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NON LINEAR SELF CONSISTENCY OF MICROTEARING MODES.

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ABSTRACT- The self consistency of a microtearing turbulence is studied in non linear regimes where the ergodicity of the flux lines determines the electron response . The current which sustains the magnetic perturbation via the Ampère law results from the combined action of the radial electric field in the frame where the island chains are static and of the thermal electron diamagnetism . Numerical calculations show that at usual values of β_{pol} in Tokamaks the turbulence can create a diffusion coefficient of order $\nu_{th}\rho_i^2$ where ρ_i is the ion larmor radius and ν_{th} the electron ion collision frequency . On the other hand , collisionless regimes involving special profiles of each mode near the resonant surface seem possible .

I Introduction

Recently, a serious degradation of confinement has been observed on most tokamaks during additional heating. The microtearing modes / 1,2,3,4,5/ could provide a possible explanation for this experimental fact. As the large scale tearing modes due to $\frac{dI}{dr}$, they mainly consist of magnetic perturbations

$$\delta B_r = \frac{1}{r}\frac{\partial}{\partial\theta}\,\delta A \;.\; \delta B_\theta = -\frac{\partial}{\partial r}\,\delta A$$

derived from a parallel potential vector

$$\delta A = A_l(r)\exp i(l\theta + m\varphi - \omega t) + cc \qquad (1)$$

which does not cancel on the magnetic surface $r=r_l$ where $m+\frac{1}{q(r_l)}=0$. They are thus able to destroy the magnetic surfaces and thereby to degrade the confinement of circulating electrons. They differ from the $\frac{dI}{dr}$ modes since they are driven at much larger l, m by the electron thermal energy rather than by the poloidal magnetic energy. The parallel current density $\delta I = I_l(r)\exp i(l\theta + m\varphi - \omega t) + cc$ which sustains the magnetic perturbation δB is induced very close to the resonant surface by diamagnetic effects.

A crucial question is to know whether that current is able to sustain the mode through the Ampère law

$$\mu_0 I_l = -\frac{\partial^2 A_l}{\partial r^2} + \left(\frac{l}{r}\right)^2 A_l \qquad (2)$$

from which one can derive

$$\frac{1}{\mu_0}\int \left(\left|\frac{\partial A_l}{\partial r}\right|^2 + \left(\frac{l}{r}\right)^2 |A_l|^2 \right) dr = \int I_l A_l^* \, dr \qquad (2a)$$

$$\frac{1}{\mu_0}\left|\frac{l}{r}\right|^2 \int A_l \cdot dr = \int I_l \, dr \qquad (2b)$$

In the linear case one finds that those equations can be satisfied for unstable modes in collisional regimes only. In fact the linear analysis is correct only when the magnetic island topology does not influence the response of electrons to the electromagnetic perturbation. The goal of the present article is to investigate the stability of the microtearing modes in the realistic regimes where one must take account of such non linear effects to derive the electron current density δI. We compute that current in the presence of a set of modes (1) producing overlapping magnetic islands. Cancelling the current component in quadrature with each $A_l(r)$ determines the velocity $\frac{\omega}{l/r}$ of the island chains in the plasma frame where no average radial electric field exists, or equivalently the average radial electric field $-\frac{\overline{\partial U}}{\partial r} = \frac{\omega}{l/r} B_0$ in the frame $\mathfrak{R}$ where the island chains are static. It is then possible to calculate the current component $I_l(r)$ in phase with $A_l(r)$, which appears as a combined effect of this radial electric field and the thermal diamagnetism. We find that self sustaining microtearing turbulence is possible in presence of collisions, but the computations also suggest collisionless regimes with attenuated values of $A_l(r)$ near the resonance surfaces $r=r_l$. The threshold values of β_{pol} for those process are estimated.

