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SUPPRESSION OF INTERNAL PLASMA OSCILLATIONS BY TRAPPED HIGH ENERGY NUCLEI

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ABSTRACT

The suppression of macroscopic oscillations (so-called "sawteeth") of the central region of a magnetically confined plasma column is related to the effect of magnetically trapped energetic nuclei produced by the injection of rf waves or neutral beams. We evaluate the threshold for, and describe the process involved in, the transition of $m^0 = 1$ modes, which trigger the crash phase of sawtooth oscillations, to a regime that is shown to be stable.

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In experiments carried out by the JET machine, the internal electron temperature relaxation oscillations (the so-called "sawtooth" oscillations), which are associated with the excitation of a mode with $m^0 = 1$, $n^0 = 1$ poloidal and toroidal mode numbers, have been suppressed for periods of up to 3.2s during high power injection of ion-cyclotron frequency waves (with coupled power $P_{RF} \gtrsim 4MW$) and/or neutral beams (with $P_{NBI} \gtrsim 7MW$) [1]. This observation cannot be explained within the framework of ideal or resistive mhd models, even when the stabilising effects of the ion diamagnetic frequency and of the electron drift wave frequency are considered [2]. In fact, according to these models, stability would occur either when the value q_o of the q parameter on the magnetic axis exceeds unity, or for sufficiently small values of the plasma poloidal beta and low collisionality. Instead, Faraday rotation measurements [3] indicate that $q_o \sim 0.7$, while high poloidal betas $\beta_p(r_o)$ often in excess of the ideal mhd instability threshold [4] β_p^{mhd} are reached during sawtooth-free periods in JET. Here, $\beta_p(r_o) \equiv [8\pi/B_p^2(r_o)]\ [\langle p(r_o)\rangle - p(r_o)]$, where B_p is the poloidal magnetic field, r_o is the mean radius of the $q = 1$ surface, p is the plasma pressure, and $\langle p(r_o)\rangle$ is its average value within the $q = 1$ volume. Specifically we find [5] $\beta_p(r_o) > 0.2$ at moderate plasma currents ($I_p \sim 2MA$), while taking into account the effects of plasma shaping [5,6], $\beta_p^{mhd} \sim 0.1\text{-}0.2$.

JET discharges which exhibit sawtooth suppression are characterised by the presence of anisotropic energetic nuclei, accelerated by radio frequency fields, and/or produced by energetic neutral beam injection (with injection energy $\mathcal{E}_{inj} \simeq 80keV$). In preliminary analyses [7,8] it was pointed out that the energetic trapped nuclei can indeed play a stabilising role. Possibly the best experimental clue is provided by "ICRH switch off" experiments [9]. In these experiments, the stable period is terminated by a sudden interruption of the applied rf fields (no neutral beams are injected in these cases). A time delay of the order of 100ms is observed between the rf switch off time and the collapse of the central electron temperature. This time delay is a finite fraction of the fast nuclei slowing down time, therefore suggesting that the loss of these nuclei is the dominant destabilising effect [10].

Plasma stabilisation by energetic particle populations was proposed and analysed for the case of Astron [11] and ion ring devices [12], and for the ELMO bumpy torus [13]. In axisymmetric toroidal configurations, stability against ballooning modes was theoretically shown to improve in the presence of magnetically trapped fast nuclei (see e.g. Ref. 14). For the case of the $m^0 = 1$ mode, energetic nuclei were initially considered [15,16] in order to show that they can destabilise a branch of the $m^0 = 1$ dispersion relation. In fact, following the analysis of Ref. [2], the relevant dispersion relation in the absence of suprathermal particles is

