A}{a^2} \int_b^a S \, dx} = \frac{N \text{ed}}{\frac{2S_a A}{a} \int_b^a S \, dx}.
\]

provides a useful measure for quantifying the ergodic divertor performance. In this form \( T_p \text{ed} \) is easily evaluated when it is known the source term, \( S = n_0 \overline{v_e} \text{ion} \) is localized within the \( \text{ed} \) layer. Under this assumption the \( H_\lambda (\lambda = 6563 \text{ A line}) \) intensity (I) provides a direct measure of \( \int_b^a S \, dx \) where \( I \) is simply expressed as the ratio of \( \overline{v_e} \text{exc} \) to \( \overline{v_e} \text{ion} \) times \( \Gamma_0 B_{32} (4\Pi)^{-1} \) and \( B_{32} \) is the branching ratio from the excited state to either the \( n_2 \) or \( n_1 \) ground state. \( N \text{ed} \) (the number of particles in the \( \text{ed} \) layer) is determined by density profile measurements. The source localization assumption is justified if the density and temperature in the \( \text{ed} \) layer are high enough that the ionization mean free path (\( \lambda_0 \)) of the neutral flux from the walls satisfies \( \lambda_0 < \Delta \) where \( \Delta = a - b \) is the width of the \( \text{ed} \) layer. For low \( \text{ed} \) layer density or temperature the denominator may be replaced with an expression of the form \( \int_b^a \overline{v_e} \text{ion}/v_0 \, dx \), where \( \Gamma_{\text{exc}} \) is an externally imposed flux and \( R \) is the recycling coefficient. The exponential attenuation factor depends, in the slab approximation, primarily on the electron density profile across the \( \text{ed} \) layer (located with its outside edge at \( r = a \)) and the ratio of the surface areas \( S_b/S_a \) on each side of the cylindrical \( \text{ed} \) layer. Thus it is necessary to measure the \( n_e \) profile and the neutral particle flux, \( \Gamma_0 = \overline{v_e} \text{exc} \) through surface \( r = a \) to obtain \( T_p \text{ed} \). The same approach may be developed for impurity transport studies and correlated with perturbation experiments such as: gas and metal impurity injections experiments. Gas puff, laser blow-off, or pellet injection hardware on Tore Supra provide the tools for these experiments. Finally, this approach can be used to quantify runaway electron confinement with an \( \text{ed} \) layer. The high energy electrons provide a method of isolating parallel transport properties from perpendicular transport more typical of the heavy ions.

4. IMPURITY TRANSPORT AND CONTROL. The dynamics and production of impurities with an \( \text{ed} \) layer is a key technical questions to be
addressed in the Tore Supra experiments. Initially experiments will
study transport, using the models described above. The second goal is
to study production. Therefore, it is essential that we test the
hypothesis that a cold, high density layer, forming near the wall,
will radiate strongly and cool the outflowing plasma particles before
they hit the facing components. If achieved, this will reduce
sputtering losses from the walls and lower the concentration of high-Z
impurity throughout discharge. Diagnostics designed specifically for
these studies have been implemented on Tore Supra. Finally, the drag
effect [5] is believed to be important for sweeping impurities out of
the discharge and will be studied in detail. The ed layer model given
in section 3 and a heat transport model as in ref. [5] are used to
develop a physical understanding of impurity transport and production.
These will ultimately be linked to more complex models for
impurity transport and production with an ergodic layer but first
experimental data must be acquired and the basic models tested.

5. SUMMARY The Tore Supra ergodic divertor program will concentrate
on particle confinement and impurity control experiments during a
first series of experiments. Numerical and analytical models are
producing results which are useful for guiding the experiments and
which will be used to analyze the measurements. There are a wide
variety of heat and particle transport experiments underway and MHD
studies will follow shortly. The results from these programs will
provide a more complete assessment of ergodic divertor benefits for
high temperature, long pulse fusion reactor plasmas and will yield new
physical insight into the complex domain of plasma surface
interactions.

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MAGNETIC FIELD STRUCTURE AND TRANSPORT INDUCED
BY THE ERGODIC DIVERTOR OF TORE SUPRA

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Introduction
The problem of extracting large power fluxes at the edge of tokamak plasmas with tolerable production of impurities, as well as of improving the plasma pumping efficiency, may be alleviated if the plasma interacts with the wall with a large density, small temperature. Such a configuration should be obtained by magnetic connection between the hot plasma and the wall, if the parallel energy flux induces a temperature decrease along the flux lines, while the pressure nT tends to remain constant. This is achieved with axisymmetric divertors [1,2]. On Tore Supra, we have planned to test that possibility with the ergodic divertor [3], where the magnetic connection is due to a resonant magnetic perturbation which creates a stochastic layer between the hot plasma and the wall. The aim of this paper is to present the expected transport performances of the implemented device in the framework of the quasilinear theory.

Magnetic topology
The coils creating the magnetic perturbation consists of 6 identical modules 60° apart in the toroidal direction (fig.1). To take account of the toroidal effects, it is convenient to change the coordinates from the usual r, φ, R, to a system r', φ', R', where 2πB_ø is the toroidal flux labelling each unperturbed magnetic surface, φ' and R' being the intrinsic poloidal and toroidal coordinates which exhibit a constant 7x, and finally where the average length along flux lines per toroidal rotation is 2πR'. The perturbation of the trajectory of the flux lines, and generally of pure v_ø particles, by a vector potential A_ø^p(r', φ', R') is then calculated as in the cylindrical case. This statement holds when an electric potential perturbation U(r', φ', R') is present.

Fourier analysis of A_ø^p in φ', R' produces a spectrum A_ø^p(M, N). The wave number M along φ' is distributed in an interval 24 ± 5 at the plasma edge r' = 0.8m (M = 10 ± 2 at r' = 0.6m). The wave number N along R' takes values : N=6, 12 or N=3, 9 according to whether parallel or opposite voltages are applied to successive modules. The predominant resonance occurs for q_{edge} = M/N = 4 (N=6) or 2.66 (N=9). For the case q_{edge} = 4 and N=6 the position of the resonance surfaces M/N = 24/6, 23/6, etc., are shown on fig.1. Figure 1 also displays the Chirikov parameter 2B_ø/1d_{ls} where
\[ S_{\text{is}} = \left( \frac{\tilde{\mathcal{A}}(H,N)}{B_{\text{res}}(dq/qdr)} \right)^{1/2} \]

is half the island width, and \( d_{\text{isl}} \) is the radial distance between successive island chains. The island chains strongly overlap within the 0.1 m edge layer, the number of resonant surfaces being \( \sim 5 \), which is large enough to ensure a stochastic diffusion of the field lines. The quasilinear diffusion coefficient of the flux lines

\[ D_{\text{qll}} = \sum_{\mathcal{M}} \frac{\tilde{\mathcal{A}}(H,N)/B_0}{(H/r')(\tilde{\mathcal{A}}(H,N)/B_0)} \]

on a given resonance surface \( \mathcal{M} \) is corrected to account for weak overlapping in the plasma core, yielding the effective value of the diffusion coefficient, \( D_{\text{qll}} \), shown on fig. 2. Note the strong decrease towards the plasma center.

**Heat conduction**

A radial perturbation of the magnetic field allows a radial transport of energy through the parallel motion of the electrons. When the ergodic divertor is activated, the thermal diffusion coefficient \( \chi_{\text{erg}} \) thus combines the stochastic transport process \( D_{\text{qll}} \) and the parallel transport. In the plasma core, the latter is collisionless, and thus

\[ \chi_{\text{erg}} = D_{\text{qll}} \frac{n(2T/m_e)^{1/2}}{\sqrt{\pi}}. \]

It is here essential that \( D_{\text{qll}} \) vanishes towards the center so that the ergodic heat transport does not prevail over the actual turbulent heat transport \( (\chi/n \sim 1 \text{ m}^2/\text{s}) \). In the ergodic layer at temperature \( \lesssim 100 \text{ eV} \), a collisional regime applies. The heat conduction is then

\[ \chi_{\text{erg}} = D_{\text{qll}} \chi_{/qR} h, \]

where \( \chi_{/qR} = 2.10^{22} \text{m}^2/\text{s}^{-1} \), and where \( h \) is a correcting factor depending on \( \log(\chi_{/qR}) \) \cite{4}. Values of \( h = 0.3 \pm 2 \) are expected from 3D calculations \cite{3}. The energy flux \( \phi \sim 0.1 \pm 0.3 \text{ MW/m}^2 \) escaping from the plasma core should thus sustain a strong thermal gradient, \( dT/dr \sim -\phi/\chi_{\text{erg}} \), which is necessary to achieve a density accumulation at constant pressure.

**Plasma convection**

The ergodic layer will maintain a constant pressure profile provided the radial convection has the form

\[ \Gamma_{\text{erg}} = -D_{\text{erg}} \frac{\partial T}{\partial r} \]

with a large enough stochastic diffusion coefficient \( D_{\text{erg}} \), namely \( D_{\text{erg}} \gg D_\perp \), \( D_\perp \) being the diffusion coefficient due to microturbulence. This is the case when the ED acts alone and the relation holds with \( D_{\text{erg}} \sim D_{\text{qll}} \). However when a self consistent electric field, i.e. a radial component \(-\partial U/\partial r\) and a fluctuating component \(-V(\partial U)\) is taken into account, the transverse electric drift motion contributes to \( \Gamma_{\text{erg}} \) and one expects the particle flux induced by the ergodic field lines to be modified. Since constant pressure is at the crux of the effects expected from the ergodic divertor it is important to characterize this change. The electric potential is such that there is no parallel electric current \( J_p \) carried by electrons, thus using the Braginskii coefficient \( \beta_e \):
\( \rho J e / T = - (\nabla P) / (n - e) \) \( (\nabla U) / T - \alpha (\nabla T) / T = 0 \); \( \alpha = 1 + \beta_v = 1.7 \div 2.5 \) (3)

This expression is split into fluctuating (\( \tilde{n} \), etc.) and non-fluctuating (\( n \), etc.) parts. In the quasilinear framework the plasma continuity and dynamical equation along the field lines are linearized and one obtains for each Fourier mode \((M,N)\):

\[
-2nT \left[ i k_x(- + i + \frac{\partial}{\partial T} + \frac{\partial}{\partial \tau'}) \right] - i n k_x B_0 \frac{\partial U}{\partial \tau'} \tilde{\nabla}_n \tilde{\nabla}_n = 0 \quad (2a)
\]

\[
-i k_x \tilde{\nabla}_n \tilde{\nabla}_n - \frac{k_x}{B_0} \left( \frac{\partial}{\partial T} - \frac{\partial}{\partial \tau'} \right) \tilde{\nabla}_n \tilde{\nabla}_n = 0 \quad (2b)
\]

where \( k_x = m_i \omega / c \), \( k_x = (M + N)/R \), where \( m_i \) is the ion mass, and where the tilde symbols now means the Fourier components \((M,N)\). The coefficients \( \gamma_v, \gamma_n \) are positive and reflect viscosity and density damping at small scale. They are assumed small. The quasilinear radial flux \( \delta B / B = \langle n - \nabla \phi \rangle + \frac{\delta \bar{E} \bar{B}}{B^2} \) is obtained from (3a,3b) in the form:

\[
\Gamma_{\text{erg}} = \sum_{M,N} \frac{i k_x n}{B_0} \left( k_x \tilde{\nabla}_n \tilde{\nabla}_n \right) + \text{c.c.} = - D_{\text{erg}} \frac{\partial (n T)}{\partial \tau'} \quad (3)
\]

\[
D_{\text{erg}} = \sum_{M,N,\xi=\pm 1} \pi \psi \delta(k_x^2 + \xi \kappa_x^2) \left( k_x \tilde{T} / (\tilde{T} / \partial \tau') \right) \left[ \frac{\tilde{\nabla}_n}{\partial \tau'} \right] | \omega T / (\omega T / \partial \tau') |^{1/2} \quad (4)
\]

Equation (4) implies a dependence of \( \Gamma_{\text{erg}} \) on the pressure gradient. The perturbation \( k_x \tilde{T} + k_x \tilde{\nabla}_n / n \tilde{\nabla}_n \) drives the resonances is the derivative \( \nabla \phi \tilde{T} \) along the stochastic flux lines. Neglecting \( k_x \tilde{T} \) yields \( D_{\text{erg}} \approx D_{\text{el}} \). However the actual values of \( \nabla \phi \tilde{T} \) experience an attenuation and a radial localization on the resonant surfaces, which involve \( \log (x / x_p) \).

**Conclusion**

We have performed a comprehensive quasilinear analysis of transport in stochastic fields, including the electric perturbation, which will be tested by the ergodic divertor experiments on Tore Supra.

**References**

Figure 1: Geometry of a coil module.

Figure 2: Radial position of resonant magnetic surfaces \( M (N = 6, q_{edge} = 4) \), Chirikov parameter \( S \), and quasilinear diffusion coefficient of flux lines \( D_{QLC} \) corrected to take account of small island overlapping towards the center.
THERMAL EQUILIBRIUM OF THE EDGE PLASMA WITH AN ERGODIC DIVERTOR

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1 Introduction

The knowledge of the possible thermal equilibria of a radiating edge plasma layer is of first importance to understand the thermal plasma stability, which in turn controls various edge plasma phenomena (attached-detached plasma, density limit, etc.). The aim of this paper is to study an ergodic divertor (ED) layer in presence of radiating impurities. The main expected effect of the divertor is to build up a layer with a constant radial pressure profile except in a thin layer where neutral interactions take place. We assume the electron density to be sufficiently high so that the thickness of that neutral layer can be neglected. The number of steady equilibria is discussed in terms of the incident power, impurity species and concentrations. The possibility of bifurcation between equilibrium states is discussed (S shaped diagrams).

In the case of radiation in the coronal limit, we show that the standard situation exhibits only one steady state.

2 Model

We consider the peripheral plasma region $r_t < r < r_Q$ where the ergodic divertor regime is established. Assuming a slab geometry the heat equation is given by:

$$\frac{d}{dr} \Phi(r) = -C_n^2 F(T) ; \Phi(r) = -\chi \frac{dT}{dr}$$

(1)

where $C_n$ is the impurity concentration, $n$ is the plasma density and where $F(T)$ describes the radiation losses. In the collisional regime [1]

$$\chi = 2 \times 10^{25} \frac{T^{5/2}}{D} \frac{q R_0 (SB/B)^2}{\nu}$$

with $D = 2 \pi q R_0 (SB/B)^2$ ; $q$ is the safety factor, $R_0$ the major radius and $SB$ the magnetic perturbation created by the ED coils.

The number of equilibria is related to the boundary conditions. In the present case the heat equation has to be solved for a given heat flux from the plasma core $\Phi(r=r_t) = \Phi_t$, while $\Phi(r=r_Q) = \Phi_Q$ satisfies the sheath conditions:

$$\Phi_0 = \gamma_1 \Gamma_0 T_0 ; \quad \Gamma_0 = \gamma_2 n_0 (2 T_0 / m)^{1/2}$$

(2)

where $\Gamma_0$ is the plasma flux flowing to the wall and where $n_0$, $T_0$ are respectively the density and temperature at the wall. The factor $\gamma_1$ incorporates both the effect of the electrostatic sheath potential and the energy exchange with the neutrals. Similarly $\gamma_2$ takes account of the
density and velocity variation across the neutral layer (the sheath effect yields the sound velocity dependence). The particle flux escaping from the plasma is assumed null. Finally the neutral flux from the wall is taken as the sum of two contributions $\Gamma_0 + \Gamma_{ext}$, where $\Gamma_{ext}$ is an independent source and $R$ the wall recycling coefficient. The particle balance within the neutral layer then determines the flux $\Gamma_0 = \Gamma_{ext} / (1 - R)$.

3 Equilibrium of a radiating layer

Eq. 1 allows a straightforward change to an equation of a moving body in a given potential $W(u)$:

$$\frac{1}{2} \left( \frac{du}{dr} \right)^2 + W(u) = \frac{1}{2} \phi^2$$

$$u = \int \chi \, d\Gamma, \quad W(u) = (n_0 T_0)^2 \chi_{\max} \int_u^{\max} \frac{1}{T^2} \phi(T) \, du$$

The cutoff, $u_{\max}$, which corresponds to a temperature of $\sim 100 \text{eV}$, is due to the strong decrease of $\chi/T^2$ at high temperature when the non-collisional regime is reached ($\chi$ scales as $T^{1/2}$). Finally we restrict the study to equilibria such that $u(r_i) > u_{\max}$, so that $W(u(r_i)) = 0$.

In order to solve Eq. (3), we will calculate $\phi$, and thereby $T$, in view of Eq. 2, for a given control parameter $\phi$, as done below. For a given couple $(\phi, T)$, eq. (3) then turns into an initial value problem for which the solution is unique [2]. The number of solutions of Eq. 3 for a given $\phi$ is thus determined by the number of initial conditions $\phi$ to the problem. Using Eqs (2, 3), the ratio $Q_0 = \phi / \phi_i$ is solution of the following nonlinear equation:

$$q^2 = 1 - \alpha^2 q_0 \, S(q_0) \quad ; \quad S(q_0) = \int_{q_0}^{\max} \int_{T^*}^{T^*} F(y T^*) y \, dy$$

Here $T^*$ is a normalization temperature and the parameter $\alpha$ and $\beta$ are given as follows in MKS units, and normalized to Tore Supra characteristic length

$$\alpha = 5.3 \times 10^{22} (T^*/10)^{0.5} \, \alpha^* \quad \beta = 1.1 \times 10^{23} \beta^* / (T^*/10) (r_i / 0.6) (R_0 / 2.37)$$

with $\alpha^* = \chi_0 / \gamma^{1/2} \gamma_0^{1/2}$ and $\beta^* = (P / 10) / (G_{\gamma} / 3.5 \times 10^{18})$.

Note that Eq (4) is formally independent of $\delta B / B$ and then can be applied to study non ED situations such as those encountered in diverted plasmas [3].

4 Bifurcation of equilibria with different impurity species

We choose the coronal expression for $F_n(T)$ [4] and solve Eq (4) numerically. For carbon impurity only, the solution $q(\beta^*)$ is plotted on fig. 1 for $\alpha^* = 3$. Two bifurcation points are found for $\beta^* = 0.86$ and $\beta^* = 1$. In
the plane \((\alpha^*, \beta^*)\) the curves shown on fig.2 reproduce the bifurcation points of the function \(q(\beta^*)\) for different \(\alpha^*\) values. The two branches of this curves separate the \((\alpha^*, \beta^*)\) plane in two regions: inside the two branches Eq.(4) admits three solutions, outside only one solution exists. The intersection point, \(A\), of the two branches plays the role of a critical point below which no bifurcation can occur when \(\alpha^*\) or \(\beta^*\) are varied. For instance, if \(\alpha^*\) is less than 0.5 no bifurcation occurs whatever the power. The behaviour of an impurity mixture composed with carbon and oxygen where the ratio of concentrations \(c^o/c^e=0.075\), is plotted on fig.3 and 4, showing a more complicated structure. The curve \(\alpha^*(\beta^*)\) exhibits two distinct domains where the heat equation admits 3 steady solutions. For higher \(\alpha^*\) and \(\beta^*\) values areas with 5 solutions can be found. With \(\gamma_1=6\) and \(\gamma_2=1\), high impurity concentration \((c^e/\gamma=20.)\) and high power \(P\geq 24.F_o/3.5 10^{22}\) are needed to get a bifurcation.

6 Conclusion

Within the coronal hypothesis, the main conclusion is that a strong impurity concentration is required for bifurcation to appear when the energy flux is varied. Of great importance in this statement is the high \(\gamma_1\) value imposed by the sheath physics more than the coronal hypothesis. With lower \(\gamma_1\) value, lower impurity concentration will be necessary. Thus the heat equation admits generally only one equilibrium state. This practically guarantees the thermal stability of the layer except in cases of oscillating modes.

References


Figure captions

Fig.1: Variation of \(q^o(\beta^*)\) for \(\alpha^*=2.\) with carbon impurity.

Fig.2: Bifurcation points of \(q^o(\beta^*)\) for different \(\alpha^*\) value and with carbon impurity only. Inside the two branches, the heat equation admits three steady solutions: \(A\) is a critical point below which no bifurcations occur.

Fig.3: Variation of \(q^o(\beta^*)\) for \(\alpha^*=7.45\) and for an impurity mixture of carbon and oxygen (with a relative impurity concentration \(c^e/c^o=0.075\))

Fig.4: Bifurcation points of \(q^o(\beta^*)\) for different \(\alpha^*\) value, in the case of the carbon-oxygen mixture, \(c^e/c^o=0.075\).
NEUTRAL CONFINEMENT IN PUMP LIMITER WITH A THROAT

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1 Introduction

In quasi-steady tokomak operation, active control of recycling is achieved by pumping a fraction of the plasma outflux with pump limiters. In order to optimize the efficiency of this apparatus one can either act on the properties of the plasma SOL, or one can improve the particle collection with the help of the throat leading from the nozzle to the neutralizer plate [1,2,3]. This reduces the amount of escaping neutrals owing to plasma-neutral-sidewall interaction. In this paper, we derive the characteristic lengths of this effect, and we compute the neutral density build-up of pump limiters with a long throat (for a large ratio length L over the transverse dimension a).

2 Geometrical effect of the throat on neutral outflux

The basic throat effect is to reduce the amount of neutrals escaping from the pump limiter without bouncing on the sidewall. This effect depends on the ratio of the transverse velocity (thermal velocity v) to the parallel velocity u and on the ratio of the distance to the sidewall a_w over the distance to the nozzle d=L-x. For shifted Maxwellian distributions the probability P_G to reach the nozzle on a ballistic trajectory is:

\[ P_G(u/v,a_w,d) = \frac{1}{2} \left[ \frac{1+\text{erf}(u/v,a_w/d)}{1+a_w^2/d^2} \right] \left( 1+\text{erf} \left( \frac{u/v}{1+a_w^2/d^2} \right) \right) \]  

For current values of u/v, only the neutrals with a birth point close to the nozzle have to be considered, fig.1. One can then assume that a_w~a/2 and u~0 for neutrals resulting from charge exchange (vanishing plasma average velocity at the nozzle) and that a_w~a and u~V_s (V_s is the sound velocity) for neutrals undergoing neutral-sidewall interaction. The flux of neutrals in the state of ballistic flight to the nozzle is then:

\[ \Gamma_{b_{\perp}} = \int_0^L dx \{ n_p \langle ov \rangle \rangle x \ P_c(u/v=0,a/2,L-x) + n_P (V_s/v,a,L-x) \} n(x) \]
where \( n_p \langle \sigma v \rangle_{cx} \) is the rate of charge exchange, and where \( v \) is the frequency of neutral-sidewall interaction. One finds that the neutrals on ballistic trajectories are created at a distance \( \leq 2a \) from the nozzle. Hence \( \Gamma_{ba1} \) can be neglected when the neutral density is vanishing at \( x \leq L - 2a \), i.e. when the plugging regime, described below, is reached on distances smaller than \( L - 2a \).