II Basic equations

To simplify the computations, we assume that there exists a frame of reference $\mathfrak{R}$ (rotating around the major axis), where the magnetic configuration is static. This is the case for a set of perturbations (1) with the same $m = m_0 \cdot \omega$. In the frame $\mathfrak{R}$, the magnetic pertubation is

$$\delta A = \exp i m_0 \varphi \sum_l A_l(r) \exp i l \theta + cc \quad (3)$$

There exists also a static electric field $-\vec{\nabla}(U(r,\theta,\varphi))$, that we will restrict to its average radial component $-\frac{\overline{\partial U}}{\partial r}$ in the considered layer. The role of the fluctuation $-\vec{\nabla}(\delta U(r,\theta,\varphi))$ superimposing to $-\frac{\overline{\partial U}}{\partial r}$, which in fact prevents microtearing perturbations except at small radial scales, will not be discussed in this note. Each component l in (3) creates an island structure near the resonant surface $r=r_l$. The magnetic surfaces are destroyed and the field lines wander stochasticly if the islands overlap. This is achieved if the island half width $\delta_I = \left|\frac{8A_l L_s}{B_0}\right|^{1/2}$ is larger than $\delta_r/2$, where $\delta_r = \frac{1}{qR}\frac{L_s}{1/r}$ is the distance between two neighbouring resonant surfaces r_l, r_{l+1}. We assume Chirikov parameters $\frac{2\,\delta_I}{\delta_r}$ of order unity and $\frac{1}{r}\frac{\delta_I}{} \ll 1$. The correlation length of any perturbation along the flux lines then scales with respect to the shear length L_s as $\mathcal{L} = \frac{L_s}{\frac{1}{r}\delta_I}$ /6/. This means that the magnetic islands exclude values of $K_{/\!/} \simeq 1/\mathcal{L}$ smaller than $\frac{1}{r}\frac{x}{L_s}$ with $x = \delta_I$. We will assume that only the electrons contribute to the parallel current δI. That current is derived from the Fokker-Planck equation for electrons written in the phase space $\left(r,\theta,\varphi,v_\perp^2,v_{/\!/}\right)$

$$v_{/\!/}\nabla_{/\!/}f + v_{E\theta}\cdot\frac{\partial f}{r\partial\theta} - \frac{e}{m}\nabla_{/\!/}U\frac{\partial f}{\partial v_{/\!/}} = C(f) + D(v^2)\Delta f + S(r,v^2)\;(4)$$

$$\delta I(r,\theta,\varphi) = \int f\left(r,\theta,\varphi,v_\perp^2,v_{/\!/}\right)\, e v_{/\!/} d_3 v$$

where $v_{E\theta} = \frac{\partial U}{B_0\partial r}$ is the electric drift velocity along θ, $\nabla_{/\!/}$ isthe derivative along the actual field lines. S is a source. C(f) is the $v_{/\!/}$ collision operator and $D\Delta f$ is a small space diffusion operator . These operators prevent the singularities due to the stochasticity of field lines and guarantee that the distribution function is not far from a Maxwellian varying smoothly in r. The nonlinear effects associated with the islands topology will play an important role if the transverse diffusion coefficient D verifies

$$D\frac{\mathcal{L}}{v_{/\!/}} \ll \delta_I^2 \qquad (5a)$$

Also the drift velocity $v_{E\theta}$ must not extract an electron from an island structure before a parallel transit time $\frac{\mathcal{L}}{v_{\parallel}}$ .This means that

$$\frac{1}{r}\frac{\partial U}{B_0 \partial r} \ll \frac{v_{\parallel}}{\mathcal{L}} \qquad (5b)$$

Anticipating that $\frac{e}{T}\frac{\partial U}{\partial r} \sim \frac{\partial n}{n \partial r}, \frac{\partial T}{T \partial r} \sim \frac{1}{L_n}$ .this is equivalent to the condition $\frac{L_s}{L_n}\rho_e \ll \delta_I$. If it is fulfilled , the electron response may be developped in powers of the electric drift velocity $v_{E\theta}$. We are then led to handle the system ( ) in two steps,namely to write $f=f_0+f_1$, where f_0 is the solution of the equation (4) without the term $v_{E\theta}\cdot\frac{\partial f}{r\partial\theta}$

$$T(f_0) = v_{\parallel}\nabla_{\parallel} f_0 - \frac{e}{m}\nabla_{\parallel} U \frac{\partial f_0}{\partial v_{\parallel}} - D\Delta f_0 - C(f_0) = S \quad (6)$$

and then f_1 is given by

$$T(f_1) = - v_{E\theta}\cdot\frac{\partial f_0}{r\partial\theta} \qquad (7)$$

Let us notice that in the linear case , the condition (5b) would imply $\omega \ll K_{\parallel} v_{\parallel}$ and is not fulfilled in the current sheet . The expansion (6),(7) relies on the fact that the magnetic islands produce $K_{\parallel}$ which may be everywhere larger than $\omega/v_{\parallel}$.

