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ABSTRACT

The presence of energetic particles is shown to qualitatively modify the stability properties of ideal as well as resistive magnetohydrodynamic (MHD) modes in tokamaks. Specifically, we demonstrate that, consistent with high-power ICRF heating experiments in JET,[1] high energy trapped particles can effectively stabilize the sawtooth mode, providing a possible route to stable high current tokamak operation. An alternative stabilization scheme employing barely circulating energetic particles is also proposed. Finally, we present analytical and numerical studies on the excitations of high-n MHD modes via transit resonances with circulating alpha particles.

I. STABILIZATION OF SAWTOOTH OSCILLATIONS

Consider a large-aspect-ratio ($\epsilon=a/R_o < 1$) tokamak with circular magnetic surfaces. The plasma is taken to consist of a resistive MHD background plasma and a lower density, hotter (h) trapped particle component. The dispersion relation of the m=1 resistive internal kink mode has been derived variationally[2-4] and is given by

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

where $\Lambda = -i[\omega(\omega-\hat{\omega}_{*e})(\omega-\omega_{*i})]^{1/3}/\gamma_R$, $\gamma_R = S^{-1/3}\omega_A$ is the resistive growth rate, S is the magnetic Reynolds number, ω_A is the shear Alfvén frequency, $\omega_A = v_A/(\sqrt{3}R_o\, rq')$ with v_A the Alfvén velocity, and $q' = dq/dr$ with q the safety factor. The ω_* terms are diamagnetic frequencies of the background plasma 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 inclusion of the diamagnetic terms was carried out by Bussac et al.,[5] and Ara et al.,[6] generalizing the work of Coppi et al.[7] The expression $\delta W_f \sim O(\epsilon^2\beta_{pc})$ is the minimized ideal variational energy for the internal kink, first calculated by Bussac et al.,[8] giving the ideal growth rate $\gamma_I = -\omega_A\delta W_f$, and δW_k is the kinetic contribution coming from the trapped particle distribution F,

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

with $[y] = (2\int yrdr)/r_s^2$, 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 dispersion relation thus depends parametrically on the trapped particle beta β_h and six frequencies ω_A, ω_d, ω_{*e}, ω_{*i}, γ_I, ω_R, as well as on the form of the distribution function F. We have examined the solution to Eq. (1) for particle distributions chosen to model neutral beam injection and ICRF heating. Numerically generated Monte-Carlo distributions as well as simplified models were used. Qualitatively, the results are insensitive to the form of the distribution function. A model slowing-down distribution with a single bounce angle and energy dependence $F \sim E^{-3/2}$ allows analytic evaluation and suffices to display the properties of the solution. Results for Maxwellian and other distributions are similar and will be reported elsewhere. We choose parameters appropriate for JET to illustrate the solutions, but the numerical results should be taken only as an indication of qualitative features. Detailed comparison with experiment must be done using appropriate distribution functions. The full dependence of the solutions on all the parameters is too involved to discuss here. Considerations are limited to that domain which allows complete stabilization, and may be of relevance for JET and other experiments. At high temperatures there are two branches of interest. The resistive internal kink branch, responsible for the sawtooth, and which for $\beta_h=0$ and $\gamma_I > \omega_{*i}/2$ has a real frequency near $\omega_{*i}/2$ and growth rate near γ_I, and the fishbone branch. Increasing β_h decreases the growth rate of the kink branch if $\langle\omega_d\rangle > \gamma_I$, and at sufficiently high temperatures the mode is stable for β_h sufficiently large. Stabilization of the ideal mode was known previously,[3,9] but stabilization of the resistive mode occurs only if the arguments of the Euler gamma functions in Eq. (1) are large enough that the mode is essentially ideal. This condition gives, for $\gamma_I > \omega_{*i}/2$

$$\gamma_I \gg (\gamma_R/\pi)\,\ln(\langle\omega_d\rangle/\gamma_I) \quad . \tag{3}$$

Thus stabilization occurs only in a sufficiently high β plasma. On the other hand, if $\gamma_I < \omega_{*i}/2$, a small amount of resistivity can significantly change the results. The dissipative effects associated with the ion resonance are not sufficient to overcome the resistivity and the mode cannot be stabilized.

The fishbone branch is, on the other hand, destabilized by the hot particles.[2] It has a real frequency approximately equal to $\langle\omega_d\rangle$, and a threshold in β_h given approximately by $\beta_c \approx \langle\omega_d\rangle/\omega_A$.