3 Plugging length, effect of plasma-neutral interaction

We are interested in the plugging regime where the neutral flux exhibits an exponential decay from the neutralizer to the nozzle with an e-folding length, the plugging length \( L_p \), smaller than the throat length \( L \). This length \( L \) has been derived precisely but assuming no neutral pumping [3]. A simple interpretation of this expression is that the neutrals experience a random walk process with a step \( L_a \) defined by the distance covered before a momentum change, i.e. \( L_a^{-1} = L_{cx}^{-1} + L_p^{-1} \), where \( L_{cx} \) is the charge exchange mean free path and where \( L_p \) is the characteristic length of momentum \( p \) loss on the sidewall \( L_p \approx p/\partial_p \). The number of steps before ionization is \( L_i/L_n \) where \( L_i \) is the ionization mean free path. The plugging length \( L_p \) is thus:

\[
L_p^2 = (L_1/L_n)^2 L_a^2 \quad ; \quad L_p = (L_1 L_n)^{1/2} (1 + L_p/L_{cx})^{-1/2}
\] (3)

To incorporate the effect of neutral pumping in the vicinity of the neutralizer we consider the continuity and mechanical equations:

\[
\begin{aligned}
\partial_x n_n(x) + \partial_x n_p(x) &= P \left[ n_p(x=0) u_p(x=0) \right] ; \quad \partial_x u_p(x) = - \frac{n_n \langle \sigma v \rangle_{cx}}{L_{cx}}. \\
\partial_x (n_n T_n + m_n u_n^2) &= - n_p \langle \sigma v \rangle_{cx} u_p + n_p \langle \sigma v \rangle_{cx} (u_n - u_p) - n_p \langle \sigma v \rangle_{cx} (u_p - u_n) - \frac{n_n \langle \sigma v \rangle}{T_n/m}^{1/2} L_p^{-1}.
\end{aligned}
\] (4)

where \( P \) is the pumped fraction of the plasma flux reaching the neutralizer plate. This set of equations accepts a solution with an exponential decay of the neutral density \( \exp(-x/L_n) \) while the neutral temperature \( T_n \), and average velocity \( u_n \), the plasma density \( n_p \), temperature \( T_p \), and average velocity \( u_p \) are constant. One thus obtains the following decay law and the plugging length:

\[
\begin{aligned}
\frac{\partial x^2 n_n}{x^2} - \frac{p}{L_{cx} T_n} \left( \frac{2T_p}{T_n} \right)^{1/2} \left( \frac{n_p(x=0)}{n_p(x=L)} \right) - \frac{\langle \sigma v \rangle_{cx}}{L_1 L_p} &= 0 \\
L_n^{-1} &= L_p^{-1} + \left( L_n^{-2} + L_p^{-2} \right)^{1/2} ; \quad L_p = \frac{2L_{cx}}{P(2T_p/T_n)^{1/2} (n_p(x=0)/n_p(x=L))}.
\end{aligned}
\] (5)

This expression reduces to (3) when \( P = 0 \) and increases with the pumping coefficient \( P \). This increase is a consequence of the larger momentum exchange between neutrals and plasma due to the unbalanced particle
fluxes. Furthermore the plugging length depends on the plasma density decrease along the throat, which induces a nonlinear dependence on the pumping coefficient. The temperature $T_n$ is derived from a balance between the energy gain by charge exchange and the energy loss on the sidewall hence $T_n/T_p \sim L_k/L_{cx}$, where $L_k$ is the characteristic length of energy loss due to neutral-sidewall interaction ($L_k \sim L_a$ is the smallest length scale). At first order in $P$, one then finds taking $n_p(0)/n_p(L) \sim 1/3$ [3]:

$$L_n = L_n \left(1 + \frac{P L_k L_t}{3 \sqrt{2} L_{cx} L_{cx}} \right)^{1/2}$$  \hspace{1cm} (6)

4 Neutral density build-up

The neutral density build-up at the neutralizer, derived in [3], can be interpreted in terms of a balance between the pressure gradient and the momentum loss on the sidewalls.

$$\partial_z (n_n T_n) \sim \frac{n_n(0) T_n(0) m n_p(0) u_p(0)}{L_n} \sim \frac{n_n(0) u_p(0)}{\tau}; \quad n_n(0) u_p(0) \sim n_p(0) u_p(0)$$  \hspace{1cm} (7)

The time scale of neutral momentum loss is $\tau = L_p/(T_n/m)^{1/2}$, which yields the approximate value of the neutral density at the neutralizer plate:

$$n_n(0) \sim \frac{1}{3} \left(\frac{L_k L_t L_{cx}}{L_{cx} L_p} \right)^{1/2} \frac{T_p}{T_n(0)} \frac{Q}{\left(T_p (T_p/m)^{1/2} \right)}$$  \hspace{1cm} (8)

Here the SOL is determined by the plasma temperature and the available energy flux $Q \sim T n_p(T_p/m)^{1/2} T_p$. The compression for both open and closed configuration are displayed on fig.2. Two domains of improved efficiency of closed configuration with respect to open configuration are found: one at low temperature $T_p \lesssim 10 \text{ eV}$, yielding the highest neutral build-up and one at high temperature $T_p \gtrsim 100 \text{ eV}$.

5 Conclusion

In this paper we show that pump limiters with a long throat can be operated in the plugging regime which strongly enhances the neutral density build-up, especially at low plasma SOL temperature. The efficiency of the apparatus should thus be increased in cold and dense edge plasmas.

References:


Figure 1: The escape probability via ballistic trajectories to the nozzle for neutrals created at $x$ with average velocity $u$ and thermal velocity $v$. $L_G$ is the characteristic length such that $P_G(L-L_0)=0.1$ ($L = 10 \text{ a}_w$).

Figure 2: Neutral density at the neutralizer plate at given heat flux flowing to the limiter, here $Q=10\text{MW/m}^2$, as a function of the plasma SOL temperature.
FASTIONS LOSSES
DURING NEUTRAL BEAM INJECTION IN TORE-SUPRA

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INTRODUCTION
Future neutral beam injection experiments in the Tore-Supra tokamak have two characteris-
tics of major importance concerning fast ions confinement: a high toroidal field ripple and a qua-
si-perpendicular injection angle. We present a Fokker-Planck calculation of the fast ions losses
occurring in such a configuration. Then, we test the sensitivity of these losses to various plasma
parameters, and we determine the most favorable heating scenarios.

NBI AND TOROIDAL FIELD RIPPLE IN TORE-SUPRA
The Tore-Supra neutral beam injection system [1] will have a total power capability of 9 MW
($D^0$, 100 keV) or 4 MW ($H^0$, 80 keV). Two boxes are in co and one in counter-injection. Each
box provides two beams in a quasi perpendicular injection geometry: Tangency major radius of
the beams $R_T=1.24$ m ($R_{axis} = 2.44$ m, $\alpha_{axis}=60^\circ$). The beam ions born between the outer edge
(plasma minor radius $a=0.80$ m) and $r=0.25$ m execute banana orbits.
The toroidal field created by the 18 superconducting coils has a high toroidal modulation going
from $\delta=2 \times 10^{-4}$ up to $\delta=7 \times 10^{-2}$ at the outer edge. As a consequence, for standard discharges
($I_p=1.7$ MA), there is an important zone inside the plasma where local magnetic wells exist
along field lines. In this region, the ripple is "effective" ($\delta^* > 0$) and the beam ions on banana
orbits can entrap the local magnetic wells [2] and leave the plasma before their complete
thermalisation. Figure 1 shows the flux surfaces, the contour $\delta^*=0$ where the modulation along
a field line vanishes and the position of the banana tips of fast ions first orbits.

FOKKER-PLANCK CALCULATION OF FAST ION LOSSES
We have developed a Fokker-Planck calculation in order to test $\delta^*$ sensitivity of the ripple
losses to plasma and beam parameters. A complete description of this code, a comparison with
Monte-Carlo simulations and with TFR experimental results can be found in [3]. The present
version of the code includes improvements concerning the magnetic configuration description
(Shafranov shift, horizontal and vertical displacements) and the modelling of the ripple losses
operator. Our calculation is based on the stationary Fokker-Planck equation for fast ions:
where $\tau_5$, $\tau_i$ are the slowing down and ion-ion collision times, $v_c$ is the critical velocity and $\mu = v/v_c$. The first and second terms represent collisional effects: drag and angular scattering in the Legendre operator approximation (the energy diffusion is neglected). The third term is the source and the fourth is an operator modelling the ripple losses. In the previous version of the code [3], it is simply written as $R(f) = -f/\langle T_r \rangle$ where $\langle T_r \rangle$ is an averaged ripple loss time. This operator applies the losses to the whole distribution function, whereas only the trapped ions are concerned: this leads to overestimate the losses, especially in the case of tangential injection. To overcome this difficulty, we now write the operator as $R(f) = W(\mu, \mu_{\text{trap}})f/\langle T_r \rangle$, where $W(\mu, \mu_{\text{trap}})$ is defined by $W=1$ for $\mu<\mu_{\text{trap}}$ and $W=0$ for $\mu>\mu_{\text{trap}}$, and applies the losses to the trapped ions only. In this case, the solution of the Fokker-Planck equation is obtained after an expansion of $R(f)$ on the Legendre polynomials and a matrix inversion. An iterative procedure is then used in order to insure the coherence of $\langle T_r \rangle$ with the solution $f$.

The characteristic loss time is defined by $1/\langle T_r \rangle = 1/\tau_1 + 1/\tau_2$ where the averages are taken over the trapped ions distribution $\langle f \rangle = \int f \, d^3v / \int d^3v$.

- $\tau_1$ is a characteristic loss time due to trapping in the local wells. This loss term is largely dominant in the effective ripple zone. To compute $\tau_1$, we use the expression given in [2]: $\tau_1 = \tau_b/F$, where $\tau_b$ is the bounce period and $F$ is the trapping rate per bounce taking into account collisional and collisionless trapping mechanisms.

- $\tau_2$ is a loss time that models the effect of banana drift diffusion: This term dominates the transport in the zones where the ripple is not effective. We define it as $\tau_2 = \Delta r^2 / 2D$. $D$ is the ripple-plateau diffusion coefficient and $\Delta r$ is the averaged distance between the banana tip and the contour where the modulation vanishes.

RESULTS FOR A STANDARD CASE

Figure 2 shows the computed ripple losses for a Tore-Supra standard case: Deuterium plasma, $R_0=236\text{cm}$, $a=80\text{cm}$, $I_p=1.7\ \text{MA}$, $B=4T$, parabolic density profile ($\langle n_e \rangle = 5 \times 10^{19} \text{m}^{-3}$), parabolic squared temperature profile, $T_e(0)=3\ \text{keV}$, $100\ \text{keV}$ (full energy) $D_0$ neutrals. The local ripple lost fraction $g_r(r)$ (defined as the power lost on a given magnetic surface divided by the power injected on this surface) increases dramatically from the center to the edge. The losses increase rapidly around $r=40\ \text{cm}$: Above this radius the banana tips stay in the effective ripple zone and the ions are lost before thermalisation because $\langle T_r \rangle < \tau_5, \tau_i$. For smaller radii the beam ion transport is due to banana drift diffusion. This mechanism feeds the poor confinement zone with fast ions born in the central part of the plasma and is responsible for substantial losses in the region $35\text{cm}<r<50\text{cm}$. The global losses $G_r(r)$, computed with a classical beam deposition code, are also displayed: They represent the power lost fraction integrated up to a given radius. The total fraction lost is very important for the standard case: $G_r(a)=50\%$. 

\[
\frac{1}{\tau_5 v_c^2} \frac{\partial}{\partial v} \left( v^3 + v_c^3 \right) f + \frac{1}{2\tau_i} \frac{\partial}{\partial \mu} \left( 1 - \mu^2 \right) \frac{\partial}{\partial \mu} f + \hat{n}_b \delta(v-v_c)\delta(\mu-\mu_0)/2\pi v_0^2 + R(f) = 0
\]
MINIMISATION OF THE LOSSES

Two main approaches can be examined in order to reduce these dramatic losses.
- for a given plasma geometry \((R_{\text{axis}}, a, B_\phi, B_0)\), one may try to minimize the beam ions deposition in the poor confinement region: the most sensible parameter in this case is the plasma density (average and profile)
- for given plasma profiles, one may try to expand the good confinement zone: in this case the plasma position inside the vacuum chamber plays the major role.

**Effect of plasma density:**

Figure 3 shows the total lost fraction (ripple and shine-through) as a function of the volume averaged density for three different peaking factors \(n_e(0)/<n_e>=1.9, 2.6, 4.0\) (the other plasma parameters are the standard ones). For a fixed peaking factor, the losses can be reduced by a decrease in the density down to \(2.5 \times 10^{19}\) that gives a more central fast ions deposition. A further decrease in the density results in a dramatic increase of the losses due to the shine-through fraction. For a fixed average density, the losses can always be decreased by peaking the density profile. However, for these plasmas with standard minor radius \(a=80\) cm, the losses remain important: 25% for \(<n_e>=2.5 \times 10^{19}\) and \(n_e(0)/<n_e>=4\).

**Displacements of the plasma column inside the vacuum vessel:**

If the plasma minor radius is reduced and the plasma pushed on the inner wall, the good confinement zone becomes more important and this leads to reduced losses. Figure 4 displays the effect of plasma position (the plasma parameters are the standard ones excepted \(a=70\) cm). For a plasma with no vertical displacement (\(\Delta z=0\) cm) the total ripple loss fraction is \(G_r(a) = 32\%\) instead of 50% in the standard case. Beam ions are captured and lost in the zone \(z>0\) because they drift towards the top of the vacuum vessel: if the plasma is displaced in the vertical direction towards the bottom of the vacuum vessel, the good confinement zones are no more symmetric with respect to the equatorial plane and the upper one becomes bigger. As shown in figure 4, for \(\Delta z=-10\) cm we get \(G_r(a) = 23\%\). To further reduce the losses, one may combine the favorable effects of plasma position and density profile peaking: for \(\Delta z=-10\) cm, \(n_e(0)/<n_e>=4.0, a=70\) cm losses are reduced to \(G_r(a) = 15\%\).

**CONCLUSION**

Fokker-Planck calculations predict important fast ions losses during NBI in Tore-Supra (up to 50% of the injected power for standard operation). These losses are mainly due to beam ions on banana orbits with tips in the effective ripple zone. Plasma position inside the vacuum vessel as well as density (average and profile) are sensible parameters. For smaller, inner wall limited plasmas (\(a=70\) cm) and for peaked density profiles, losses can be reduced to smaller values (\(<25\%)\). A vertical displacement of the plasma column in the opposite direction of the ion drift velocity can also improve fast ion confinement.
REFERENCES


Fig. 1: \( \delta^* = 0 \) contours and position of the first orbits banana tips (D\(^0\) 100 keV)

Fig. 2: Ripple losses (local and integrated) for a T-S standard case

Fig. 3: Effect of plasma density on total losses (ripple + shine-through) for three peaking factors \( n_e(0) / \langle n_e \rangle = 1.9, 2.6 \) and 4.0

Fig. 4: Effect of plasma position on ripple losses (local and integrated)
MICROTEARING TURBULENCE AND HEAT TRANSPORT

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INTRODUCTION

Up to now, experimental results do not allow to decide whether electrostatic or magnetic turbulence is the source of anomalous transport in tokamaks. Magnetic instabilities such as microtearing modes are often dismissed, for the transport they induce is usually estimated to be small. These modes are linearly unstable in collisional regimes for reasonable

\[ \beta_p = \frac{2\mu_0 nT}{B^2} \]

corresponding to a potential vector \( A \) which is radially constant near the resonant surface. However, they induce in non linear regimes a small electron heat diffusivity \( D_e \). Indeed, the collisional impedance requires that \( \frac{D_e}{\delta^2} \ll \nu \) where \( \delta \) is the mode radial scale, while a mode is stabilized by electric fluctuations unless \( \delta \) is smaller than the ion Larmor radius \( \rho_i \). In collisionless regimes, microtearing modes are linearly stable. However, preliminary results [1] show that the modes may be driven unstable by the radial diffusion they induce in non linear regimes where magnetic islands overlap. To study this effect, we use a model where the non linear action of modes on a given Fourier component of the perturbed distribution function is represented by a velocity dependent diffusion coefficient operator. This introduces a new scale \( \delta_e \alpha D_e^{1/3} \) which must be larger than the linear current width \( \delta_e \) so that non linear destabilization takes place, and smaller than \( \rho_i \) to avoid electric stabilization. Unstable modes are found for \( \beta_p \approx 1 \) associated with non constant \( \lambda \) profiles. The dependence of the marginal threshold \( \beta_p \) on the parameters \( K_0 \delta_p, \delta_e / \delta_p \) and \( \rho_i / \delta_p \) is studied in this work.

I BASIC EQUATIONS AND NUMERICAL CALCULATION

We study in a cylindrical equilibrium a perturbed potential vector \( \delta A \) and a perturbed electric potential \( \delta U \)

\[ \delta A = -A(r-r_i) \exp(i(\omega+\nu)\omega t)) + cc \]

and

\[ U(r-r_i) \exp(i(\omega+\nu)\omega t)) + cc \]

which induce perturbed current and charge densities
\( S_j \) and \( S_p \). The Ampère equation and the electroneutrality constraint may be written in the linear case under a variational form whose functional
\[
\mathcal{L} = -\int d^3 x \frac{|\nabla A|^2}{2\mu_0} + \int d^3 x \left( j_A^* - \rho U^* \right),
\]
extremum with \( A^* \) and \( U^* \), is derived from the usual electromagnetic action. The densities \( j \) and \( \rho \) are deduced from a Vlasov equation for each species with non linear Landau terms replaced by a linear diffusion operator [3] with a diffusion coefficient \( D(v_f) = D_r |v_f| \) \( v_f = \sqrt{2T/m} \).

It is easier to compute the particle responses in the Fourier space
\[
\langle j(r-r'),\rho(r-r') \rangle = \int_{-\infty}^{+\infty} \frac{dK}{2\pi} \langle j(K),\rho(K) \rangle \exp\{iK(r-r')\}
\]
leading to the following expression of \( \mathcal{L} [2] \)
\[
\mathcal{L} = -\frac{1}{\nu_s} \int_{-\infty}^{+\infty} \frac{dK}{2\pi} K^2 A(K) A^*(K) + \sum_{s=i,e} \mathcal{L}_n (1)
\]
\[
\mathcal{L}_n = \frac{n_s e^2}{T_s} \int_{-\infty}^{+\infty} \frac{dK}{2\pi} U(K) U^*(K) + 
\]
\[
+ \frac{i n_s e^2}{T_s} \int_{-\infty}^{+\infty} \frac{dK}{2\pi} \int_{-\infty}^{+\infty} \frac{dK'}{2\pi} \frac{dv_f'}{v_f^2} \exp\left\{-\frac{v_f^2}{v_f'^2} \right\} \int_{0}^{+\infty} dh \exp\left\{-h \right\} 
\]
\[
\left\{ \omega_n \frac{\partial}{\partial T_s} - \frac{3}{2} \right\} \left\{ \frac{K' - K}{v_f'} \right\} \exp\left\{ i \frac{\omega}{v_f'} (K' - K) - \frac{D_s(v_f')}{6} (K'^2 - K^2) \right\}
\]
\[
J_o \left( K' \rho_s \sqrt{h} \right) J_o \left( K \rho_s \sqrt{h} \right) \left( \frac{v_f'}{v_f} + \frac{v_{ts} h_1}{2 v_f'} \right) (U(K) - v_f \Lambda(K))
\]
\[
J_o \left( K' \rho_s \sqrt{h} \right) J_o \left( K \rho_s \sqrt{h} \right) \left( \frac{v_f'}{v_f} + \frac{v_{ts} h_1}{2 v_f'} \right) (\Lambda^*(K') - v_f \Lambda^*(K'))
\]
where \( \omega_n, \omega^* \) are the diamagnetic frequencies.
\[ K_0 = 1/r \quad ; \quad K^2 = K^2 + K_0^2 \quad ; \quad \dot{\chi'} = K_0 L_0 \]

and \( Y \) is the Heaviside function. The Bessel averaging involving the safety factor \( q \) takes account of the drift surface shift and is correct if \( K \cdot r \cdot dq / r \cdot dq \). This variational form has been implemented in a code [2] which finds the modes by scanning over the parameters.

\[
\beta_p^* = \frac{n_e T_e}{E_p/2} \left( \frac{L_m}{L_n} \right)^{1/2} \text{ and } \omega. \]

II STABILITY AND ANORMAL TRANSPORT

In the linear case, i.e. \( \delta_e = \omega / \dot{\chi'} T_e \), the first term in (1) is imaginary and vanishes if \( \omega = \omega^{*} + \frac{1}{2} \omega^{*} \), whatever \( A(K) \) is. The following term, of order \( \delta_e / \delta_0 \), is found destabilizing for profiles \( A(r-r_f) \) decreasing as \( \exp(-|K_0 (r-r_f)|) \) outside the resonant layer and exhibiting a hollow shape around the resonant surface (figure 1a).