One may show that the current δI_0 associated with f_0 is in quadrature with δA, i.e., it produces imaginary $I_{0l}(r)A_l^*(r)$.Similarly, the current δI_1 relying on the distribution f_1 given by (7) is in phase with δA. The self consistency imposes that δI_0 cancels and that the $I_l(r)$ generated by δI_1 verify (2). A small error on f_0 should not influence the currents δI_0 and δI_1. As f is not far from a Maxwellian, we may therefore replace in (6) the term $-\frac{e}{m}\nabla_{\parallel} U\frac{\partial f_0}{\partial v_{\parallel}}$ by $\frac{e}{T}v_{\parallel}\nabla_{\parallel} U f_0$ .Similarly , the term $-\frac{e}{m}\nabla_{\parallel} U\frac{\partial f_1}{\partial v_{\parallel}}$ may be dropped in equation (7). One may still derive from (5),(7) the very important following statement: in a collisionless regime, if one makes the traditional assumption of a constant $A_t(r)$ over the current profile, the integrated correlation $\int I_t(r)A_t^*(r)dr$ is purely imaginary, so that (2) cannot be satisfied. A microtearing turbulence is only possiblein the collisional or in the non constant $A_t(r)$ cases.

III Computation scheme

The phase space is divided in small volumes d_3v_k centred around $(v_{\perp k}, \pm v_{\parallel k})$, $v_{\parallel k} > 0$, each couple defining a group of particles labelled by k. For each k, we define the quantities n_k and $\bar{n}_k e\, v_{\parallel k} U_k$, respectively the density and the parallel current density of the group k, $\bar{n}_k$ being the space averaged value of n_k over the layer. We use a Crookes collision operator based on a proper collisional frequency for the ν_k for the group k. The equations (6) and (7) then determine n_{ok}, U_{ok} and n_{1k}, U_{1k} associated with f_o and f_1 for each group k.

III - 1) Equilibrium gradients

Putting

$$N_{ok} = \frac{n_{ok}}{\bar{n}_k} + \frac{e}{T} U(r)$$

the equation (6) provides

$$\begin{cases} \nabla_{\parallel} U_{ok} = d_k \Delta N_{ok} \quad s_k d_k S_o \\ \nabla_{\parallel} N_{ok} = d_k (\Delta - b_k) U_{ok} \end{cases} \qquad (8)$$

where

$$d_k = \frac{D_k}{v_{\parallel k}}, \quad b_k = \frac{\nu_k}{D_k}, \quad s_k d_k S_o = \frac{2\, S\left(r, v^2\right) d_3 v_k}{\bar{n}_k \quad v_{\parallel k}}$$

and $S_o(r)$ reflects a normalized source and sink localized at the edges of the studied layer. We dispose of a numerical code performing the inversion of (8) through a Fourier analysis along (r,θ,φ) from which we derive

$$N_{ok} = s_k\ \mathcal{F}_o\left(\vec{x}, b_k, d_k\right) \qquad (9)$$

$$U_{ok} = s_k\ \mathcal{G}_o\left(\vec{x}, b_k, d_k\right)$$

the current density calculated at this step,

$$\delta I_o = e \sum_k n_k\ v_{\parallel k}\ s_k\ \mathcal{G}_o\left(\vec{x}, b_k, d_k\right)$$

is in quadrature with δA and must be cancelled in view of the Ampere equation (2). By differentiating (9) along r and averaging across the layer, the source coefficients s_k may be linked to the averaged gradients $\frac{\partial n}{n\partial r}, \frac{\partial T}{T\partial r}, \frac{e}{T}\frac{\partial U}{\partial r}$. The collisionless regime corresponds to the situation where $\nu_k\ \mathcal{L}/v_{\parallel k} = b_k d_k \mathcal{L} \ll 1$. It appears that the functions

