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ABSTRACT.

When the external injection of power is sufficiently high a tokamak ohmically relaxed state with maximum entropy can no longer be sustained and the plasma becomes resistively unstable. The instability is the result of a peculiar interaction between the edge resistivity, the current and the thermal transport at the edge of the current channel. The thresholds of the instability depend on the existence of a current pedestal, on the detachment of the current channel from the material limiters combined with a sufficiently high ratio between the auxiliary power and the toroidal magnetic field and a sufficiently low toroidal current. In the paper the thresholds are expressed in a form which allows the comparison with the experiments. The theoretical predictions are generally consistent with the observations on the H-transitions in tokamaks, in particular with the scalings observed in DIII-D and JIPP T-IIU.

1. INTRODUCTION

A resistive instability of the tokamak discharge was recently discussed (MINARDI, 1988) which results from a peculiar interaction between transport and resistive effects at the edge of the confinement zone. Indeed, in a situation in which the thermal diffusivity increases strongly beyond a critical q-value at the edge of the discharge (for instance, for $q \geq 2$) the

critical q = const surface is also approximately a surface with constant temperature. This means that when the surface q = const is moved by any slight resistive perturbation of the current, the temperature also moves in such a way as to remain constant on this surface. This fact implies a link between the T = const and the q = const surfaces near the edge, or, which is the same, between the local resistivity and the local current. This link may involve a resistive instability (circuit instability, see MINARDI, 1984, 1986) whose existence and time scale strongly depends on the local steepness of the current density profile and on the existence of a pedestal. Since, at the same time, the physical conditions of the plasma at the edge will depend, among other factors, on the incoming heat flux from the mixing region and on the applied power, as well as on the readiness of the system to build up a temperature and current density pedestal, a close relationship will exist between these quantities and the threshold values of the resistive instability, whose parametric dependence we are going to discuss in the present paper. We shall find that, under certain conditions, the parametric dependence of the thresholds is consistent with that observed in the H-transition. This fact seems to imply a situation in which the plasma assumes and maintains resistive states with the physical properties of the H-states provided that it has been able in its previous history to build up the threshold conditions and in particular it was ready, as a consequence of its transport properties, to build up a current density pedestal required by the new resistive state. In fact, properly speaking, the theory to be presented, is not a theory of the pedestal formation, but of the existence and maintenance of resistive states with pedestal. In the resistively unstable situation these states are amplified, while in the stable situation the perturbation forming the pedestal tends to disappear. The description of the detailed mechanisms producing the threshold conditions necessary for the development of the new resistive states is beyond the scope of the present theory, but the development of these states is certainly related to a strong disturbance of the resistive equilibrium of the discharge, due to a number of different causes and involving time scales shorter than the resistive time.

2. TRANSPORT-DISSIPATION PROBLEM IN A SLIGHTLY RESISTIVE TOKAMAK

Let us divide the poloidal flux $A(r,t)$ of a cylindrical tokamak in a part $A_o(r)$ related to an ohmically relaxed current density configuration $j_o(r) \propto T_o^{3/2}(r)$ (Z_{eff} is supposed to be uniform) and in a time dependent part $A_1(r,t)$ which describes the resistive effects generated by the temperature change $T(r,t) - T_o(r)$:

$$-\frac{c}{4\pi}\frac{1}{r}\frac{d}{dr}\left(r\frac{dA_o}{dr}\right) = j_o(r), \qquad j_o(r)\ \eta(T_o(r)) = E \tag{1a}$$

$$-\frac{c}{4\pi}\frac{1}{r}\frac{\partial}{\partial r}\left(r\frac{\partial A_1}{\partial r}\right) = E\left(\frac{1}{\eta(T)} - \frac{1}{\eta(T_o)}\right) - \frac{1}{\eta(T)c}\frac{\partial A_1}{\partial t}, \qquad \eta(T)j = E - \frac{1}{c}\frac{\partial A_1}{\partial t} \tag{1b}$$

In order to calculate the temperature $T(r,t)$ one should couple the Eqs.(1) with the equation for the energy balance, assuming, as discussed above, that the temperature remains constant on the critical q-surface at the boundary. However, the problem can be considerably simplified by exploiting the tendency of a slightly resistive plasma to adhere to the structure of an ideally conductive equilibrium, characterized by a specific functional relation between A and q, provided that the time dependence of the total current remains approximately linear. Indeed, rigorously speaking, in the presence of dissipation, the relation between A and q should be modified by an explicit time dependence $A(r,t) = A(q,t)$. However, it was shown, on the ground of analytical and numerical arguments (MINARDI, 1986, 1988) that, under the condition above for the total current (which can be satisfied for a significant fraction of the resistive time), one can still express $A(r,t)$ in terms of $q(r,t)$ only, namely $A(r,t) = A(q(r,t))$. Moreover, in view of the link between the q = const and the T = const surfaces imposed by the boundary condition, $A(r,t)$ can also be expressed in terms of the temperature only, namely $A(r,t) = A(T(r,t))$, at least in the neighbourhood of the critical q-surface.

The relation $A = A(T)$ is easily found from the condition that, at zero order, it must be identical to the static Ohm's law $\eta(T_o) = E/j_o(A_o)$ ($j_o(r)$ can be expressed uniquely in terms of $A_o(r)$ because $\partial A_o/\partial r \neq 0$ for $r > 0$). Thus, one must have in general $\eta(T) = E/j_o(A)$.

