HIGH BETA, SAWTOOTH-FREE TOKAMAK OPERATION
USING ENERGETIC TRAPPED PARTICLES

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Abstract

It is shown that a population of high energy trapped particles, such as that produced by ion cyclotron heating in tokamaks, can result in a plasma completely stable to both sawtooth oscillations and the fishbone mode. The stable window of operation increases in size with plasma temperature and with trapped particle energy, and provides a means of obtaining a stable plasma with high current and high beta.

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28

A sawtooth-free domain of tokamak operation, also stable to fishbone oscillations, could prove very desirable for reactor operation, permitting a stable equilibrium with the safety factor q significantly below one on axis. In this work we find conditions such that a high energy trapped particle population completely stabilizes the resistive internal kink mode without exciting the fishbone mode.

Experiments on JET¹ have perhaps already entered this domain, but the high energy population is maintained for too short a time to be certain.

Recently, the effect of an energetic trapped particle population on magnetohydrodynamic modes in a tokamak has been explored with the use of a variational formalism.²⁻⁶ The plasma is considered to consist of two components, a warm background treated with resistive magnetohydrodynamics, and a lower density hot component, treated with a gyrokinetic formalism. The usual branch of the ideal internal kink,⁷ unstable for plasma β (the ratio of plasma pressure to magnetic pressure) greater than a threshold value, is stabilized by a trapped particle population as long as the average toroidal precession rate of the particles is greater than the mode growth rate.

On the other hand, it was found that for β near the internal kink threshold value the trapped particles resonantly destabilize a second branch of the internal kink mode, with a real frequency given by the average precession frequency of the particle distribution, provided that the trapped particle beta, β_h , exceeded a threshold value. This branch is responsible for the fishbone oscillation. The dispersion relation describing both branches of this mode was generalized to include resistive effects in Ref. 6. The dispersion relation takes the form

$$\delta W_c + \delta W_k - \frac{8i \Gamma((\Lambda^{3/2}+5)/4) (\omega(\omega-\omega_{*1}))^{1/2}}{\Lambda^{9/4} \Gamma((\Lambda^{3/2}-1)/4) \omega_A} = 0, \quad (1)$$

where $\Lambda = -i[\omega(\omega - \hat{\omega}_{*e})(\omega - \omega_{*1})]^{1/3}/\gamma_R$, $\gamma_R = S^{-1/3}\omega_A$ is the resistive growth rate, S is the magnetic Reynolds number, ω_A is the shear Alfvén frequency, $\omega_A = v_A/(\sqrt{3}Rq')$ with v_A the Alfvén velocity, R and r the major and minor radii, respectively, and $q' = dq/dr$ with q the safety factor. The ω_* terms are diamagnetic frequencies of the background plasma with $\omega_{*1} = -(c/neBr)(dp_i/dr)$, $\omega_{*e} = (c/neBr)(dp_e/dr)$, and $\hat{\omega}_{*e} = \omega_{*e} + 0.71(c/eBr)(dT_e/dr)$. The term in Eq. (1) involving the Γ functions arises from the inertial layer, so all expressions are evaluated at the $q = 1$ surface. The inclusion of the diamagnetic terms was carried out by Bussac *et al.*,⁸ and Ara *et al.*,⁹ generalizing the work of Coppi *et al.*¹⁰ The expression δW_0 is the minimized ideal variational energy for the internal kink, first calculated by Bussac *et al.*,⁷ and δW_k is the kinetic contribution coming from the trapped particle distribution F ,

$$\delta W_k = \frac{2^{3/2}}{B^2} \pi^2 \left[\int d(\alpha B) \int \frac{dTE^{5/2} K_2^2 \omega(\partial/\partial E + \hat{\omega}_*/\omega_d) F}{K_b(\omega_d - \omega)} \right] \quad (2)$$

with $[y] = (2 \int y r dr)/r_s^2$, r_s the $q = 1$ radius, $\alpha = v_{\perp}^2/v^2$, $\hat{\omega}_*$ a differential operator associated with the hot particle diamagnetic drift frequency, and K_2 and K_b are elliptic functions arising from bounce averaging. The integrand contains a resonance at a frequency equal to the trapped particle toroidal precession frequency, given for deeply trapped particles approximately by $\omega_d = Eq/(mrR\omega_0)$ with E the particle energy, m the mass, and ω_0 the gyro frequency. Details of the derivation of this expression are found in Refs. 3 and 5.