$\frac{\mathcal{F}_o\left(\vec{x}, b_k, d_k\right)}{d_k}$ and $\frac{\mathcal{G}_o\left(\vec{x}, b_k, d_k\right)}{d_k}$ (for a given magnetic configuration and δ_o) are then independent of b_k, d_k over all transverse scales greater than the fine structure scale $\delta = (d_k\mathcal{L})^{1/2}$ assumed $\ll \delta_I$ in (5a). Cancellation of δI_o is then insured under the simple quasilinear condition

$$-\frac{e}{T}\frac{\partial U}{\partial r} = \frac{1}{n}\frac{\partial n}{\partial r} + \alpha\frac{1}{T}\frac{\partial T}{\partial r} \qquad (10)$$

with $\alpha = \frac{1}{2}$. A similar calculation gives $\alpha = \frac{5}{2}$ in the collisional case $\nu_k\mathcal{L}/v_{\parallel k} = b_k d_k \mathcal{L} \gg 1$. One may stress that the Ampère law, by imposing that δI_o cancels, imposes also the cancellation of the average radial electron flux. Indeed, the latter is proportional to the correlation $\delta I . \delta B_r = \delta I . \frac{\partial \delta A}{r\partial\theta} = \delta I_o . \frac{\partial \delta A}{r\partial\theta}$, since δI_o is the component of the parallel current δI in phase with δA. This situation is achieved by adjustment of $\frac{\partial U}{\partial r}$: the Eq.(15) expresses that the electrons tend to appear in thermodynamical equilibrium in the frame of reference where the islands are static. Equivalently, in the plasma frame, the islands rotate at an electron diamagnetic velocity.

<u>II -2) Parallel current in phase with δA</u>

The density n_{1k} and the parallel flux

$\bar{n}_k U_{1k} v_{\parallel k}$ for the group k, associated with f_1 verifying(7), are still given by equations of type (8) except that the source term in the continuity equation is now $S_p = -\dfrac{v_{E\theta}}{v_{\parallel}k}\dfrac{1}{r}\dfrac{\partial N_{ok}}{\partial\theta}$ and a source of parallel

momentum $S_M = -\dfrac{v_{E\theta}}{v_{\parallel k}}\dfrac{1}{r}\dfrac{\partial U_{ok}}{\partial\theta}$ appears in the dynamical

equation. $V_{E\theta}$ being now determined by (10). We obtain

$$U_{1k} = \frac{v_{E\theta}}{d_k v_{\parallel k}} s_k\ \mathcal{G}_1\left(\vec{x}, b_k, d_k\right)$$

with $\mathcal{G}_1 = \mathcal{G}_1' + \mathcal{G}_1''$,the functions $\mathcal{G}_1'$ and $\mathcal{G}_1''$ being induced by the particle source S_p and the momentum source S_M, respectively.The flux U_{1k} is responsible for the current δI_1 in phase with δA. The numerical code provides, at each parallel velocity $v_{\parallel k}$, the coefficients

$$K_l'(b_k, d_k) = \int dr a_l^*(r)\frac{\mathcal{G}_{1l}'(r, b_k, d_k)}{d_k\dfrac{\partial\mathcal{F}_0\left(\vec{x}, b_k, d_k\right)}{\partial r}}, K_l''(b_k, d_k) = \ldots.$$

where

$$a_l = \frac{A_l}{B_o},\ \mathcal{G}_{1l}' = \int \frac{d\theta}{2\pi}\frac{d\varphi}{2\pi}\mathcal{G}_1' \exp-i(l\theta+m\varphi), \mathcal{G}_{1l}'' = \ldots$$

Using the notations:

L_n = density gradient length

L_s = shear length

$$\eta = \frac{d\log T}{d\log n}$$

$$\beta_p = 2\eta(2+\eta)\,\mu_0\frac{nT}{B_o^2}\frac{L_s^2}{L_n^2}$$

and putting

$$H_l = \sum_k \frac{\bar{n}_k}{n}\left(4 - \frac{m}{T}v_k^2\right)\left(K_l'(b_k, d_k)+K_l''(b_k, d_k)\right) \qquad (11)$$

the Ampère constraint (2a) may be expressed for each mode l as

$$\beta_p = \frac{8\int dr\left(\left|\frac{\partial a_l}{\partial r}\right|^2 + \left(\frac{1}{r}\right)^2 |a_l|^2\right)L_s^2}{H_l} \qquad (12a)$$

On the other hand , one deduces from (3b)

$$\beta_p = \frac{8\int -\left(\frac{1}{r}\right)^2 a_l\, L_s^2}{H_l} \qquad (12b)$$

if $K_l^{'}$ and $K_l^{''}$ are built up from $\int \mathcal{G}_{1l}^{'}dr$ and $\int \mathcal{G}_{1l}^{''}dr$ rather than from $\int a_l^*\mathcal{G}_{1l}^{'}dr$ and $\int a_l^*\mathcal{G}_{1l}^{''}dr$.

