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SAWTOOTH STABILIZATION BY ENERGETIC TRAPPED PARTICLES

By

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Abstract

Recent experiments involving high power radio-frequency heating of a tokamak plasma show strong suppression of the sawtooth oscillation. A high energy trapped particle population is shown to have a strong stabilizing effect on the internal resistive kink mode. Numerical calculations are in reasonable agreement with experiment.

Experiments on JET using high power ICRF on-axis heating of a minority ion species[1] show efficient electron heating with the occurrence of long sawtooth-free periods of up to 1.6 sec. This giant sawtooth regime is accompanied by an increase in energy confinement and in density. The experiments involved heating of either He^3 or H in a background plasma of D or He^4. In this letter we show that a high energy trapped ion population such as that produced by the heating in JET can significantly decrease the growth rate of the resistive internal kink mode, leading to an increase of the sawtooth period.

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 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

$$\delta W_c + \delta W_k + \frac{8S^{-1/3}\Lambda^{-9/4}\Gamma\left((\Lambda^{3/2}+5)/4\right)}{\Gamma\left((\Lambda^{3/2}-1)/4\right)}\left[\Omega\left(\Omega+i\,\frac{\omega_{*i}}{\omega_R}\right)\right]^{1/2} = 0 \quad , \qquad (1)$$

where $\Omega = -i\omega/\omega_R$, $\Lambda = \left[\Omega(\Omega + i\,\hat{\omega}_{*e}/\omega_R)(\Omega + i\,\omega_{*i}/\omega_R)\right]^{1/3}$, $\omega_R = S^{-1/3}\omega_A$ is the resistive frequency, S is the magnetic Reynolds number, ω_A is the shear Alfvén frequency

$$\omega_A = \frac{v_A}{\sqrt{3}\ Rrq'} \tag{2}$$

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 with $\omega_{*i} = -(c/neBr)(dp_i/dr)$, $\omega_{*e} = (c/neBr)(dp_e/dr)$, 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_c 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,

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

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 diamagnetic drift frequency, and K_2 and K_b are elliptic functions arising from bounce averaging. Details of the derivation of this expression are found in Refs. 3 and 5. The trapped particle population contributes to the dispersion relation only in the domain $q < 1$ in this formalism, which involves an expansion to lowest order in the inverse aspect ratio r/R. This is because the minimizing eigenfunction in this approximation is the usual internal kink step function, consisting of a single harmonic with poloidal and toroidal mode numbers equal to one. A more complete calculation, with several poloidal harmonics, would include significant contribution from the entire trapped particle population, also outside the q = 1 surface. Numerical codes for the solution of the resulting dispersion relation for arbitrarily shaped cross-sectional geometry are being developed,[11] but until they allow a detailed analysis of the full trapped particle contribution, comparison with experiment must be only approximate. Equation (1), without the diamagnetic terms, was used in Ref. 6 to calculate resistive modification of the fishbone threshold, given approximately by $\beta_h \approx \langle\omega_d\rangle/(\pi\omega_A)$ in the ideal limit, where $\langle\omega_d\rangle$ is the average precession frequency of the trapped particle population.

In the present work we examine the effect of a trapped particle population on the resistive internal kink branch, which is purely growing in the absence of kinetic effects. Without a trapped particle population, $\delta W_k=0$, this branch of the solution to Eq. (1) possesses two well-known limits. We neglect the diamagnetic frequencies ω_{*e} and ω_{*i} for this discussion. They have the effect of decreasing the growth rate and giving the mode a real frequency; but, they complicate the algebra, and the qualitative behavior of the solution can be understood without them. The numerical results presented include the diamagnetic terms.

For large S (high temperatures) the large argument limit of the Γ functions can be used to obtain

$$0 = \delta W_c - i\omega/\omega_A \quad , \tag{4}$$

giving a purely growing mode for $\delta W_c < 0$, the ideal internal kink mode. Note that this limit requires $\Omega \gg 1$, or $|\delta W_c|S^{1/3} \gg 1$. In the second limit, with

$|\delta W_c| S^{1/3} << 1$, the dispersion relation is satisfied for $\Gamma((\Omega^{3/2}-1)/4) = \infty$, or $\omega = i\omega_R$. This is the usual m = 1 tearing mode, which is of interest to us in the present work, and in the following we will take $\delta W_c = 0$. The stabilization of the ideal internal kink by trapped particles was noted in Ref. 5.