Substitution of $\eta(T_o)$, $\eta(T)$ into (1b) gives the following self-consistent equation for the flux dissipation:

$$\frac{c}{4\pi}\frac{1}{r}\frac{\partial}{\partial r}\left(r\frac{\partial A_1}{\partial r}\right) + j_o(A_o + A_1) - j_o(A_o) = \frac{j_o(A_o + A_1)}{Ec}\frac{\partial A_1}{\partial t} \tag{2}$$

In the following we shall limit ourselves to the linearized version of this equation:

$$\frac{c}{4\pi}\frac{1}{r}\frac{\partial}{\partial r}\left(r\frac{\partial A_1}{\partial r}\right) + \frac{dj_o}{dA_o}A_1 = \frac{j_o(A_o)}{Ec}\frac{\partial A_1}{\partial t} \tag{3}$$

We consider a confinement region $\lambda a \leq r \leq a$ and a boundary region $a \leq r \leq b$ without current and with high thermal conductivity, where b is the location of the material limiters. Assuming an exponential behaviour in time, one is led to study the eigenfunctions defined by the equation

$$\frac{1}{x}\frac{d}{dx}\left(x\frac{dA_n}{dx}\right) + (\nu a)^2 A_n = \lambda_n A_n \tag{4}$$

where $x = r/a$ and

$$\nu^2 \equiv \frac{4\pi}{c}\frac{dj_o}{dA_o} \tag{5}$$

for $\lambda < x < 1$ and $\nu^2 = 0$ for $1 \le x \le b/a$. In general $\nu^2(x)$ is a slow function of x so that it can be approximated by a constant. Taking the maximum or the minimum value of $\nu^2(x)$, the negative or the positive sign of λ_n is respectively a sufficient condition for stability or instability.

The marginal solution

$$A_1 = hJ_o(\nu r) + kY_o(\nu r) \qquad (\lambda a \le r \le a) \tag{6}$$

separates the stable ($\lambda_n < 0$) from the unstable ($\lambda_n > 0$) eigenfunctions. Two sets of boundary conditions can be considered:

$$A_1(b) = 0, \qquad (\partial A_1/\partial r)_{\lambda a} = 0 \tag{7a}$$

or

$$A_1(b) = 0, \qquad A_1(\lambda a) = 0 \tag{7b}$$

The condition of $r = b$ expresses the assumption that the material limiters act as perfect conductors. At the mixing radius $r = \lambda a$ one can take $B_{p1}(\lambda a) \equiv -(\partial A_1/\partial r)_{\lambda a} = 0$, assuming that the poloidal magnetic field in λa and the current inside the sawtooth region $r \le \lambda a$ remain the same as in the zero order, while the process described by $A_1(r,t)$ develops in the outer region. These modes, which were considered in the previous papers, give rise in the nonlinear regime to a motion of the inner boundary of the intermediate region and to an expansion or a shrinking of the core of the current channel. Here we shall study the second family of modes, which keep the inner boundary of the confinement region fixed and conserve the poloidal flux in the region $r \leq \lambda a$. A discontinuity of B_p will eventually appear in $r = \lambda a$ associated with a toroidal sheet current.

The two conditions (7b) involve an algebraic linear system of equations whose compatibility condition gives the marginal stability curve

$$\ell g \frac{b}{a} = \frac{1}{\nu a} \frac{J_o(\nu a)Y_o(\lambda\nu a) - J_o(\lambda\nu a)Y_o(\nu a)}{J_1(\nu a)Y_o(\lambda\nu a) - J_o(\lambda\nu a)Y_1(\nu a)} \tag{8}$$

This equation has an infinite number of branches in the space $(\nu a,\ b/a)$ for a given λ. An example of stability diagram for different values of λ is given in Fig.1 (first branch) and in Fig.2 (second branch). The deformation of the marginal current density profile $j(r) = j_o(A_o + A_1) = j_o(r) + (c\nu^2/4\pi)A_1$ produced by modes of the first and of the second branch, is illustrated in Fig.3. The modes of the first branch describe a global heating (or cooling, depending on the sign of A_1) with a broadening (or a shrinking) of the current density profile. The modes of the second branch, which possess a zero in the interval $\lambda a < r < a$, describe a central cooling and a peripheral heating of the discharge (or vice versa).

The time dependence of the modes results from the two terms in the left hand side of (3). The first is the ordinary diffusive term and implies a characteristic time of order $4\pi a^2/c^2\eta(T)$. The second term contains the effect of the link between the local resistivity and the local current and involves a characteristic time $\tau_\nu = 4\pi/c^2\nu^2\eta(T)$. In general $(\nu a)^2 \gg 1$ so that the time scale of the second effect is much shorter than that of the ordinary resistive diffusion. The parameter $(\nu a)^2$ can reach local values of 50 or more. Thus, for T = 300ev and a = 1 m, τ_ν can be of the order of 0.1s or less.

The time scale is modified by nonlinear effects. Indeed one can write, up to second order

$$j_o(A_o + A_1) - j_o(A_o) = \frac{c}{4\pi} \nu^2 A_1 + \frac{d^2 j_o}{dA_o^2} \frac{A_1^2}{2} \tag{9}$$

One sees from this relation and from Eq. (2) that when A_1 and d^2j_o/dA_o^2 have the same sign, the time scale of the process is shortened and the process tends to take a disruptive character. At the contrary, in the opposite case, the nonlinearity quenches the instability. In this case the system tends to settle in a neighbouring saturated equilibrium (see MINARDI, 1988, 1986).

3. ZERO ORDER MAXIMUM ENTROPY STATES

We shall discuss the dependence of the marginal curves on the physical parameters of the discharge taking, at zero order, those ohmically relaxed steady states of the tokamak which correspond to a vanishing entropy production in the presence of ohmic and auxiliary heating. Thus the magnetic configuration remains stationary under the external action of the heating. These states are then expected to be most likely realized in practice because they correspond to the same maximum value of the entropy as that of the isolated system. The requirement above leads, for a cylindrical tokamak, to the following equation for the axial current density j_o (MINARDI, 1988):

$$\frac{1}{r} \frac{d}{dr} \left(r \frac{dj_o}{dr}\right) + \mu^2 j_o = - \mu^2 \frac{p_A}{E} \tag{10}$$

where p_A is the auxiliary power density, E is the axially induced electric field and μ^2 is a parameter which depends on an "a priori" arbitrary power E^n of the electric field. The dependence of the electric field, of the temperature and of the global confinement time on the auxiliary power and on the ratio $\hat{q}/q_a$ between the safety factor at $r = \lambda a$ and $r = a$, occurs through the biunivocal dependence of these quantities on the parameter μ (see below). Considerations of consistency with the Goldston scaling of the global confinement time imply a cubic dependence $\mu^2 \propto E^3$ and that μ is independent of the toroidal magnetic field B_T and of the total current I (see previous papers).