If $\langle\omega_d\rangle \gg \omega_{*i}$, i.e., if the trapped particle population is sufficiently hot, the stabilization of the kink branch can occur for $\beta_h < \beta_c$, so that complete stabilization is achieved. There then exists a triangular domain in the γ_I, β_h plane in which both these branches are stable.[10] The stable domain is shown in Fig. 1 for an approximate JET equilibrium with R = 296 cm for two different values of the magnetic Reynolds number, $S = 10^7$ and 10^8. The particles were taken to be a slowing-down distribution with an average energy of 700 keV. The toroidal field was B=24 kG. The shear Alfvén frequency was $\omega_A = 2\times10^6$/sec, and the diamagnetic frequencies were $\hat{\omega}_{*e} = -3\times10^4$/sec and $\omega_{*i} = 2\times10^4$/sec. A triangular stable domain also exists for Maxwellian distribution functions. Although it is tempting to associate this domain with the occurrence of sawtooth-free operation on JET, it is necessary to make more detailed comparisons using noncircular tokamak equilibria, correct distribution functions, and reasonably accurate current density profiles.

Recently, Eq. (1) has also been generalized to include a highly anisotropic ($P_\perp \gg P_\parallel$) but circulating energetic component.[11] In this case, we have

$$\delta W_f \sim O(\varepsilon\beta_{ph,u}) > 0 \qquad (4)$$

and $\delta W_k = 0$ since no trapped particle is involved. Equation (1) then suggests the interesting possibility of significant stabilization by employing "barely" circulating energetic particles mainly within the q = 1 surface. Note since $\delta W_f S^{1/3} > 1$, this scheme, in addition to achieving ideal stability, could also significantly suppress the sawtooth oscillations.

II. ALPHA-PARTICLE EXCITATIONS OF HIGH-N MODES

In proposed tokamak ignition experiments such as CIT, fusion α particles typically have velocities $v_\alpha \sim v_A$ and instabilities via transit resonances between shear Alfvén waves and circulating α particles must be considered. In the n (toroidal mode number) >> 1 limit, the toroidicity-induced shear Alfvén "gap" mode,[12] which suffers little damping due to phase mixing, could therefore be readily excited.

To illustrate the physics, we consider a large-aspect-ratio tokamak with circular magnetic surfaces. The corresponding shear Alfvén eigenmode equation is then given by, in the ballooning-mode representation,

$$\left[\frac{d}{d\theta}\, p(\theta)\, \frac{d}{d\theta} + \left(\frac{\omega}{\omega_A}\right)^2 p(\theta)\, (1+2\varepsilon\cos\theta) + \Delta g_o(\theta)\right]\delta\phi \tag{5}$$

$$- 4\eta_\alpha \langle\, \omega_d \delta g\rangle_\alpha = 0 \quad ,$$

where $p(\theta)=1+\hat{s}^2\theta^2$, $\hat{s}=rq'/q$, $\Delta=\beta' q^2 R_o$, $\omega_A=v_A/qR_o$, $g_o=\cos\theta+\hat{s}\theta\sin\theta$, $\eta_\alpha = \pi\omega e_\alpha q^2 R_o^2/c^2 k_\theta^2$, $k_\theta=nq/r$, $\omega_d=\hat{\omega}_d g_o$, $\hat{\omega}_d=k_\theta\,(v_\perp^2/2+v_\parallel^2)/R_o\omega_c$, ω_c is the cyclotron frequency, $\langle\ldots\rangle \equiv\int d^3\underline{v}(\ldots)$, and δg_α satisfies

$$\left(\frac{d}{d\theta} - iW\right) \delta g_\alpha = i\bar{\beta}_\alpha g_o \delta\phi \quad , \tag{6}$$

and $w = (\omega-\omega_d)/\omega_t, \omega_t = v_\parallel/qR_o$, and

$$\hat{\beta}_\alpha = [(e/m)\ (\omega\partial/\partial E+\hat{\omega}_*)\ F_o\hat{\omega}_d/\omega\omega_t]_\alpha .$$

