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* *See Appendix 1*

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THERMAL EFFECTS IN DRIFT-TEARING MODES

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ABSTRACT - The stability of drift-tearing modes in the collisional and semi-collisional regimes is studied analytically, when thermal effects are included. We use a fluid description of a cylindrical plasma and assume a variable plasma resistivity, depending on the local power input. The linear growth rates for the thermally modified tearing modes ($\Delta' > 0$) and for the thermal instabilities ($\Delta' < 0$) are derived. The saturation mechanisms for these instabilities are also discussed.

1. INTRODUCTION

Tearing instabilities are commonly considered as key processes in tokamak discharges and, in recent years, the influence of thermal effects in the growth and saturation of such instabilities has received increasing attention (Sykes and Wesson, 1980; Steinolfson, 1983; Rebut and Hugon, 1984; Scott and Hassam, 1987; Kim et al, 1988).

On the other hand, the kinetic theory was able to show that tearing modes with large poloidal numbers, $m \gg 1$, also called microtearings, could eventually be destabilised by temperature gradients (Haseltine et al, 1975; Drake and Lee, 1977). This would have an important impact on the magnetic confinement, leading to ergodisation of magnetic field lines and enhanced thermal transport (Drake et al, 1980). However, the kinetic theory seems to have overestimated the growth rate (Hassam, 1980) and it is not clear that in an actual tokamak discharge, where toroidal effects play an additional stabilising effect (Finn et al, 1983; Hahm and Chen, 1986) these modes will be unstable.

In the present work we consider the influence of thermal effects on drift-tearing and microtearing modes, using a hydrodynamical approach and treating the plasma resistivity as an additional variable which depends on the amount of power locally deposited in the plasma. We will study both collisional and semi-collisional regimes, in a cylindrical geometry. In particular, we will show that the variation in power deposition can supply an alternative mechanism for driving unstable microtearing modes. In some sense we generalise here the work of Steinolfson (1983), where the low m tearing modes were considered in a slab geometry and where the diamagnetic drifts were disregarded, and the work of Rebut and Hugon (1984) where the steady state of large thermal islands were considered.

In Section 2, we state the basic starting equations, which are used in Section 3 to obtain a general analytical expression for the linear dispersion relation, in the collisional regime. The case of a semi-collisional plasma, which is perhaps more relevant to present-day high temperature tokamaks, will be discussed in Section 4. The saturation processes leading to steady state magnetic islands will be discussed in Section 5 and the conclusions are stated in Section 6.

2. BASIC EQUATIONS

Let us consider a tokamak plasma, in the cylindrical approximation, with radius a and axial periodicity length $L = 2\pi R$. It is well known that any helical perturbation in the magnetic field $\vec{B}$ and in the plasma velocity $\vec{v}$ can be described by a flux function ψ and a stream function ϕ, such that:

$$\vec{B} = \hat{z} \wedge \nabla\psi \qquad \vec{v} = \hat{z} \wedge \nabla\phi \tag{1}$$

The evolution equations for these two functions can be written, in a-dimensional form as (White, 1979; Biskamp, 1978):

$$\left[\begin{array}{l} \frac{\partial\psi}{\partial t} + (\vec{v} + \vec{v}^*) \cdot \nabla\psi = \eta j - E_0 \\ \\ \frac{\partial W}{\partial t} + (\vec{v} + \vec{v}_i^*) \cdot \nabla W = \hat{z}\nabla\psi\wedge\nabla j \end{array}\right. \tag{2}$$

where:

$$\nabla^2\psi = j - \frac{2ka}{m} B_z \qquad \nabla^2\phi = W$$

B_z in the constant magnetic field, $k = n/R$, m and n are the poloidal and toroidal mode numbers. Here, the electron and ion diamagnetic effects were included in the simplest way compatible with an exact formulation: $\vec{v}^* = \vec{v}_e + \vec{v}_i^*$ and $\vec{v}_{e,i}^*$ are the particle diamagnetic drift velocities.

The plasma resistivity η, appearing in eq. (2) is a function of the temperature, $\eta = \eta(T)$, and T will be determined by:

$$\frac{3}{2} n_0 (\frac{\partial}{\partial t} + \vec{v}\cdot\nabla)T - n_0\chi_{\perp}\nabla_{\perp}^2 T = P \tag{3}$$

where n_0 in the mean electron density, $\chi_{\perp}$ the transverse thermal conductivity, and P is the total power input. In the simple case of a discharge with no auxiliary heating, P can be described as a balance between the power absorbed by ohmic heating and the losses by radiation:

$$P = P_{ohm} - P_{rad} = \eta J - n_0^2 \rho T^{\alpha} \tag{4}$$

where ρ is only a function of radial position and the exponent α can be 1/2 for bremstrahlung, 2 for cyclotron radiation and a negative value for radiation due to impurities.