In the case of a uniform p_A the solution of (10) has the form

$$j_o(r) = HJ_o(\mu r) + KY_o(\mu r) - \frac{p_A}{E} \quad (11)$$

where H and K are determined by the conditions $j(\lambda a) = \hat{j}$, $j(a) = j_a$. The further conditions

$$B_p(\mu a) = \frac{aB_T}{Rq_a}, \qquad B_p(\lambda\mu a) = \frac{\lambda a B_T}{R\hat{q}} \quad (12)$$

applied to the solution for $B_p(\mu r)$ involve a relationship between the dimensionless parameter $p \equiv p_A/E\hat{j}$ and μa for given values of $\hat{q}/q_a$, $\hat{j}_a/j$ and λ:

$$p(\mu a) = \frac{\hat{q}/q_a - \lambda^2 - (2/\mu a D)[(2/\mu a\pi)(1 + j_a/j) + \lambda\Lambda + j_a G/j]}{(2/\mu a D)(4/\pi\mu a + \lambda\Lambda + G) + \lambda^2 - 1} \quad (13)$$

where

$$D = J_o(\lambda\mu a)Y_o(\mu a) - J_o(\mu a)Y_o(\lambda\mu a)$$

$$\Lambda = J_o(\mu a)Y_1(\lambda\mu a) - J_1(\lambda\mu a)Y_o(\mu a)$$

$$G = J_o(\lambda\mu a)Y_1(\mu a) - Y_o(\lambda\mu a)J_1(\mu a) \quad (14)$$

Correspondingly the parameter ν^2 (defined by (5)), through which the marginal stability values of the zero order equilibrium are determined in accordance with (8), can also be parametrized in terms of μa, for given values of $\hat{q}/q_a$, $\hat{j}_a/j$ and λ:

$$(\frac{\nu}{\mu})^2 = - \frac{2}{\lambda\mu a D} [p(\Lambda + \frac{2}{\pi\lambda\mu a}) + \Lambda + \frac{2j_a}{\pi\lambda\mu a\hat{j}}] \qquad (15)$$

Note that this relation is independent of $\hat{q}/q_a$, a fact which will be of importance in the next section. The following properties can be checked numerically:

- For fixed $\hat{q}/q_a$, $j_a/\hat{j}$ and λ, $p = p(\mu a)$ and $X \equiv p(\mu a)(\mu a)^{2/3}$ are decreasing functions of μa.
- $\nu(\mu a)$ is a decreasing function of μa and an increasing function of p and X.

The results above constitute the basis for the discussion, in the next section, of the threshold of the instability in terms of the characteristic parameters of the zero order equilibrium.

4. THRESHOLDS FOR H-TRANSITIONS

4.1 Case of $\hat{q}/q_a$, $j_a/\hat{j}$ variable and p, μ, E, λ fixed

Let us consider a marginal point $(b/a)_m$, λ_m, ν_m on the marginal curves of Figs.1 and 2. We specify the zero order equilibrium by giving the electric field E (related to μa by $E \propto (\mu a)^{2/3}$) and the parameter p, which can be written in the form

$$p = \frac{p_A}{E\hat{j}} = p_A \frac{2\pi R}{EcB_T} \tag{16}$$

where the relation $\hat{q} = cB_T/2\pi R\hat{j} = 1$ was used. We solve (15) with respect to $j_a/\hat{j}$ obtaining the marginal value of the pedestal for the equilibrium specified above:

$$\frac{j_a}{\hat{j}} = -\frac{\pi\lambda_m\mu a}{2} [p(\Lambda + 2/\pi\lambda_m\mu a) + \Lambda + (\lambda_m\mu aD/2)(\nu_m/\mu)^2] \tag{17}$$

Introducing this expression into (13) one can solve for $\hat{q}/q_a$ and obtain the values of $\hat{q}/q_a$ corresponding to the marginal value of the pedestal for the chosen values λ_m, ν_m, E(or μa) and p. Note that in view of $\hat{q} = 1$ one has $\hat{q}/q_a \propto I/B_T$, where I is the current flowing in the cylinder with radius r = a. If now ν_m varies along the marginal curve as the function of the detachment parameter b/a (with λ_m fixed) $j_a/\hat{j}$ and $\hat{q}/q_a$ will also change as functions of b/a and one obtains a representation of the marginal curve in the spaces $(j_a/\hat{j}, b/a)$ or $(\hat{q}/q_a, b/a)$. Examples of this representation are given in Figs.4 and 5 for different values of the parameter $\mu a \propto E^{3/2}$. One sees from the figures that a decrease of $\hat{q}/q_a \propto I/B_T$ (accompanied by a decrease of the pedestal), as well as a decrease of E (or μa), destabilizes the system, for fixed b/a. Similarly, an increase of b/a, for fixed values of E and I/B_T, destabilizes the system. These results are consistent with the experimental observations on the tokamak JIPP TIIU (TOI, 1989). In the plots of Figs.4 and 5 the parameter μa is changed by keeping fixed $p(\mu a)^{2/3}$. Since $(\mu a)^{2/3} \propto E$, this implies recalling (16), that p_A/B_T has the same value for the different curves. Also p_A/B_T is constant along each marginal curve with given E. It follows that the threshold value of the auxiliary power is directly proportional to B_T when one moves along a marginal curve by changing

B_T with I fixed, but remains constant when one moves along the marginal curve by changing I with fixed B_T.

The variations of the parameters discussed above constitute one way for producing a marginal situation. It is not, however, the only way. One can create a marginal situation with a different arrangement of the set of parameters, as we shall discuss below.