To proceed further analytically, we assume $\hat{s},\Delta<1$ to allow two-spatial-scale expansions. Letting $\delta\psi=(1+\hat{s}^2\theta^2)^{1/2}\delta\phi$, we shall concentrate on the toroidicity-induced shear Alfvén waves by assuming

$$\delta\psi = A(\theta_1)\cos(\theta_o/2) + B(\theta_1)\ \sin(\theta_o/2) \tag{7}$$

such that $\theta_1-1/\hat{s}>>\theta_o-1$. Equation (5) then yields the following coupled equations between $A(\theta_1)$ and $B(\theta_1)$,

$$f_+A=-B'+(i\beta_2-\Delta\hat{s}\theta/2p)B \quad , \tag{8}$$

$$f_-B = A'-(i\beta_2+\Delta\hat{s}\theta/2p)A \quad , \tag{9}$$

where

$$f_\pm = \Gamma_\pm - \hat{s}^2/p^2 \pm \Delta/2p \quad ,$$

$$\Gamma_\pm = (1\pm\varepsilon)(\omega/\omega_A)^2 - \frac{1}{4} - \beta_1 \quad , \tag{10}$$

$$\beta_1 = \langle\pi\hat{\beta}\ \hat{\omega}_d\ w_o\ [(\frac{1}{4} - w_o^2)^{-1} + (\frac{9}{4} - w_o^2)^{-1}]\rangle_\alpha \quad ,$$

$$\beta_2 = \langle\pi\hat{\beta}\hat{\omega}_d[(\frac{1}{4} - w_o^2)^{-1} - 3(\frac{9}{4} - w_o^2)^{-1}]\rangle_\alpha/2 \quad ,$$

and $w_o = \omega/\omega_t$. Equations (8) and (9) can be solved via asymptotic matching analyses yielding the following dispersion relation

$$\Gamma_+ = -\Gamma_-(\frac{\pi}{4})^2(\Delta/\hat{s} - \hat{s})^2 \quad ;$$

that is, we have approximately

$$\omega_r = \omega_A/2 \quad , \tag{11}$$

$$\omega_i/\omega_A = \text{Im}\beta_1 \sim f_\alpha k_\theta \rho_\alpha \beta_\alpha R_o/L_{p\alpha} \quad , \tag{12}$$

with $f_\alpha \sim O(1)$ being the fraction of resonant α particles and $L_{p\alpha}$ being the α-particle scale length. Noting that electron Landau damping provides the main stabilization and that $\max(k_\theta \rho_\alpha) \sim \epsilon v_A/v_\alpha$, we then have a critical β_α given as

$$\beta_{\alpha,c} \sim (v_\alpha/v_A f_\alpha)\ (\beta_e\ m_e/m_i)^{1/2} \sim O(10^{-3}) \quad . \tag{13}$$

The analytical results already mentioned are supported by numerical results from a comprehensive kinetic linear eigenfrequency eigenfunction calculation.[13] This calculation is fully electromagnetic and includes all of the relevant kinetic effects for the magnetically trapped and untrapped particles for each species, including bounce-frequency resonances for the trapped particles, transit-frequency resonances (Landau damping) for the untrapped particles, magnetic (gradient and curvature) drift frequency resonances, and full finite Larmor radius effects. The calculation employs the ballooning formalism to lowest order in 1/n, and accordingly the calculation is local to a chosen magnetic surface and involves a system of three one-dimensional integrodifferential equations along the unperturbed magnetic field lines. Here, a version of the design for the proposed CIT is considered which has (vacuum) major radius R_0=1.75 m. Radial profiles of $n_j(r)$ and $T_j(r)$ for each species and also of the safety factor q(r) and the pressure p(r) are obtained from a BALDUR transport code run for this design. In general, up to four particle species are included: j = e for background electrons, j = i for background hydrogenic ions, j = He for the thermalized helium ash, and j = α for the hot α-particles which are in the process of slowing down. Accordingly, the so-called slowing-down distribution function $F_{SD} \propto n_\alpha(r)/(v^3+v_c^3)$, where v_c is the so-called critical velocity, is used for the α-particles and the distribution functions for the other species are taken to be Maxwellian. The MHD equilibrium is calculated numerically, since the CIT device has a strongly shaped cross section, with a vertical elongation factor of about two and significant triangularity. Further details are given in Ref. 14. Numerical results are presented here for a magnetic surface with (average) minor radius r ≈ a/2, which is somewhat outside the q=1 surface. In these results, β is varied artificially from its BALDUR value, $\beta=\beta_{CIT}$=5.82%, by varying all of the $n_j \propto \beta$, at fixed n_j/n_e, T_j, B_0, and n. For each different β value, the entire pressure profile is multiplied by β/β_{CIT} and the numerical MHD equilibrium is recomputed, for a fixed q(ψ) profile.