In writing eqs. (2) and (3) we have assumed that parallel thermal conductivity is so high that it prevents the existence of any parallel temperature gradients.

Let us now follow the standard procedure and assume a perturbation of the form:

$$\psi_1(\vec{r},t) = \psi_1(r) \exp i(kz - m\Theta - \omega t) \tag{5}$$

around some equilibrium state, specified by $\psi_0(r)$, $\eta_0(r)$ and $\phi_0 = 0$. The resulting equations for the perturbations are:

$$\left[\begin{array}{l} -i(\omega - \omega^*)(\psi_1 + \frac{m}{r_s}\frac{\psi_0'}{\omega}\phi_1) = \eta_0\nabla^2\psi_1 - \frac{3}{2}\frac{\eta_0 j_0}{T_0} T_1 \\ -i(\omega - \omega_i^*)\nabla^2\phi_1 = i\frac{m}{r_s}\psi_0'\nabla^2\psi_1 - i\frac{m}{r_s}j_0'\psi_1 \\ -i\omega T_1 - \frac{2}{3}\chi_\perp\frac{\partial^2 T_1}{\partial x^2} = a_1\nabla^2\psi_1 + a_2 T_1 + i\frac{m}{r_s}T_0'\phi_1 \end{array}\right. \tag{6}$$

where $r = r_s$ defines the position of the rational surface where the safety factor is $q(r_s) = m/n$ and $x = r - r_s$ is the appropriate radial variable to study the local perturbation. The operator ∇^2 means here:

$$\nabla^2 = \frac{\partial^2}{\partial x^2} - [(\frac{m}{r_s})^2 + k^2]$$

The diamagnetic frequencies are defined by $\omega^* = \omega_e^* + \omega_i^*$ and $\omega_{e,i}^* = (m/r_s)v_{e,i}^*$, where the unperturbed values of $v_{e,i}^*$ are assumed. In order to calculate the perturbed diamagnetic drift velocities we have used electron and ion pressure perturbations given by:

$$p_1 = \frac{m}{r\omega}\phi_1\frac{\partial}{\partial r}p_0$$

In deriving eq. (6) we have also assumed a Spitzer like resistivity, which allowed us to write, for the resistivity perturbation:

$$\eta_1 = -\frac{3}{2}\frac{\eta_0}{T_0}T_1$$

Furthermore, we have assumed that the total power input P is some function of the temperature and the current, $P = P(j, T)$. This explains the coefficients a_1 and a_2 which are defined as:

$$a_1 = \frac{2}{3n_0}\left(\frac{\partial P}{\partial j}\right)_0 \qquad a_2 = \frac{2}{3n_0}\left(\frac{\partial P}{\partial T}\right)_0 \tag{7}$$

In the particular case of eq. (4), we can get:

$$a_1 = \frac{4}{3n_0 j_0} P_{ohm} \qquad a_2 = -\frac{1}{n_0 T_0}\left(P_{ohm} + \frac{2\alpha}{3} P_{rad}\right) \tag{8}$$

We see that a_1 is always positive and a_2 is almost always negative.

It is important to note that, far from the resistive layer, where the ideal MHD can be used, $\eta_0 = 0$, we see from eq. (6) that the evolution of ψ_1 and ϕ_1 is decoupled from thermal effects. This means that the external value for the logarithmic discontinuity Δ' can be calculated in the usual way. In the particularly simple model of $\psi_0' = \tanh x$ and for $k = 0$, we would get the well known result:

$$\Delta' \equiv \frac{1}{\psi_1(0)}\left[\left(\frac{\partial\psi_1}{\partial x}\right)_{0+} - \left(\frac{\partial\psi_1}{\partial x}\right)_{0-}\right] = 2\left[\frac{r_s}{m} - \frac{m}{r_s}\right] \tag{9}$$

In the following, we will be mostly interested in the high m modes, for which $\Delta' < 0$.