4.2 Case of $\hat{q}/q_a$, $j_a/\hat{j}$, ν, λ variable and p, μ, E fixed

As seen from the marginal curves (Figs.1 and 2) the marginal values ν_m depend only on λ for a given value $(b/a)_m$ of the detachment parameter. Then the same holds for the marginal values of $j_a/\hat{j}$ (given by (17)) and of $\hat{q}/q_a$ (given by (13)). Examples of the dependence of these quantities on λ are given in Figs.6 and 7. Note that these marginal equilibria exist in a limited range of $\hat{q}/q_a \propto I/B_T$, $j_a/\hat{j}$ and λ and that the λ-dependence is rather independent of the value chosen for p. Just as above, the quantity p_A/B_T is invariant along these marginal equilibria, so that the threshold value of the auxiliary power is directly proportional to the toroidal field (with fixed E) not only when b/a is varying with a fixed λ, but also for a fixed b/a with a varying λ, $\hat{q}/q_a$ and $j_a/\hat{j}$.

A direct proportionality between the threshold value of the power and the toroidal field has been observed in DIII-D (CARLSTROM, 1989).

In our theory the proportionality factor depends on the loop voltage of the ohmic equilibrium. Introducing the total auxiliary power $P_A = p_A 2\pi^2 a^2 kR$ (where k is the elongation) one can write the marginal relation in the form

$$P_A(\mathrm{MW}) = 5p_m V(\mathrm{Volts}) B_T(\mathrm{Tesla}) ka^2(\mathrm{m})/R(\mathrm{m}) \qquad (18)$$

where p_m is a marginal value of p.

4.3 Case of p, μ, ν, E variable and $\hat{q}/q_a$, $j_a/\hat{j}$, λ fixed

Let us introduce the variable

$$X = p(\mu a)^{2/3} = p_A \frac{2\pi R(\mu a)^{2/3}}{cB_T E} \qquad (19)$$

or also $X = P_A(\mathrm{MW})\cdot F$ where

$$F = 0.2(\mu a)^{2/3} R(\mathrm{m}) / V(\mathrm{Volts}) B_T(\mathrm{Tesla}) ka^2(\mathrm{m}) \qquad (20)$$

Since $E \propto (\mu a)^{2/3}$, the variable X is directly proportional to P_A/B_T for any value of μa. If P_A is increased, keeping B_T fixed, X increases and the same holds for ν because ν is an increasing function of X when the equilibrium parameters above are fixed. As is seen from the marginal curves (Figs.1 and 2), the increase of ν from the marginal point with fixed b/a and fixed λ drives the system into the unstable region.

The space (X, b/a) allows a convenient representation of the marginal curve for fixed values of $\hat{q}/q_a$ (or I/B_T), $j_a/\hat{j}$ and λ. Fig.8 shows an example of this representation based on the first and on the second branch of the marginal curve, in the absence of current pedestal ($j_a = 0$). Fig.9 shows the destabilizing effect of the current pedestal in a representation based on the second branch of the marginal relation. It should be noted, however, that the current pedestal of ohmically relaxed states is not independent of the auxiliary power. Thus, in the case of ohmic relaxation (i.e. $j_o(r) \propto T_o^{3/2}(r)$) the description of the power thresholds in which the pedestal is taken as an independent variable, can have no more than a qualitative significance.

4.4 Case of $\hat{q}/q_a$, $j_a/\hat{j}$, p variable and μ, E, λ, b/a fixed

By varying

$$p = P_A(MW) \frac{0.2\ R(m)}{V(Volts) B_T(Tesla) ka^2(m)} \quad (21)$$

while the parameters above are fixed, one obtains, according to (13) and (17), marginal equilibria of the form illustrated in Fig.10, where the values of $\hat{q}/q_a \propto I/B_T$ are related to the pedestal $j_a/\hat{j}$. The region of instability corresponds to a decrease of the current away from the marginal equilibrium related to a decrease of the pedestal, for fixed values of B_T and of the parameters above.

5. FINAL REMARKS

We have considered ohmically relaxed steady states of the tokamak defined by vanishing entropy production in the presence of ohmic and auxiliary heating. When the external injection of power is sufficiently high, namely above the threshold given by (18), the ohmic relaxation in a steady state can no longer be sustained and the plasma becomes resistively unstable as a consequence of a peculiar interaction between the edge resistivity, the current and the thermal transport occurring at the edge of the discharge. The ohmic relaxation is then broken and the plasma enters in a new (generally time dependent) resistive state, provided that specific plasma conditions, which are necessary for the existence of the new resistively unstable states, are realized in the discharge as a result of its previous history.

The realization of the threshold conditions required for the development of the new resistive states can be due to a number of different mechanisms, all producing a strong departure from the resistive equilibrium of the discharge and involving time scales shorter than that of the resistive diffusion. The explanation of these mechanisms constitutes a separate problem. The threshold conditions involve the existence of a current pedestal, the detachment of the current channel from the material limiters combined with a sufficiently high ratio between the auxiliary power and the toroidal magnetic field and a sufficiently low toroidal current (i.e. $q_a \gtrsim 2$). The theoretical model is supported by experimental observations in DIII-D and JIPP T-IIU.

Besides a better understanding of the mechanisms producing the threshold conditions (in this connection see HINTON, 1985), a consistent identification of our resistively unstable states with the H states needs the demonstration of the improved particle and energy confinement after the transition. For treating this problem, which is related to the change of the temperature and current density profiles after the transition, the reduced Eq.(2) is too simplified and one should revert to the complete set of the Ohm-Maxwell-Transport equations, taking into account however, the boundary conditions resulting from the peculiar link, discussed above, between a critical value of q and a constant temperature on the same critical q-surface, at the edge of the confinement zone.