The results of varying β in this way are shown in Fig. 2 for the kinetically calculated MHD ballooning mode (kMHDBM). Also shown is the growth rate curve for the simplest ideal MHD ballooning mode equation, without ω_{*pi} and without α-particles or helium, for reference. Including ω_{*pi} and the many other kinetic effects, but with no α-particles or helium, lowered the growth rate curves substantially to those for the kMHDBM. The corresponding kinetic β-critical, β_{ci}^{kin}, can be either raised, as for n = 19, or lowered, as for n =

58, from the simple ideal MHD β_{cl}^{MHD} = 8.25%, due to ion magnetic drift resonances and other kinetic effects. Including the α-particles for n = 19 raises the growth rate substantially for $\beta < \beta_{cl}^{MHD}$ and, in fact, lowers β_{cl}^{kin} to below 2%! The addition of helium ash for $\beta = \beta_{CIT}$ = 5.82% has only a small additional destabilizing effect.

The corresponding real frequencies ω_r are shown in Fig. 3. They are all in the ion diamagnetic direction. The real frequency for the simple ideal MHD ballooning mode without the ω_{*pi} is just $\omega_r = 0$. For the kMHDBM for n=19, adding the α-particles is seen to substantially increase $|\omega_r|$. Various characteristic frequencies are shown on the vertical axis. With the α-particles included, ω_r is of the same order as the average untrapped α-particle transit frequency $\bar{\omega}_{t\alpha}$, and is 0.5 to 0.7 times the Alfvén frequency $\omega_A \propto n_j^{-1/2} \propto \beta^{-1/2}$. This allows the strong destabilization of the mode by the α-particle transit-frequency resonances that has been mentioned.

ACKNOWLEDGMENT

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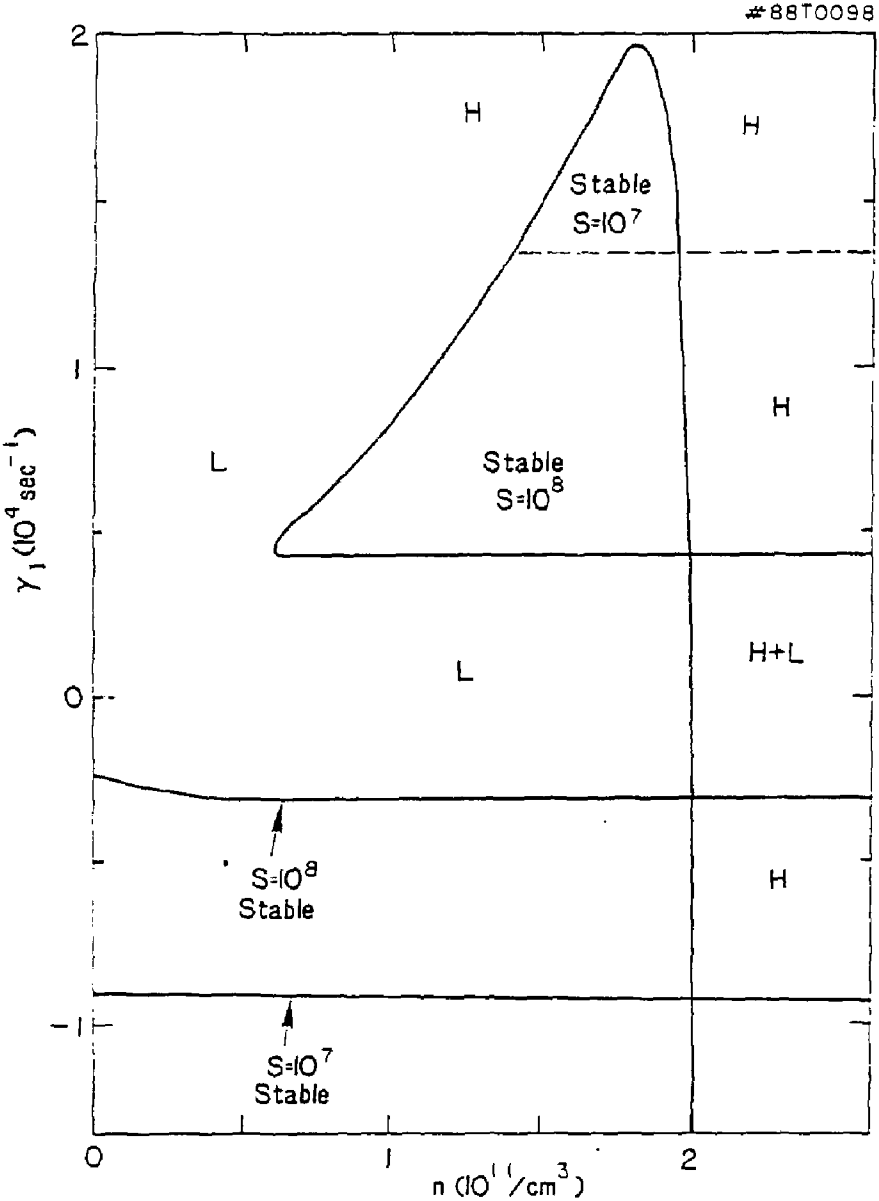
FIGURE CAPTIONS