However, the associated electric potential has a stabilizing influence which depends on the ion resonant response. If the mode is strongly localized, i.e. \( \delta_i \ll \delta_0 \), the mode is nearly adiabatic and the modes are weakly stabilized. Increasing the ratio \( \delta_i / \rho_i \), the mode do not distinguish between electrons and ions and the threshold \( \beta_p^* \) increases, as shown on figure 1b. This figure allows to derive for a given \( \beta_p^* \) the transport coefficient

\[
D = 6 \cdot K_e L_m \text{ with } K_e = 0.25 \quad \text{ and } \quad \frac{d\log T_e}{d\log n} = 2. \]

For instance, \( D = 0.04 \rho_1^2 \frac{v_{Te}}{L_n} \) with \( K_e \rho_1 = 0.25 \), \( \frac{L_m}{L_n} = 20 \), \( m_e/m_i = q=2 \), and \( d\log T_e/d\log n = 2 \).
CONCLUSION

A simple model shows that microtearing modes can be unstable in non linear regimes for reasonable values of physical parameters. The bounds obtained on the heat transport coefficient do not allow to neglect such a turbulence and encourage to perform more accurate calculations.

REFERENCES


Figure 1a: Marginal profile of a magnetic mode (U=0) for $\beta_p^*=2.9, K_b \delta_b = 0.1$
The solid and dashed lines represent the real and imaginary parts of $A$ (the scale is $\delta_b/\rho_i = 0.1$).

Figure 1b: Marginal threshold $\beta_p^*$ versus $\delta_b/\rho_i$ for $q=2, \omega_T/\omega_e^*=2, K_b \delta_b = 0.1$
, $\delta_b/\delta_e = 5$. 

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INTRODUCTION

An important problem which has been addressed since the beginning of exploitation of Tokomak devices is the precise understanding of the MHD activity. Many calculations have been performed to cover the laminar phases where the modes maintain integrability of flux lines [1,2,3], as well as the disruptive phases where stochastic effects play a leading role. The theoretical analysis we present here reconsider the former case in the spirit of a confrontation with the observed poloidal structures, growth rates and frequencies of the modes. Such data can in fact provide useful informations on the equilibrium characteristics which determine the radial structure of the various poloidal components of each unstable mode in the plasma bulk between the resonant layers, and also on the irreversible processes which, together with plasma rotation and diamagnetic effects, control the active coupling within the resonant layers.

We consider a magnetic perturbation derived from a potential vector parallel to the unperturbed toroidal field $B_0$: $\delta B = \frac{Ro_t}{R_0} \delta \psi \frac{B_0}{RB_0}$, $\delta \psi = \psi(r, \theta, \phi) \exp(i \omega t + \text{c.c})$ involving several helicities: $\|f = \sum_m m \psi_m(r) \exp(i m \theta + n \phi)$. The mode is resonant on the magnetic surfaces $r = r_m$ where $q = m/n$, creating around each of them a thin island chain of half width $w_m = \frac{\sqrt{q} m \psi_m(r_m) L_m}{B_0}$. The Ampere's law is equivalent to state that the functional:

$$F(\psi, \psi^*, \omega) = \frac{1}{2} \int \frac{\| \nabla \psi \|^2}{R^2} d^3x + \frac{1}{2} \int \frac{\| \psi^* \|^2}{R} d^3x$$

is an extremum in $\psi^*$. Here the parallel current response $J(r, \theta, \phi) \exp(i \omega t + \text{c.c})$ is considered as a known functional of $\psi(r, \theta, \phi)$. As the functional $F(\psi, \psi^*, \omega)$ is linear in $\psi^*$, its extremum value in $\psi^*$ cancels. This provides the mode frequency $\omega$ for a given geometrical structure $\psi(r, \theta, \phi)$. On the other hand, for a magnetic perturbation applied at a real frequency $\omega$, the quantities $2 \omega \text{Im}(F)$ and $2n \text{Im}(F)$ are respectively the power and the $\phi$ momentum rate coupled to the plasma.
TOROIDAL COUPLING AND OUTER SOLUTION

We split the functional $F$ as given by (1) into $F_{\text{bulk}} + F_{\text{res}}$, arising from the plasma intervals outside the resonant layers and from the resonant layers, respectively. In the bulk plasma and vacuum intervals, the MHD response $J$ applies and the functional $F_{\text{bulk}}$ is just equivalent to the variational form of the linearized energy principle derived from the toroidal quasi-equilibrium equation [4], namely to:

$$-\delta W = -\int \int \int |\nabla \psi|^2 / R^2 d^3 x + R_{\phi} \int \int \int \psi \nabla \cdot (\mu_0 R J_\phi) / R B_0 d^3 x$$

(2)

where $J_\phi$ is the equilibrium current density and $U(r, \theta, \phi, t)$ is the scalar potential of the velocity field defined by: $\nabla / U = \nabla \psi$. The toroidal terms in $\delta W$ proportional to $\beta$ [5] have been ignored because, they are strongly divergent near the resonant surfaces $r_m$, and in fact belong to $F_{\text{res}}$. The fact that $\delta W$ is real means that it contributes to no secular transfer of energy or $\phi$ momentum to the plasma. The extremal value of $\delta W$ gives the slope jumps $[\nabla \psi]$ of $\psi$ across the resonant layers $r_m$; we have:

$$\delta W = \sum_m \int |\psi^* \cdot \nabla \psi| / R^2 d\sigma$$

(3)

where $\psi$ is the solution of the Euler equation of (2) and $\hat{n}$ is the unit vector normal to the equilibrium resonant surfaces $r_m = \text{cst}$.

In principle, the extremal value (3) could depend on all the values $\psi_m(r_m)$ $m \neq m'$; however, it is found that linear relations exist between $[\psi_m'(r_m)]$ and $[\psi_m'(r_m)]$, just leaving $\delta W$ as a functional of $\psi_m(r_m)$ only:

$$-\delta W = \sum_{m, m'} T_{mm'} \psi_m(r_m) \psi_{m'}(r_{m'})$$

(4)

In cylindrical geometry $T_{mm'} = A_m \delta_{m,m'}$, $A_m$ being the classical logarithmic slope jump. The matrix elements $T_{mm'}$ specify through (3) the slope jumps corrected by the toroidal effects. Expanding $\delta W$ to the second order in the small parameter $\varepsilon = r / R$ leads to analytical values $T_{mm'}$, with $m' = m \pm 1$, involving the cylindrical tearing profiles $\psi_m(r)$ and the usual characteristics of the toroidal equilibrium: $\beta_{\text{pol}}$, $\Lambda$, etc.

RESONANCES AND INNER SOLUTION

The $\psi$ solutions within the outer regions of the plasma are connected to the resonant inner layer solutions through alternate expressions of the slope jumps derived from integration of the resonant parallel current responses $I$. These resonant currents determine the contribution

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$\int J^* / R \, d^2 x$ of each resonant layer $r_m$ in $F_{res}$. They are calculated in the frame rotating in the toroidal direction at the same velocity $R_e(\omega)/N$ as the islands, supposing a non linear regime where the electrons reach thermodynamical equilibrium over each perturbed magnetic surface. This regime applies for observable modes, for which the parallel transit frequency $K_m w_m V_{th}/RL_m$ is larger than the diffusion rate $D_e w_m^2$ assumed larger than the mode frequency $\omega$. This situation imposes the detailed density, temperature and electrostatic potential profiles in terms of the unperturbed diamagnetic frequencies $\omega^e_n, \omega^i_T$ and of the island velocity $R_e(\omega)/N - \Omega$ in the plasma frame. The parallel current $I$ is then determined from the charge continuity equation:

$$\text{div}(J) = -\text{div}(J_L)$$

where the expression of the transverse current is $J_L = e(\Phi_\Phi - \Phi_i) + J_L$, $\Phi_\Phi$ and $\Phi_i$ being the diffusion electron and ion transverse fluxes due to the small scale turbulence present in the plasma, while $J_L$ is due to the drift curvature and to the ionic effects of inertia, F.L.R. and viscosity [6]. To complete the determination of $I$, so that it represents the deviation from the Rutherford current $I_R$, a further constraint on average value of $I$ over each perturbed magnetic surface must be imposed: the inductive electric field, proportional to the mode growth rate $\gamma$ must balance the ohmic friction in the average. If we assume ambipolarity of the local diffusion, i.e. $\Phi_\Phi = \Phi_i$, then the deviation of $I - I_R$ is simply $J_L$ and the electron diamagnetism does not influence the mode frequency. However, if the turbulent modes which determine $\Phi_\Phi$ and $\Phi_i$ exchange momentum with the plasma over a radial range larger than the island width, the ambipolarity condition must be removed, resulting in a largely dominant contribution of $e(\Phi_\Phi - \Phi_i)$ compared to $J_L$. The calculations [7] then give:

$$F_{res} = \delta_{mm} K_m \psi_m(r_m) \psi^*_m(r_m)$$

where $K_m$ is an explicit complex non-linear function of already defined quantities: $D_e \omega^e_n, \omega^i_T$ for both electrons and ions and $\gamma$, $R_e(\omega) - \Omega(r_m), \psi_m(r_m)$. The mode consistency is expressed by matching the outer and inner solutions. This is equivalent to extremalize $F_{bulk} + F_{res}$ in $\psi_m(r_m)$ where, $F_{bulk} = -\delta W$ and $F_{res}$ are given by (4) and (6) respectively. This yields the following non linear system of equations for the growth rate $\gamma$, the frequency $R_e(\omega)$ and the normalized amplitudes $\psi_m(r_m)$:
\[ \sum \{ T_{m'} m'' - K_{m'} (\gamma, R_q(\omega), \psi_{m'}) \delta_{m'm''} \} \cdot \psi_{m''}(r_{m''}) = 0 \quad (7) \]

PRELIMINARY RESULTS

As a first application of the above theory we analyze the saturated Rutherford regime \( \Omega = \omega^* = R_q(\omega) = 0 \) which corresponds in (7) to \( K_{m'} = \gamma R_{m'} R_m \approx -l \psi_m^{1/2} \). Starting from an unstable tearing mode \( m = 2, n = 1 (\Lambda_2 > 0) \) we find that: if the \( m = 3, n = 1 \) harmonic is cylindrically unstable, then the two harmonics \( m = 2 \) and \( m = 3 \) are mutually pumped with a relative phase 0 or \( \pi \) depending on the sign of the coupling coefficient \( T_{2,3} \); if the \( m = 3, n = 1 \) harmonic is stable then it can be destabilized by toroidal coupling with the same phase, at a normalized amplitude up to \( a/R_0 \). Generally, the frequency \( R_q(\omega) \) compromises the friction effects within the two resonance surfaces \( m = 2,3 \) (and of course within the resistive wall). Assuming \( T_{3,3} < 0, T_{2,2} > 0 \) one finds, in the above Rutherford situation, only one positive root \( \gamma = \gamma_R \). For non-vanishing \( \Omega, \omega^* \), the mode may bifurcate from a state \( \gamma = T_{2,2}/R_2 \) to a state \( \gamma = \gamma_R \) with a frequency influenced by the frictions at \( r = r_2 \).

CONCLUSION

Generally the model shall provide by solving (7) the relative amplitude \( \psi_m(r_m) \) from which the external mode structure may be deduced and compared to the magnetic probe data. The model provides also the structure of the pumped internal \( m = 1 n = 1 \) mode by its neighbouring \( m = 2 n = 1 \), which may be then compared with the soft X-ray data analysis. On the other hand the model should allow to explain the mode frequency \( R_q(\omega) \) in terms of diamagnetism and plasma rotation at the resonant layers. Bifurcation in \( R_q(\omega) \) could for instance explain the mode locking.

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COLLISIONLESS FAST IONS DYNAMICS IN TOKAMAKS

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INTRODUCTION

In thermonuclear plasmas of the next Tokamak generation, the dynamics of fusion products needs to be understood as it could well influence the reactor performances. In such a low collisionality regime, a careful analysis of the associated transport and relaxation processes is required. For this program to be completed, the effect of a perturbing field on fast particle individual motion must be precisely evaluated. This imposes to address the question within the Hamiltonian mechanics. In this framework, resonances of particle trajectories with low frequency fields arise naturally as a potential mechanism for transport and relaxation. The latter processes result from the destruction of adiabatic invariants when a stochasticity threshold is reached. Then, the quasi-linear theory is applicable, and diffusion coefficients can be derived. In the following, the method will be briefly reviewed and two reactor relevant situations analysed: heating of energetic ions under the effect of a compressionnal wave and alpha-particle losses induced by the toroidal field ripple.

HAMILTONIAN DESCRIPTION OF COLLISIONLESS TRANSPORT

In an axisymmetric unperturbed tokamak magnetic configuration, the particle trajectory is integrable and described by a set of three action-angle variables \( \{ J_k, \Omega_k \} \). Its Hamiltonian \( H_0(J_k) \) depends only on the action variables and the motion is quasi-periodic: the actions remain constant and the angles rotate linearly \( \Omega_k = \omega_k(J_k)t + \Phi_{0k} \). Essentially, the effect of low frequency perturbations is to disrupt this adiabatic behaviour in regions of the phase space where an efficient coupling (i.e. resonances) between fields and particles occurs. In order to identify those regions, the Hamiltonian perturbation is expanded as a Fourier series:

\[
\delta H = \sum_{n_k} h_{n_k} \exp(\text{in}_k \Phi_k + i \text{o}t) + \text{complex conjugate}
\]

Resonant interaction takes place when the phase of the perturbation is stationary along the unperturbed trajectory. For a given triplet of integers \( \{ n_k \} \), the location of the resonance surfaces is defined by \( \Omega = n_k \omega_k(J_k) + \text{o} = 0 \). Around these surfaces an island like structure emerges and above a stochastic threshold \( S > 1 \) [1], resonance overlapping destroys this regular behaviour; \( S \) is the product of the island width \( 2\Delta \Omega_l = 4\sqrt{2\alpha_h \alpha n_k} \), \( (\alpha = n_k \frac{\partial^2 H_0(J)}{\partial J_k \partial J_l}) \)

by the density of resonant surfaces. The stochastic behaviour legitimates the random phase approximation and the quasi-linear theory.
COMPRESSIONNAL WAVE HEATING OF ENERGETIC IONS

In the case of ICRF, wave particle resonance occurs when the cyclotron frequency of the particle matches the RF frequency at some points along its trajectory. An irreversible transfer of energy (heating) is possible if the phase of the particle get randomised in between two successive resonance crossings. For low energy ions, the randomisation is insured by both collisions and intrinsic stochasticity [2]. When the wave spectrum contains only one toroidal number N, the fast ion population recovers its adiabatic behaviour, but when a full toroidal wave spectrum is considered, stochastic heating reappears.

To demonstrate that point, we focus on passing particles. Their motion is represented by the following set of action-angle variables for the cyclotronic, poloidal and toroidal motions:

\[
J_1 = -\frac{m}{e} \mu \\
J_2 = e \Phi (J_3/e) + q (\Psi_p) R_0 v_\parallel \\
J_3 = e \Psi_p + m R_0 v_\parallel
\]

\[
\Phi_1 = \phi_c + \frac{R_0 \Omega \phi_c}{R_0 \Omega_2} \sin(\theta) \\
\Phi_2 = \theta \\
\Phi_3 = \phi
\]

with the magnetic poloidal and toroidal fluxes \(\Psi_p, \Phi;\) major radius \(R_0;\) poloidal and toroidal angles \(\theta\) and \(\phi;\) safety factor \(q(\Psi_p)\) and for the particle : charge \(e;\) mass \(m;\) guiding centre radius \(R_c;\) magnetic moment \(\mu;\) parallel velocity \(v_\parallel;\) cyclotronic phase and averaged pulsation \(\phi_c, \Omega_c.\)

The compressional field is taken as a radially running wave with one toroidal harmonic N and :

\[
\delta H = e \nu \Phi_2 \exp(i(\omega t + kr) \exp(iN \Phi_2) + c.c.
\]

The fundamental harmonic resonance \((n_1 = 1)\) gives rise to a family of resonances in the phase space, labelled by \((1, n_2, N),\) of density \(1/\nu \Omega_2.\) The stochasticity threshold is :

\[
S = 4 \frac{q}{\nu \Omega_2} \left| N + \frac{n_2}{q} \right| \sqrt{2 \frac{2 \Omega_2 \nu_\parallel}{m} \sqrt{\frac{2}{\pi x}}}
\]

where \(2x = 2 \nu \Omega_2 (\Omega_c / \nu \Omega_2)\) is the number of effective resonances.

If for the low energy particles, \(S\) is greater than 1 and stochastic heating is effective, it is no longer the case for highly energetic ones as \(S\) scales with \(\nu^{-1/4}.\) For a field strength of 50 V/cm, stochasticity disappears for alpha-particles of energy typically greater than 100 keV.

However for a wave toroidal spectrum, the resonance families are superimposed leading to a higher resonance density and to a destruction of the regular trajectories (Fig.1). Applied to \(\alpha\)-particles in a reactor, ICRH continues to take place, allowing the use of the quasi-linear theory. This gives confidence in the use of 2-D ICRF modelling even for high energetic particles and to the further Fokker Planck calculations.

STOCHASTIC INSTABILITY OF ENERGETIC TRAPPED IONS IN THE T.F. RIPPLE

In the transport regime relevant for \(\alpha\)-particles in a reactor, i.e. the ripple banana drift regime, the neoclassical diffusion coefficient scales as the particle collisionnality and is consequently very low. As already pointed out [3,4], collisionless effects may dominate the \(\alpha\) transport. In the hamiltonian analysis, the trapped \(\alpha\)-particle motion in presence of a ripple perturbation...
$\delta B \cos(N_c \phi)$ is strongly influenced by the existence of resonant surfaces. For this class of particles, many resonances appear because the bounce frequency $\omega_b = \partial H_0 / \partial J_2$ and the precession frequency $\omega_p = \partial H_0 / \partial J_3$ are of the same order. Here $J_2$ and $J_3 = e \psi_\rho(r)$ ($r = r_0$) are the invariants associated respectively to the bounce ($\Phi_2$) and precession ($\Phi_3$) motions. The Fourier expansion of the perturbation writes:

$$\delta H = \sum_{n=\infty}^{+\infty} \frac{1}{2} \mu \delta B J_{n_2} (N_c q \theta) \exp (iN\Phi_3 + in_2\Phi_2) + \text{complex conjugate}$$

where $J_{n_2} (N_c q \theta)$ reflects the averaging over the bounce motion ($\theta$ is the banana tip poloidal angle). The resonant surfaces are given by $\Omega = N_c \omega_3 + n_2 \omega_2 = 0$ and their radial positions for a typical reactor case are shown on Fig.2. Using first order $\varepsilon$ expansions ($\varepsilon = r/R_0$) of the action-angle variables, the stochasticity threshold takes the following form:

$$S = f(s, \theta) \frac{N_c q^2}{\varepsilon} \frac{\rho_c}{r} \sqrt{\frac{\delta(r)}{J_{n_2} (N_c q \theta)}}$$

where $\rho_c$ is the Larmor radius, $\delta(r) = \delta B / B$ the ripple value and $f$ a function of the shear parameter $s = r (dq/dr)$ and of the banana tip poloidal angle $\theta$ ( $f$ is of order of one and for strongly trapped particles $f = 4 \left(3s/2 - 1\right)^{1/2}$ ). The critical value of $S$ scales as $\rho_c^2$ and can be very low for energetic particles. This scaling is quite different from the one proposed in [4]. However, for 3.5 MeV $\alpha$-particles in an INTOR-like device, we obtain the same critical value ($\delta = 10^{-3}$). This expression has been successfully confirmed by Poincaré maps (fig.2). Above this threshold, i.e. above a critical minor radius, the motion becomes stochastic and leads to a radial diffusion at constant energy and magnetic momentum. The diffusion coefficient is given by the quasi-linear theory. For typical trapped particles ($\theta = 1$ and $N_c q \theta > n_2$):

$$D_{QL} = \frac{\sqrt{2}}{4} N_c \delta \rho_c^2 \frac{\varepsilon}{\theta} \frac{q^2}{\varepsilon^{3/2}}$$

Under typical reactor conditions, values of order $10$ m$^2$s$^{-1}$ are easily reached for fusion products with $\delta = 10^{-3}$, leading to fast losses.

**CONCLUSIONS**

We have investigated two reactor relevant situations where collisionless processes dominates the dynamics of fast ions population:

- Concerning ICRF, provided that a realistic wave spectrum is considered, intrinsic stochasticity allows for the use of the quasi-linear theory even for very energetic ions.

- Concerning the TF ripple, we have shown that this perturbation dominates the trapped $\alpha$-particles transport as soon as $\delta = 10^{-3}$.

Such kind of stochastic behaviour is of major importance for the next tokamak generation and gives constraints on the forthcoming fusion reactor designs.
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Figure 1: Poincaré maps representing $J_1(\Phi_2)$ for 500 keV $\alpha$-particles in presence of an ICRF field. Resonant surfaces are labelled by $(n_2,N)$. Left: with a single $N=20$ harmonic, no resonance overlapping is achieved. Right: a second $N=21$ harmonic is superimposed, leading to overlapping and stochastic areas.

Figure 2: Poincaré map for $\alpha$-particles in a reactor case with low ripple ($R_0/a=4$, $\rho_c/a=0.04$, $\theta=1$, $q_{\text{edge}}=3$, $N_c=12$, $\delta_{\text{edge}}=0.12\%$, $\delta(r) \approx R^{N_c}$): the radial positions of the particles (related to the invariant $J_3$) are plotted as a function of their precession angle $\Phi_3$ when they cross the equatorial plane. The figure also displays the resonances and the ripple values. Two different regimes can be clearly identified: stochastic ($r/a>0.8$) and adiabatic ($r/a<0.7$).
VARIATIONAL DESCRIPTION OF LOWER HYBRID WAVE PROPAGATION AND ABSORPTION IN TOKAMAKS

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ABSTRACT:
A variational approach to a global solution of the lower-hybrid wave propagation and absorption in tokamaks is presented. It is based on the use of trial functions which are locally solutions of the homogeneous problem ("eikonal trial functions"). Preliminary results in circular cylindrical geometry are reported, showing the radial power deposition profiles for various toroidal wavenumbers excited by the r.f. source.

1- INTRODUCTION
Variational techniques provide an efficient scheme to describe both propagation and absorption of electromagnetic waves in tokamak plasmas. Such a scheme has been successfully implemented for the ion cyclotron frequency range but, due to the large number of degrees of freedom of the wave field, an extension to higher frequencies seems prohibitive.