IV Results

As explained in section II - b , in the constant $A_l(r)$, collisionless regimes.the r-integrated current carried by each group k , proportional to $\mathcal{G}_1^{'} + \mathcal{G}_1^{''}$, vanishes. This means that the integrals $K_l^{'}$ and $K_l^{''}$ cancel one another in H_l . This situation , numerically verified , implies that (12) is impossible . It is instructive to first consider the case of a single island chain.The system (5) then simply reflects that the particles pass round the bulk of the islands through the X points. This providesa relatively smooth source of particles $S_p \propto \frac{1}{r}\frac{\partial N_{ok}}{\partial\theta}$, which, by short circuiting along the flux lines, produces a current $\propto \mathcal{G}_1^{'}$ decreasing smoothly outside the islands. On the other hand, the source of momentum $S_M \propto \frac{1}{r}\frac{\partial U_{ok}}{\partial\theta}$ is localized very close to the separatrix. Acting as a ponderomotive force, it produces a current $\propto \mathcal{G}_1^{''}$ which drops abruptly outside the islands.Without collisions, the functions $\mathcal{G}_1$ and $\mathcal{G}_1^{''}$ are maintained within the islands by electron viscosity (proportional to the diffusion coefficient D) at a constant value determined by their profile just outside. They then exhibit opposite $K_l^{'}$ and

K_l''. The value of K_l' is positive and plays a destabilizing role in (12), while K_l'' is negative, stabilizing. In collisional regime, the parallel friction effect, now competing with electron viscosity, reduces the function $\mathcal{G}_1''$, more localized in the islands, at a larger extent than $\mathcal{G}_1'$, resulting in positive destabilizing $K_l'+K_l''$. On the other hand, a simple geometrical reduction of the island width obtained by attenuating $A_l(r)$ within the chain has a similar effect.

A similar behaviour is found in the general stochastic case where several overlapping island chains are present. Again, the destabilizing current $\propto \mathcal{G}_1'$ is less localized than the stabilizing current $\propto \mathcal{G}_1''$ near the resonant surfaces and is less affected by parallel friction. The computations are done for a given geometry characterized by a weak value of $\frac{1}{r}\,\delta_I \sim 0.1$. various ratios $\frac{2\delta_I}{\delta_r}$ of order unity and a shear length L_s generating a longitudinal correlation scale $\mathcal{L} = \frac{L_s}{(1/r)\,\delta_I}$. It appears that the quantity $\frac{H_l\delta_I}{|a_l|^2 L_s^2}$ which enters in (1) depends on the geometrical parameters L_s, $\delta_I \sim \delta_r \ll \frac{r}{1}$, mainly through $\frac{2\delta_I}{\delta_r}$. For collisional regimes $\frac{\nu_k\mathcal{L}}{v_{\neq k}} = b_k d_k \mathcal{L} \gg 1$, the function $\mathcal{G}_1'$ behaves as $d_k^3 b_k$ and $\mathcal{G}_1''$ as $\frac{d_k}{b_k}$, so that the ratio $\frac{\mathcal{G}_1''}{\mathcal{G}_1'}$ varies as $\frac{1}{d_k^2 b_k^2} \sim \left(\frac{v_{\neq k}}{\nu_k}\right)^2$, decreasing abruptly with collisionality. The summation over k in (11) may therefore be restricted to values of k for which the contribution associated to $\mathcal{G}_1''$ is less than half the one associated to δ_1 . This procedure produces the threshold β_p given on Fig.1 for various $\frac{2\delta_I}{\delta_r}$, in function of the collisionality of the electron assembly .The latter is represented by the parameter $\frac{\delta_I^2 \nu_{th}}{D_{act}}$ where ν_{th} is the collision rate for electrons at the velocity $v_{th} = \left(\frac{2T}{m}\right)^{1/2}$ and $D_{act} = D + D_{erg}$ is the computed collisionless maxwellian averaged diffusion coefficient of electrons in the considered layer.This parameter scales as $\frac{\nu_{th}\mathcal{L}}{v_{th}}$ in the stochastic cases $\frac{2\delta_I}{\delta_r} \neq 0$, where $D_{act} \simeq D_{erg}$. In fact the corresponding curves on Fig.1 are comparable to the curve for the non stochastic case $\frac{2\delta_I}{\delta_r} = 0$ obtained with a single island chain and $D_{act} = D$ .One will note that the Fig.1 allows to estimate,for a given β_p,the diffusion coefficient D_{act} normalized to $\nu_{th}\,\delta_I^2$ which can be produced by collisional microtearing modes. A further study shows that the scale δ_I is limited to small values of the order of the ion Larmor radius ρ_i by the effect of the potential fluctuations.For realistic β_p values of order unity,the Fig.1 then implies values of D_{act} of order $\nu_{th}\delta_I^2$ scaling as $\nu_{th}\rho_i^2$.