It is instructive to consider the ideal limit, $\Lambda \rightarrow \infty$, but to keep the effects of ω_{*1} , using a model slowing-down distribution for the hot trapped

particles

$$F(E, \mu) = n(r) \delta\left(\frac{\mu}{E} - \alpha\right) E^{-3/2} \quad E < E_m \quad (3)$$

with μ the magnetic moment, $dn/dr = 1/\epsilon$, and E the energy. The parameter $\epsilon = a/R$ is the inverse aspect ratio. Equation (1) then becomes

$$\delta W_c = \frac{1[\omega(\omega - \omega_{*1})]^{1/2}}{\omega_A} + \frac{\beta_h \omega}{\epsilon \omega_{dm}} \ln\left(1 - \frac{\omega_{dm}}{\omega}\right) = 0, \quad (4)$$

where ω_{dm} is the precession rate of a particle with $E = E_m$. The condition for threshold ($\omega = \text{real}$) gives for $\omega_{*1} < \omega < \omega_{dm}$ an equation for the threshold frequency

$$\gamma_I = \frac{[\omega(\omega - \omega_{*1})]^{1/2}}{\pi} \ln\left(\frac{\omega_{dm}}{\omega} - 1\right) \quad (5)$$

with $\gamma_I = -\delta W_c / \omega_A$ the ideal kink mode growth rate, and the associated value of beta is $\beta_h = \epsilon \omega_{dm} (1 - \omega_{*1}/\omega)^{1/2} / (\pi \omega_A)$, a monotone increasing function of ω . The ideal growth rate is related to the background plasma. For example, for a quadratic q profile and circular flux surfaces γ_I was found by Bussac et al.⁷ to have the form

$$\gamma_I = \frac{\omega_A r_s^2 3\pi}{R^2} \left(1 - q(0)\right) \left(\beta_p^2 - \frac{13}{144}\right), \quad (6)$$

where β_p is the poloidal beta of the background plasma. In general γ_I depends on the full q -profile and plasma shape.

The right-hand side of Eq. (5) is zero for $\omega = \omega_{*1}$ and $\omega = \omega_{dm}/2$, and is positive and real for $\omega_{*1} < \omega < \omega_{dm}/2$. For γ_I small and positive there are

two solutions, the lower frequency one corresponding to the stabilization of the kink mode branch and the higher frequency one to the destabilization of the fishbone. We are interested in the case in which these threshold frequencies, and associated values of β_h , are widely separated, so we consider $\omega_{*1} \ll \omega_{dm}/2$. This condition puts a lower bound on the trapped particle energy. If also $\gamma_I \ll \omega_{dm}/(2\pi)$, the threshold frequencies are given approximately by $\omega_1 = \omega_{*1} + \pi^2 \gamma_I^2 / [\omega_{*1} \ln^2(\omega_{dm}/\omega_{*1})]$ and $\omega_2 = \omega_{dm}/2 - \frac{\pi}{2} \gamma_I$, and the corresponding values of β_h indicate that a stable gap can exist between the two modes. The stabilization of the ideal internal kink, which has also been pointed out by Pegoraro et al.,¹¹ occurs for $\beta_h = \epsilon \omega_{dm} \gamma_I / (\omega_A \omega_{*1} \ln(\omega_{dm}/\omega_{*1}))$. For $\omega_{*1} = 0$ the stabilization occurs at $\beta_h = \epsilon \omega_{dm} / (\pi \omega_A)$, as pointed out in Ref. 5.