$$\omega(\omega - \omega_{di}) = - \gamma^2_{mhd}, \qquad (1)$$

where ω is the mode frequency (in the frame of reference where the equilibrium $\underline{E}x\underline{B}$ drift vanishes at the $q = 1$ surface), $\omega_{di} \equiv [- c(dp_i/dr)/enBr]_o$ is the bulk ion diamagnetic frequency at $r = r_o$ and γ_{mhd} is the growth rate found by the ideal mhd approximation [4]. When $\gamma_{mhd} < |\omega_{di}|/2$, Eq. (1) yields two marginally stable roots. By introducing the effects of finite electrical resistivity, the one with the lower frequency becomes unstable [2] and is believed to correspond to the mode responsible for the crash-phase of sawtooth oscillations. The higher frequency mode ($\omega \approx \omega_{di}$) instead can be driven unstable [15] by an "effective viscosity" arising from the resonance between this mode and the trapped energetic nuclei with bounce averaged magnetic drift frequency $\omega^{(o)}_{Dh}$ equal to ω. This resonant interaction is most effective when $\omega_{di} \sim \bar{\omega}_{Dh}$, with $\bar{\omega}_{Dh}$ the characteristic value of $\omega^{(o)}_{Dh}$ over the distribution of the energetic nuclei.

The "sawtooth-free" regimes in JET are generally characterised [6] by $|\omega_{di}| \lesssim \gamma_{mhd} < \bar{\omega}_{Dh}$, while the ratio between the transverse pressure of the trapped energetic nuclei $p_{\perp h}$ and the bulk pressure p can be as high as the inverse aspect ratio $\epsilon_o \equiv r_o/R_o$. In this case, as shown by the derivation in the second part of this Letter, the dispersion relation (1) for modes with frequency $\omega < \bar{\omega}_{Dh}$ is modified into

$$[\omega(\omega-\omega_{di})]^{1/2} = i(\gamma_{mhd} - H\,\omega), \tag{2}$$

where

$$H \approx \frac{\pi}{3s_o}\,\epsilon_o\,\beta_{ph}\,\frac{\omega_A}{\bar{\omega}_{Dh}}\ , \tag{3}$$

$$\beta_{ph} \approx -\frac{8\pi}{B_p^2(r_o)}\int_0^{r_o} dr\left(\frac{r}{r_o}\right)^{3/2}\frac{d}{dr}\left[\left(\frac{r}{r_o}\right)^{1/2} p_{\perp h}\right].$$

$s_o = r_o q'(r_o)$ is the dimensionless shear parameter, $\omega_A = v_A/R_o\sqrt{3}$, and $v_A = B/(4\pi m_i n_i)^{1/2}$ is the Alfvén velocity. Acceptable roots of Eq. (2) must satisfy the condition $\gamma_{mhd} - H\,\mathrm{Re}(\omega) > 0$ for the corresponding eigenfunction to be spatially regular [15]. Purely oscillatory modes are obtained from (2) for

$$H \geq H_{cr} \equiv \frac{\gamma_{mhd}}{\omega_{di}}\left(1 - \frac{\omega_{di}^2}{4\gamma_{mhd}^2}\right). \tag{4}$$

At $H = H_{cr}$, $\omega = \gamma_{mhd}/(1 + H_{cr}^2)^{1/2} \lesssim \omega_{di}$ and $\gamma_{mhd} - H\omega = \gamma_{mhd}\,[1 - (1 + H_{cr}^{-2})^{-1/2}]$. As H is increased above γ_{mhd}/ω_{di}, only the lower frequency root satisfies the condition $\gamma_{mhd} - H\omega > 0$. Stabilisation arises because energy must be spent to perturb the energetic trapped nuclei when their magnetic drift velocity is larger than the mode phase velocity. With [4] $\gamma_{mhd} \sim \omega_A \epsilon_o^2 [\beta_p^2 - (\beta_p^{mhd})^2]$, the stability criterion (4) for $\beta_p > \beta_p^{mhd}$ can be rewritten as

$$\beta_{ph} > \alpha \frac{s_o r_o}{R_o} \beta_p^2 \frac{\bar{\omega}_{Dh}}{\omega_{di}} \qquad (5)$$

where α is a numerical factor of order unity which depends primarily on the q profile. Condition (5) depends on the density of the energetic nuclei and on their spatial profiles, and, since $\bar{\omega}_{Dh} \propto \mathcal{E}_h$, is independent of their mean energy $\mathcal{E}_h$. It is consistent with the plasma parameter values ($\beta_{ph} \sim 0.05$, $\mathcal{E}_h \sim 100keV$, $\bar{\omega}_{Dh}/\omega_{di} \sim 3$, $\beta_p \sim 0.2$-0.3, and $R/r_o \sim 6$-7) of the regimes where sawtooth oscillations are suppressed in JET. In addition, when off-axis rf heating is applied, the profile of $p_{\perp h}$ may not be sufficiently peaked to fulfill condition (5) and this can explain why sawteeth are difficult to suppress in this case.