3. COLLISIONAL REGIME

In order to drive a linear dispersion relation for the drift-tearing modes in the collisional regime, we make the three usual approximations: First, we assume that the flux function remains mainly constant across the resistive layer, $\psi_1(x) = \psi_1(0)$. Second, we take $\nabla^2 = d^2/dx^2$, which is only valid for moderately high values of m and n. Third, we consider a linear variation of ψ_0' around $x = 0$. We then get, from eq. (6):

$$\left[\begin{array}{l} \dfrac{d^2\bar{\phi}_1}{dx^2} = -\left(\dfrac{m}{r_s}\right)^2 \dfrac{x}{\omega\,(\omega - \omega_i^*)}\,\dfrac{d^2\psi_1}{dx^2} \\[2ex] \dfrac{d^2\psi_1}{dx^2} = \dfrac{\omega - \omega^*}{\tilde{\eta}}\left[\psi_1(0) + (x - \tilde{x})\,\bar{\phi}_1\right] + i\,\dfrac{j_0}{T_0}\,\dfrac{\chi_\parallel}{\omega - ia_2}\,\dfrac{\eta_0}{\tilde{\eta}}\,\dfrac{d^2T_1}{dx^2} \end{array}\right. \tag{10}$$

where $\bar{\phi}_1 = (m/r_s\omega)\phi_1$. We have also introduced an effective resistivity $\tilde{\eta}$ and a thermally induced displacement $\tilde{x}$, such that:

$$\left[\begin{array}{l} \tilde{\eta} = \eta_0 \left(1 - \frac{3}{2}\frac{j_0}{T_0}\frac{ia_1}{\omega - ia_2}\right) \\ \\ \tilde{x} = \frac{3}{2}\frac{j_0}{T_0}\frac{\eta_0 \omega T_0^1}{(\omega - \omega^*)(\omega - ia_2)} \end{array}\right. \tag{11}$$

In the following we will neglect this small displacement and will assume that $\chi_\perp = 0$. In the next section we will show that the contribution of the term in $\chi_\perp$ is negligibly small, at least in the semi-collisional approximation. From eq. (10) we can now get an equation for $\bar{\phi}_1$, which takes the form:

$$\frac{d^2\bar{\phi}_1}{dx^2} + \left(\frac{m}{r_s}\right)^2 \frac{1}{\omega(\omega - \omega_i^*)\,\tilde{\eta}} \left[x^2\bar{\phi}_1 + x(\omega - \omega^*)\psi_1(0)\right] = 0 \tag{12}$$

Integration of eqs. (12) and (10a) through the resistive layer (White, 1979) will then lead to the following form of the general dispersion relation:

$$\omega(\omega - \omega_i^*)(\omega - \omega^*)^3 \left[1 + \frac{a_0}{a_2 + i(\omega - \omega^*)}\right]^{-3} = i\gamma_T^5 \tag{13}$$

where $a_0 = 3j_0 a_1/2T_0$ and γ_T in the linear growth rate for the collisional tearing mode:

$$\gamma_T = \left[\frac{\Gamma(1/4)}{\pi\Gamma(3/4)}\Delta'\right]^{4/5} \eta_0^{3/5} \left(\frac{m}{r_s}\right)^{2/5} \tag{14}$$

From eq. (13) we get, in the case $a_0 = \omega^* = 0$ the tearing mode limit, $\omega = i\gamma_T$, even if $a_2 \neq 0$. This means that local power input dependence on the current is the dominant thermal correction. If $a_0 = 0$, and $\omega^* \neq 0$, eq. (13) will reduce to the well known drift-tearing dispersion relation (Biskamp, 1978). But the new result contained in eq. (13) is that, for $a_0 \neq 0$ we can eventually obtain unstable modes when $\Delta' < 0$. To show this we write $\Delta' = |\Delta'|e^{i\rho}$, where $\rho = 0$ (for $\Delta' > 0$) or $\pm\pi$ (for $\Delta' < 0$). Neglecting, to simplify, the diamagnetic effects and using $\gamma \equiv -i\omega = \gamma_r + i\gamma_i = |\gamma|\, e^{i\Theta}$, we can write, from eq. (13):

$$\begin{cases} |\gamma|^{5/4} |1 - \frac{a_0}{\gamma - a_2}|^{-3/4} = \gamma_T^{5/4} \\ \frac{5}{4}\theta - \frac{3}{4}\alpha = \rho \end{cases} \qquad (15)$$

where α is defined by:

$$\alpha = \tan^{-1} \frac{\gamma_i a_0}{\gamma_i^2 + (\gamma_r - a_2)(\gamma_r - a_0 - a_2)} \qquad (16)$$