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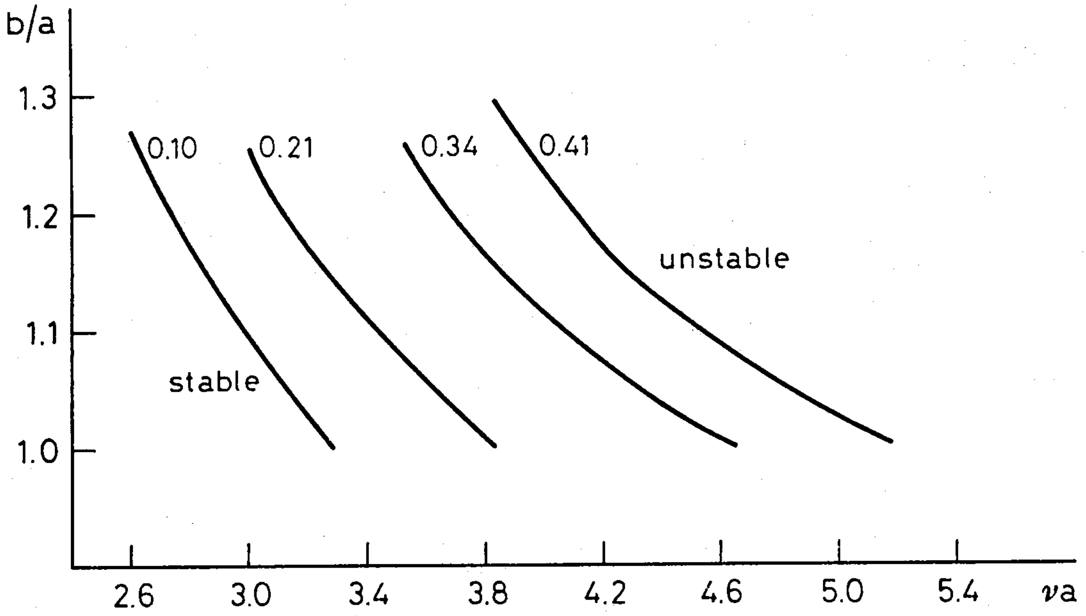

Fig.1 First branch of the marginal curve (8), for different values of the parameter λ.

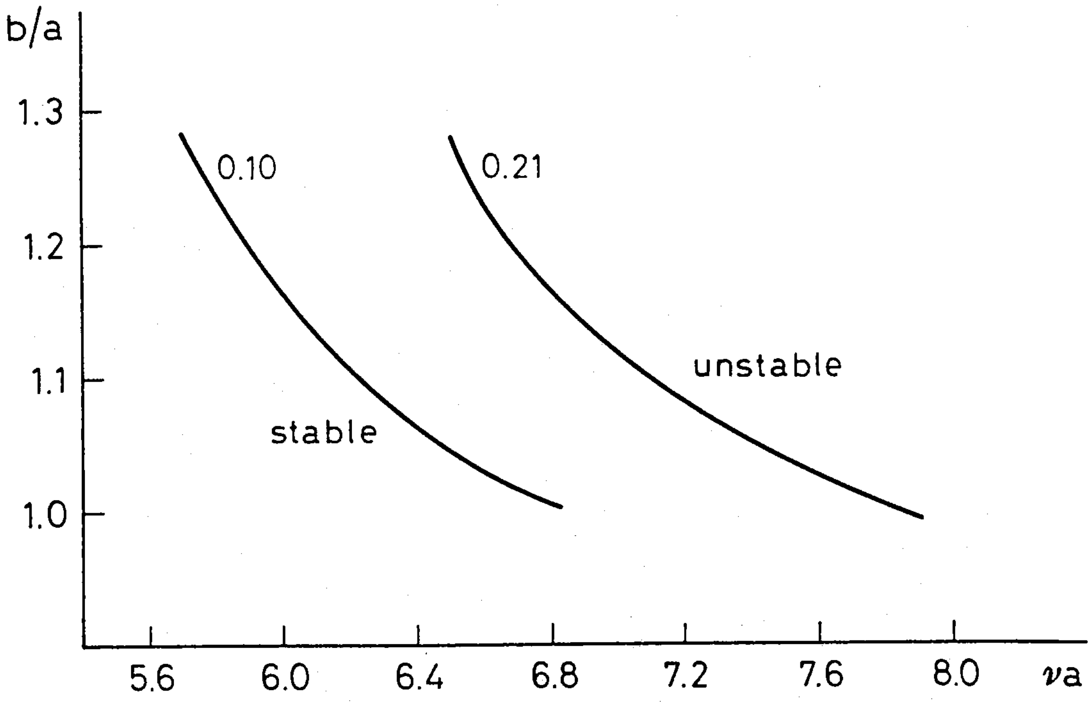

Fig.2 Second branch of the marginal curve (8), for different values of the parameter λ.