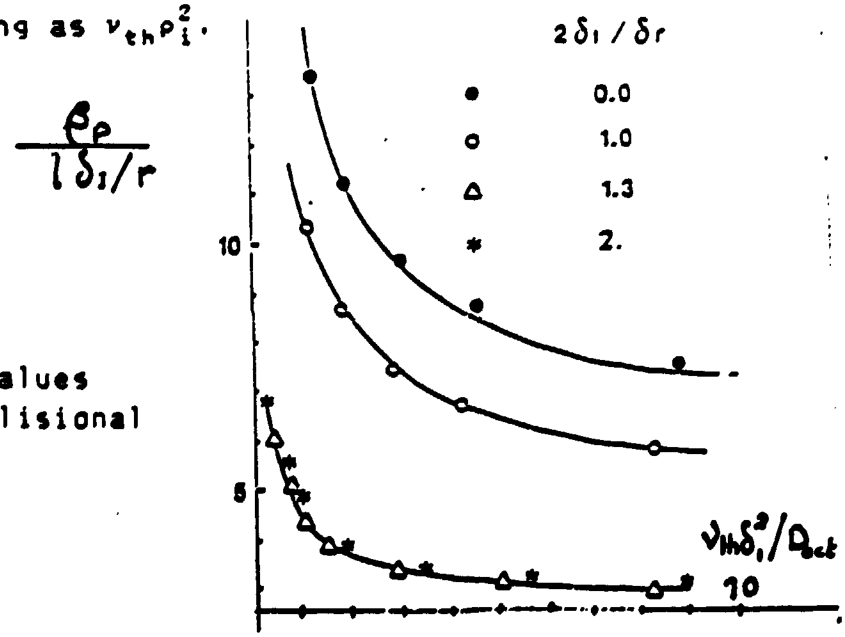

Fig 1:
Threshold values of β_p in collisional regimes

Non constant $A_l(r)$, collisionless regimes have been studied by attenuating the profiles $A_l(r)$ near the resonance surface $r = r_l$ by a coefficient

$$\left(1 - \frac{2}{3} \exp - \left(C \frac{r-r_l}{\delta_I}\right)^2\right)$$

. The function $\mathcal{G}''_{1l}$ is more localized near the surface $r=r_l$ than $\mathcal{G}'_{1l}$, so that the total parallel current $I_l(r)$ due to $\mathcal{G}'_{1l} + \mathcal{G}''_{1l}$ has the same sign as $-\frac{\partial^2 A_l}{\partial r^2}$ on the surface $r = r_l$ and at some distance from that surface . It is in fact possible to adjust the constant C between 0.6 and 1 to fit the shape of the profiles of $I_l(r)$ and $-\frac{\partial^2 A_l}{\partial r^2} + \left(\frac{l}{r}\right)^2 A_l$ in the Ampere law . The condition (12a) then gives the threshold beyond which the imposed variations $\frac{\partial A_l}{\partial r}$ near the resonant surface can be created by $I_l(r)$.On the other hand,the condition (12b) gives the threshold $\frac{\beta_p}{\frac{1}{r}\delta_I}$ above which the r-integrated value of $I_l(r)$ can reconnect the vacuum values of $A_l(r)$ on each side of the resonant surface.This latter condition can be satisfied only for $\frac{2\delta_I}{\delta_r} \leq 1.5$.For larger $\frac{2\delta_I}{\delta_r}$, the function $\mathcal{G}''_{1l}$ is still more localized near the resonant surface than $\mathcal{G}'_{1l}$, but exhibits a large stabilizing peak value on that surface , preventing to satisfy (12b). The obtained results are given on the following table .