The second of eqs. (15) will give the condition for the instability to occur. The mode will be unstable for $\cos\theta > 0$. When $\Delta' > 0$, the mode will be unstable for $\gamma_i = \theta = \rho = 0$. Let us now examine the case of $\Delta' < 0$. If the mode is almost a purely growing mode, $\gamma_r \gg \gamma_i$, we can get from eq. (15b) the following condition:

$$\frac{\gamma_i}{\gamma_r} \approx \frac{5\pi}{4} [1 - \frac{3}{5} \frac{\gamma_r a_0}{(\gamma_r - a_2)(\gamma_r - a_0 - a_2)}]^{-1} \ll 1 \qquad (17)$$

which leads to a growth rate of the order of a_0 or a_2. But, as illustrated by eq. (8), a_0 and a_2 are of the order of the inverse of the energy confinement time. This means that the mode satisfying condition (17) will have a very small growth rate, or will not even exist due to the impossibility of satisfying simultaneously eq. (15a). If we now consider γ_i of the order of (or larger than) γ_r, we arrive at similar conclusions: the angle α must be always significantly different from zero, which means that a_0 and a_2 must be of the order of γ_r.

The inclusion of the diamagnetic effects does not seem to modify these conclusions, which are somewhat negative in what concerns the physical implications of the thermal effects on the collisional tearing modes, for $\Delta' < 0$. However, only a numerical integration of eqs. (6), with less restrictive approximations, would eventually lead to a definitive answer.

4. SEMI-COLLISIONAL REGIME

In high temperature plasmas, for which the electron collision frequency ν_c becomes of the order of (or lower than) $k^2 V_{the}^2/\omega^2$, the previous MHD analysis cannot be applied and we enter the so-called semi-collisional regime. We know that in this new regime the electrostatic effects and the plasma radial displacements can be considered negligible (Drake and Lee, 1977). We can then neglect the influence of the stream function ϕ in eqs. (6). This leads to:

$$\left[\begin{aligned} \frac{d^2\psi_1}{dx^2} &= \left(\frac{\omega - \omega^*}{i\eta_0} + p\right)\psi_1 + \frac{3}{2}\frac{j_0}{T_0}T_1 \\ T_1 &= \frac{ia_1}{\omega - ia_2}\left(\frac{d^2\psi_1}{dx^2} - p\psi_1\right) + i\,\frac{2}{3}\,\frac{\chi_\perp}{\omega - ia_2}\,\frac{d^2T_1}{dx^2} \end{aligned}\right. \tag{18}$$

where $p = (m/r_s)^2 + k^2$. From these equations we then obtain the following equation for ψ_1:

$$-\frac{2i}{3}\frac{\chi_\perp}{\omega - ia_2}\frac{d^4\psi_1}{dx^4} + \left[a_3 + \frac{2}{3}\frac{i\chi_\perp}{\omega - ia_2}\left(\frac{\omega-\omega^*}{i\eta_0} + p\right)\right]\frac{d^2\psi_1}{dx^2} = \left[\frac{\omega-\omega^*}{i\eta_0} + pa_3\right]\psi_1 \tag{19}$$

where we have used:

$$a_3 = 1 - \frac{3}{2}\frac{j_0}{T_0}\frac{ia_1}{\omega - ia_2}$$

Eq. (19) can be simplified because $\chi_\perp$ is a very small quantity. This means that, to the lowest order in $\chi_\perp$, we can write:

$$\frac{d^2\psi_1}{dx^2} = \left[\frac{\omega - \omega^*}{i\eta_0 a_3} + p\right]\psi_1 \tag{20}$$

If we differentiate twice this equation and use the result to estimate the two terms proportional to $\chi_\perp$ in eq. (19), we see that these terms exactly cancel each other. This means that eq. (20) remains valid even for $\chi_\perp \neq 0$. Let us now integrate eq. (20) across the resistive layer, assuming again that ψ_1 is nearly constant. The result can then be written as:

$$\Delta' = \left[-i\,\frac{\omega - \omega^*}{\eta_0 a_3} + p\right]\delta \tag{21}$$

where $\delta = (\omega\nu_c)^{1/2}L_s/k_sV_{the}$ is the width of the resistive layer (Drake and Lee, 1977) and $L_s = [B_z^{-1}(\partial B_\theta/\partial x)]^{-1}$ is the shear length. This can be written as an explicit dispersion relation for ω using:

$$\begin{aligned} b &= \left(\frac{\Delta'}{\delta} - p\right)\eta_0 + a_2 \\ c &= \left(\frac{\Delta'}{\delta} - p\right)(a_2 + a_0)\eta_0 \end{aligned} \tag{22}$$

and neglecting the weak dependence of δ on the frequency, we obtain:

$$\omega_{\pm} = ½(\omega^* + ib) \pm ½ \sqrt{(\omega^{*2} - b^2 + 4c) + 2i\omega^* b} \tag{23}$$