Now consider resistive modification of these results. For the gap to continue to exist it is necessary and sufficient to require that the arguments of the Γ functions be large (ideal limit) at the threshold locations ω_1, ω_2 . If the arguments of the Γ functions are small, it is easy to show that the growth rate tends asymptotically to zero for $\beta_h \rightarrow \infty$, so that complete stabilization does not occur. Using ω_1 the condition $\Lambda \gg 1$ gives

$$S \gg \frac{16}{\pi^2} \frac{\ln^2(\omega_{dm}/\omega_{*1}) \omega_A^3}{\gamma_I^2 (\omega_{*1} - \omega_{*e})} \quad (7)$$

Since $\omega_2 \gg \omega_1$, this condition automatically insures that $\Lambda \gg 1$ at $\omega = \omega_2$. It is desirable to achieve a stable gap for S not too large. From Eq. (7) we see that this requires taking γ_I large, consistent with there being two solutions to Eq. (5), i.e., γ_I must remain small compared to $\omega_{dm}/(4\pi)$. The stable domain was not observed in a previous calculation¹² because δW_c was taken to be zero, i.e., $\gamma_I = 0$, which from Eq. (7) gives $S = \infty$.

We can now write necessary and sufficient conditions for the existence of a range of hot particle density between the sawtooth stabilization value and that for fishbone destabilization for $\gamma_I > 0$. The limits on the ideal growth rate γ_I are given by Eq. (7) and by the maximum value permitting a solution to Eq. (5). This gives

$$\frac{4\omega_A}{\pi} \left(\frac{\omega_A}{S(\omega_{*1} - \omega_{*e})} \right)^{1/2} \ln \left(\frac{\omega_{dm}}{\omega_{*1}} \right) < \gamma_I \leq 0.09\omega_{dm} \left(1 - \frac{2.4\omega_{*1}}{\omega_{dm}} \right) \quad (8)$$

It seems perhaps strange that to achieve stabilization using trapped particles the kink mode must be above its ideal threshold ($\gamma_I > 0$), but this is understandable in that it is precisely this instability which preserves the ideal character of the mode. In the resistive limit the mode cannot be stabilized, and for small values of S the gap described by Eq. (8) vanishes. The upper bound of Eq. (8) is the point at which the kink branch and the fishbone branch coalesce in the complex ω plane. For γ_I larger than this there are no thresholds for any value of β_n . The kink and fishbone branches exchange roles; the kink branch is destabilized by the trapped particles and takes on a large real frequency, and the fishbone branch is stable.

The range of trapped particle beta in which both modes are stable is given approximately by

$$\frac{\omega_{dm}\gamma_I}{\omega_A\omega_{*1}\ln(\omega_{dm}/\omega_{*1})} < \frac{\beta_n}{\epsilon} < \frac{\omega_{dm}}{\pi\omega_A} \left(1 - \frac{\omega_{*1}}{\omega_{dm}} - \frac{\pi\gamma_I\omega_{*1}}{2\omega_{dm}} \right) \quad (9)$$

Since ω_{dm} is proportional to the particle energy, Eq. (9) is seen to be a condition on the density of the hot trapped species. The upper limit in Eq. (9) is the fishbone threshold, and the lower limit is the internal kink stabilization point. The stability domain in the β_n, γ_I plane is

approximately triangular. Outside this triangle one or both of the branches are unstable. If γ_I is below the lower limit in Eq. (8) and β_H/ϵ above the upper limit of Eq. (9), sawtooth and fishbones can occur simultaneously.

For $\gamma_I < 0$ only the fishbone branch is ideally unstable, so below the fishbone threshold instability requires $\Lambda \leq 1$. Numerical and analytical examination shows that there is a low frequency ($\omega \leq \omega_{*1}$) unstable mode in the range

$$|\gamma_I| \leq \frac{4}{\pi} \ln\left(\frac{\omega_{dm}}{\omega_{*1}}\right) \frac{\omega_A^{3/2}}{S^{1/2}(\omega_{*1} - \omega_{*e})^{1/2}} \quad (10)$$

for both signs of γ_I . For very large S the mode is purely growing with growth rate $\gamma = \gamma_R^3/|\omega_{*1}\omega_{*e}|$. This mode cannot be completely stabilized by energetic trapped particles, although the growth rate can be significantly reduced.¹² For γ_I even more negative there is a second domain of stability. Without energetic trapped particles the domain corresponds to region 6 of Ref. 8, and energetic particles at densities below the fishbone threshold do not destabilize any branch of the dispersion relation in this domain.