As the value of H is increased and the marginal stability condition is fulfilled, a term representing the mode particle resonance $\omega = \omega_{Dh}^{(o)}$ has to be added on the r.h.s. of Eq. (2). Specifically H is to be replaced by $H[1 + i\, g(\omega_R/\bar{\omega}_{Dh})]$, where $\omega_R \equiv Re(\omega)$ and $g(\omega_R/\bar{\omega}_{Dh})$ depends on the velocity distribution of the energetic nuclei. For the model given by Eq. (13), g reduces to $(\omega_R/\bar{\omega}_{Dh})^{3/2}$. With the addition of the resonant contribution, the lower frequency root of Eq. (2) is damped.

However, the restriction to modes with frequency $\omega < \bar{\omega}_{Dh}$ ceases to be relevant if H is increased above a second threshold value. A "macroscopic stability window" is thus found [7]. In fact, beyond this second threshold, an unstable mode with $\omega \sim \bar{\omega}_{Dh}$ can be excited [16]. The second threshold depends on the equilibrium distribution of the energetic nuclei, but generally it involves values of H considerably larger than H_{cr} in Eq. (4).

The conclusions given above do not take into account the effect of the small but finite electron resistivity. When this is included, the relevant dispersion relation is (see Refs. [2] and [15])

$$[\omega(\omega-\omega_{di})]^{1/2} = i\,[\gamma_{mhd} - H\omega(1 + ig)]\,\frac{Q^{3/2}}{8}\,\frac{\Gamma[(Q-1)/4]}{\Gamma[(Q+5)/4]} \tag{6}$$

where $Q^2 \equiv i\omega(\omega-\omega_{di})(\omega-\hat{\omega}_{*e})/(\omega_A{}^3\epsilon_\eta)$, with $\epsilon_\eta \equiv \eta c^2\, s_o/(4\pi\, r_o\, \omega_A)$ the inverse magnetic Reynolds number. In the experiments of interest, ϵ_η is typically $10^{-7} - 10^{-8}$. The parameter γ_{mhd} can become [5] larger than the characteristic growth rate $\omega_A\, \epsilon_\eta{}^{1/3}$ of resistive $m^0 = 1$ modes, while $\omega_A\, \epsilon_\eta^{1/3} \gtrsim |\hat{\omega}_{*e}| \gtrsim |\omega_{di}|$, with $\hat{\omega}_{*e} = [(T_e c/eBrn)(dn/dr)\ (1+1.71\eta_e)]_o$ and $\eta_e = d\ln T_e/d\ln n$. The dispersion relation (2) is recovered from (6) to lowest order by taking $|Q| > 1$. In this limit, treating the effect of resistivity as a perturbation, we find that the mode-particle resonance overcomes the destabilising effect of resistivity when

$$2(\gamma_{mhd} - H\omega_R)\, H\omega_R\, g > (5/2)\, \epsilon_\eta\, \omega_A^3(\omega_R - \hat{\omega}_{*e})^{-1} \tag{7}$$