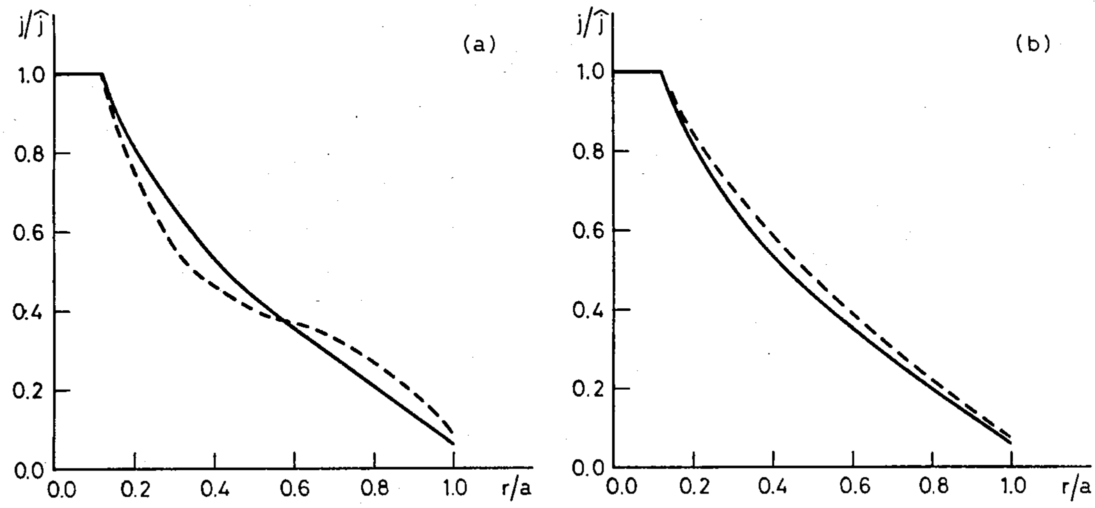

Fig.3 Example of a zero order current density profile (solid line, $\hat{q}/q_a$ = 0.33, $j_a/\hat{j}$ = 0.06, λ = 0.12, μa = 0.35) and of the profile distortion (dashed line) produced by a marginal mode. (a), mode of the second branch; (b) mode of the first branch.

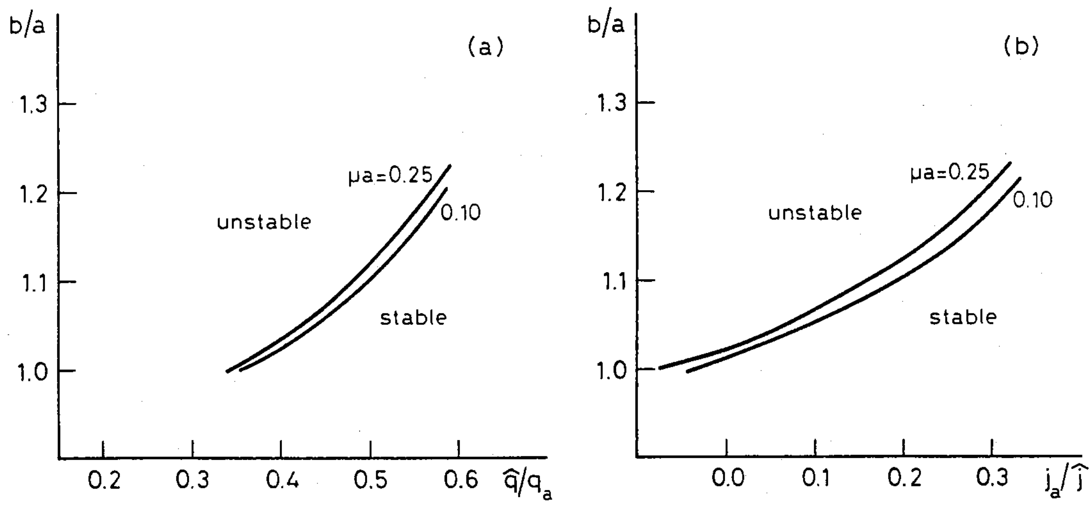

Fig.4 Examples of marginal curves corresponding to the first branch for given p, μa (or E) and λ; $p(\mu a)^{2/3}$ = 2.00, λ = 0.2944. The marginal equilibria can be labelled equivalently by $\hat{q}/q_a$ (plots (a)) or by the related values of the pedestal $j_a/\hat{j}$ (plots (b)).

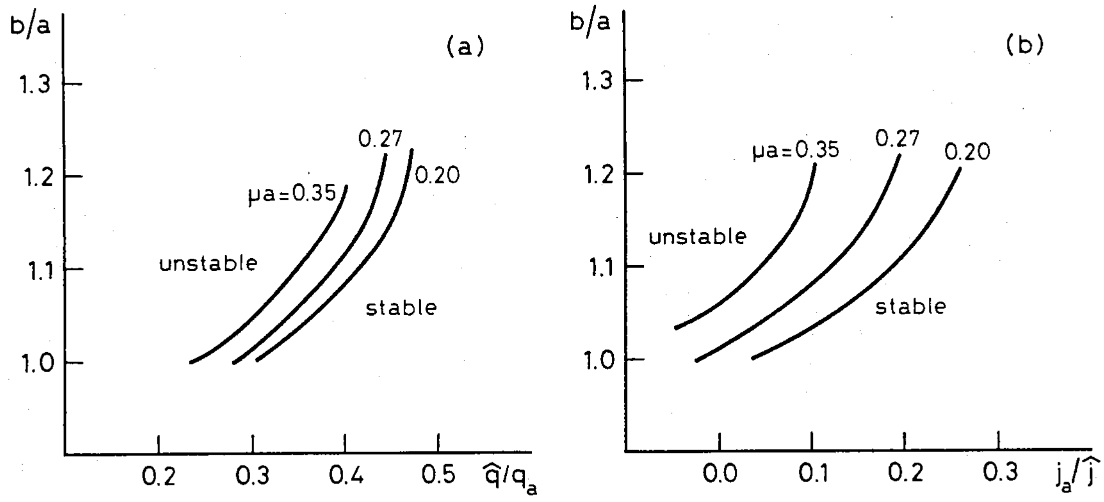

Fig.5 Same as in Fig.4 for marginal curves corresponding to the second branch; $p(\mu a)^{2/3} = 4.5952$, $\lambda = 0.12667$.