A few points are worth noting. First, the $\gamma_I > 0$ window of operation is accessible with present-day tokamaks. The condition $\omega_{dm} \gg 2\omega_{*1}$ requires a fairly high energy trapped particle population, but the 1 MeV particles produced in JET easily satisfy this requirement. Beam injection experiments on PDX and other machines involved particle energies not sufficiently high to observe this effect. The internal kink stability threshold is also readily accessible. Second, the lower bound of this stability domain decreases with increasing temperature, giving a stable window suitable for reactor parameters. For example, the energy of the trapped particles can be chosen so that the upper bound on γ_I approximates the Troyon limit. Third, the effect

of accidental excursions to values of background plasma β or trapped particle density which are too large would be to produce, respectively, a lowering of the background β through increased transport inside the $q=1$ surface caused by a sawtooth cycle, or an ejection of part of the trapped particle population through a fishbone burst. Further, if a sawtooth cycle were desirable (e.g., for impurity ejection), it could be produced by a temporary lowering of the trapped particle density. The stable domain with $\gamma_I < 0$ is also interesting, although it requires lower plasma beta, or perhaps particular pressure or current profiles which give γ_I sufficiently negative at high beta.

Finally we note that although we have used a model particle distribution function in this analysis, previous experience with the fishbone mode⁵ indicates that the results are rather insensitive to the exact form of the distribution function, although some dependence exists and profile shaping could perhaps be utilized to increase the stability domain. A naturally occurring stabilization of the sawtooth by fusion-product alpha particles is a possibility.

Results of a numerical solution of Eq. (1) are shown in Figs. 1 and 2 for a hydrogen minority species in JET. A root finding procedure using a Stokes plot¹³ was used to follow unstable roots. Shown in Fig. 1 are the complex frequencies, and the growth rates, as a function of trapped particle density for an approximate JET equilibrium with $R = 296$ cm, for two different values of the magnetic Reynolds number, $S = 10^6$ and 10^7 , and $\gamma_I = 1.4 \times 10^4 \text{ sec}^{-1}$. Both the resistive internal kink branch and the fishbone branch are shown. The trapped particle density n ranges from zero to $2.3 \times 10^{11} / \text{cm}^3$. The particles were taken to be a slowing-down distribution with an average energy of 700keV. The toroidal field was $B = 24$ kG and the averaged trapped particle precession rate was $\langle \omega_d \rangle = 1.3 \times 10^5 / \text{sec}$. The shear Alfvén frequency was $\omega_A =$

$2 \times 10^6/\text{sec}$, and the diamagnetic frequencies were $\hat{\omega}_{*e} = -3 \times 10^4/\text{sec}$ and $\omega_{*i} = 2 \times 10^4/\text{sec}$.

As seen, the $S=10^7$ case is almost ideal, with $\gamma \approx \gamma_I$ and $\omega_r \approx \omega_{*i}/2$ at $n=0$. There is a range of trapped particle density in which both the kink mode branch and the fishbone branch are stable. For smaller values of S the resistive modification of the kink branch makes it more unstable and eliminates this stable gap. For $S = 10^6$, the decrease of the growth rate with increasing hot particle density, results in values of $\Lambda \approx 1$. The behavior for large β_h is then dominated by the small argument of the r functions, and the growth rate only asymptotically approaches zero for $\beta_h \rightarrow \infty$, as found previously,¹² and complete stabilization does not occur. The fishbone branch is not noticeably modified by resistive effects, because although the growth rates of the two branches are comparable, the real frequency of the fishbone branch is much larger, giving $\Lambda \gg 1$. In Fig. 2 are shown the stability domains for these same parameters. Outside of the stable domains one would expect unstable modes of low (L) or high (H) frequency or both (H+L). Only the small $|\gamma_I|$ limits depend on the Reynolds number, and two values are shown.

In conclusion, we find that for typical JET parameters there exist domains of hot trapped particle density in which both the sawtooth mode and the fishbone mode are stable. One stable domain requires that the plasma be ideally unstable to the internal kink mode ($\gamma_I > 0$), i.e., operation at fairly high plasma beta, and that the trapped particle population have a sufficiently high energy ($\omega_d \gg \omega_{*i}$). This domain of stable operation appears to be suitable for use in a fusion reactor.

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FIGURE CAPTIONS

Fig. 1 The complex frequency plane, showing kink and fishbone branches for a range of hot trapped particle density.

Fig. 2. The stable domain in the γ_I , n plane, for JET parameters.

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Fig. 1

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Fig. 2

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