At $H = H_{cr}$, Eq. (7) reduces to $\omega_{di}/(\epsilon_\eta{}^{1/3}\,\omega_A) \gtrsim (\bar{\omega}_{Dh}/\omega_{di})^{1/2}$ for $g(\omega_R/\bar{\omega}_{Dh}) \sim (\omega_R/\bar{\omega}_{Dh})^{3/2}$. For larger values of $\omega_{di}/(\epsilon_\eta\, \omega_A)$ we find a range in H where both roots of Eq. (2) are stable. This is illustrated in Fig. 1 where the two roots, as obtained from the numerical integration of the eigenvalue problem leading to Eq. (6), are followed in the complex frequency plane as a function of H for two different values of $\omega_{di}/(\epsilon_\eta^{1/3}\ \omega_A)$. The oscillation frequency is reduced when H is increased above H_{cr}. As a consequence, the mode damping due to the resonance becomes weaker. A similar effect occurs for values of γ_{mhd} below $\omega_A\ \epsilon_\eta^{1/3}$, and thus the dispersion relation (6) does not yield complete stabilisation when $\gamma_{mhd} \to 0$. However, in this limit and for $\omega_{di}/\omega_A > 2\epsilon_\eta^{1/3}$, the viscous dissipation resulting from bulk ion collisions effectively replaces the resonant damping and ensures stability [17], independently of the presence of the trapped energetic nuclei.

A problem remains for more realistic values of JET experimental parameters, corresponding to the dashed curves of Fig. 1. In this case the instability is not completely suppressed even though the ideal mhd growth rate is strongly depressed. However one should bear in mind that in this case the occurence of magnetic reconnection is necessary and that the linearised mode description breaks down for rather small amplitudes, corresponding to a width of the magnetic island of the order of the mode transition layer. Since this layer becomes of the order of the ion gyroradius, the mode amplitude for which the linear stability analysis breaks down is very modest. Another point to be considered is that the set of equations that we have adopted to describe the possible occurrence of resistive modes may not be adequate, given the low rate of collisionality of the regime of interest.

The main analytical steps which lead to the conclusions we have given are now outlined. We assume the ordering $p_{\perp h} \sim \epsilon_o p$. In addition, since we limit our considerations to a population of trapped energetic nuclei, we have $p_{\parallel h} \sim \epsilon_o\, p_{\perp h}$. The dominant component of the Lagrangian displacement is represented by $\tilde{\underline{\xi}} = \hat{\underline{\xi}}(r)\ \exp(-\ i\omega t + iP)$ where $P = \zeta - \theta$, ζ and θ are the poloidal and toroidal angles, respectively. The total momentum conservation equation is written in the form

$$nm_i \frac{d^2\underline{\tilde{\xi}}}{dt^2} = - \nabla\cdot\underline{\underline{\Pi}} + \frac{1}{c} (\underline{\tilde{J}}x\underline{B} + \underline{J}x\underline{\tilde{B}}), \qquad (8)$$

where the pressure tensor $\underline{\underline{\Pi}}$ contains the contributions of both the background plasma and the energetic nuclei. Equation (8) can be reduced to an equation in $\tilde{\xi}_r$ using a standard procedure [18]. Then we expand $\tilde{\xi}_r$ in powers of ϵ_o. To lowest order in ϵ_o we find $\hat{\xi}_r = \Theta(-x)\ \hat{\xi}_o$, where $\hat{\xi}_o$ is a constant, $x = (r-r_o)/r_o$, and $\Theta(x)$ is the Heaviside function. To the next relevant order, we find [15,18] for $x \ll 1$, i.e. approaching the $q = 1$ surface,

$$\frac{d\hat{\xi}_r}{dx} = - \frac{\lambda_H + \lambda_K(\omega)}{\pi s_o x^2}\ \hat{\xi}_o.$$

The parameter $\lambda_H \sim O(\epsilon_o\ \beta_p)^2$, which is derived e.g. in Refs. [4], [6] and [19], is proportional to the mhd energy functional δW and is such that $\lambda_H > 0$ when $\beta_p > \beta_p^{mhd}$. The parameter $\lambda_K \sim O(\epsilon_o\ \beta_{ph})$ represents the contribution of the energetic nuclei and is expressed [15,20] in terms of their perturbed perpendicular pressure $\tilde{p}_{\perp h}(\omega)$