This paper describes a method for attempting a full wave description of r.f. fields in the lower hybrid frequency range, in order to study current generation and profile control via Landau absorption. The basic idea is to use eikonal trial functions which locally satisfy the cold dispersion relation. This allows to use a mesh larger than the wavelength and reduces the number of unknowns.

2- VARIATIONAL PRINCIPLE AND TRIAL FUNCTIONS
Since mode conversion between the cold propagation branches occurs commonly in lower hybrid current drive experiments, it is important to start with a full wave equation which describes both the electron plasma wave and the whistler wave, ie from a vectorial wave equation which takes into account the electromagnetic and the electrostatic components of the fields. As opposed to the magnetosonic case, this leads us to use the three components of the electric field as our variables rather than to perform a reduced description in terms of two components of the vector potential.

Maxwell's equations are equivalent to finding the extremum of the following form:
\[ L(E, E^*) = \int dr \left[ (\nabla \times E(r)).(\nabla \times E^*(r)) - \frac{\omega^2}{c^2} E^*(r).K(r).E(r) - \frac{4i\pi\omega}{c^2} E^*(r).J_0(r) \right] \]  

(1)

for all variations \( \delta E^*(r) \) subject to the condition \( n \times \delta E^*(r) = 0 \) on the metallic boundary.

Because of the small wavelengths involved in the problem, performing a discretization of the variational form on a scale length which is smaller than the wavelength would require a prohibitive processing time. As a consequence we must use trial functions which can approximate the solution over several wavelengths. In practice we divide the plasma in a number of radial cells whose width is a given fraction \( \eta \) of the inhomogeneity scale length, i.e.

\[ \eta = \text{Max} \left( \frac{\Delta n_e}{n_e}, \frac{\Delta r}{r} \right). \]  

(2)

The simplest functions which can approximate both the analytic behaviour of the fields near the singularity at \( r = 0 \) ("whispering gallery" effect), and the propagative nature of the wave away from the magnetic axis, are of the form \( r^\lambda \) where \( \lambda = \alpha + i \beta \). \( \lambda \) is a complex number which can take four values \( \lambda_i \) (i=1,4) to be determined from the local wave equation.

In each cell the \( \lambda_i \)'s (i=1,4) are the eigenvalues associated with a 4x4 matrix and they physically correspond to the slow and fast waves propagating both inward and outward. Associated with each of these four values, there is a polarization vector for the electric field \( E_i \) and the trial function in a particular radial cell can be written:

\[ E(r,t) = \sum_{i=1}^{4} \alpha_i E_i r^{\lambda_i} \exp \left( i n \phi + i m \theta - i \omega t \right) \]  

(3)

where \( r, \theta, \phi=z/R_0 \) are the cylindrical coordinates, \( R_0 \) being the major radius of the torus, \( n \) and \( m \) are toroidal and poloidal wavenumbers respectively, and \( \omega \) is the pulsation of the r.f. field. In this paper we restrict ourselves to the cylindrical approximation of the tokamak so that \( n \) and \( m \) are integers which can be specified together with the frequency of the r.f. source.

3-PRINCIPAL BOUNDARY CONDITIONS AND LAGRANGE MULTIPLIERS

Having chosen the trial functions described above, the variational form (1) has a matrix representation:

\[ L(E, E^*) = X^+.A.X - X^+.S \]  

(4)

where \( X \) is the vector whose components are the \( \alpha_i \)'s (eq.2). \( A \) is a square matrix of dimension \( 4N - 2 \), \( N \) being the total number of cells, (in the central cell two unphysical solutions for \( \lambda_i \) must be eliminated).
Our choice for the trial functions entails that $A$ is diagonal by blocks of dimension $4 \times 4$, as cells are completely decoupled. Therefore, for the problem to be well-posed, it is obvious that some continuity constraints have to be imposed to supplement the extremum principle.

In fact it can be shown that, imposing the continuity of the tangential components of the electric field at the border of each cell is both necessary and sufficient for the uniqueness of the solution, provided that there is some dissipation (even infinitesimal) at every point in the volume considered. Such boundary conditions are usually called principal boundary conditions as opposed to the natural boundary conditions (continuity of $\nabla \times \mathbf{E}$) which are automatically satisfied when the variational form is extremum [1].

The constraints inherent to the principal boundary conditions can be written:

$$B \cdot X = 0$$  \hspace{1cm} (5)

with our notations, where $B$ is a large rectangular matrix. We therefore introduce a set of Lagrange multipliers, the vector $\Lambda$, and look for a vector $X$ which, for a particular choice of $\Lambda$, satisfies simultaneously

$$B \cdot X = 0$$  \hspace{1cm} (6)

and

$$\delta X^+ \cdot (A \cdot X - S + B^+ \cdot \Lambda) = 0$$  \hspace{1cm} (7)

with no restriction on the variation $\delta X^+$.

If uniqueness is assessed (absorption at every point), it can be shown that $A$ is regular and that $B \cdot A^{-1} \cdot B^+$ is also regular. Then it is straightforward to write the solution as:

$$X = A^{-1} \cdot [S - B^+ \cdot (B \cdot A^{-1} \cdot B^+)^{-1} \cdot (B \cdot A^{-1} \cdot S)].$$  \hspace{1cm} (8)

4-POWER DEPOSITION PROFILES IN CYLINDRICAL GEOMETRY

As first results, we present some radial deposition profiles which were obtained assuming parabolic dependences for the ion and electron densities $n_i(r)$, $n_e(r)$, the electron temperature $T_e(r)$, and the safety factor $q(r)$. Both the collisional absorption and the resonant Landau damping of the waves are taken into account by including the corresponding anti-hermitian component into the local dielectric tensor. In these calculations the plasma parameters on axis are $n_e(0) = 5 \times 10^{19} \text{ m}^{-3}$, $T_e(0) = 5 \text{ keV}$, $q(0) = 1$, and at the edge $q(a) = 3$. The toroidal magnetic intensity is $3.4 \text{ T}$, the frequency $3.7 \text{ GHz}$ and the major and minor radii of the tokamak, $R_Q = 3 \text{ m}$ and $a = 1 \text{ m}$ respectively. The poloidal mode number was fixed at $m = 10$. Various toroidal mode numbers ranging from $n = 450$ to $n = 850$ have been selected thus showing the influence of the parallel wavenumber on the power deposition (fig. 1). In addition, we have considered two distinct values of the discretization parameter, $\eta_1 = 0.1$ and $\eta_2 = 0.05$, corresponding to 86 and 171 radial cells respectively (figs. 2, 3 and 4). In both cases, the central cell has a fixed radius of $1 \text{ mm}$ and the outer one which extends itself from the antenna layer ($r = 0.96 \text{ m}$) to the edge is also fixed.
In cases of weak absorption ($n = 450$) the Bessel-like structure of the wave appears clearly near the centre (fig.5). The discontinuous character of the density profile shows up by giving rise to local resonances (figs.3 and 4) in some particular cells whose width matches the wavelength of the field. These resonances have a precise physical origin and therefore it seems possible to avoid them by going to a finer mesh when the damping is too weak. The stability of this variational scheme requires further investigation and will be reported later.

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CURRENT PROFILE CONTROL BY ELECTRON-CYCLOTRON AND LOWER-HYBRID WAVES IN TORE SUPRA

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Introduction. Radial control of the tokamak current is of primary importance for sawtooth stabilization and access to the second stability regime. Specifically, off-axis RF current drive for constant plasma current has the effect of raising the value of q on axis, which provides a path to the second stability regime. In principle, this can be achieved by driving current by high-phase-velocity lower-hybrid (LH) waves. However, it is known that a major drawback of this method of current drive is the difficulty of efficient and flexible spatial control of the LH phase velocity spectrum. This limitation results mainly from the complicated mechanisms responsible for the formation of the wave spectrum within the plasma. Combination of electron-cyclotron (EC) and LH waves seems a more appropriate RF system for remote control of the radial current profile. By the appropriate choice of the injection angles, the EC wave power deposition can be shifted both in the real and the velocity space, which allows a flexible tailoring of the plasma current profile. The long pulse operation and good confinement properties of Tore Supra offer the opportunity of extensive studies of current control by LH and EC waves. Specifically, a current I > 0.5 MA will be driven at central density and temperature near $5 \times 10^{13}$ cm$^{-3}$ and 3 keV, respectively, by a LH wave power $P_{LH} = 3 - 6$ MW distributed in a spectrum $1.8 < n_p < 3.5$, where $n_p$ is the parallel retardation index. This power will generate and sustain a long superthermal tail which provides the target for EC wave absorption. In this work, an extensive theoretical study of the current drive process by simultaneous injection of EC and LH waves in Tore Supra is presented. Top launching of EC waves of frequency $f = 110$ GHz is considered, both for $f < f_e$ and for $f > f_e$, $f_e$ being the EC frequency. The EC wave damping along the ray trajectories is computed by means of a toroidal ray-tracing code, incorporating the evaluation of the relativistic dielectric tensor for an arbitrary electron distribution, e.g., the LH-driven electron tail. The latter is determined by means of a 3-D bounce-averaged Fokker-Planck code, which evaluates also the evolution of the electron distribution function during the absorption of high-power EC waves ($P_{EC} > 2$ MW).

Current drive by LH and EC waves. The kinetic equation describing the process of current generation by LH and EC waves contains a Fokker-Planck collision term, a quasilinear parallel diffusion term related to the absorption of the LH waves and a perpendicular diffusion term rel-
ated to EC waves. For the LH wave term we adopt a very simple model, i.e. we assume a constant diffusion coefficient in the interval \( p_e < p_n < p_i \), where \( p_n = p_B^2 / B \), \( p \) is the electron momentum, \( B \) is the tokamak magnetic field and \( p_n, p_i \) are related to the boundaries of the wave \( n \) spectrum. For parameters typical of the Tore Supra LH experiment (\( a=0.7 \) m, \( R=2.25 \) m, \( P_e = 3 \) MW, \( p_v \sim 0.3 \) mc, \( p_e \sim 0.6 \) mc) we obtain a current \( I_{\text{EC}} = 0.465 \) MA, with the profile shown in Fig. 1 (dashed line). The distribution function carrying this current is characterized by a long parallel tail in the range \( p_e < p_n < p_i \), which is the target for the absorption of the EC waves. We first consider the case of top launching of an extraordinary wave beam propagating perpendicularly to the toroidal magnetic field, for \( B = 4.5 \) T on axis. The wave beam is absorbed close to the plasma center by electrons at \( p_n \approx p_e \). The effect of the wave absorption is to decrease the collision rate of those electrons, which causes a significant enhancement of the parallel tail for \( p_n > p_e \), thus in the relativistic range. This fact that for \( f \)-sustained electron system can be used to simulate, in a low temperature plasma, the Maxwellian tail of a hot plasma in order to investigate EC current drive in conditions similar to those existing in the reactor regime. For a wave beam of angular half-width \( \Delta \psi = 5^\circ \) and power \( P_{\text{EC}} = 2 \) MW, the change in the current density profile is shown in Fig. 1 (solid line). The total absorbed power is \( 1.4 \) MW and the additional driven current is \( I_{\text{EC}} = 0.130 \) MA. This launching configuration is thus suited for obtaining EC peaked current density profiles. By an appropriate choice of the launching angles, the EC wave power can be deposited and the current density profile can be modified at different radial locations. For instance, if the EC wave beam is launched at an angle \( \phi = 15^\circ \) with respect to the vertical (towards the low magnetic field side) and with an angle \( \psi = 100^\circ \) with respect to the magnetic field, the wave power is mainly absorbed at \( r \sim 0.3 \) m and the current density profile is modified as shown in Fig. 2. In this case the wave power is absorbed by electrons at \( p_n \sim p_e \). Moreover, \( I_{\text{EC}} = 0.145 \) MA and \( P_{\text{EC}} = 2 \) MW. The same current drive scenario can be investigated at constant plasma current. This can be obtained, e.g., by decreasing the power input in the LH waves as the current driven by the EC waves increases. The results of the numerical simulation of this process are shown in Fig. 3, for \( P_{\text{EC}} = 2 \) MW (a), and \( P_{\text{EC}} = 4 \) MW (b). Note the significant drop in the central current density: the value of \( q \) on axis is 0.95, 1.21 and 1.42 for \( P_{\text{EC}} = 0, 2 \) MW and 4 MW, respectively. This shows that off-axis current drive can be used in Tore Supra to investigate the approach to high-g regimes, which may provide a stable path to the second stability regime.

We finally discuss EC current drive at upshifted frequency, i.e., \( f > f_e \). This scenario is very attractive for a steady-state reactor, since it minimizes the deleterious effects of electron trapping, but it requires high electron temperatures. Again, the ECHR system offers the opportunity of testing this method in a low temperature experiment. The current density generated by a 2 MW wave beam injected at \( \psi = 75^\circ, \phi = -5^\circ \), for \( B = 3.8 \) T on axis, is shown in Fig. 4. Note the broad profile of the additional EC-driven current, due to the beam angular spread and to the fact that for \( f > f_e \) the wave power deposition is strongly sensitive to the value of the \( \psi \) angle. In this case we obtain \( I_{\text{EC}} = 0.165 \) MA and \( P_{\text{EC}} = 2 \) MW.
Fig. 1: radial profile of the current driven by the LH waves alone (dashed line) and by the ECH system (solid line). 
$B=4.5$ T; extraordinary mode, $P_{EC} = 2$ kW, $\psi = 90^\circ$, $\phi = -5^\circ$.

Fig. 2: as in Fig. 1, for $\psi = 100^\circ$, $\phi = 15^\circ$. 

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Fig. 3: as in Fig. 2, but for constant plasma current.
(a) $P_{EC} = 2$ MW,
(b) $P_{EC} = 4$ MW.

Fig. 4: as in Fig. 1, but for the ordinary mode, $\psi = 75^\circ$,
$\phi = -5^\circ$; $B = 3.8$ T.
1. Introduction.

Two identical antennae delivering 3.5 MW each for quasi-CW operation at 3.7 GHz are being installed on Tore Supra. Each antennae is composed of 16 evacuated modules using cnc H-plane hybrid junction (HJ) and 6 E-plane bijunctions as power dividers. Two modules of the same column are fed by a 500-kW klystron via a hybrid junction terminating the pressurized transmission line. The balance port of the HJ is connected to a load.

The large number of waveguides—32 columns of 4 waveguides—allows to excite a narrow, strongly asymmetric, \( \lambda/2 \) spectrum for current drive studies. The \( \lambda/2 \) peak value of the spectrum can be tailored by adjusting the power amplitude and phase of each klystron when the phase between each module of the same line is varied from -90° to 90°, the \( \lambda/2 \) peak value is shifted from 1.46 to 2.40.

Two main issues are investigated: the power reflected towards the 8 klystrons feeding one antenna and the power launched into the plasma. The reflected power has to be low enough for operation with no circulators and the launched power has to be very directive (i.e., large ratio of the power in the main \( \lambda/2 \) peak) for large current drive efficiency. For this purpose, the exact power division in amplitude and phase, which depends on the reflection due to the plasma, has to be known.

2. Main features of the antenna.

In order to fit the plasma shape, the mouth of the RF modules (2x4 wg) were machined with a poloidal curvature radius of 820 mm and consequently the length of the second E-plane bijunction is different of 43° for the external and internal waveguides of the same column (fig. 1). Such a difference was compensated with a phase shifter of 90°+43° located in the hybrid junction to obtain RF fields with the same phasing at the mouth. The position of the short-circuit of the balance port was individually determined from low RF power measurements to obtain a balanced power injection from the HJ. A passive waveguide was added on each side of the grill. Lengths of the 2 bijunctions and the passive waveguides were optimized with the numerical code SWAN /3/ which computes the intercoupling with the plasma of 32 waveguides of infinite height in the same row. Electrical length of the 2 transmission lines from the HJ to the RF window were experimentally adjusted to obtain at the mouth a phasing of -90° of the upper modules with respect to the lower ones.

The 10X10 scattering matrix has been measured on a prototype module with the same procedure than in /4/. It was also computed from the two 180° and 90° MJ scattering matrices obtained from the SWAN code. The same symmetries are observed indicating that only 6 S parameters are linearly independent and are needed to define the whole module. For the modules of the antenna, the symmetry is lost due to the poloidal shape. For example, on fig. 2, the variation of the phase of the S13 parameter describing the coupling through the L1=180° multifunction is given as a function of the short-circuit position: a good agreement between the theoretical and experimental values is obtained.

For the 2 types of multifunctions of 4 waveguides, the SWAN code was run in the case of 8 semi-modules at different densities with different phasings of the modules (θ = 0°, 90°, 90°): in case of good matching (h = 10° m⁻¹) the coefficient of reflection R of the lateral modules does not exceed 3.5% and the value of R averaged on the 8 semimodules is lower than 2% (fig. 3). The calculated values of reflected fields in phase and amplitude allow to compute the power division of the hybrid junctions and the reflected power P for the entire module (fig. 4). The procedure was iterated with the new values of power division. Finally, taking into account the electrical lengths of the 2 transmission lines, the RF power reflected towards each of the 8 klystrons was calculated.
4. Results.

Due to the low coefficient of reflection of the multijunctions, only one iteration was needed to obtain the exact power division. In case of good matching ($N=0.6 \times 10^3 \text{ m}^{-3}$), the unbalance of the power injected from the HJ into the two semi-modules is at most .55. At lower density ($N=0.3 \times 10^3 \text{ m}^{-3}$), the strongest unbalance is .45. However, the excited spectrum is just slightly modified, less than 1% for the directivity $D=<(N_{+})(P(N_{+})-P(-N_{+}))dN_{+}/N_{+}^2$, and the coefficient of reflection is not significantly changed. It was found that the inner rows of waveguides with $L_1=180^\circ$ had a lower directivity for negative phasing of the modules (at $0.6 \times 10^3 \text{ m}^{-3}$, 62% against 67.5% for $\phi_{\text{inj}}=90^\circ$).
where as for positive phasing of the modules, the directivity of the inner waveguides was higher (53% against 42% for $\Delta \phi_{m}=90^\circ$). The coefficient of reflection of the entire module is very similar for the 2 modules of the same column; lower than 1% for $\Delta \phi_{m}=0$ and than 2.5% for $\Delta \phi_{m}=-90^\circ$ (fig.5) at $0.6 \times 10^{18} \text{ m}^{-3}$. Moreover the difference of phase of the reflected fields recombining at the pressurized HJ is very low: less than 6° at $0.6 \times 10^{18} \text{ m}^{-3}$ and 15° at $0.3 \times 10^{18} \text{ m}^{-3}$. Therefore most of the reflected power is injected into the load: less than 0.01% at the medium density and than 0.2% at lower density is reflected towards the klystron. Actually, these very low figures do not take into account: 1) the directivity of the HJ ($\sim 30$ dB) and the return loss of the line ($\sim 25$ dB) 2) the accuracy of measurements of the electrical length of the transmission lines ($\sim 5$°). It can be reasonably expected that the coefficient of reflection seen by the klystrons will be lower than 0.5%.

![Graph](image)

**Fig.5.** Coefficient of reflection for the 2 rows of 8 modules

**Conclusion**

The full power division of the lower hybrid antenna, featuring 120 power dividers and 128 waveguides facing the plasma, has been calculated. It was found that 2 rows of 32 waveguides, using the same type of multijunctions, excite a very similar $N$ spectrum, but 2 rows with different types of multijunctions have different directivity. Due to the symmetry of the system, reflected fields recombining in the HJ are very close in phase and in amplitude, and the coefficients of reflection toward the klystrons are very low ($<0.1\%$). Even with departures from the ideal system, the reflection is low enough to consider a CW operation of the klystrons, able to sustain a VSWR of 1.5 ($R=1.7\%$), with no circulators.

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MAGNETODRIFT TURBULENCE AND DISRUPTIONS

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A qualitative model has been recently reported [1] for sawtooth relaxation in Tokamaks. Here we also discuss minor and major disruptions.

MAGNETODRIFT TURBULENCE

As proposed in that model, magnetodrift motions (md) are a particular type of motions with exceptional properties. They can exist only in configurations having along B lines a sufficient large scale asymmetry of the local magnetic shear.

In order that the kink motion can fully develop, the mhd flux conservation constraint has to be broken by some parallel electric field. In a hot plasma, the electron pressure term is dominant since the ratio of $V_\parallel / p_e / e_n$ by $\eta J_\parallel$ varies as $T_3$.

In a low $\beta$ plasma any perpendicular fluid velocity field is almost divergence free, and so $\nabla \times V \neq 0$ almost everywhere. Then during the disruption, strong turbulence must develop in spite of the magnetic shear. Thus the corresponding motions must be free of that constraint, which means that $\delta (J \times B)$ must be weak. Therefore we would have $d/dt (\nabla \times V) \sim 0$ and the perpendicular spectrum will be quasi-isotropic. From this property we deduce $k_\parallel \sim 2\pi s*/qR$ where $s*$ is a local value of the shear parameter.

An asymmetry of the local magnetic shear produces a variation of $k_\parallel$ along the B lines. When this asymmetry is sufficient, $\omega$ exists such that:
- In the region where $s*$ is weak, $\omega k_\parallel > V_e$, the electron continuity equation is dominated by its perpendicular part, and mhd dynamics is valid with $\delta E_\parallel \sim 0$.
- In the region where $s*$ is strong, $\omega k_\parallel < V_e$, the electron continuity equation is dominated by its parallel part, and drift wave dynamics is valid with $\delta E_\parallel \sim V_\parallel / p_e / e_n$. These motions can tap the mhd free energy ($Vp$ versus curvature) and become free of the global magnetic shear constraint through disposal of several wavelengths before toroidal closure.