In the drift limit, $\omega^* >> b$ and c, we can write, for $\omega = \omega_r + i\nu$:

$$\left[\begin{array}{l} \omega_{r\pm} = ½(\omega^* \pm \omega^*) \pm \frac{1}{2\omega^*}(c - b^2) \\ \nu_{\pm} = \frac{b}{2}(1 \pm 1) \end{array} \right. \tag{24}$$

The high frequency mode, corresponding to the solution ω_+, will be unstable for $b > 0$. Assuming $a_2 < 0$, the instability will occur for:

$$\Delta' > [(\frac{m}{r_s})^2 + k^2 + \frac{a_2}{\eta_0}]\delta \tag{25}$$

This instability will correspond to the usual semi-collisional tearing mode, modified by thermal effects. The thermal corrections introduce an instability threshold and lead to a decrease in the growth rate, as shown by eqs. (22) to (25).

In the opposite limit of $\omega^* << b$ and c, we will have:

$$\nu_{\pm} = \frac{b}{2} \pm ½\sqrt{b^2 - 4c} \tag{26}$$

Apart from the previous tearing mode, which will still be present, for $b > 0$, we have now a thermal instability, which will occur for $b < 0$ and $c < 0$. For $\Delta' < 0$ and $a_2 < 0$, this corresponds to the condition:

$$|a_2| < \frac{3}{2}\frac{j_0}{T_0} a_1 \tag{27}$$

Using eqs. (8) we can rewrite this condition as $P_{ohm} > (2\alpha/3)P_{rad}$. If the radiation losses are due to impurities ($\alpha < 0$) this inequality is always satisfied. If bremstrahlung is the dominant process ($\alpha = 1/2$) the instability range remains very large, $P_{rad} < 3P_{ohm}$.

It is important to note that the growth rate for this mode will strongly increase with the poloidal mode number m, because not only $|\Delta'|$ increases with m but also p depends quadratically on m. This instability mechanism can then eventually explain the formation of microtearing modes in the semi-collisional regime.

5. MODE SATURATION

We have seen previously that the perpendicular heat conductivity $\chi_\perp$ gives a negligible contribution to the linear growth rate. We will show here that, in spite of that, $\chi_\perp$ can play a significant role in the mode saturation and in the determination of the steady state magnetic island structure. We will follow here an approach similar to that used by Rebut and Hugon (1984), but adapted to our perturbative analysis.

Let us return to the equation of heat transport, eq. (3) and assume an already formed magnetic island. We will assume that the island width, w is small and very far from its saturation value, but already large enough to allow thermal isolation of the island with respect to the surrounding plasma. This means that $w \gtrsim \rho_i$, where ρ_i is the ion Larmor radius, and will strictly not apply to very high m modes. However, even for $w < \rho_i$ some thermal isolation will remain and some general conclusions can still be drawn.

If we integrate eq. (3) over the island volume V, between the planes z and $z + \Delta z$, we obtain:

$$-i\,\frac{3}{2}\,\omega \int_V n_0 T_1 dV - \int_S m_0 \chi_\perp \nabla_\perp T_1 \cdot d\vec{s} = \int_V P_1 dV \tag{28}$$

where s is the lateral island surface. Integration can be easily performed if we assume that T_1 and P_1 are nearly constant inside the island, and we get:

$$\left(\frac{3}{4}\, i w \omega + \chi_\perp\right) T_1 = \frac{P_1 w^2}{2 n_0} \tag{29}$$

where we have used $\nabla_\perp = \partial/\partial x = -2/w$. If we now relate the perturbation in the input power P_1, to the perturbations j_1 and T_1, we arrive at the following result:

$$T_1 = \frac{a_1 j_1 w^2}{{}^4/_3 \chi_\perp - (i\omega + a_2) w^2} \tag{30}$$

At this level we have to remind that w is a function of the flux function perturbation:

$$w^2 = 16 \frac{\psi_1(0)}{\psi_0''}$$

If we now replace eq. (29) in eq. (18a) we can obtain a closed equation for ψ_1, which is formally identical to eq. (20), but with a_3 replaced by:

$$\tilde{a}_3 = 1 + \frac{3}{2}\frac{j_0}{T_0}\frac{a_1 w^2}{{}^4/_3\chi_\perp - (i\omega + a_2)w^2}$$