The behavior of the tearing mode branch of the dispersion relation can be found by using a model slowing-down trapped particle distribution, which permits analytic evaluation of the kinetic contribution. Introducing the distribution

$$F(E,\mu) = \delta(\mu-\mu_o)/E \qquad E < E_m \tag{5}$$

with μ the magnetic moment and E the energy, the kinetic contribution becomes[3]

$$\delta W_k = i\beta_h \Omega A \ln(1+i/(A\Omega)) \tag{6}$$

with β_h the hot trapped particle beta, and $A = \omega_R/\omega_{dm}$ the ratio of the resistive growth rate to the maximum toroidal precession rate. The precession rate ω_{dm} is approximately given by the value for deeply trapped particles,

$$\omega_{dm} = E_m q/(mrR\omega_0) \tag{7}$$

with ω_0 the gyrofrequency, $\omega_o = ZeB/(mc)$, and for this distribution the average precession rate is given by $\langle\omega_d\rangle = \omega_{dm}/2$. Thus the dispersion relation for the tearing mode becomes

$$8\Gamma((\Omega^{3/2}+5)/4)/\Gamma((\Omega^{3/2}-1)/4) + i\beta_h S^{1/3}\Omega^{9/4} A \ln(1+i/(A\Omega)) = 0 \quad . \tag{8}$$

The trapped particles play an important role only for

$$\beta_h S^{1/3} A > 1 \quad , \tag{9}$$

which is, within a factor of π, the fishbone threshold condition. Note that this is independent of the Reynolds number S. The combination $\beta_h A$ is proportional to the hot trapped particle density and charge, and independent of its mass and energy. Thus, as with the fishbone mode, it is the trapped particle density which is the important parameter in determining the nature of the solutions to the dispersion relation. Two limits are of interest. For small $\beta_h AS^{1/3}$ the solution can be obtained by expanding about $\Omega = 1$, which is the solution for $\beta_h = 0$, giving

$$\Omega = 1 - 2iA \ln(1+i/A)\beta_h S^{1/3}/(3\sqrt{\pi}) \quad . \tag{10}$$

The trapped particles are initially destabilizing, and more so for $A \approx 1$, i.e., when the precession frequency is resonant with the resistive rate. For large $A\beta_h S^{1/3}$ there are two roots, with $\Omega \to \infty$ and $\Omega \to 0$. For the small root, letting $\alpha = 1/|\Omega|$ we find an equation determining α

$$\alpha^{9/4}\ln(\alpha) = \Gamma(3/4)A\beta_h S^{1/3}/(2\Gamma(5/4)) \tag{11}$$

with $\Omega = (1/\alpha)e^{-i2\pi/9}$. This gives for the mode $\omega = (\omega_R/\alpha)\, e^{i5\pi/18}$, and it is strongly stabilized for $A\beta_h S^{1/3} >> 1$. This asymptotic value depends only on the trapped particle density, and is independent of energy. Note, however, that the analysis leading to the dispersion relation, Eq. (1), assumed that the temperature of the hot species is of order $(R/r)^2$ greater than that of the background plasma, i.e., trapped ions of at least 50 keV for the JET discharges.

The second root can be found by using asymptotic expressions for the Γ functions, giving

$$\Omega = \beta_h S^{1/3} \quad . \tag{12}$$

As long as $\langle\omega_d\rangle > \omega_R$, the first root, Eq. (11), is the tearing mode root, and Eq. (12) describes the fishbone mode. If instead $\langle\omega_d\rangle < \omega_R$, the trapped particles destabilize the tearing mode and its asymptotic limit is given by Eq. (12).

Thus we find that to stabilize the tearing mode it is necessary to produce a trapped particle density which is large enough to destabilize the fishbone. The occurrence of the fishbone is, however, also contingent upon proximity to the internal kink threshold.[3] Depending on plasma β and other equilibrium parameters, we must, in general, expect either sawtooth stabilization or fishbone oscillations, with fishbone oscillations dominating for high β operation. The two could exist together only if the fishbone failed to eject the trapped particles.

We have examined the solutions to Eq. (1) using a numerical code developed for the investigation of the fishbone mode.[5] Numerical values of temperatures, densities, diamagnetic frequencies, etc., were chosen to approximate the experiments done on JET. The kinetic contribution δW_k is generated by a Monte Carlo procedure. For minority species ion cyclotron heating experiments the hot trapped particle distribution is well approximated by

$$F(E,\mu,r) = n(r)e^{-E/T}\delta(\mu/E-\alpha) \quad . \tag{13}$$

The value of n is typically of the order of a few percent of the average plasma density, and the temperature T ranges from 70 to 150 keV depending on the ICRF power. For the experiments performed on JET using on-axis heating the approximation of a single value of μ/E corresponds to a single value of the bounce angle, $\theta_b \approx \pi/2$. In the present work, as well as in the case of the fishbone mode, the details of the solution are not particularly sensitive to small changes in the particle distribution function.