MAGNETIC RELAXATION

Magnetodrift turbulence will produce fast transport. Braiding of magnetic lines (mhd + non mhd $\delta B$) and non equilibrium forces due to $\nabla p \leftrightarrow 0$ will result in a tendency of $J/B$ to become constant in the turbulent domain.
SAWTOOTH DISRUPTION

Magnetodrift turbulence is the main ingredient of the model proposed for the sawtooth relaxation D1 (disruption near and inside q=1). It triggers a fast transition from resistive kink to quasi-mhd kink and it allows a catastrophic increase and a full development of this convective motion. This effect of the turbulence is obtained by relaxing asymmetrically the current density spikes associated with the $B_\phi(r)$ field distorted by compression-expansion. An increase of the shear parameter $s_1$ near $q=1$ results. It produces an enhancement of the $\xi_{22}$ $(m/n=2/1)$ harmonic of the kink motion and thus an increase of the instability. Therefore the system enters a positive feedback loop which constitutes the trigger mechanism.

When the symmetry is sufficiently broken by the helical deformation, a strong burst of this turbulence produces the collapse of the density and temperature profiles and flattens the current density profile in some annular domain.

Various experimental aspects of D1 (cold bubble, crescent shape of the core, twisting, precursor and successor oscillations, partial sawteeth, snake motion) fit naturally in that model.

Experimental measurements of the specific fluctuations associated with the sawtooth relaxations in TFR plasmas [2] show several similarities with the proposed magnetodrift turbulence.

ROTATION AND HELICAL DEFORMATIONS

The experimental observation that the various particle and energy transport coefficients have similar numerical values whatever the particle mass and energy, strongly suggests that fluctuations of the electric drift velocity is the main turbulent transport process. Since collisional transport is much larger for ions than for electrons, the plasma will generally develop a negative charge density. The resulting radial electric field produces a poloidal rotation, which is slower in the central part because of the sawtooth relaxations which reduce the potential gradient.

Helical current pertubations can develop around the resonant magnetic surfaces with the most simple ratios $m/n=q$ due to tearing instability. When the resulting islands of neighbouring resonances interact, their spatial phases are no more independent. The potential energy $W$ of the configuration depends on their relative phase. The value of this phase which minimize $W$ will be imposed to the system through non equilibrium forces acting on the relative toroidal rotation.

As an exemple, a strong 1/1 deformation inside a 2/1 "cavity" will have a toroidal phase such that the 1/1 displacement directed outwards takes place in the poloidal plane where the flux surfaces are stretched along the equatorial direction by the 2/1 deformation. This particular phase permits the 1/1 displacement to be stronger in the direction of low magnetic field than in that of high field. In this situation the poloidal and toroidal rotations combine in such a way that the coupled helical deformations rotate toroidically at the same frequency.

Another noteworthy exemple is the coupling of two island systems, one due to plasma tearing current pertubations and one due to symmetry imperfections in the external currents. The common toroidal frequency for strong coupling is zero, which means that for a sufficiently large amplitude the plasma helical deformation will appear as blocked.
SOFT MINOR DISRUPTION

When the maximum of \( J'' \) (\( = \partial^2 J / \partial r^2 \)) is near the resonant radius \( r_2 \) (\( q = 2 \)) the \( J(r) \) profile is favorable to 2/1 tearing instability. The plasma evolves to this situation either by an increase of density (atomic losses displaced inwards) or by an increase of the current (\( r_2 \) displaced outwards).

The island helical deformation of the configuration produces an asymmetry of the local magnetic shear \( s^* \) which for \( r \leq r_2 \) is weakened near the hyperbolic axis and is enhanced near the ventral segment. On these magnetic surfaces near the separatrix, while the relative poloidal rotation of magnetic lines with respect to the island system is accelerated beside the ventral segment, it is reduced near the X point where they accompany the hyperbolic axis for several turns. There the poloidal rotation of lines is more regular than in the unperturbed equilibrium, and a weak ballooning of the perturbation is sufficient to tap the mhd free energy. An eventual enhancement of this effect can be produced by a 4/2 mhd harmonic induced by nonlinearity. Indeed this harmonic produces variations of the island angle that can result in an enlargement of the stagnation azimuthal domain.

We suppose that, from these geometrical properties, \( m \)d motions or \( d \) (drift wave) motions coupled to the mhd free energy (depending on the actual asymmetry of the local shear) become unstable in a small annular domain near and inside (\( r \leq r_2 \)) the inner branch of the separatrix.

The resulting turbulence that initiates the D2 disruption (near and inside \( q=2 \)) will not be azimuthally localized as in the case of D1 because the local shear asymmetry is weaker and the poloidal period is smaller. Density and temperature transport, and current density relaxation will result and favour an inward spreading of turbulence. This inward 2/1 deformation of the magnetic configuration will induce by toroidal coupling a 1/1 displacement of the central core.

HARD MINOR DISRUPTION

If the 1/1 displacement is sufficiently pronounced it will permit the development of \( m \)d turbulence near and inside \( r_1 \) mainly localized in the wake of this displacement. Then the resulting magnetic relaxation allows a fast growth of the internal kink motion.

Meanwhile the two main helical distortions 2/1 and 1/1 generate second order and high order islands by toroidal and nonlinear couplings. Thus transition to global stochasticity of magnetic lines [3] which are hamiltonian trajectories in real space can be reached.

Transport of electron energy and current relaxation will result producing an evolution of the magnetic structure. The time varying stochasticity can be considered as a particular sort of low frequency turbulence. Density and ion temperature will not be much affected because electric drift velocity is not a basic element of that turbulence. On the contrary, the electron temperature profile is expected to shift down as a whole for \( r < r_2 \), and the current density profile to become somehow flattened inside \( r_2 \) and stepened outside.
If the energy loss is important, the toroidal equilibrium is lost and the plasma ring contracts producing an increase of the total current and a decrease of the voltage, by inductance reduction. Scraping off by the inner limiter will enhance the edge steepening of the $J(r)$ profile.

**MAJOR DISRUPTION**

If the current density profile is sufficiently enlarged, external mhd kink can become unstable [4, 5]. Considering a situation with $2 < q_a < 3$, the $3/1$ unstable kink will develop a three petal-like deformation of the plasma edge. Scraping-off by the limiter will steepen the $J(r)$ profile near the edge and thus make the kink more unstable.

When the plasma edge, transferred inwards by scraping-off reaches the $q=2$ surface, the $2/1$ external kink will in turn develop. The final step will be the $1/1$ external kink propelling the remaining plasma to the wall.

More direct routes to major disruption can be given by any process generating an enlarged and edge steepened current density profile. Skin effect due to fast current rise, ballooning mode turbulence near the $\beta$ limit, strong plasma wall interaction, are three possible mechanisms.

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SPECIFIC TURBULENCE ASSOCIATED WITH SAWTOOTH RELAXATIONS IN TFR PLASMAS.

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INTRODUCTION

Recently we reported [1, 2] the existence of a specific turbulence during the sawtooth relaxation which superimposes the quasi-stationary (QS) turbulence.

Experimentally it corresponds to bursts of fluctuations observed on the high frequency side of the QS spectrum. The bursts appear synchronously with the relaxation and globally last 60 μs. It has been shown that the specific fluctuations cover the 0.4 to 3 MHz domain depending on the wave number and, at least, the 3 to 17 cm⁻¹ observed wave number range. In this range, the k spectrum follows a k⁻⁵±₁ rapidly decaying variation. The phase velocity of this turbulence is typically 8.10³ m/s, i.e. 4 times higher than the one of the QS turbulence.

A CO₂ coherent scattering diagnostic is used to perform the experiments. The scattering volume consists on a vertical chord which is radially scanned. The presently reported experiments, performed on TFR Tokamak, concern k = 5 or 7 cm⁻¹. A careful analysis of the temporal amplitude variation lets us define three stages in the turbulent process (Fig.1), in agreement with the soft X Ray signals. Each stage typically lasts 20μs. The first stage is associated with a displacement of the hot core. The increase of fluctuations amplitude is clear but modest. The second stage coincides with the main crash of the emissivity signal. Its amplitude is large. Finally, a third stage coincides with the reorganisation process during which the central region of the plasma recovers its symmetry. During this phase, the amplitude and spectrum of the turbulence are comparable with the ones of the first phase. The following reported data is uniquely concerned with the second phase and our attention is focused on the spatial localization of the specific turbulence.
SPATIAL LOCALIZATION

The specific turbulence power $\Delta S(\omega)$ is defined as the difference between the instantaneous spectrum and the average QS spectrum. By radially scanning the scattering volume, the frequency integrated specific turbulence power $\Delta n^2$ is measured and gives us Fig.2. From the scattered dots a first conclusion is drawn: the specific fluctuations only exist at the interior of the $q=1$ surface. The large dispersion of the points lets us suspect that a hidden parameter plays an important part.

Azimutal Localization

By following the m=1 precursor mode, observed on the soft X ray signals just at the interior of the inversion radius ($r < r_1$), the azimutal position of the cold side at the beginning of the first phase (fast displacement of the hot core) is noticed as $\theta$. We define a radius $r_0$ where the turbulence is supposed to be maximum. From the raw experimental points on Fig.2 we took $r_0 = 3$cm. Then $\theta_0$ is the azimutal angle of the $r_0$ circle intersection with the axis of the scattering volume. The amplitude of the specific fluctuations $\Delta n^2$ is symmetrical in $(\theta - \theta_0)$ and then Fig.3 can be presented as a function of $(\theta - \theta_0)^2$. The amplitude drastically depends on $(\theta - \theta_0)^2$, following an approximately exponential law. It is maximum for $\theta = \theta_0$ with an experimentally observed half width of $\sim 60^\circ$. We conclude that the turbulence is azimutally localized on the cold side.

Radial Localization:

If we add in Fig.2 the points obtained by extrapolating to the maximum of the fluctuations (i.e. for $\theta = \theta_0$), we get an envelop (the dashed line on Fig.2). From this azimutal and radial analysis we conclude that the fluctuations exist in a cell localized around a ring of radius $r_0 - 2.5$ cm $- r_1/2$.

Numerical Model:

The experimentally observed signals are the result of several convolutions. To estimate with more care the spatial parameters of the turbulence cell, we try to adapt a spatial level model decaying in both the $r$ and $\theta$ directions with gaussian profiles. Several hypothesises concerning the $k$ spectrum of the turbulence have been tested. These are: isotropic turbulence, radial $k_r$ turbulence, and azimutal $k_\theta$ turbulence. The results of the analysis are compared with the experiment and we conclude:

- The turbulence cell is located on the cold side.
- Its azimutal half width is $60 \pm 8^\circ$ (for $k = 5$ cm$^{-1}$)
- An isotropic spectrum in an asymmetrical cell (asymmetrical gaussian function) with $r_0 = 2$ cm, $\Delta r_i = 0.3$cm in the interior direction and $\Delta r_e = 2$ cm in the exterior direction gives satisfactory comparison; however, whereas a pure $k_r$ spectrum is definitively discarded, a $k_\theta$ spectrum with $r_0 = 3$ cm, $\Delta r_i = \Delta r_e = 2$ cm is also possible.
Heat Flux

The heat flux outside the q=1 surface is not azimuthally symmetrical. Analysis of the soft X rays signals lets us define the azimuthal direction in which the heat flux appears "sooner and faster". Assuming a constant rotation velocity of the hot core, the heat flux appears to be directed 130° away of the hot side i.e. apparently at one corner of the turbulence cell. If we share Andreoletti's model [3], the plasma rotation would be modified by the turbulent kink. The edge part (r < r_{1}) rotation slows down whereas the one of the most central part speeds up, producing a twisted state. Under these conditions, the direction of the "sooner and faster" heat flux agrees with the central position of the turbulence cell (Fig. 4).

With the same approach we observe a temporal and spatial localization of some turbulence in conjunction of the q=2 minor disruptions.

CONCLUSION

A cell of turbulence, localized in the interior of the q=1 surface with an azimuthal full width extension of about 120° has been detected and analysed. Its life time coincides with the collapse of the sawtooth. These observations are in fair agreement with the so called "magnetodrift" turbulence in the sawtooth disruption model proposed by Andreoletti.

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Fig. 1 Temporal analysis

Inside

Outside

20 µs

X = 0

k = 5 cm⁻¹

Fig. 2 Radial analysis

Inside

Outside

k = 7 cm⁻¹

Fig. 3 Azimuthal analysis

Fig. 4. Spatial localization during the collapse

(TFR plasma rotation is fast. Thus twisting is expected to be strong: ~ 50°)

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DETAILED STRUCTURE OF THE $q$ PROFILE AROUND $q=1$ IN JET

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1- INTRODUCTION

The limitation of ablation on rational surfaces has been shown to be an efficient mechanism of striation formation during pellet ablation [1]. In JET, a very large striation is observed when the pellet crosses the $q=1$ surface [2]. This paper presents a thorough analysis of the pellet ablation in this region and shows that an extended shearless zone around $q=1$ is necessary to reproduce the experimental signal. Such a feature is likely to be an essential ingredient in the understanding of internal disruptions.

2 - DESCRIPTION OF THE MODEL:

The basic idea is that dark striations correspond to locations where ablation is lower than average, i.e. that the available electron energy is lower there than at other radii. This can be understood as follows for electrons of parallel velocity $v_{\|}$: the pellet, surrounded by its neutral cloud (of effective diameter $\Phi_p$), moves across the discharge with a velocity $v_p$. It interacts during $\delta t = \Phi_p / v_p$ with every toroidal shell of infinitesimal thickness. Therefore, all the electrons located at a distance smaller than $\delta t v_{\|}$ from the pellet (following the field line) are intercepted by the neutral cloud. If the pellet is on a resonant magnetic surface, (corresponding to a rational $q$ value $q_R = m/n$), the field line has a finite length $L(q_R)$. If $L(q_R) < \delta t v_{\|}$, the energy flux on the neutral cloud vanishes after $\delta t' = L(q_R) / v_{\|}$, leading to the following reduction of the local ablation rate:

$$\alpha_R = \frac{L(q_R) v_p}{v_{\|} \Phi_p} \quad (1)$$
The influence of a rational surface extends until the difference of poloidal rotation (after m toroidal turns) reaches the angular sector sustained by the neutral cloud. Inside this region, the shape of the striation is linked to the q profile by the relation:

\[ \alpha = \alpha_R \left( 1 + \left( \frac{\Phi}{\Phi_p} \right) \left( \frac{2\pi m r}{q - q_R} \right)^{1/3} \right) \]  \hspace{1cm} (2)

provided the width of the striation is larger than \( \Phi_p \).

Inside \( \Delta r \), the observed \( H_\alpha \) signal is assumed to be proportional to the matter deposition \( \Delta M \):

\[ \Delta M = 4\pi r_p^2 \int_0^\infty \frac{dE}{\Delta r} \]  \hspace{1cm} (3)

where \( r_p \) is the pellet radius and \( \eta \) the ice density. The expression of \( \dot{r}_p \) given in [3] can be revised to take into account the shape of the striation \( \alpha \) and the Maxwellian distribution of the electrons \( f(E) \); leading to:

\[ \dot{r}_p = 5.14 \times 10^{-9} \int_0^\infty \frac{dE}{\Delta r} \left( \int_0^\infty \frac{dE}{\Delta r} \right)^{1/3} \]  \hspace{1cm} (4)

where \( n_e \) is the local electron density.

This model can be used to determine the details of the q profile around a resonant surface displaying a striation at this location, provided \( q_R \) is known. Indeed, \( \Phi_p \) (the only free parameter of the model) can be determined from the experimental value of \( \Phi_R \) by equations (1), (3) and (4); and then the q profile computed from the experimental shape factor \( \alpha \) by equations (2), (3) and (4).

3- THE q PROFILE AROUND q=1 IN JET:

An example is given for JET shot 9228 the main parameters of which are: \( n_e(0) = 3.7 \times 10^{19} \text{ m}^{-3} \), \( T_e(0) = 3.6 \text{ keV} \), \( q_\psi(a) = 4.8 \), and \( q=1 \) at \( r \simeq 51 \text{ cm} \) in the equatorial plane. The pellet has an initial radius \( r_{p0} = 2.2 \text{ mm} \) and is injected with a velocity \( \nu_p = 10^3 \text{ ms}^{-1} \).

On figures 1a to d, comparisons between model and experiment are shown for different q profiles around q=1: for a smooth q profile (\( \nabla q = 0.25 \text{ m}^{-1} \), fig.1a), the computed striation is too shallow and considerably narrower than the observed one. The large dip of about 14 cm wide exhibited by the experimental signal implies the presence of a flattening of roughly the same extent on the q profile [4]. This flattening is modelled with a constant q gradient: \( \nabla q = 0.03 \text{ m}^{-1} \) (fig.1b, and curve labelled b on fig.2) and a zero q gradient (fig.1c and 2, curve c). These two cases emphasize the extreme sensitivity of \( \alpha \) to the details of the q profile. Finally, the best fit of the experimental data is displayed on figure 1d and the corresponding q profile on
The diameter of the neutral cloud deduced from the maximum depth of the striation $\alpha_R$ is $\Phi_p = 0.7$ cm. Outside the striation, the residual discrepancy between the computation and the experimental signal is $\leq 20\%$ and local variations of the latter do not exceed this value. Expecting a similar accuracy inside the striation leads to $\delta q/q \leq 2.5 \times 10^{-4}$. Therefore, the extremely flat region ($Vq = 0.003 \text{ m}^{-1}$) of about 8 cm wide around $q=1$ is not an artifact of the computation.

4- DISCUSSION:
Two effects are neglected in this model: the radial heat and matter transport and the density structure of the neutral cloud which is not - stricto sensu - a perfectly absorbing sphere. However, for a given $q$ profile, they both tend to smooth the resultant striation. Therefore, to take them into account will strengthen further the strong $q$ flattening at the center of the striation, which gives yet more confidence in the reality of this feature.

The existence and persistence of an extended shearless annular region in the vicinity of the $q=1$ surface should be taken into account in any theory of the internal disruption. It is likely that this zone acts as a control to release the energy accumulated at the center of the discharge. The precise mechanism can be studied using pellet ablation as a shear diagnostic since an hypothetical $q=1$ island would be distinguishable by this method, even if local temperature gradients are very low. A first step is to study this shoulder on the $q$ profile at several times in the sawtooth period, to check if its evolution is consistent with normal resistive diffusion.

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Figure 1: Comparison between experimental (thin line) and computed (thick line) signals for different q profiles around q=1:
a. Smooth q profile (without flattening), b. Constant q gradient, c. Zero q gradient, d. Optimized q profile to best fit data.

Figure 2: (q - 1) profile around q=1 in the three cases b, c and d. For case d the error bar is indicated in the upper left corner of the figure.
TURBULENCE PROPAGATION DURING PELLET INJECTION

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ABSTRACT
It has been observed in the TFR tokamak that an injected pellet was preceded by a cold front propagating at about twice its velocity. The aim of this paper is to show that the observed phenomena can be interpreted as a manifestation of the ballooning structure of modes destabilised by the pellet.

I. INTRODUCTION
Details about experimental data during pellet injection in TJr can be found in [1,2]. The most striking features are:
- The "cold front" (position where the electron temperature starts decaying) propagates inwards with velocities ranging between \(10^3\) and \(2\times10^3\) m/s outside the \(q=1\) surface (increasing with decreasing minor radius), and much larger inside the \(q=1\) surface.
- Between this front and the pellet which moves more slowly (0.6 \(10^3\) m/s), there is a zone with enhanced transport (\(\chi_e = 10-100\) m²s⁻¹) that broadens in time up to a width of 0.1 m (half the plasma minor radius).
- When the pellet is fully ablated the heat transport goes back everywhere to its previous value in less than 10 \(\mu\)s.
- The density fluctuations (measured by scattering) are greatly enhanced. The spectrum shape is similar to the one without pellet. A striking feature is that the increase in the fluctuation level starts 30-40 \(\mu\)s after the pellet injection and stops 50 \(\mu\)s after the enhanced transport period.

During fast heat transport which lasts typically 150 \(\mu\)s, it seems that the particle transport is only weakly affected. Such a difference in the behaviour of \(n_e\) and \(T_e\) profiles has been reported by other authors [3,4]. It results that there is a turbulent zone propagating beyond the pellet.

A question addressed by these experiments is how the excited electromagnetic perturbations propagate radially. Two possible mechanisms already proposed are the spreading
of a turbulent zone by non-linear interaction [5] or linear destabilisation by local gradient at the turbulent front [2,6]. However, it should be pointed out that a radial propagation is a natural consequence of the toroidal mode coupling, as discussed in section II. In section III, we propose a scenario based of this model to explain the experimental facts.

II. A LINEAR MECHANISM FOR RADIAL PROPAGATION OF MICROINSTABILITIES

Heat transport is believed to be due to microinstabilities. They can be described by eigenfunctions localized around resonant magnetic surfaces coupled by toroidal effects. When the plasma properties are assumed to be invariant by radial translation, it can be shown that the general solution of the equations defining the instability can be written as [7]:

$$\delta E, \delta B = \sum h(\tau-n) \exp[i(\theta-n\Phi-\omega t+1\delta)]$$

where $h(\tau-n)$ are functions localized around resonant surfaces $\eta$ and $\delta$ is a phase factor. These surfaces are separated by a distance $d = (n dq/d\tau)^{-1}$. This expression is equivalent to the ballooning representation. Generally one considers solutions with $\delta = 0$. All the modes $h(\tau-n)$ are oscillating in phase with a frequency $\omega_0$. However, the general solution corresponds to arbitrary values of $\delta$ and $\omega$ is a function of $\delta$. At scales larger than $d$, the perturbation behaves like a travelling wave with a wave vector $K = \delta/d$ and a frequency $\Omega = \omega(\delta) - \omega_0$. When a perturbation is excited on a magnetic surface its energy propagates at the group velocity $v_g = \partial Q/\partial K$ (like phonons in a crystal). When the plasma is weakly inhomogeneous, WKB theory can hold for such waves. When the mode propagates in a zone where it is stable (growth rate $\gamma < 0$), the evanescent length is given by $L = v_g/|\gamma|$. In the general case $\omega(\delta)$ can only be computed numerically [7]. However a dimensional analysis shows that $v_g = \alpha \epsilon \nu^* / s$ where $\epsilon$ is the inverse aspect ratio representing the mode coupling, $\nu^*$ is the diamagnetic velocity, $s = (r/q)(\partial q/\partial r)$ the shear parameter and $\alpha$ is a constant.