$$\lambda_K(\omega) = - \frac{4\pi^2 i}{B_{po}{}^2 s_o\ \xi_o} \int_o^{r_o} dr\ r^2 \int_{-\pi}^{\pi} \frac{d\theta}{2\pi} [\bar{e}_{||} \times \bar{\kappa}\cdot\nabla\tilde{p}_{\perp h}(\omega)]\ \exp(i\omega t - iP), \qquad (9)$$

where $\bar{e}_{||} = \bar{B}/B$ and $\bar{\kappa} = (\bar{e}_{||}\cdot\nabla)\bar{e}_{||}$. The solution for $\hat{\xi}_r$ given above, which is valid outside a small layer around $x = 0$, in a region where the inertial and finite ion Larmor radius being contained in Eq. (6) are unimportant, can be matched with the relevant asymptotic solution within this layer. Then we arrive at the dispersion relation

$$[\omega(\omega-\omega_{*i})]^{1/2} = i\omega_A[\lambda_H+\lambda_K(\omega)] \qquad (10)$$

of which Eq. (2) is a special case, showing that $\gamma_{mhd} = \omega_A \lambda_H$.

We derive $\tilde{p}_{\perp h}(\omega)$ from the perturbed distribution function $\tilde{f}_h$ neglecting effects related to the presence of a radial equilibrium electric field. Solving the linearised Vlasov equation in the drift kinetic approximation, $\tilde{f}_h$ can be conveniently split in two parts: $\tilde{f}_h = \tilde{f}_h^{ad} + \tilde{f}_h^{nad}$. The adiabatic part is $\tilde{f}_h^{ad} = -\tilde{\underline{\xi}}_\perp \cdot \nabla F_{oh}$, with F_{oh} the equilibrium distribution function of the energetic trapped nuclei. For $\omega < \omega_{bh}$, with ω_{bh} the fast nuclei bounce frequency, and neglecting finite banana orbit effects, the non-adiabatic part to lowest order in ϵ_o is $\tilde{f}_h^{nad} = \hat{f}_h^{nad} \exp(-i\omega t + iS)$, where $S = (\zeta - q\theta)$ and

$$\hat{f}_h^{nad} = \frac{\hat{\xi}_r}{R_o} \frac{\omega - \omega_{*h}^T}{\omega - \omega_{Dh}^{(o)}} \mathcal{E} \frac{\partial F_{oh}}{\partial \mathcal{E}} [\cos(q\theta)]^{(o)}, \tag{11}$$

with $(\partial F_o/\partial \mathcal{E}) \omega_{*h}^T = -\bar{e}_{||} x \nabla F_{oh} \cdot \nabla P/(m_h \Omega_h)$, and $\mathcal{E} = m_h v^2/2$. The superscript "(o)" indicates particle orbit averaging. The phases of $\tilde{f}_h^{ad}$ and $\tilde{f}_h^{nad}$ differ by a factor $S-P = (1-q)\theta$, as in the considered frequency range $\tilde{f}_h^{nad}$ must be constant to lowest order along the trapped particle orbit, i.e. $\bar{e}_{||} \cdot \nabla \tilde{f}_h^{nad} = 0$.

The parameter λ_K is also split in two terms,

$$\lambda_K(\omega) = \lambda_K^{ad} + \lambda_K^{nad}(\omega)$$

where

$$\begin{Bmatrix} \lambda_K^{ad} \\ \lambda_K^{nad} \end{Bmatrix} = \frac{2\pi}{B_{po}^2 s_o R_o \xi_o} \int_o^{r_o} r dr \int_{-\pi}^{\pi} d\theta \begin{Bmatrix} \hat{p}_{\perp h}^{ad} & \cos\theta \\ \hat{p}_{\perp h}^{nad} & \cos(q\theta) \end{Bmatrix} \tag{12}$$

and $\hat{p}_{\perp h} = (m_h/2) \int d^3v\, v_\perp^2\, \hat{f}_h$.