Results of a Monte Carlo simulation are shown in Figs. 1 and 2 for a hydrogen minority species in JET. Shown in Fig. 1 are the trajectories in the complex frequency plane for an approximate JET equilibrium with R = 296 cm, for two different values of the magnetic Reynolds number, $S = 3 \times 10^6$ and 10^7. The trapped particle density ranges from zero to $10^{12}/cm^3$. The particle distributions were of the form given by Eq. (13), and had temperatures of 50 and 150 keV. The toroidal field was B = 24 kG and the average trapped particle precession rates were $\langle\omega_d\rangle \approx 2 \times 10^4$/sec and 6×10^4/sec, respectively. The shear Alfvén frequency was $\omega_A = 2 \times 10^6$/sec, and the diamagnetic fequencies were $\hat{\omega}_{*e} = - 3 \times 10^4$/sec and $\omega_{*i} = 2 \times 10^4$/sec. Without the trapped particles, the mode is almost purely growing but with a growth rate almost an order of magnitude smaller than the resistive value ω_R, a well-known effect of the diamagnetic terms. The trapped particle population is destabilizing for small density if $A \approx 1$, becomes stabilizing if $A\beta_h S^{1/3} > 1$ provided that $\langle\omega_d\rangle > \omega_R$, and for large density the stabilization is more effective than predicted by Eq. (11). We find this to be true for a wide range of values of ω_{*e}, ω_{*i}. A more complete study of the dependence of this effect on the diamagnetic frequencies will be reported in a separate publication.

For $S = 3 \times 10^6$ and 50 keV, $\langle\omega_d\rangle \approx \omega_R$ and the mode is barely stabilized. If either S or the energy is decreased, the mode is destabilized by the particles. Increasing the particle energy above 150 keV does not significantly change the results, which are independent of energy for $\langle\omega_d\rangle \gg \omega_R$. In Fig. 2 is shown the mode growth rate as a function of trapped particle density for the trajectories of Fig. 1. The growth rate is decreased by a significant amount for values used in JET, where trapped ion densities are estimated to be greater than $10^{12}/cm^3$. From these curves it is clear that a significant lengthening of the sawtooth period is possible, assuming that the sawtooth period scales as the inverse of the linear growth rate. We note that the trapped particles also stabilize the ideal internal kink,[5] and thus the qualitative features of this result are unchanged if a short time scale ideal model of the sawtooth is used.

The nature of this mechanism makes it clear that the most effective stabilization should be achieved using heating on-axis, where the density is largest and the biggest trapped particle population can be expected to be obtained.

Two different antenna settings were used in the JET experiments. It was observed that an in-phase antenna setting produced giant sawteeth with either a hydrogen or He^3 minority species, but an out-of-phase antenna setting produced giant sawteeth only with a He^3 minority. This is understandable in that the factor A is proportional to the charge Z. Thus for equal density and energy, a He^3 trapped ion population is more stabilizing than one consisting of hydrogen.

During neutral beam heating of JET at lower toroidal field values (21 kG) and higher beta, magnetic signals associated with the fishbone mode were observed.[12] The threshold value of β_h for the fishbone was calculated[13] to be 2×10^{-3}. Because of the efficient ejection of trapped ions by the fishbone, the trapped particle population would be effectively limited to approximately this value, which corresponds, for the neutral beam energy of 70 keV, to a trapped particle density of $3 \times 10^{11}/cm^3$, too low a value to produce stabilization of the sawtooth. In general, since the stabilization condition implies that the fishbone threshold be exceeded, it appears, perhaps unfortunately, that at high beta operation the occurrence of the fishbone will not allow sawtooth stabilization with trapped particles. To our knowledge, fishbone oscillations have not yet been produced during ICRF heating, so the nature of the transition from sawtooth stabilization to fishbone is still unexplored experimentally. The signature for this transition would be the occurrence of ion bursts with an associated drop in neutron production during the sawtooth period.

For the same reason it is understandable that during high β operation using neutral beam injection when PDX was strongly unstable to the fishbone mode, the phenomenon of giant sawteeth was not observed.