FIG. 1 The stable domain in the γ_I, n plane, for JET parameters.

FIG. 2 Linear growth rates γ versus (local) $\beta \propto n_j$, for fixed T_j and B_o, for the CIT parameters, in units of 10^5 sec^{-1}.

FIG. 3 Real frequencies ω_r corresponding to Fig. 2. Characteristic single-particle frequencies are indicated on the vertical axis.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2
1
0
-1
γ_1 (10^4 sec^{-1})
H
H
Stable
S=10^7
H
L
Stable
S=10^8
L
H+L
S=10^8
Stable
H
S=10^7
Stable
0
1
2
n (10^{11}/cm^3)

Figure 1

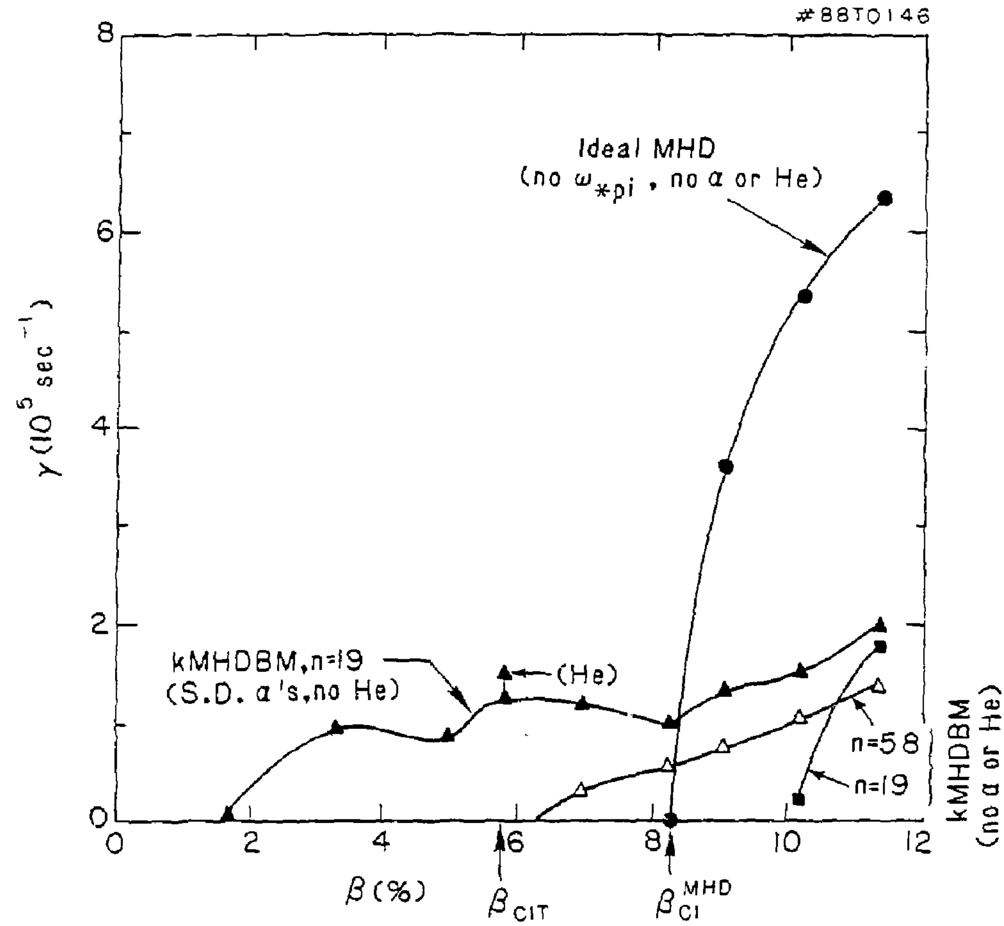
#88T0146
Ideal MHD
(no ω*pi , no α or He)
kMHDBM,n=19
(S.D. α's,no He)
(He)
n=58
n=19
kMHDBM
(no α or He)
γ (10^5 sec^-1)
β (%)
β CIT
β CI MHD
0
2
4
6
8
10
12