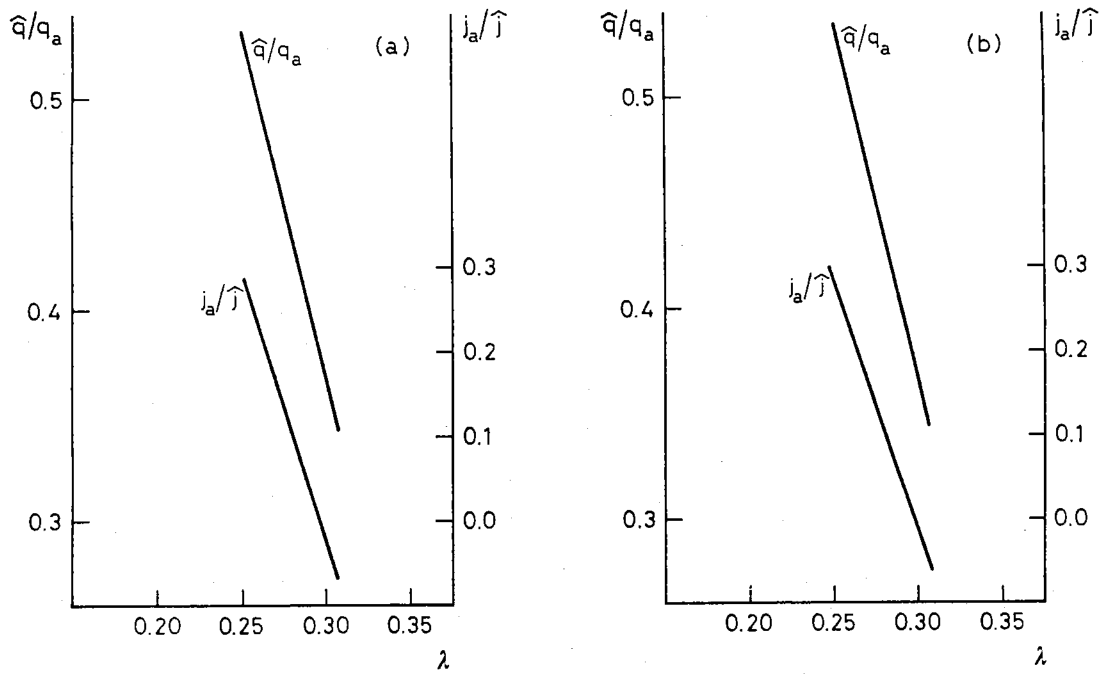

Fig.6 Marginal equilibria corresponding to the first branch for given p, b/a and μa (or E) and variable $\hat{q}/q_a$, $j_a/\hat{j}$ and λ. (a) p = 3.51, b/a = 1.02, μa = 0.12; (b) p = 10.207, b/a = 1.02, μa = 0.10.

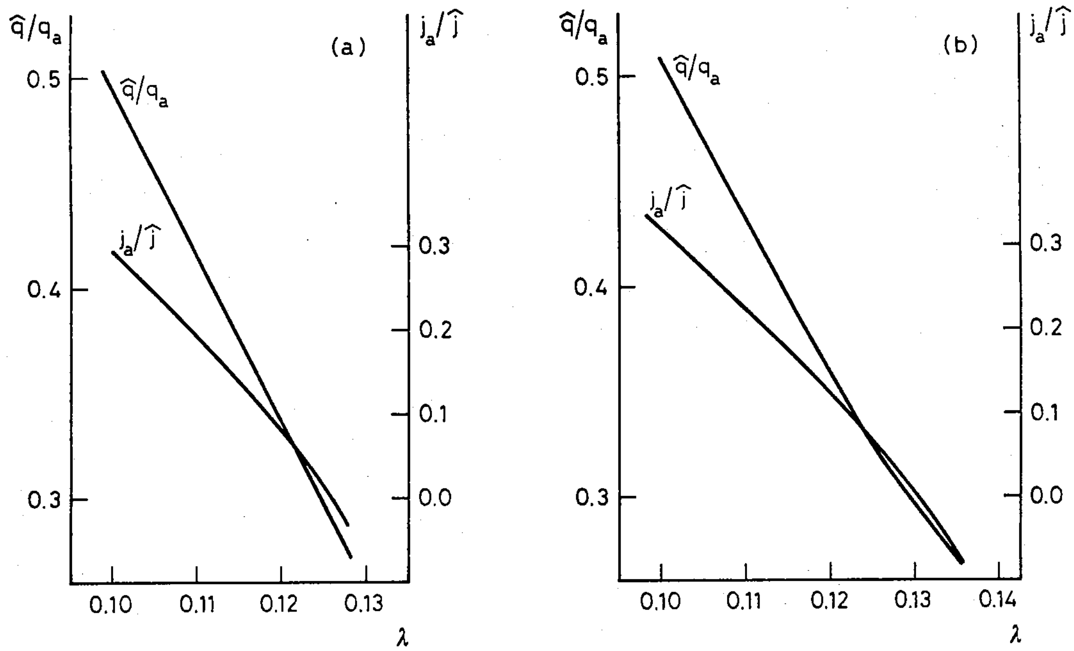

Fig.7 Same as in Fig.6 for marginal equilibria corresponding to the second branch. (a) p = 5.0495, b/a = 1.00, μa = 0.35; (b) p = 11.00, b/a = 1.00, μa = 0.22.