In conclusion, we find that the presence of a high energy trapped ion population introduces a stabilization of the sawtooth in a tokamak. Numerical calculations of the effect are in reasonable agreement with experiments on JET, giving almost an order of magnitude increase in the period. Although the stabilization of the sawtooth mode in low beta discharges is encouraging, at higher β the same trapped particle population should destabilize the fishbone branch, as has been observed to happen with neutral beam injection. The fishbone should be then expected to limit the trapped particle population to a value too low to provide sawtooth stabilization.

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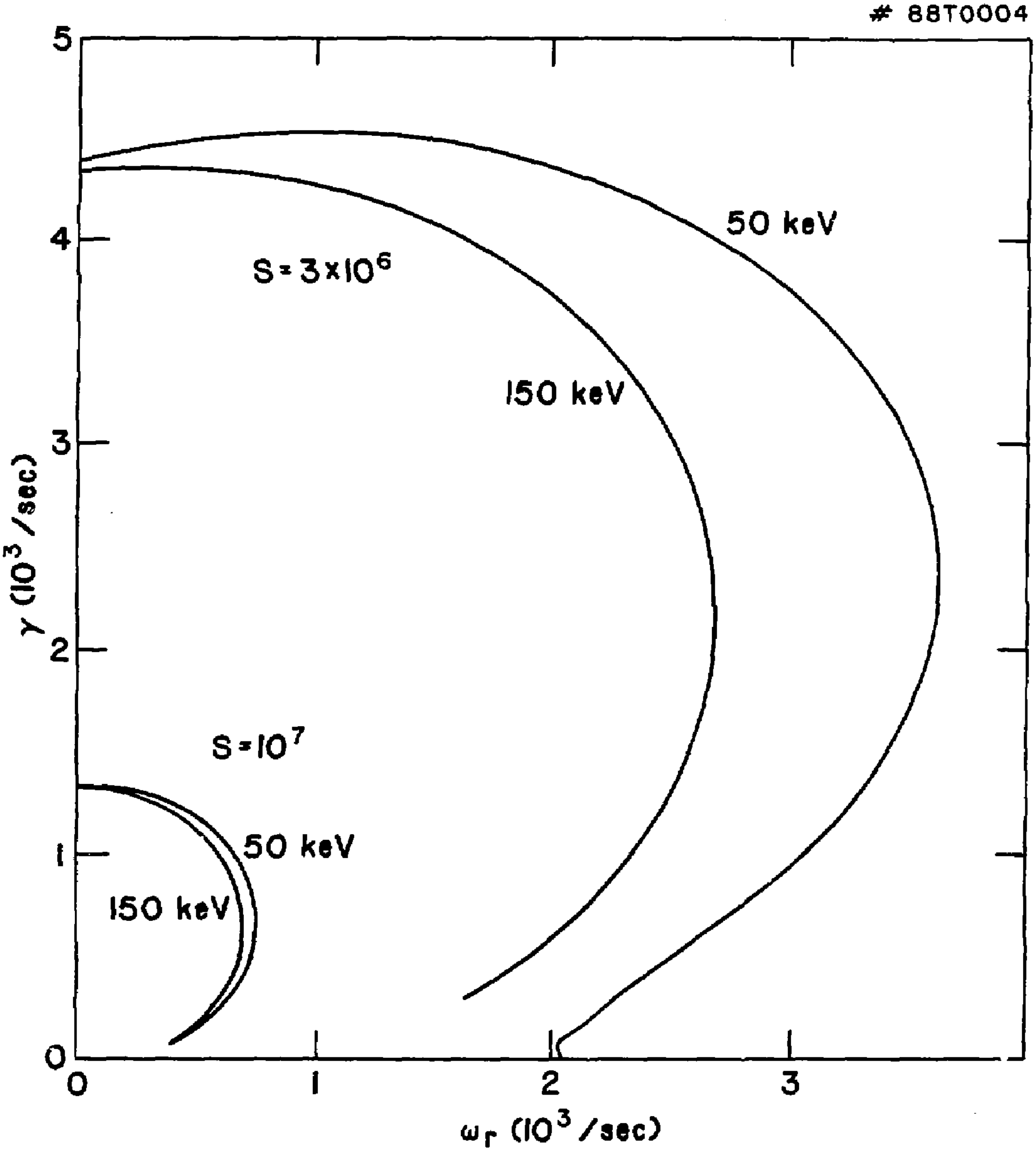
FIGURE CAPTIONS