III. A POSSIBLE SCENARIO TO EXPLAIN THE PLASMA BEHAVIOUR DURING PELLET INJECTION

a. Mode excitation

When the pellet reaches a flux tube, it creates a strong perturbation that lasts a time of order $\tau = r_{\text{pellet}}/v_{\text{pellet}} = 1\mu s$ [8]. When the crossed magnetic surface is rational ($q = l/n$) this excitation stays localized around a closed magnetic field line and the shear limits its radial extension $\Delta r$ to values given by:

$$\Delta r < \left(2 \frac{\partial q}{\partial \tau} \right)^{-1} = \frac{d}{n}.$$ 

The effect is maximum for $\Delta r = r_{\text{pellet}}$ i.e. $l = \sqrt{\frac{q r}{s r_{\text{pellet}}}}$. 

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These perturbations are short ($\tau \omega < 1$) and well-localized ($\Delta r < r_{\text{th}}$, $L = 30$). They have significant projections on the plasma eigenmodes. These modes may already exist in the plasma; they may be stable: in the second case the pellet creates forced oscillations.

The magnetic field is perturbed because the current is expelled out of the cold flux tubes (resistive perturbation) or because the field lines are perpendicularly displaced by the impact of the pellet (shear Alfvén waves). The density is increased drastically in the crossed flux tubes (electrostatic drift waves). However, in the latter case, the toroidal equipartition time limited by the ambipolarity constraint is a few tens of microseconds.

From the experimental point of view, the difference between the heat and particles transports and the fact that the density perturbations are delayed suggests that magnetic perturbations are responsible for the enhanced transport.

Cold front propagation

According to our model, these perturbations should propagate radially. The observed velocity ($1000 \text{ ms}^{-1}$) is compatible with our previous estimate taking $\alpha = 4$. (On TFR, $v^* = 2500 \text{ ms}^{-1}$, $\varepsilon = 0.2$ and $s = 2$ at the edge). The acceleration of the cold front is compatible with the decrease of $s$ along its trajectory. Since transport stops as soon as the pellet disappears, the excited modes are damped. The fact that the turbulent zone extends up to $L = 10 \text{ cm}$ beyond the pellet imposes an upper bound for the damping rate $|\gamma| < \frac{v_g}{L} = 10^4 \text{ s}^{-1}$.

A noticeable feature is that the fast transport stops everywhere within $10 \mu s$. This implies that the information propagates at a velocity of $10^4 \text{ ms}^{-1}$, no more compatible with the precedent group velocity. However, a possible explanation is that behind the cold front the plasma is strongly turbulent: in such a situation, non-linear mode interaction could dominate the toroidal coupling. In the expression of the group velocity, $\varepsilon$ (representing the toroidal effects) should be replaced by a larger factor.

In JET, a cold front propagating faster than the pellet has not been observed. This could be explained by our model in two ways:

- In large tokamaks the group velocity can be less than the pellet velocity, because the diamagnetic velocity is lower and the pellets are injected at higher speed.
- The evanescent length $L = \frac{v_g}{\gamma}$ is much smaller than the plasma minor radius. In this case, although the turbulence may propagate faster than the pellet, the cold front is observed to stay at a constant distance $L$ from the pellet. This interpretation corresponds to the experimental data on JET [9] with $L = 5$ to $10 \text{ cm}$.
IV. CONCLUSION

Toroidal coupling of modes is the simplest mechanism allowing a radial propagation of low frequency modes that otherwise would remain localized near resonant surfaces. Such a phenomenon should have a strong influence on anomalous transport. For example, this scheme can explain how edge turbulence affects the plasma bulk and why the saturation levels at different radii are coupled.

The study of radial propagation of perturbations can be an useful way to test the properties of microturbulence. First, it can provide a macroscopic evidence of the ballooning structure and then gives quantitative informations about micromodes from the experimental values of $L$ and $v_g$. Pellets, although not very selective for the excited modes are an useful experimental tool for such investigations.

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I Introduction

It is reasonable to believe that plasmas in thermonuclear reactors will have to be steady state. In the case of tokamaks this requires a large amount of recirculated power to maintain the plasma current (a few hundred of megawatts in the present state of the art). However, a part of this current, the bootstrap current, is expected to be generated by the plasma itself. We investigate here the concept of a reactor with 100% bootstrap current obtained by monitoring the $n_e$ and $T_e$ profiles.

II A tokamak discharge with 100% bootstrap current

We consider the case of a discharge with equal ion and electron temperature ($T_i$ and $n_i$). In the collisionless case, the expression of the bootstrap current in a tokamak is:

$$J_B(r) = \frac{2n_e(r)T_e(r)}{B_\phi(r)} \left[ \frac{r}{R} \left( \frac{2.3}{L_N} + \frac{.17}{L_T} \right) \right]$$

where $B_\phi$ is the poloidal field, $R$ the major radius, and $L_N, L_T$ the gradient lengths. Our starting point is that $J_B$ mainly depends on the density gradient and may be increased by decreasing $L_N$, while maintaining the pressure gradient below the MHD instability threshold. We consider here the case of a discharge of minor radius $a$, major radius $R$, toroidal field $B_t$ with the following characteristics:

a) Inside a given radius $r=r^*$, the safety factor profile is flat with $q=1$. The pressure profile $p(r)$ is also flat to avoid pressure-driven modes i.e. $L_N = -L_T$. Equation (1) defines the $n_e$ profile which is peaked at the plasma centre, while the $T_e$ profile is hollow (see Fig. 1):

$$n_e(r) = n_{eo} e^{-\frac{1.25}{\beta_t} \left( \frac{r}{R} \right)^3}$$

with $\beta_t = p(0) / \left( \frac{B_t^2}{2 \mu_0} \right)$
b) Outside the \( r = r^* \) (i.e. \( q = 1 \)) surface the pressure \( p(r) \) decreases from \( p(r^*) = p(0) \) to \( p(a) = 0 \) at the plasma edge. The ratio \( \frac{L_H}{L_T} = \eta \) is assumed to be constant. The pressure gradient is limited by a simplified ballooning stability criterion [1]:

\[
2 \mu_0 \left( \frac{\partial q}{\partial r} \right) = -\frac{0.6}{R q(r)^{\frac{3}{2}}} \quad \text{where} \quad s = \frac{r}{q} \frac{\partial q}{\partial r} \quad (3)
\]

Combining equations (1) and (3) one obtains a differential equation for the profile of the safety factor

\[
\frac{1}{q(r)} \frac{\partial q}{\partial r} = \frac{2 r^{-\frac{5}{2}}}{r^{-\frac{5}{2}} + \left(\frac{0.69 + 0.048 \eta}{1 + \eta}\right)} R^{-\frac{5}{2}} \quad (4)
\]

Giving the value of the safety factor at the edge \( q(a) \) determines the solution of the preceding equation. It finally appears that \( R/a, q(a) \) and \( \eta \) determine \( p(r), q(r) \), the volume averaged beta and the ratio \( r^*/a \). This is also the case for \( n_e(r)/n_{eo} \), \( T_e(r)/T_{eo} \) and the "hollowness parameter" \( \lambda = T(r^*)/T(0) \). The equi beta and equi-hollowness profiles are drawn in the \( q(a) - R/a \) plane for \( \eta = 3 \) in figure 2. The beta value given by the Troyon limit which can also be represented this plane is drawn in figure 2b. When \( \eta \) is small the beta values and the hollowness parameter for a given \( q(a) \) are smaller since the density gradient which creates the bootstrap current is larger.

II Discussion of the results; relaxation of the current profile

The above simple model shows that driving a discharge with 100% bootstrap current requires unusual \( n_e \) and \( T_e \) profiles. We tentatively examine how these profiles could be obtained in a thermonuclear reactor:

- \( n_e(r) \): The peaked \( n_e \) profile can be obtained by repeated pellet injections. The particle confinement time is expected to be large enough so that the pellet injection rate will be reasonable. However an usual pellet will deposit its particles at the plasma edge. A specific technology has to be developed. It could consist in shielded pellets, ablated on magnetic axis by a laser beam. Such a development looks quite reasonable compared to what is needed for other current drive methods and could in fact solve the center fuelling problem.

- \( T_e(r) \): The \( T_e \) profile will naturally become hollow if the pellet is ablated at the plasma center. The duration of the dip at the center will depend on transport phenomena and will be increased by the hollow thermonuclear power profile and the peaked bremsstrahlung (and possibly of a small amount of heavy impurity) radiation profile.
The obtained $j(r)$ profile will not be flat inside the $q=1$ surface because the $n_e$ and $T_e$ profiles obtained by pellet injection will not correspond exactly to expression (2). Moreover the expression (1) of bootstrap current is not valid in a small zone around the magnetic axis. One can escape to these difficulties by using the fact that for large enough bootstrap effect inside the $r=r^*$ surface, the current profile will tend to become hollow. The $q$ values will be slightly above and below 1 at $r=0$ and $r=r^*$ respectively. Such profiles are unstables [2,3]. One may accordingly hope repeated sawtooth-like redistribution of the current profile towards a lower energy state flat $q=1$ profile. Such a mechanism has already been suggested in [4]. One may then show that the bootstrap current $J_B$ has no longer to be locally equal to the actual current $J_p$. Rather the bootstrap electromotive force $J_B/\rho$ ($\rho$ is the plasma resistivity) must now simply balance the resistive friction $J_p/\rho$ in average in the $r=r^*$ surface. The question of course arises whether the particle and energy losses unavoidably associated to the considered disruptive relaxations are acceptable or not.

It should be noted that in recent JET discharges with pellet injection, hollow $q$ profiles have been obtained with values around 1.5. They were terminated by a sawtooth-like event associated to a m/n=3/2 mode [5].

### Compatibility with a thermonuclear reactor

Coming back to the model of section II, for given $R,a,q(a),\eta, B$, it is possible to define a reactor discharge. It is chosen to have average beta and density values below the Troyon and Murakami limits. The energy confinement time is twice the Kaye-Golston scaling. We find for example that the following circular cross section discharge is ignited and delivers a fusion power of 3.3 GW (wall load 2.45 MW m$^{-2}$):

- $R=9$ m
- $a=3$ m
- $r^*=1.1$ m
- $q(a)=2.55$
- $I=15$ MA
- $B=7.1$ T (13 T at the conductor)
- $Z_{eff}=1$
- $\eta=10^{20}$ m$^{-3}$
- $T=12.5$ keV
- $\beta=1.9\%$

The hollowness parameter $\lambda$ is 3.1. The corresponding $T_e$ and $n_e$ profiles are shown in figure 1.

### Conclusion

A concept of reactor with 100% bootstrap current using pellet injection and dynamo effect has been presented. Its viability depends on a low particle and heat transport in the plasma center and of the possibility that plasma current will relax repeatedly around a state of flat $q=1$ profile in the central region. It is difficult to make predictions concerning these two points since they are not fully understood in existing plasmas. However current experimentation in tokamaks should improve our understanding of these two points.
References:

FIG. 1: Temperature and density profiles in the reactor proposed in section IV.

FIG. 2 a: Contours of constant hollowness parameter λ in the q(a)-R/a plane for η = 0.3. The cross represents the position of the reactor of section IV.

FIG. 2 b: Contours of constant average beta (solid line) for η = 0.3. The dotted line represents the Troyon β limit for η = 0.3 with q = 2.8.
OPTIMIZATION OF A STEADY STATE TOKAMAK DRIVEN BY LOWER HYBRID WAVES

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1 - INTRODUCTION

As suggested by experimental results obtained on many tokamaks steady state operation appears feasible with the plasma current \( I_p \) driven by lower hybrid waves (LHW). By this way it is expected to increase reactor reliability and availability. In the present paper, we consider a steady state device like NET/ITER [1] driven by LHW. The power multiplication factor \( Q \) will be an essential parameter; so we have determined optimum plasma characteristics for maximizing \( Q \) value under the restriction of the LHW penetration limitations. In all expressions, MKS units are used with temperature \( T \) in units of 10 kev, current in megamperes (MA), power in megawatts (MW) and densities in units of \( 10^{20} \) e.m.

2 - CHOICE OF THE RF PARAMETERS

2.1 According to the quasi-linear theory of current drive by LHW [2], the current merit factor \( \eta \) is given by [3]:

\[
\eta = \frac{n_e R I_{HF}}{P_{HF}} \approx 2 \cdot S(N_{H}) \cdot \frac{W(N_{H}, T_e, Z_{eff})}{N_{H}} \cdot \frac{4}{5+Z_{eff}}
\]

with:

\[
S(N_{H}) = \left[ \frac{1}{2} \left( \frac{N_{H}}{N_{Hm}} \right)^2 \right] \cdot \sqrt{2} N_{Hm} \cdot \log \left( \frac{N_{H}}{N_{Hm}} \right)
\]

where \( R \) represents the major radius, \( I_{HF} \) the toroidal current driven by LHW and \( P_{HF} \) the HF power injected into the torus. In equation (1b), \( N_{Hm} \) and \( N_{Hl} \) are the upper and the lower refractive index which reach the central part of the plasma. The improvement in current drive efficiency \( W \) has been calculated by KARNEY and FISCH [2]. We use an approximate expression given by EHST and EVANS [4]:

\[
W(N_{H}, T_e, Z_{eff}) = 1 + \left\{ \frac{4.9 \cdot 10^{-3} (5+Z_{eff}) + 3.2 \cdot 10^{-3} (5+Z_{eff})}{2.6 \cdot 10^{-3} (5+Z_{eff}) N_{Hm} T_e / Z_{eff}} \right\} \cdot \frac{N_{H}}{N_{Hm}}
\]

2.2 We see, from equations (1a) (1b) (1c), that the current drive merit factor \( \eta \) can be maximized by:

i) working at the lowest \( N_{H} \) (limited by the accessibility condition \( N_{H} \gg N_{Hacc} \)) with a spectrum width \( \Delta N_{H} = N_{H} - N_{Hm} \approx N_{Hm} \), as small as possible. Here the \( N_{H} \) -spectrum excited by the coupling structure is determined by its lower \( N_{Hm} \) and upper \( N_{Hl} \) limits.

ii) increasing the fraction \( \frac{M_{abs}}{M_{acc}} \) of the injected power which is absorbed by the electrons (\( \eta_{HF} \)). We have [3]:

\[
\frac{M_{abs}}{M_{acc}} = \frac{P_{HFE}}{P_{HF}} = \frac{M_{D}}{M_{acc}} \cdot \frac{1 - \frac{\Xi}{M_{p}}}{}
\]

where \( \mu_{D} < 0.8 \) represents the directivity and \( \frac{M_{acc}}{M_{p}} \) the accessible part of the \( N_{H} \) -spectrum injected into the torus [3]. In order to minimise the HF absorption by background ions and \( \alpha \)-particle (\( \Xi/M_{p} \)) it is necessary to choose an appropriate frequency \( f \). Wave damping on fuel ions is avoided when [3]:

\[
73
\]
Wave damping on α-particle can be avoided simply making the perpendicular wave phase velocity, (c/Nj, with Nj perpendicular refractive index) faster than the α birth velocity, i.e. Nj > 23.3. From the cold dispersion relation, without toroidal effects, we obtain:

\[
\frac{f_3}{f_1} \approx \left( \frac{\alpha \nu}{\nu} \right)^{1/2} \left[ \frac{1 + 2.33 \nu/2}{1 + 10.3 \nu} \right]^{1/2} \left[ 1 + \left( \frac{4 \nu^2 (1 + 10.3 \nu)}{1 + 10.3 \nu} \right)^{1/2} \right]
\]

For a given value of Nj > Njacc the maximum frequency is determined by the accessibility condition:

\[
f_{\lambda} \leq f_3 \approx 2.1 \left( \frac{\alpha \nu}{\nu} \right)^{1/2} \left[ 1 - \frac{N_{jacc}^2}{N_{jacc}^2 + 6.4 \nu^2 N_j/B_T} \right]^{1/2}
\]

In practice we always have \( f_{\lambda} < \frac{\nu}{\nu} \). On the figure 1 we have represented the evolution of the frequency limits \( f_1 \) and \( f_3 \) as a function of the density for a toroidal magnetic field \( B_T = 5 \) teslas, \( N_{jacc} = 2 \) and \( \Delta \nu = N_{jacc} - \nu \approx 0.25 \). We see that in order to avoid wave damping on α-particle until \( \nu < 10^{29} \) e.m. a frequency higher than 4.5 GHz seems necessary.

iii) Optimizing the wave penetration to a hot plasma core. Complete power absorption by quasi-linear electron LANDAU damping occurs for a given ratio between the parallel phase of velocity (c/Nj) and the thermal speed of the electrons (kTe/me)\(^1/2\). From \( /5/ \) we obtain the maximum penetration temperature given by \( (k = b/a = \text{plasma elongation}) \):

\[
N_{\nu_{Te}} \text{Te} \leq 2.5 \times 10^{-3} \frac{\nu^2}{\nu_{Te}} R \times \nu \times \left( \frac{\nu_{Te}}{\nu} \right)^{1/2} \frac{\Delta \nu / \nu_{Te}}{N_{\nu_{Te}} / \nu_{Te}} \frac{T_{\nu_{Te}}}{T_{\nu_{Te}}}
\]

The evolution of \( N_{\nu_{Te}} \) as a function of density is shown in figure 2, for \( (\Delta \nu / \nu_{Te}) / \nu = 10^{-3} \nu_{Te} \) and \( \nu = \nu_{Te} \). We see that wave penetration can be largely improved by choosing small densities, narrow \( N_{jacc} \)-spectra, large input power. Finally for \( 1.5 < \nu < 2 \) and considering the wave penetration limited by quasi-linear LANDAU damping, we can determine the value of the enhancement factor \( \nu \) given by equation (1c) and shown in figure 2. Typically \( \nu \approx 1.8 \nu_{Te} \).

2.3 We have calculated the current merit factor \( \gamma \) considering profiles in the form \( \rho = \rho_0 (1 - \xi) \) where \( \rho_0 = \text{ne}, T \). Thus average density (ne) and density-averaged temperature (T) are given by \( \langle \text{ne} \rangle = \text{ne} / (1 + \alpha \nu) \) and \( \langle T \rangle = T (1 + \alpha \nu) / (1 + \alpha \nu + \alpha \nu) \). The coupling structure of the grill type is composed of a network of \( N_{\nu_{Te}} \) waveguides along the toroidal direction with a phase difference between the HF electric fields in adjacent waveguides optimized for each value of density in order to maximize \( \gamma \). For \( 4 < B_T < 6 \) Teslas and \( 3 < \nu < 5 \) GHz, the results are plotted in figure 3. For \( \nu \approx 0.4 \times 1.5 \times 10^{20} \) m \(^{-3} \), we can derive a simple analytic expression:

\[
\gamma \approx 0.1 N_{\nu_{Te}}^{1/3} B_T / [1.5 + 2.5 + (4 + 3) \nu_{Te}] = 9.18 N_{\nu_{Te}}^{1/3} B_T / [1.5 + 2.5 + (4 + 3) \nu_{Te}^{1/3} \text{ne}]
\]

3 - LOWER HYBRID DRIVEN STEADY STATE TOKAMAK

3.1 In steady state operation \( Q = PFUS / PHFT \), where PFUS represents the total fusion power delivered by the plasma, will be maximized. Since HF power for current drive heats the plasma, the HF current drive is strongly coupled with the power balance of the plasma. We use the following zero-dimensional global power balance equation:
where the input power $P_{\text{input}} = P_{\text{fo}} + P_{\text{HFT}}$ and $P_{\text{fo}}$ is the total α-power delivered by the plasma. Here $W_p$ is the plasma energy and the global energy confinement time $\tau$ given by GOLDSTON:

$$\tau = 0.037 H I_p \frac{R}{a^2} \left( \frac{1}{1.5} \right)^{1/2} \frac{P_{\text{HFT}}}{P_{\text{input}}}$$

where $H$ represents the enhancement factor of the L-mode confinement. The plasma current includes the bootstrap current contribution $I_{\text{BS}}$ and we have:

$$I_{\text{BS}} = I_p (1 - \chi)$$

$\chi = I_{\text{BS}} / I_p$ has been determined by M.Y. HSIAO et al. /7/, and:

$$I_p = \frac{5 \cdot B_T [1.17 - 0.65 (\frac{\rho}{\rho_T})]}{\frac{H p}{q} \left( \frac{1 - \gamma}{\rho_T} \right)^2}$$

3.2 The numerical applications are made using $Z_{\text{eff}} = 1.5$, $Z_p = 1.2$ and $\gamma = 0.5$ at $\gamma = 1$. Considering the ITER parameters for the technological phase /1/ ($A = 1.2$, $a = 0.3$, $q = 3$) we have plotted in figures 4a and 4b for various values of $B_T$ the evolution of $Q$ and the neutron wall loading $W_n$ as a function of $\langle n_e \rangle$ when $H = 1$ and $H = 1.2$. In order to ensure the penetration of the L.H.W until the plasma core it is necessary to limit the density $\langle n_e \rangle$ (given by the relation (6) when $N_{\text{LMW}} \sim N_{\text{LMW}}$). The results are indicated on the following table.

<table>
<thead>
<tr>
<th>$B$ (T)</th>
<th>$Q$</th>
<th>$P_{\text{EVS}}$</th>
<th>$P_{\text{HFT}}$</th>
<th>$\frac{P_{\text{HFT}}}{W_{\text{ne}}} (\frac{\text{MW}}{\text{m}^2})$</th>
<th>$\frac{\langle n_e \rangle}{\text{cm}^{-3}}$</th>
<th>$&lt;T&gt;$ KEV</th>
<th>$\chi$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$H = 1$</td>
<td>5</td>
<td>2.65</td>
<td>398 MW</td>
<td>150 MW</td>
<td>0.78</td>
<td>10.7</td>
<td>0.14</td>
</tr>
<tr>
<td></td>
<td>5.5</td>
<td>3.6</td>
<td>483</td>
<td>134</td>
<td>0.58</td>
<td>0.73</td>
<td>12.6</td>
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<tr>
<td></td>
<td>6</td>
<td>5</td>
<td>586</td>
<td>117</td>
<td>0.69</td>
<td>0.68</td>
<td>14.9</td>
</tr>
<tr>
<td>$H = 1.2$</td>
<td>5</td>
<td>5.2</td>
<td>405 MW</td>
<td>78 MW</td>
<td>0.48</td>
<td>0.56</td>
<td>14.9</td>
</tr>
<tr>
<td></td>
<td>5.5</td>
<td>8</td>
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<td>63</td>
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<tr>
<td></td>
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<td>14.3</td>
<td>589</td>
<td>41</td>
<td>0.69</td>
<td>0.4</td>
<td>25.4</td>
</tr>
</tbody>
</table>

4 - CONCLUSIONS

We have determined the main characteristics of the RF system (frequency $f > 4.5$ GHz, $W_{\text{ne}} / N_{\text{LMW}} < 10^{-2}$) in order to ensure the wave penetration until the plasma core and to maximize the current merit factor $P_{\text{MF}}$. A simple analytic expression of $\chi$ has been obtained. It seems possible in ITER, using LHW alone, to reach $Q \sim 8$, $W_n \sim 0.7$ MW/m$^2$ with $\langle n_e \rangle = 0.5 \cdot 10^2$ m$^{-3}$ and $P_{\text{HFT}} = 60$ MW, if a moderate value of the enhancement factor $H = 1.2$ can be achieved.