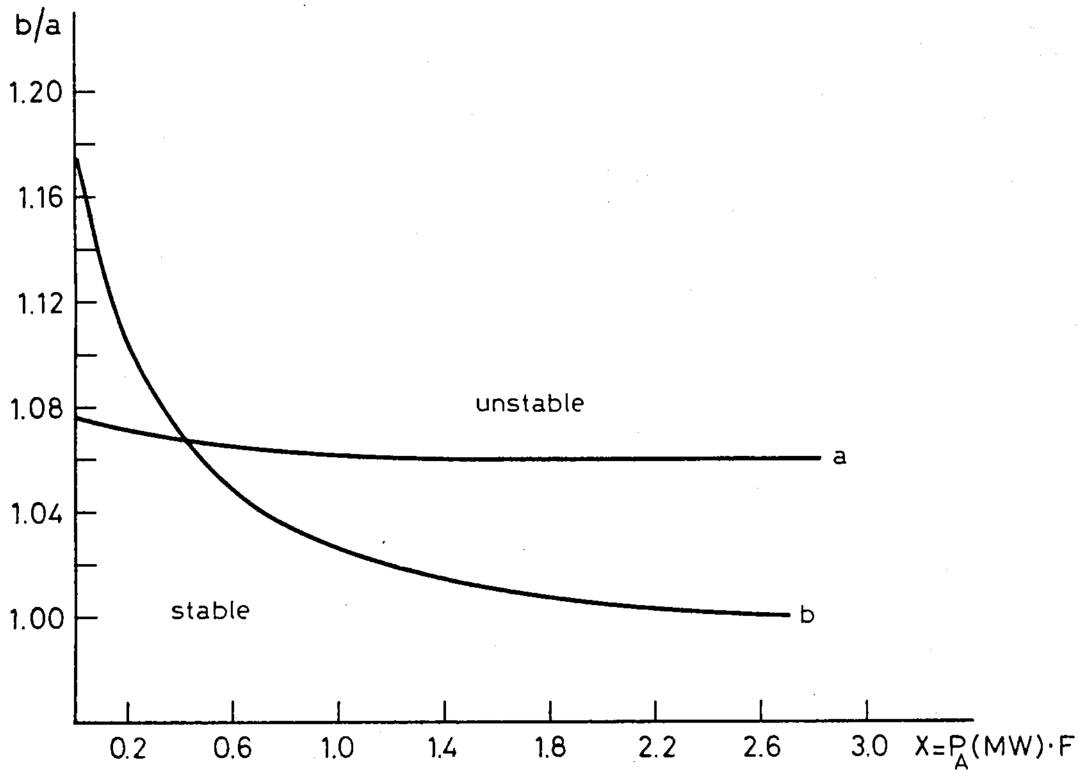

Fig.8 Examples of marginal curves in the space (P_A, b/a) for given $\hat{q}/q_a$, λ and $j_a = 0$. (a) $\hat{q}/q_a = 0.40$, $\lambda = 0.30$, first branch; (b) $\hat{q}/q_a = 0.30$, $\lambda = 0.10$, second branch. Recalling that $E \propto (\mu a)^{2/3}$ one has that the factor F is a constant.

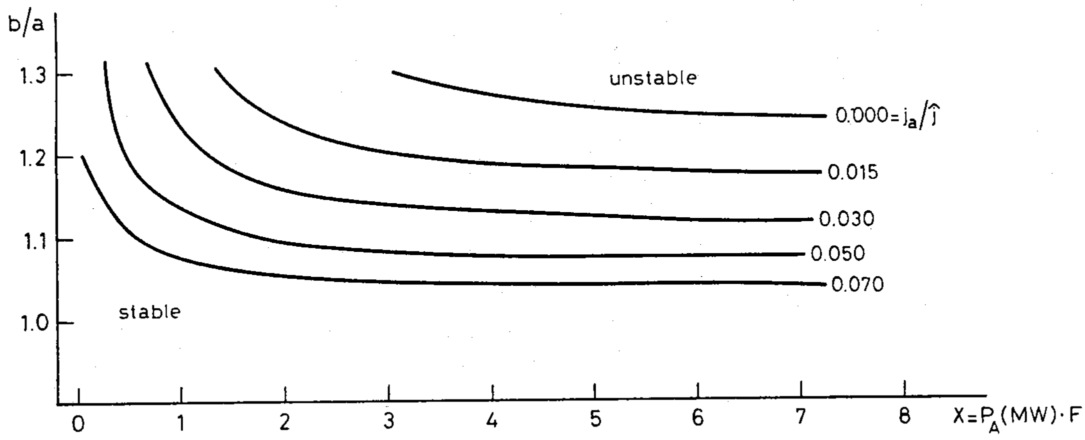

Fig.9 Examples of marginal curves in the space (P_A, b/a) for given $\hat{q}/q_a$, λ and different values of the pedestal; $\lambda = 0.10$, $\hat{q}/q_a = 0.40$, second branch.

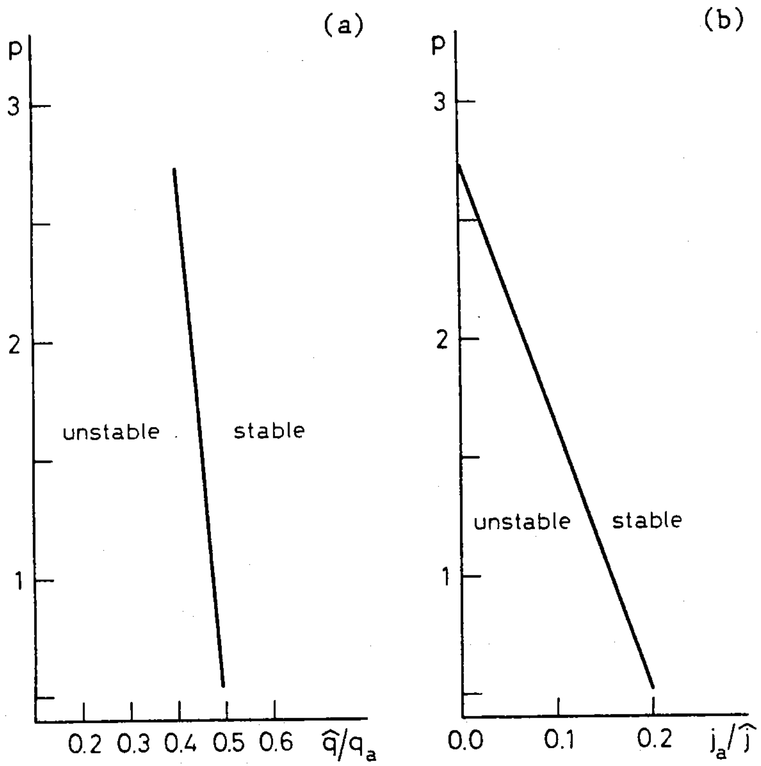

Fig.10 Typical examples of marginal equilibria for $\hat{q}/q_a$, $j_a/\hat{j}$ and p variable; $\mu a = 0.70$, $\lambda = 0.25$, $b/a = 1.02$, first branch. The equilibria can be labelled equivalently by $\hat{q}/q_a$ (plot (a)) or by the related value of the pedestal (plot (b)).