Intergration of this equation leads then to a nonlinear dispersion relation, where the growth rates depend on the actual island width w. The new dispersion relation will be formally identical to eq. (24), but with b and c replaced by:

$$\left[\begin{array}{l} \tilde{b} = b - \frac{4}{3}\frac{\chi_\perp}{w^2} \\ \\ \tilde{c} = c - \frac{4}{3}\frac{\chi_\perp}{w^2}\left(\frac{\Delta'}{\delta} - p\right)\eta_0 \end{array}\right. \tag{31}$$

For the thermal instability, determined by eqs. (26)-(27) this nonlinearity leads to the appearance of a stationary state ($\nu_+ = 0$), determined by $\tilde{c} = 0$, or equivalently, by:

$$\frac{3}{2}\frac{j_0}{T_0} a_1 = |a_2| + \frac{4}{3}\frac{\chi_\perp}{w^2} \tag{32}$$

But a closer look at these equations clearly shows that such a stationary state cannot be attained to a growing island, because the growth rate ν_+ and $|\tilde{c}|$ can only increase, for an increasing w. Or, in other words, this equilibrium state is not accessible for an island which starts near $w \gtrsim 0$.

Let us now return to eq. (20), with a_3 replaced by $\tilde{a}_3$, and follow the usual quasi-linear procedure (Rutherford, 1973). We can then derive an equation for the time evolution of the island width:

$$\frac{dw}{dt} = \frac{\eta_0}{\pi}|\Delta'|\,\tilde{a}_3(w) \tag{33}$$

where Δ' is calculated using the derivatives of ψ_0 at $x = \pm\, w/2$. On the other hand, we can say that:

$$\frac{dw}{dt} = \frac{\nu}{2}\, w$$

We then conclude that:

$$\nu = \frac{2}{\pi}\, \eta_0 \;\; |\Delta'| \;\; \tilde{a}_3(w) \qquad (34)$$

This result clearly shows that the unstable mode will saturate for $\Delta' = 0$, or for $\tilde{a}_3 = 0$. The first condition corresponds to the usual quasi-linear saturation process, when the source of magnetic energy driving the instability comes to an end. The second one exactly corresponds to the non-linear saturation state defined by eq. (32), which is not accessible to a growing island. From this we can conclude that, even for thermally modified tearing modes, the saturation mechanism will be the quasi-linear one. However, a different saturation mechanism is also suggested by eq. (32), when the unstable island overlaps with nearby islands, leading to the formation of stochastic field lines. In this case, it is known that the heat conductivity $\chi_\perp$ will grow very rapidly with w, near the island surface. Numerical studies (Hugon et al, 1988) suggest that $\chi_\perp$ can grow like $(w - w_c)^\beta$, where $\beta = 2.5$ and $w = w_c$ corresponds to the overlapping condition. This means that, in this case, the last term in eq. (32) will eventually grow with w, instead of decreasing, and the non-linear stationary state defined by $\tilde{a}_3 = 0$ will become accessible.

6. CONCLUSIONS

We have studied the stability of drift-tearing modes when thermal effects are retained, in the collisional and semi-collisional regimes, using a purely analytical approach. We have used plasma fluid equation, where the electron and ion diamagnetic effects were included, in a cylindrical geometry.

We have shown that a new mechanism for driving the instability can exist, associated with the variations in the local power deposition, due to the formation of a magnetic island structure. In particular, this mechanism is very likely to explain the generation of microtearing islands in the semi-collisional regime, which can have an important influence in the perpendicular heat transport. However, only a careful numerical analysis of the growth rates, and comparison with the temperature gradient effects not considered here, can eventually say if this is a plausible explanation for magnetic turbulence in specific experimental conditions.

We have also discussed in a somewhat qualitative way, the saturation processes occurring in these tearing-thermal instabilities, and we have shown that, apart from the well known quasi-linear saturation mechanism, there is one specifically thermal mechanism due to the strong increase in the perpendicular heat conductivity when nearby unstable island overlap. However, this mechanism is far from being elucidated by the present work and a self-consistent theory, where stability and transport are linked together, is required.

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APPENDIX 1.

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