The perturbed perpendicular pressure $\tilde{p}_{\perp h}$ of the energetic nuclei depends rather sensitively on their equilibrium distribution function F_{oh}. For the present analysis, we take

$$F_{oh} = \left(\frac{m_h}{2\mathcal{E}_h}\right)^{3/2} \frac{n_h(r)}{K(1/2)\sqrt{\pi}} \exp\left(-\frac{\mathcal{E}}{\mathcal{E}_h}\right) \delta\left[\left(\frac{R}{2r}\right)^{1/2}(\Lambda - 1)\right], \qquad (13)$$

where $\Lambda \equiv (v_\perp/v)^2 B_o/B$ is the pitch angle variable in velocity space, B_o is the value of B on the axis, $B_o/B \approx 1 + (r/R)\cos\theta$, and K(x) is the elliptic integral of the second kind ($K(1/2) \approx 1.85$). For simplicity $\mathcal{E}_h$ is taken to be a constant value $\mathcal{E}_h \sim 100$keV. This distribution function represents trapped energetic nuclei with a magnetic turning point at $\theta_o = \pi/2$. It can be used as a rough model for the anisotropic energetic nuclei tail which is formed during ICRH heating in a toroidal magnetic confinement configuration, since trapped nuclei having the tips of their banana orbits on the rf resonant layer along a vertical chord across the centre of the plasma column tend to absorb relatively more power from the launched wave compared with the other nuclei.

For $\omega/\bar{\omega}_{Dh} < 1$, $\lambda_K(\omega)$ becomes simply

$$\lambda_K(\omega) \approx \lambda_K(0) + \lambda_K' \frac{\omega}{\bar{\omega}_{Dh}} + \lambda_{K,res} + 0\left(\frac{\omega}{\bar{\omega}_{Dh}}\right)^2 \qquad (14)$$

where $\lambda_K(0) \equiv \lambda_K^{ad} + \lambda_K^{nad}(0)$ and $\lambda_{K,res}$ represents the contribution of the resonant nuclei.

We observe [21] that $|\lambda_K(0)/\lambda_K^{ad}| \lesssim 1 - q_o$. In this Letter, we treat $\lambda_K(0)$ as a correction to λ_H. When the simple analytic form of F_{oh} given in Eq. (13) is used, the coefficient of the linear term becomes $\lambda_K' = -(\pi/3s_o)\epsilon_o\beta_{ph}$ which corresponds to the expression for H entering Eq. (2).

We note that a distribution function which includes passing particles, as would be the case during neutral beam injection, yields an expression for λ_K similar to (12), where λ_K^{ad} is contributed, with opposite signs, by both passing and trapped particles, and λ_K^{nad} by trapped particles only. A near cancellation of λ_K^{ad} occurs for $p_{h\perp} \sim p_{h\parallel}$. As a result, at constant energy content of the fast nuclei, λ_K is smaller. This can explain why higher power levels are required to suppress the sawtooth oscillations with quasi-tangential beam injection.

In adopting the expansion (14) for λ_K, we have disregarded the possibility that modes with $\omega/\bar{\omega}_{Dh} \gtrsim 1$ become significant [16]. The reason for this is that the actual distribution function F_{oh} is expected, unlike expression (13), to have a spread in the pitch angle variable Λ. Preliminary results from our numerical analysis [7] indicate that the threshold value of β_{ph} for the onset of modes in this high frequency range is significantly larger than the value required to stabilise modes with $\omega < \bar{\omega}_{Dh}$. It was the discovery of this "macroscopic stability window" which induced us to first report [7,8] the possibility that sawtooth oscillations are suppressed by the presence of energetic trapped particles.