Figure 2

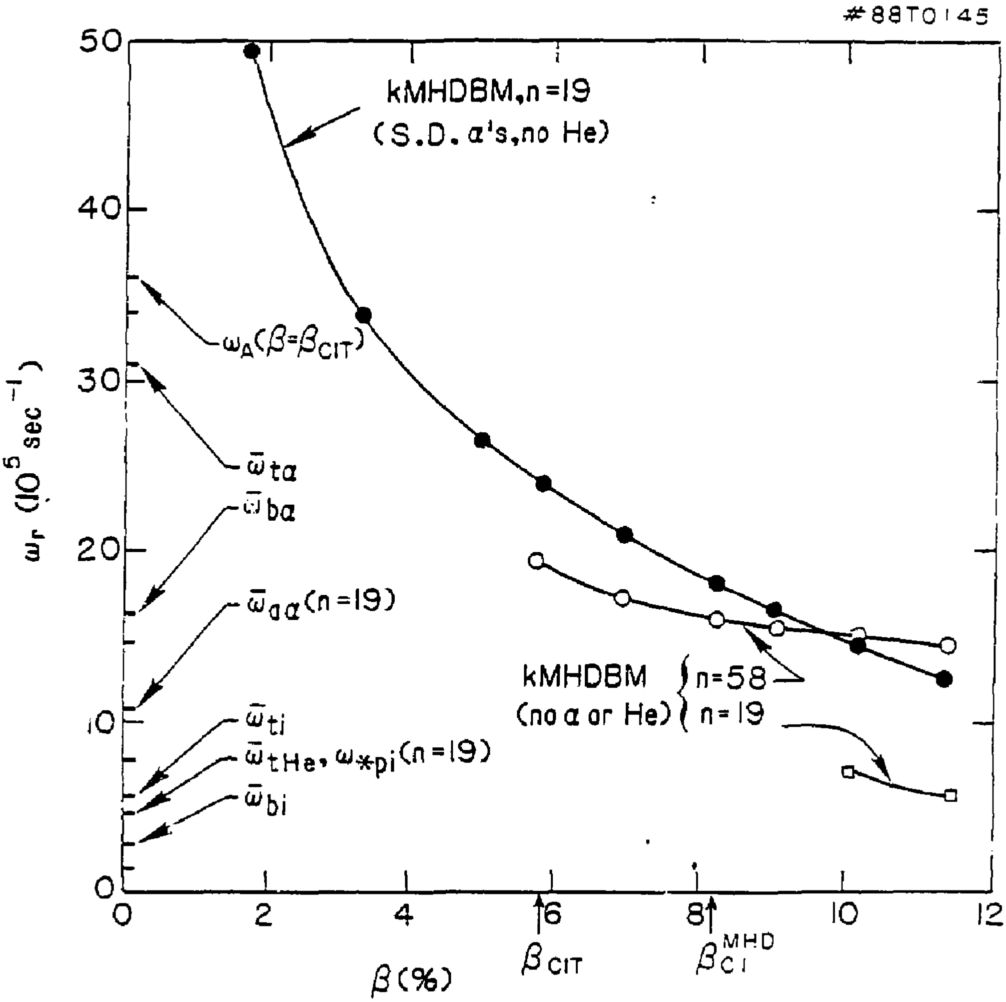
#88T0145
kMHDBM, n=19
(S.D. α's, no He)
$\omega_A(\beta=\beta_{CIT})$
$\bar{\omega}_{t\alpha}$
$\bar{\omega}_{b\alpha}$
$\bar{\omega}_{d\alpha}(n=19)$
$\bar{\omega}_{ti}$
$\bar{\omega}_{tHe}, \omega_{*pi}(n=19)$
$\bar{\omega}_{bi}$
kMHDBM
(no α or He)
n=58
n=19
ω_r (10^5 sec^{-1})
β (%)
β_{CIT}
β_{CI}^{MHD}
0
10
20
30
40
50
0
2
4
6
8
10
12

Figure 3

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