/1/ ITER Concept. Definition ITER 1 October 1988
/5/ R.W HARVEY et al - Third Topical Conf. RF Plasma Heating-PASADENA (USA) January 1978 - A 71
/6/ R. GOLDSTON - P/L= a Physics and Controlled Fusion-V26 N1 (1984) p.87
/7/ MING YUAN HSIAO et al - ANL/FPP/TM 221 May 1988
Evolution of the frequency limits \( f_2 \) and \( f_3 \) for \( B_T = 5 \) Teslas and \( \alpha = 0.4 \) as a function of \( n_e \).

![Graph 1](image1.png)

Limit \( W_{\text{lim}} \) and enhancement factor \( W \) (\( R/a = 3.1, a = 1.8n, k = 2 \))

![Graph 2](image2.png)

Evolution of normalized \( P \) as a function of \( \langle n_e \rangle \): (0 calculated points)

![Graph 3](image3.png)

Evolution of \( Q \) as a function of density \( \langle n_e \rangle \) and neutron wall loading \( J_n \) as a function of density \( \langle n_e \rangle \) for:

- \( H = 1 \) and \( H = 1.2 \)
- \( B = 6 \) T
- \( B = 5.5 \) T
- \( B = 5 \) T

![Graph 4](image4.png)
1. INTRODUCTION

It became more and more evident during the last years that the total gas balance in the fusion machines was largely dependent on the recycling properties of the graphite which is the principal component of the inner wall.

Numerous studies, reviewed e.g. in [1,2], have been done to increase the physical understanding of phenomena like retention, detrapping and diffusion of the hydrogen isotopes, which controls the recycling behaviour of the graphite.

If the graphite implanted by energetic ion beam has been extensively studied, less work was made on thermally exposed graphite [3-5] and very little information is available for low energy high density plasma exposure typical of the inner wall conditions in the Tokamaks.

Therefore, we describe in this paper the first results that we have obtained by implanting pyrolytic and polycrystalline samples of graphite with a Deuterium plasma source. The density of the plasma was $5 \times 10^{12}$ cm$^{-3}$ and the electronic temperature was $T_e = 15$ eV. After the plasma exposure, the samples were analysed by thermodesorption with different heating rates.

The thermodesorption spectra show a rather complex structure with peak desorption at lower temperature than those obtained by the authors (see e.g. [6] and ref. therein) which used energetic ion beam implant but comparable to the experiments [5] made with graphite thermally exposed to hydrogen isotopes.

The last part of the paper gives the results of the comparison between the experimental data and numerical simulation of detrapping processes and diffusion processes.

We show that the principal features of our desorption data can be explained as well by diffusion processes as by detrapping and recombination mechanism, with in the two cases activation energies between 0.2 and 1.5 eV.

2. EXPERIMENTAL

The experimental set-up will be described elsewhere with more details. It involves one plasma source and one thermodesorption device. The graphite sample can be moved without breaking of vacuum from the plasma source to the analysis apparatus.

2.1. Plasma source

It is a duopigatron source [7] with an anticathode made of polycrystalline graphite. The electronic temperature was set at 15 eV and the density is $5 \times 10^{12}$ cm$^{-3}$.

In these conditions, the electrically floating sample is implanted by Deuterium ions of about 45 eV and with a flux of $\approx 10^{18}$ s$^{-1}$.

The plasma density and temperature have been measured by a double Langmuir probe at different positions and checked by the value of the sheath potential between the anode and the anticathode. We have also verified the results by the ratio of $H_2$ and $H_\alpha$ lines.

2.2. Thermodesorption apparatus

The analysis vessel has a base pressure of $1.10^{-10}$ torr. Thermodesorption spectra are
recorded from a quadrupole mass spectrometer. The sample is heated by electronic bombardment on the side not exposed to the plasma. Temperature is measured by a thermocouple inserted inside the sample. The heating ramp, numerically driven, is perfectly rectilinear and its rate may be changed from 0 to 2.5 Kelvin s^{-1}.

The temperatures obtained were checked by an infrared pyrometer. In thermodesorption experiments, the temperature uniformity across the sample is essential. It was inspected by Infrared Camera and color analysis did not show noticeable temperature variation.

2.3. Samples

The samples were made with Carbone-Lorraine 5890PT (as used for the Tore Supra first wall) and pyrolytic graphite from Union Carbide. Dimensions are 11x11x2 mm. They are mounted on a molybdenum support. The plasma exposure is done by an aperture of 9 mm diameter in the anticathode.

Two series of experiments were made. The first one with the sample directly screwed onto the support; the second one with the sample electrically and thermally isolated from the support by little ceramic pieces. We have verified that for the second serie, the support did not degas during thermodesorption experiment by implanting a sample as usual and replacing it by another identical sample which had not been exposed to the plasma. Then, we proceeded a thermodesorption experiment during which the D$_2$ signal remained negligible. This fact rules out any influence of the support on the thermodesorption spectra.

3. RESULTS

A typical desorption spectrum is shown in Fig. 1 giving evidence for four peaks at least.

![Experimental desorption spectrum](image)

fig 1. Experimental desorption spectrum. Heating rate 1.5 Kelvin s^{-1}

--- Numerical simulation of diffusion mechanism. Desorption energies: 1.3, 1.05, 0.73, 0.60 eV. $I^2/D_0=10^{-6}$ s^{-1}. See text below.

Preliminary experiments with exposure time to the plasma from 0.4 s to 12 s have not shown differences in the amount of implanted Deuterium. That is coherent with the well-established [1] fact that graphite retains all the non reflected hydrogen until the saturation concentration is reached. However the observed height of the first peak is lower when a same fluence is obtained by one shot than by some shots separated by time intervals.

We suppose that the higher temperature rise of the sample for higher times of
continuous exposure is the reason for this lowering of this first peak at 420 Kelvin. This was confirmed by an identical lowering of the first peak when the initial temperature of the sample was 100 Kelvin above room temperature.

Series of desorption were made at heating rates of 0.5, 1, 1.5 and 2.5 Kelvin s\(^{-1}\). The spectra are reproducible and the estimated positions of peaks are given in Table 1.

<table>
<thead>
<tr>
<th>Heating rate</th>
<th>I</th>
<th>II</th>
<th>III</th>
<th>IV</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.5</td>
<td>410</td>
<td>480</td>
<td>690</td>
<td>820</td>
</tr>
<tr>
<td>1</td>
<td>430</td>
<td>500</td>
<td>720</td>
<td>840</td>
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<tr>
<td>1.5</td>
<td>444</td>
<td>510</td>
<td>750</td>
<td>860</td>
</tr>
<tr>
<td>2.5</td>
<td>520</td>
<td>765</td>
<td>895</td>
<td></td>
</tr>
</tbody>
</table>

Table 1: Peak Temperatures (Kelvin) for different heating rates

As expected, the peak positions are shifted to higher temperatures when the heating rate is increased.

The experiments done with pyrolytic graphite have given the same results for the peak positions and the general features of the spectra with a lower first peak, due to a temperature rise during the bombardment more important for the less conducting pyrolytic graphite along the c axis.

4. DISCUSSION

With high energetic ion beam exposition [6], the thermodesorption spectra show structures with peak temperatures around 1100 Kelvin; these structures are sometimes washed by the high rate of temperature rise. These experiments were, in most cases, interpreted by second order detrapping processes.

Our results are more similar to those of Atsumi and al. [5], where the interaction between graphite and deuterium was obtained by high temperature high pressure exposition of the graphite to the gas. They obtained structured spectra with peaks in the same range of temperature as our results. They have supposed that bulk and pore diffusion processes were the dominant mechanism of the desorption.

We have done two series of numerical simulation; the first with a second order process (detrapping and recombination), the other one with a diffusion model.

In the first case, the pressure \( P(t) \) is given by

\[
P(t) = \alpha n^2(t) \exp \left( -\frac{E_d}{kT(t)} \right)
\]

where \( n \) is the number of implanted species, \( E_d \) the detrapping energy of activation, \( T \) the temperature and \( \alpha \) a coefficient proportional to the preexponential term of the detrapping process and dependent on the pumping rate and geometrical features of the device.

Then, if \( v \) is the heating rate, the peak temperature \( T_p \) is given by the relation

\[
\frac{E_d}{kT_p^2} = \frac{2}{v} n(T_p) \alpha \exp \left( -\frac{E_d}{kT_p} \right)
\]

which needs simple numerical computation to be solved.

In fact, thanks to the minor variation of \( n(T_p) \) versus \( v \), numerical trials gave same results for the peak temperature variation versus the heating rate with first and second order processes. We can therefore use the well-known first order result where the slope of the quantity \( \log T_p^2/v \) versus \( 1/T_p \) gives the activation energy. If the curves thus obtained are straight lines, it is a presumption of validity for the model. These curves are shown in Fig. 2 for the four peaks that we obtained.

Therefore Fig. 2 shows that our results can be explained by a detrapping model of first or second order with low activation energies of 0.7 eV, 0.51 eV, 0.43 eV, 0.25 eV.
Fig. 2. Experimental and simulated values of \( \log \frac{T_p}{V} \) for a detrapping mechanism of 0.7, 0.5, 0.43, 0.25 eV activation energies. Heating rate in Kelvin min\(^{-1}\).

We have also simulated a diffusion process with a simple model of an initial implanted slab of uniform concentration between \( x = 0 \) and \( x = 1 \) in a semi-infinite solid \( x > 0 \). Such a simple initial concentration gives a good approximation of the experimental distribution observed by Causey and al.\([3]\) in the low temperature exposure case. During the desorption, the surface concentration is set to zero.

The solution of this model is almost analytic\((8)\), including only numerical computation for the change of time scale\((9)\) to \( I = \int_0^t D(t) \, dt \) and gives

\[
P(t) = \frac{\alpha n_0 D}{4V} \left( 1 - \exp\left(-\frac{t^2}{4V}\right) \right)
\]

Our results can be adequately described by such a model with activation energies of 1.3 eV, 1.05 eV, 0.73 eV, 0.60 eV and \( D_0 = 1 \, \text{cm}^2 \text{s}^{-1} \) for an implantation length of 30 \( \mu \text{m} \) which compares with the results of Causey and al.\((3)\).

Numerical simulation gives peaks with a tail at the left side for second order desorption and at the right side for the diffusion mechanism, but the overlapping of the peaks prevents the use of the form of the spectra to take conclusions.

In conclusion, independent measurements of the diffusion coefficients seems to be necessary for a choice of one mechanism. We are actually working on preliminary results of incoherent quasi-elastic neutron scattering experiments and we hope that we shall obtain at least an upper limit for the diffusion coefficient.

BIBLIOGRAPHY
INTRODUCTION

The injection in the plasma, of electromagnetic waves at a frequency close to the lower hybrid (LH) frequency has proven, in many tokamaks, to be very effective for current generation and plasma heating through LANDAU damping on the electron population. We report here the first results of lower hybrid current drive (LHCD) experiments in TORE SUPRA. As it is known that the density limit for wave electron interaction increases with the frequency /1/, a 3.7 GHz system has been built in order to investigate the current drive efficiency at high densities (~ $10^{20} \text{m}^{-3}$).

EXPERIMENTAL SET-UP - WAVE TRANSMISSION AND COUPLING

The LHCD experiments in TORE SUPRA /2/ were carried out under the following conditions: major and minor radii of plasma are respectively $R = 2.37 \text{ m}$ and $a = 0.77 \text{ m}$ and toroidal magnetic field is $B_t = 1.8$ Teslas.

The LH-system made of 16 klystrons, is able to deliver a total RF power of 8 MW with a duty cycle factor of 0.36 /3/. This power is injected in TORE SUPRA by two independent launchers located in two adjacent ports /4/. In fact the whole LH system is composed of two identical parts as shown on the figure 1: the 8 klystrons of each part are used to supply one antenna. For the experiments described hereunder, only one launcher has been used. This launcher is composed of 16 evacuated modules juxtaposed in such a way that the waveguide network is made of 128 waveguides (4 lines x 32 columns). One passive waveguide on the sides of the launcher is added in order to decrease the reflection coefficient of the adjacent lateral modules /5/. Two superposed modules (4 lines x 4 columns) are fed by a 0.5 MW klystron via a 3 dB hybrid junction terminating the pressurized transmission line. The measured transmission efficiency $\eta$ of the whole system ($\eta = \text{Pinjected}/\text{Pklystron}$) reaches 90%.

Each module owns internal RF splitting using a H-plane hybrid junction and a 4 waveguide E-plane multijunction /6/. The E-plane multijunction, when used in travelling waves (internal phasing $A\phi = \pi/2$), allows furthermore to decrease strongly the power reflection coefficient in the percent range as measured on PETULA /7/ and calculated using the SWAN code /8//9/. The global coupling to the plasma was investigated on a row of 2 modules fed by the same klystron for two cases: (i) with the lowest density plasma ($n_e = 10^{19} \text{ m}^{-3}$) and the launcher mouth set 3 cm behind the limiter, the coefficient of reflection was 6 to 7%. (ii) with the highest density plasma ($n_e = 2.3 \times 10^{19} \text{ m}^{-3}$) and the grill mouth moved at 1.5 cm behind the limiter, the coefficient of reflection was only 1 to 1.5%.

Assuming an e-folding length of the density in the scrape-off layer of 1.5 cm, the density at the mouth of the launcher is estimated to be around $2 \times 10^{17} \text{ m}^{-3}$ (neg ne cut-off) and $10^{18} \text{ m}^{-3}$ (neg - $<N_e>$ cut-off).
respectively in the first and second case. The measured values of the reflection seem consistent with the expected values given by numerical computations /8/. The geometrical periodicity of the launcher $\Delta = 1.05$ cm (internal waveguide width $b = 0.85$ cm, wall thickness $d = 0.2$ cm) and the large number of waveguides (32) along the toroidal magnetic field, allows to excite a narrow, strongly assymetric, $N_{\psi}$-spectrum (shown on figure 2) for current drive studies. The $\langle N_{\psi} \rangle$ peak value of this spectrum can be changed when the phase $\Delta \psi$ between two adjacent modules is properly adjusted, using the phase loop system of each klystron. For $\Delta \psi = 0$ we have $\langle N_{\psi} \rangle \sim 1.9$ and $\Delta N_{\psi} / \langle N_{\psi} \rangle \sim 0.25$. With the configuration shown on figure 1, it is possible to reduce the reflected power towards the klystrons to a level that the tube can withstand without using circulators. On Table 1, we have plotted the results obtained during a discharge, concerning the emitted power by each klystron and the corresponding reflected power towards the same klystron. We see that the V.S.W.R lies between 1 to 1.37 which is less than the nominal operating value of 1.4. The launcher was conditioned with short powerful RF pulses (10 $\mu$s every 100 $\mu$s) injected into the vacuum vessel of TORE SUPRA between two plasma shots. About 800 conditioning pulses were carried out before each 2 s-long RF shots; in less than 30 plasma shots, a RF power of 1.25 MW was injected to the plasma.

3 - L.H.C.D EXPERIMENTS

During these experiments all the modules are in phase ($\Delta \psi = 0$) and, we see from figure 2, that the mean value of the injected $N_{\psi}$-index is 1.9. The plasma current and average density are feedback controlled to be kept constant during the plasma discharge.

Typical time evolution of some main plasma parameters are plotted for two limit hydrogen plasma discharges on figure 3 and 4. In figure 3 low density ($n_e = 10^{18}$ m$^{-3}$) and relatively low current ($I_p = 430$ kA) plasma discharge is shown. During the application of RF power ($P_{RF}$ injected = 630 kW, $\Delta t = 2s$) as loop voltage decrease of 40% together with an enhancement of ECE, bolometric and $[\delta P + l/2]$ signals are observed. Any significant variation in impurity content is measured. In higher density ($n_e \sim 10^{19}$ m$^{-3}$) and higher current ($I_p = 630$ kA) plasma discharge shown on figure 4, same time evolution of plasma parameters are obtained with an injected RF power of 1.25 MW; an increase of the density is observed (until $1.4 \times 10^{19}$ m$^{-3}$). At this density where the characteristic accessible value of the parallel index $N_{\psi||} \leq 2$, the relative loop voltage drop $\Delta V_L / V_L \sim 35\%$ again. During the 1.25 MW, 2 sec long RF pulse, a strong increase of the gas pressure was measured in the tank of the launcher which was not evacuated by the $\text{Li}$ auxiliary pump /4/; the pressure rises up to 10 mPa, twice the baseline pressure and contributes probably to the increase of density which is observed. For these shots, direct electron temperature measurements was not available. But the observed increase on the $[\delta P + l/2]$ and bolometric signals during the RF pulse, can be explained by a electron heating of the plasma together with current drive. This is also supported by the evolution of the electron emission at $\mu$W and of the soft X-ray signals during the RF pulse. Then the determination of current drive efficiency by considering only the loop voltage drop appears difficult.

4 - SUMMARY

With our first multijunction-grill composed of 126 waveguides we have injected up to 1.25 MW of RF power ($P_{RF}$, Pohmic) into TORE SUPRA after less than 30 plasma shots. It was checked that the coupling, not yet optimized, is good enough for safe klystron operation without circulators. The measured value $R_p = P_{RF} / V_L$ obtained on TORE SUPRA ($B_t = 1.8$ T) is closed to one observed on PETULA-B ($B_t = 2.75$ T) at the same frequency and density /10/.
REFERENCES

/7/ - M. GONICHE et al- 13th European Conf. on Controlled Fusion and Plasma Physics-SCHLIERSEE - April 1986 vol I0C part II p. 370
/8/ - M. GONICHE et al- This Conference
/9/ - D. MOREAU and T.K. NGUYEN Report EUR-CEA 1246

Table 1

<table>
<thead>
<tr>
<th>ELEMENTS (J = 1 to 8)</th>
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<th>5</th>
<th>6</th>
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<tr>
<td>Emitted power P(\text{E}) (kW)</td>
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<td>110</td>
<td>110</td>
<td>110</td>
<td>110</td>
<td>110</td>
<td>110</td>
<td>880</td>
</tr>
<tr>
<td>Reflection coefficient R (%)</td>
<td>1.84</td>
<td>0.25</td>
<td>0.30</td>
<td>0.6</td>
<td>2</td>
<td>0.8</td>
<td>2.7</td>
<td>0.25</td>
</tr>
<tr>
<td>Reflection coefficient R (%)</td>
<td>1.84</td>
<td>0.25</td>
<td>0.30</td>
<td>0.6</td>
<td>2</td>
<td>0.8</td>
<td>2.7</td>
<td>0.25</td>
</tr>
<tr>
<td>Power [W]</td>
<td>1.28</td>
<td>1.1</td>
<td>1.16</td>
<td>1.16</td>
<td>1.16</td>
<td>1.16</td>
<td>1.16</td>
<td>1.16</td>
</tr>
</tbody>
</table>

\[ \Delta z = 1.05 \text{ cm}; b = 0.85 \text{ cm}; d = 0.2 \text{ cm} \]

Fig. 2: \( N_f \) spectra of multijunction Tore Supra launcher at various phase difference \( \Delta \Phi \) between modules: \( n_\text{E} = 0.85 \times 10^6 \text{ m} \), \( n_\text{E} = 0.85 \times 10^6 \text{ m} \), \( \Delta \Phi = 1/2 \)

Fig. 1: 3.7 GHz system schematic layout
Fig. 3: Typical time evolution of LHCD discharge: $I_p = 430$ KA and $n_e = 0.7 \times 10^{19}$ m$^{-3}$.

Fig. 4: Typical time evolution of LHCD discharge: $I_p = 630$ KA and $n_e = 1.1 \times 10^{19}$ m$^{-3}$. 
INTRODUCTION

The operation of TORE SUPRA at full power (25MW, 30s) has led to the design of a full set of actively pumped carbon limiters to remove at least 8MW and to partially control the particle balance [1,2]. An interim version is now installed, composed of 5 vertical and one horizontal outboard (OPL) pump limiters, semi-inertially water cooled. The later is a result of a collaboration between the U.S.-DoE and the Association EUR-CEA, it is fully instrumented and therefore can serve as a reference for the final design. Ohmic discharges (1.85T, 740kA, 8.5s) in helium have been used to test the thermal load on and the particle exhaust efficiency of the OPL. In these experiments the plasma is formed on the inner wall (R=232cm, a=76cm) and subsequently displaced (6cm) outward, early on the current plateau, to lean on the OPL (R=238cm, a=75cm). In addition to the limiters above, a non-pumped outboard (ONLP) limiter of identical shape to the OPL served to produce similar discharges for better comparison and determination of particle control. A comparison is made hereafter of the thermal load and particle pumping effects on the OPL when the plasma is in contact either with the OPL/ONPL alone or with the OPL and the vertical limiters together.

Thermal load on the first wall. Calorimetric measurements provide estimates of the time-integrated balance of the energy flow at the plasma edge. Radiated and charge exchange energy losses ($W_R$) are deduced from the inner vessel calorimetry (R=242cm, a=94cm), conductive/convective losses ($W_C$) onto the OPL and vertical (top and bottom) limiters (TBL) are independently measured as well as the integrated energy flow on the ergodic divertors.