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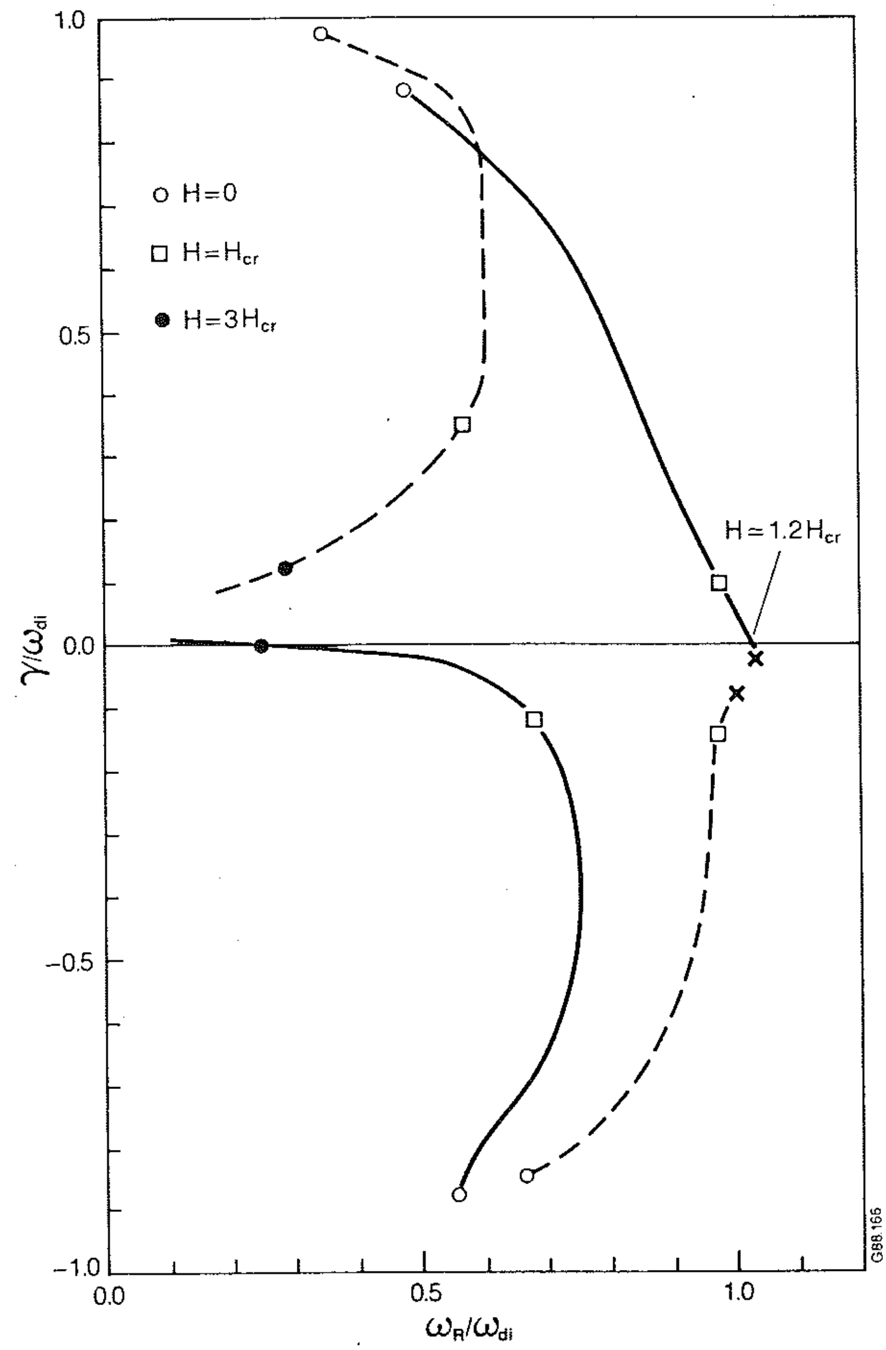

Fig. 1 Path in the complex plane of the two relevant roots of the dispersion relation (6), for increasing H and for the following set of parameters: $\gamma_{mhd}=\omega_{di}=2\hat{\omega}_{*e}/3=\bar{\omega}_{Dh}/3$ (all curves); $\omega_{di}/\omega_A=4.5\epsilon_n^{1/3}$ (solid curve); $\omega_{di}/\omega_A=1.5\epsilon_n^{1/3}$ (dashed curves). The solid curves show complete stabilisation for values of $1.2H_{cr}\leq H\leq 3H_{cr}$. The roots with higher values of ω_R violate the condition $\gamma_{mhd}-H\omega_R > 0$ (needed for their spatial regularity) at the indicated x-points.