Table 1. Threshold values of $\beta_p = 2\eta(2+\eta)\mu_o \dfrac{nT}{B_o^2} \dfrac{L_s^2}{L_n^2}$ in collisionless regimes .

$\dfrac{2\delta_I}{\delta_r}$ =	0	1.	1.3	2.
β_p deduced from (13a) =	3.	3.	2.	5.
$\dfrac{\beta_p}{\frac{1}{r}\delta_I}$ deduced from (13b) =	12.	13.	15.	stable

Conclusion

The self-consistency of a non-linear microtearing turbulence due to diamagnetic effects relies on the presence of the average radial electric field confining hot electrons in the frame of reference where the overlapping island chains are static , and more specifically on the associated drift velocity $v_{E\theta}$. If the effect of this velocity is ignored,the electron subassemblies at various energies exhibit a distorted particle distribution N_o and a parallel momentum distribution U_o,but in fact contribute by no parallel current in the Ampère law.However the electric drift $v_{E\theta}$ produces a local particle source $-v_{E\theta}\dfrac{\partial N_o}{r\partial\theta}$ which by short-circuiting along the ergodized flux lines induces a destabilizing current I'. Also the parallel momentum source $-v_{E\theta}\dfrac{\partial U_o}{r\partial\theta}$ produces a stabilizing current I'' determined by electron viscosity competing with friction effects.The stabilizing current I'' is selectively reduced by friction effects,or,at least at weak ergodization,by an attenuation of the Fourier components $A_l(r)$ near the resonant surfaces.The total current $I'+I''$ can then satisfy the Ampère law at threshold values of β_{pol} given on Fig.6 and on Table 1.It must be stressed that these thresholds simply mean that the parallel current $I^{+} + I^{-}$ can sustain existing modes, leaving open the problem of their onset.

The above mechanism would apply as well if the electric drift $v_{E\theta}$ was replaced by a guiding center velocity $v_{g\theta}$ due to a destabilizing curvature $\frac{1}{\mathcal{R}}$ of the flux lines. The thresholds β_{pol} of Fig.6 and Table 1, increased by the ratio $\frac{v_{E\theta}}{v_{g\theta}} \sim \frac{\mathcal{R}}{r}$, would then characterize situations where the curvature can maintain overlapping magnetic islands. An important difference is that the curvature drift would apply whatever the transverse scale of the turbulence. On the other hand, the drift electric velocity $v_{E\theta}$ is only effective if its action is different for ions and electrons, i.e. for very small island widths, typically of order of the ion Larmor radius ρ_i. This constraint limits the diffusion coefficients that a collisional microtearing turbulence can achieve to values scaling as the electron collision rate $\times\ \rho_i^2$. In modern tokamaks, such values can be significant at the plasma edge only. If the confinement in the core is limited by a microtearing turbulence, collisionless regimes associated with $A_\parallel(r)$ variations should play a role. Further investigations of such regimes are in progress.

References

/1/ DRAKE J.F., LEE Y.C. (1977) Phys. Fluids 20, 1341.

/2/ DRAKE J.F., GLADD N.J., LIU C.S., CHANG C.L. (1980) Phys. Rev. Lett. 44, 994.

/3/ FARENGO R., LEE Y.C., GUZDAR P.N. (1983) Phys. Fluids 26, 3515.

/4/ HAZELTINE R.D. DEBROTT D., WANG T.S (1975) Phys. Fluids 18, 1778).

/5/ MAHAJAN S.M., HAZELTINE R.D., STRAUSS H.R., ROSS D.W. (1979) Phys. Fluids 22, 2147.

/6/ RECHESTER A.B., ROSENBLUTH M.N., White R.B. (1979) Phys. Rev. Lett. 42, 1247.