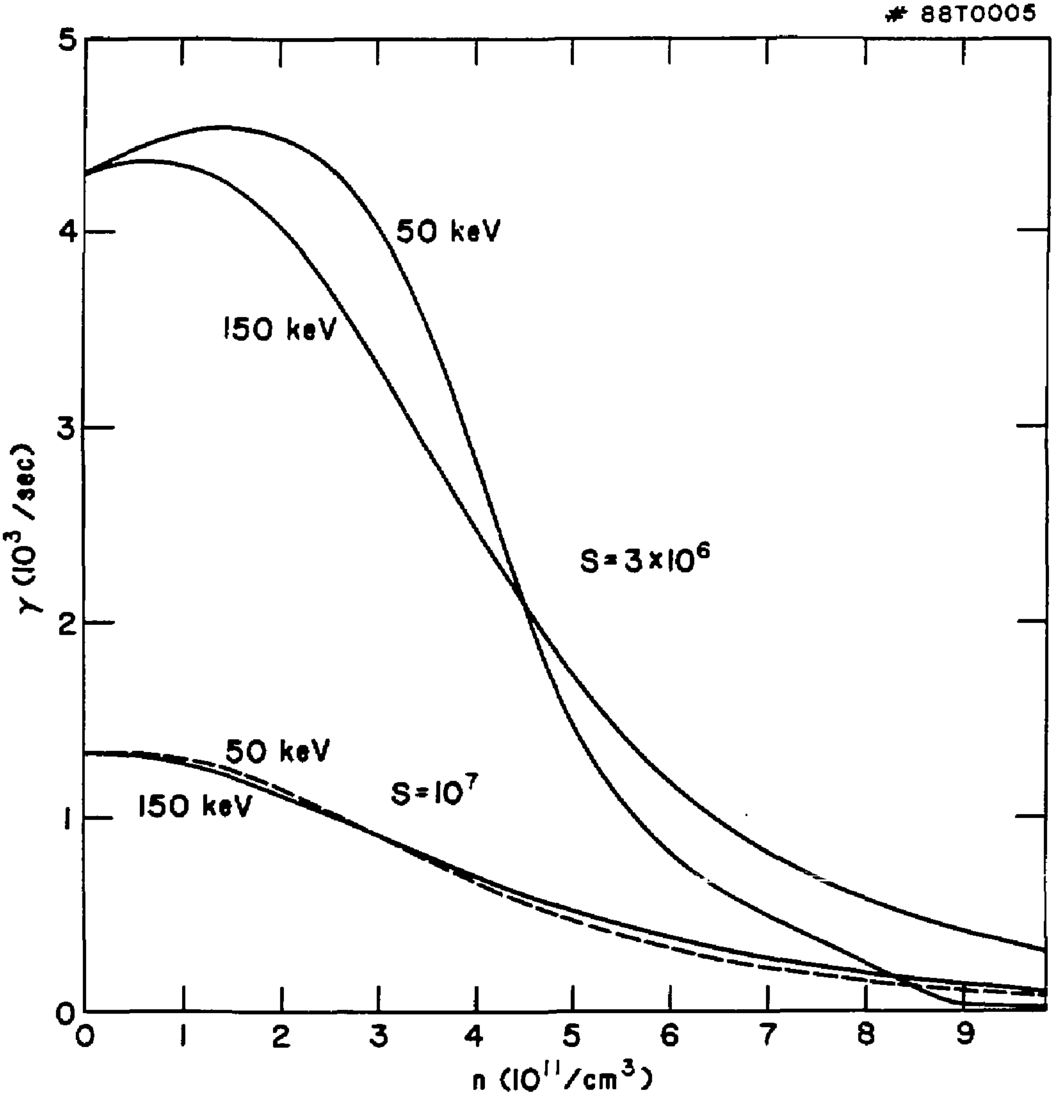
Fig. 1 The trajectories of solutions to Eq. (1) in the complex frequency plane as a function of trapped proton density, for two values of the magnetic Reynold number S, and different energies. Equilibrium parameters and Alfvén and diamagnetic frequencies were chosen to approximate JET.

Fig. 2 Growth rate as a function of trapped particle density for the trajectories of Fig. 1.



88T0004
5
4
3
2
1
0
0
1
2
3
γ (10^3 /sec)
ω_r (10^3 /sec)
50 keV
S = 3x10^6
150 keV
S = 10^7
50 keV
150 keV

Fig. 1



88T0005
50 keV
150 keV
S=3×10^6
50 keV
S=10^7
150 keV
γ (10^3 /sec)
n (10^11/cm^3)
0
1
2
3
4
5
6
7
8
9

Fig. 2

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