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which are located 5 cm in the SOL (R=238 cm, a=80 cm).

Results are shown in table I for a discharge (n=1151) with the plasma on the TBL and the OPL and a discharge (n=1158) with the plasma leaning on the OPL only. Supposing that the radiated and the charge exchange losses are spatially uniform, edge components in the SOL should receive this energy in proportion to their surface area. This is not the case for the ergodic "vertors (n=1151 and n=1158) nor for the vertical limiters (n=1158) when there are 5 cm behind the OPL location. The difference is interpreted as conductive-convective energy onto these components. The radiated energy contribution is observed to decrease (from 29 to 16%) when passing from the multilimiter to the single limiter configuration, while conductive-convective power on the OPL is nearly tripled. In the later case, the mean power flux on the OPL is 2.4MW/m².

The calorimetric results are compared with bolometric measurements of the radiated power as displayed in table II together with Ohmic power estimates. This compares well with the total energy values in table I when considering the 7 sec plateau after the outward displacement of the plasma. Although the bolometric measurements are systematically larger (30-40%) than the inner wall calorimetric measurement, they also indicate the decrease of the radiated power when moving from the multilimiter to the single limiter configuration.

An infrared imaging of the OPL was made when the plasma leans only on it. It shows that the limiter experiences the largest temperature excursions at the tip and at the leading edge. Typical temperature values at the end of the current plateau are 269°C at the tip, 265°C at the leading edge and 224°C in a large area in between. (The base temperature before shots is 180°C). A preliminary comparison with results from heat transport calculations in the limiter tend to indicate that the scrape-off heat length is between 1.0 and 1.5 cm [3].

Particle exhaust. Fig. 1 shows the central line density from interferometric measurement for the two almost successive shots: n=1137 with the ONPL and n=1139 with the OPL. The outward displacement occurs at about 1.2s after the start of the discharge. Before this displacement the densities are identical for these two shots. At the time of the displacement, the density of n=1139 departs from that of n=1137 increasingly during about 1.3 s (scoop effect) and, finally, an almost constant difference remains between the two line densities during the rest of the current plateau (the scoop being saturated). This description applies as well to the other channels of the 5-chord interferometer, indicating that the scoop effect is
experienced throughout the plasma.

A straightforward calculation of the dynamic evolution of the pressure \( p(t) \) in the pump limiter fits well the experimental values as measured near the throat inside the OPL: 
\[
p(t) = \frac{\Gamma_i^+}{C} (1 - e^{-Ct/v}),
\]
where \( \Gamma_i^+ \) is the entrance particle flux parallel to the field in front of the throat, \( C \) is the conductance of the throat (\( C \approx 4.7 \text{m}^3\text{s}^{-1} \)) and \( v \) is the volume of the OPL (\( v = 3.7 \text{m}^3 \)). The particle flux \( \Gamma_i^+ \) is given by 
\[
\Gamma_i^+ \approx n_0 T_0^{1/2} f(\lambda_T),
\]
where \( n_0 \) and \( T_0 \) are the plasma density and temperature at the limiter edge, whereas \( f(\lambda_T) \) is a geometrical expression depending on the particle flux scrape-off length \( \lambda_T \). For typical scrape-off values \( n_0 \approx 5 \times 10^{12} \text{ cm}^{-3} \), \( T_0 \approx 10 \text{ eV} \), \( \lambda_T = 2 \text{cm} \), and for the given OPL geometry, one finds 
\[
p(t) \approx 3 \times 10^{-3} (1 - e^{-1.27t}) \text{ (Torr.s)}.
\]
This expression is plotted together with the experimental pressure in the OPL on the insert of Fig. 1. The small difference on the pressure limit is easily corrected by changing slightly the values assumed for \( n_0 \) and \( T_0 \). After the scooping effect saturates, the almost linear increase of the density with time suggests a recycling coefficient slightly larger than 1.

Preliminary Langmuir probe measurements in the throat of the OPL are consistent with the values supposed for \( n_0 \) and \( T_0 \) in the scrape-off layer. In the ion drift side the ion saturation current leads to 
\[
\Gamma_i^+ \approx 3 \times 10^{20} \text{ s}^{-1} \text{ in the throat. The probe signal on the electron drift side is generally 3–5 times smaller. This is explained by shadowing (far in the SOL) by the ergodic divertor coil nearest to the OPL which is on the electron drift side of the limiter.}
\]

A zero-dimensional balance of the total number of particles in the plasma can be made: Supposing a particle confinement time \( \tau_p \approx 0.1 \text{ s} \), the particle outflux from the plasma is \( N/\tau_p \) (for shot *1137, \( N = 2.5 \times 10^{20} \)) \( : N/\tau_p = 2.5 \times 10^{21} \text{s}^{-1} \)). Using \( \Gamma_i^+ \approx 3 \times 10^{20} \text{ s}^{-1} \), a fraction \( f = \Gamma_i^+ \tau_p / N \) is pumped. In the case of shot *1139, \( f \approx 12\% \).

References.
Acknowledgements. The perseverance of C. Balorin and M. Renaud in providing bolometric and calorimetric data was appreciated as well as the participation of J. Koski and T. Lutz (SNLA) in the experiments. Part of this work was supported under U.S. DoE Contract No. DE-AC05-84OR21400

<table>
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<td>W_{TBL}</td>
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<td>540</td>
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<td>140</td>
</tr>
<tr>
<td>W_{EDC}</td>
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<td>610</td>
</tr>
<tr>
<td>WC</td>
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</tr>
<tr>
<td>WR/WT</td>
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<td>0.16</td>
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**TABLE I** CALORIMETRIC DETERMINATION OF RADIATED WR AND CONDUCTED OR CONVECTED WC ENERGY ONTO THE OUTBOARD (OPL) TOP AND BOTTOM LIMITERS (TBL) ERGODIC DIVERTORS (ED). THE RATIO OF RADIATED TO TOTAL ENERGY (WR/WT) IS GIVEN.

<table>
<thead>
<tr>
<th>P (kW)</th>
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<tbody>
<tr>
<td>P_{H}</td>
<td>1020</td>
<td>970</td>
</tr>
<tr>
<td>P_{B}(t = 2 s)</td>
<td>325</td>
<td>225</td>
</tr>
<tr>
<td>P_{B}(t = 6 s)</td>
<td>375</td>
<td>275</td>
</tr>
<tr>
<td>P_{B}/P_{H}(t = 2 s)</td>
<td>0.32</td>
<td>0.23</td>
</tr>
<tr>
<td>P_{B}/P_{H}(t = 6 s)</td>
<td>0.37</td>
<td>0.28</td>
</tr>
</tbody>
</table>

**Fig. 1** CENTRAL LINE DENSITY AS A FUNCTION OF TIME FOR SHOTS WITH THE PLASMA LEANING ON: #1137 the ONPL #1139 the OPL

IN THE INSERT IS THE EXPERIMENTAL (—) AND ESTIMATED (— — —) PRESSURE IN THE OPL

**TABLE II** OHMIC (P_{H}) AND BOLOMETRIC (P_{B}) POWER EARLY (t = 2 s) AND LATE (t = 6 s) IN THE CURRENT PLATEAU. THE RATIO P_{B}/P_{H} IS INDICATED.
EDGE PLASMA MODIFICATION INDUCED BY A RESONANT PERTURBATION DURING ERGODIC DIVERTOR EXPERIMENTS ON TORE SUPRA


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Introduction

First experiments on the TORE SUPRA tokamak were primarily devoted to edge plasma studies. We report here the ergodic divertor (ED) experiments [1,2,3]. For these experiments the Tore Supra set-up was: toroidal field \( B \sim 1.8 \) T, plasma current \( I_p \sim 0.5-1.2 \) MA, line average density \( n_e \sim 2.-3.10^{19} \) m\(^{-3}\), with Helium gas, and with two configurations.

- "EXT" configuration (\( R=2.37m, a=0.77m \)) The plasma is limited by an outer passive limiter in the mid plane 0.04 m ahead of the E.D. coils.
- "INT" configuration (\( R=2.32m, a=0.76m \)) where the plasma is limited on the inner wall (axisymmetric limiter).

Observation of resonant signatures

The aim of the ergodic divertor perturbation is to create chains of magnetic islands, of half width \( \delta(r) \), on resonant surfaces (for several values of the safety factor \( q(r) \)) with a distance \( \Delta(r) \) between these resonant surfaces. The ergodicity criteria, at radial position \( r \), is given by the Chirikov parameter \( S(r)=2\pi(r)/\Delta(r) \), the threshold being \( S(r) \geq 1 \) [2]. \( S(r) \) depends on the square root of the ED coils current and on the q-profile. Strong values of \( S \) are localized at the plasma edge.

The experimental study is devoted to the response of the plasma to such magnetic perturbations. At plasma current \( I_p \sim 0.85 \) MA, numerous resonant signatures (with a wide variety of effects) are recorded for both configurations. As an example, we examine here a particular signal resulting from a combination of \( \text{H}_\alpha \) and \( \text{He}^+(4-6) \) lines and integrated on a chord which passes relatively close to a bottom limiter. In the EXT configuration, the modification of the recycling pattern induced by the ED perturbation leads to an increase of this \( \text{H}_\alpha \) signal, Fig.1. Although the physics involved in this effect are not straightforward, a comparison can be drawn with the computed resonance on the Chirikov parameter at the...
limiter radius, $S(a)$ [2]. Both calculation and measurements exhibit a resonance at $I_p \approx 0.85$ MA, Fig. 1. No broad range of plasma currents $(0.7 - 1.1$ MA) which yield significant effects on $H_a$ (poor quality factor $I_p/\Delta I_p \sim 2$). The other $H_a$ measurements exhibit complicated (but resonant) modifications of the recycling. In particular several chords aiming at the inner wall indicate an $H_a$ decrease during ED operation in the INT configuration. Up to now, a unique resonance at plasma current $I_p \approx 0.85$ MA, $\psi_{\text{edge}} \sim 3$ has been identified. This resonance is observed on $H_a$ measurements, impurity emissivity, bolometer profiles, soft X rays, and MHD activity, in both configurations.

Effects of the magnetic perturbation on plasma parameters

**EXT configuration.** Edge power deposition integrated during a whole shot at constant perturbation level and plasma conditions are given in Table 1. With respect to non ED shots, an increase of power deposition is found on the neutralizer plates of the ergodic divertor (located between the ED current bars) and on the bottom limiter. A decrease is found on the main outer limiter and on the top limiter. Fig. 2 displays the most relevant signals for a given discharge $(I_p \sim 0.85$ MA). A large reduction (up to 50%) of $CV$ and $CS$ lines is recorded, while a lower reduction (10% to 20%) is observed on OWII and OXII lines. A small increase of the bolometric power and of the soft X ray signal at the center of the discharge is also found. The line average density $n_e$ and the bulk values of $T_e$ and $T_i$ are not affected. An outstanding feature is given by the emissivity shell of peripheral lines, especially $CV$ (ionization potential 392 eV), which are narrowed and shifted inwards [4]. As stated previously the recycling is strongly affected (Fig. 1), however on some chords viewing the inner wall no modification of the $H_a$ signal is observed.

**"INT" configuration.** Recycle studies with fiber optics and one CCD camera show that the $H_a$ signal decreases over the whole limiting inner wall. This is accompanied by a strong increase in the FeXVII metallic impurity emissivity (carbon and oxygen lines are less affected than in the EXT case). The radiated power measured by the 16 chords bolometer increases from 40 to 70% of the ohmic power. Soft X ray signal and Bremsstrahlung signal increase and can be interpreted as a 5% increase of $Z_{\text{eff}}$. An increase of $n_e$ (≈6%) and a slight decrease of $T_e$ (2-3%) are observed.

MHD activity, runaway electrons behaviour (INT and EXT). A strong reduction of MHD activity is induced by the ED perturbation (beyond a given threshold of the latter). The mode $(2,1)$ disappears or is locked whereas the mode $(3,2)$ remains unaffected. Application of the magnetic perturbation leads to a sudden decrease of the confinement time of the remaining runaway electrons.

Variation of the amplitude of the magnetic perturbation. A threshold in the intensity of the perturbation is found for most of the effects. Whether this threshold is the same for all processes has not been determined yet. Fig. 3 displays the variation of the main resonant effects with the ED perturbation amplitude. In the INT configuration, the threshold is given by $\omega \approx \langle \delta B_r \rangle / B_r(a) \approx 0.35 \cdot 10^{-3}$, no saturation of
resonant signatures are observed before the maximum value $\alpha \sim 0.87 \times 10^{-3}$. In the EXT configuration, the threshold value is about the same, $\alpha \sim 0.40 \times 10^{-3}$, and a saturation is found, $\alpha \lambda 1.40 \times 10^{-3}$.

**Discussion and conclusion**

The initial results of the ergodic divertor experiments on Tore Supra are discussed along three main lines:

**Scrape-off layer broadening.** The features of a broadening of the scrape-off layer (SOL) are found in the EXT configuration. The power and particle fluxes are redistributed deeper in the layer thus impinging retracted limiters, for instance the ED neutralizer plates which are located 0.06 m deep in the shadow of the main limiter (see Table 1). The image is different in the INT configuration where recycling occurs on a large axisymmetric limiter and decreases when the ergodic divertor is applied. Yet metallic impurity concentrations increase which could either be related to an increase of the particle confinement time or to a broadening of the scrape-off layer bringing the plasma in contact with metallic objects. This effect is currently under study.

**Edge plasma modification.** The displacement of the X ray lines [4] may be attributed to a change of the edge temperature profile as seen on TEXT [5]. Deconfinement of runaway electrons, (2,1) mode locking or disappearance, suggest that a broad zone at the edge is affected one way or the other by the magnetic perturbation (no definite answer can be given due to the lack of $n_e$, $T_e$, edge measurements).

**Bulk plasma modification.** In the INT configuration density increase and recycling decrease may suggest a particle confinement improvement. At any rate no deconfinement of the bulk plasma is observed in contrast with the TEXTOR ergodic divertor experiment [6].

These preliminary experiments have clearly shown the expected resonant effect induced by the magnetic perturbation. No degradation of the bulk of the discharge is observed. General features of the ergodic divertor discharges are a modification of recycling, a broadening of the scrape-off layer, a modification of the edge plasma and in some cases a possible improvement in the bulk plasma.

**Acknowledgements**

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Mrs BALORIN and Mr RENAUD are thanked for their technical support.

**References**


TABLE 1 - Energy and power deposition on various plasma-facing components

<table>
<thead>
<tr>
<th>Magnetic Perturbation ( a = \delta B/\delta(a) \times 10^3 )</th>
<th>0</th>
<th>0.73</th>
<th>1.05</th>
<th>1.33</th>
</tr>
</thead>
<tbody>
<tr>
<td>E. D. Neutralizer plates ( E ) ( (kJ) )</td>
<td>180</td>
<td>580</td>
<td>710</td>
<td>930</td>
</tr>
<tr>
<td>Top limiter ( (kW \ m^{-2}) )</td>
<td>105</td>
<td>58</td>
<td>74</td>
<td>47</td>
</tr>
<tr>
<td>Outer limiter ( (kW \ m^{-2}) )</td>
<td>92</td>
<td>60</td>
<td>30</td>
<td>53</td>
</tr>
<tr>
<td>Bottom limiter ( (kW \ m^{-2}) )</td>
<td>133</td>
<td>211</td>
<td>204</td>
<td>316</td>
</tr>
</tbody>
</table>

1. Calorimetry throughout the shot.
2. Power flux deduced from thermocouple measurements on two points per module.

Fig. 1: VARIATION OF THE CHIRIKOV PARAMETER AND OF THE RELATIVE INCREASE OF Hα LIGHT NEAR THE BOTTOM LIMITER WHEN A CURRENT 5kA IS PULSED IN THE ERGODIC DIVERTOR COIL.

Fig. 2: EXT CONFIGURATION: VARIATION OF Hα, C XII, O XII (TIME-EVOLUTION DEDUCED FROM SHOT # 1033). \( T_e, T_i, \text{OXX}, n_i, \text{MHD ACTIVITY AND SOFT X-RAYS INTENSITY WHEN CURRENT } I_{ED} \text{ IS PULSED IN THE DIVERTOR COIL.}

Fig. 3: RELATIVE VARIATION OF Hα EMISSIVITY \( I_\text{Fe XIX} \text{, MHD ACTIVITY } \delta \delta \delta_0, \text{Fe XIX LINE EMISSING } I_\text{Fe XIX} \text{ BOLOMETRIC POWER } P_{bolo} \text{ AS A FUNCTION OF THE MAGNETIC PERTURBATION.}
POST DEADLINE PAPER

SOFT X-RAY SPECTROSCOPY DURING TORE SUPRA ERGODIC DIVERTOR EXPERIMENTS

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1. INTRODUCTION

The initial ergodic divertor experiments on Tore Supra have been especially studied with an extreme grazing incidence spectrograph equipped with a new powerful radial scan system allowing the radial emissivity profiles of several soft X-ray lines to be recorded within a single discharge. In this paper, preliminary results on the behaviour of a few intrinsic impurity lines when applying external resonant fields are first reported. Then, the radial emissivity profiles of five light impurity lines are studied as a function of the magnetic perturbation introduced by the ergodic divertor. Perturbation effects on the radial position of the emissivity shells can be interpreted in terms of electron temperature gradient or transport modifications.

2. INSTRUMENTATION

The 2m-grating radius, extreme grazing incidence (1.5°) Schwob-Fraenkel spectrograph was developed at the Racah Institute of Physics (under CEA contract) more than 10 years ago. It has been recently equipped with two microchannelplate (MCP) detectors, both movable along the Rowland circle. Each detector consists of a MgF2 coated, funneled MCP, associated with a phosphor screen image intensifier and coupled by a flexible fiber optic conduit to a 1024 element photodiode array. Two spectral regions of roughly 40 Å width are simultaneously accessible in the range 10 - 330 Å. The spectral resolution (typically 200 mA) is given by a 20 μm entrance slit and a 600 groove/mm Jobin-Yvon holographic grating /1/.

The original feature of the instrument installed on Tore Supra lies in a new fast radial scan system. The spectrograph is repetitively moved during a plasma discharge, so that its line of sight scans the plasma radius (figure 1). The sinusoidal movement, generated by an hydraulic jack, is adjusted by a microprocessor. The scanning period can be as low as 1 second (with a plasma scanning duration of 600 ms), and the maximum spatial resolution along the plasma radius is close to 5 cm (limited by the spectrograph aperture and the integration time).

The spectrograph may be operated either in this new 'scanning mode', or in the classical 'static mode' with a chosen line of sight.
3. DETECTION OF A RESONANCE

In a first set of experiments the spectrograph was used in the static mode, integrating the line emission along the entire plasma diameter. In the discharge no 1051 the intensity of the ergodic divertor current was left constant at 1.6 kA, and the plasma current was varied following figure 2. The integrated emission of the FeXV 284.15 Å line is significantly increased for a plasma current of 840 kA, indicating the existence of a resonance for this value /2/.

Figure 2 : Time evolution of the FeXV line intensity giving evidence for a resonance. The plasma current is shown by the dotted line.

The plasma response to the magnetic perturbation has then been studied for this resonant value of the plasma current through the integrated emission of a few typical lines (CV 40.27 Å, CVI 53.74 Å, OVII 21.6 Å, OVIII...
18.97 Å, FeXV 284.15 Å, NiXVIII 292.00 Å). Two plasma configurations were
used: an external configuration where the plasma is leaning on the outer
limiter 4 cm ahead of the ergodic divertor coils, and an internal
configuration with the plasma leaning on the inner first wall. Preliminary
results show different behaviours in the two situations: in the external
configuration the intensity of the two carbon lines clearly decreases when
the ergodic divertor is applied. The intensity of the oxygen lines and the
intensity of the metallic lines remain roughly constant. In the internal
configuration a smaller effect is noticed on the carbon lines, but the FeXV
and NiXVIII line intensities increase strongly.

4. RADIAL EMISSIVITY PROFILES

The scanning mode of the spectrograph has been successfully
operated to study the modifications of the impurities emissivity profiles
with the magnetic perturbation. The scanning period of the spectrograph was
2 s (1.2 s plasma scanning duration) and the spatial resolution was close to
7 cm (100 ms integration time).

Inverted emissivity profiles of two carbon lines (CVI 33.74 Å,
CV 40.27 Å) are presented on figure 3. Inversion of the line integrated
data, taking into account the magnetic geometry [3], is obtained using least
squares minimisation techniques [4]. For each line a comparison is made
among four plasma discharges with various ergodic divertor currents (Ied =
0, 0.8, 1.5, 1.85 kA).

![Figure 3](image-url)

Figure 3 : Inverted radial emissivity profiles of the CVI (a)
and CV (b) resonance lines at 33.74 Å and 40.27 Å. The ionization potential
of each ion is indicated on the figure.

In these experiments the external plasma configuration was used. The
radial position of the CVI maximum emissivity is not modified by the
magnetic perturbation, whereas the CV maximum is clearly shifted inwards, with a corresponding narrowing of its profile. The radial emissivity profiles of three oxygen lines (OVIII 18.97 Å, OVI 173.00 Å, OV 192.80 Å) have also been recorded during the same discharges. No effect is seen for the OVIII line, but the OVI and OV profiles are slightly shifted inwards, with little shape modification. These observations are summarized in figure 4, where we have plotted the ionization potential of the different impurity ions as a function of their radial position of maximum emissivity. These data can be interpreted in terms of a temperature profile modification only if the radial impurity transport at the plasma periphery has not been modified by the ergodic divertor effects.

![Figure 4](image)

Figure 4: Ionization potential of the 4 studied impurity ions versus radial position of the ion maximum emissivity, with and without magnetic perturbation.

5. CONCLUSION

A preliminary study of soft X-ray impurity emission during ergodic divertor experiments has been undertaken. The expected resonant effect has clearly been detected. The new system of radial scan associated to the Schwob-Praenkel spectrograph has been successfully operated, and the radial emissivity profiles recorded have been carefully analyzed. The few available experimental data do not allow to choose between impurity transport and electron temperature gradient changes at the plasma edge as the cause of impurity ion emissivity profile modifications.

REFERENCES

/2/ A. GROSMAN et al., this conference.