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DISSIPATIVE EFFECTS ON FINITE-LARMOR-RADIUS MODIFIED MAGNETOHYDRODYNAMIC BALLOONING MODES

J.W. Connor,* Liu Chen, and M.S. Chance

Plasma Physics Laboratory, Princeton University

Princeton, New Jersey 08544

ABSTRACT

Finite-ion-Larmor-radius (FLR) effects provide a band of additional stability for values of β($\equiv$ plasma pressure/magnetic pressure) exceeding the limit predicted for ideal magnetohydrodynamic (MHD) ballooning modes. We examine the effect of particle collisions on the stable modes of the FLR modified ideal theory that exist in this range of β values.

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*Permanent address: Culham Laboratory, Abingdon, Oxon OX143DB, U.K. (Euratom/UKAEA Fusion Association)

INTRODUCTION

The theory of high-n ideal MHD ballooning modes in an axisymmetric torus predicts that, for given pressure and current profiles, there will be a critical value of β, i.e., β_c^{MHD}, that will cause instability [1,2,3]. When $\beta > \beta_c^{MHD}$ there will be discrete modes with frequencies

$$\omega_{MHD} = i\omega_A \, (\beta/\beta_c^{MHD} - 1)^{1/2}, \tag{1}$$

where ω_A is the Alfven frequency V_A/L_c with V_A the Alfven speed and L_c the connection length. Below this value of β lies a continuum of stable modes with $\omega_{MHD}^2 > 0$. In the simplest model [4,5] the effect of finite ion Larmor radius is to cause the substitution $\omega^2 \to \omega(\omega - \omega_{*i})$ for the ideal eigenvalue ω^2 of ideal MHD, where $\omega_{*i} = k_\perp a_i V_{Ti}/r_n$ is the ion diamagnetic drift frequency. Here $k_\perp$ is the wave number perpendicular to the magnetic field, a_i is the ion Larmor radius, V_{Tj} is the thermal speed of species j and r_n is the density scale length. Equivalently, $\omega_{*i} \sim \sqrt{b_i \beta}\, \omega_A L_c/r_n$ where $b_i = k_\perp^2 a_i^2$. Thus it is clear that ω is now given by

$$\omega = \frac{\omega_{*i}}{2} \pm (\omega_{*i}^2/4 + \omega^2_{MHD})^{1/2}, \tag{2}$$

and this corresponds to an increase in the critical β for instability to

$$\beta_c^{FLR} = \beta_c^{MHD} \left(1 + \frac{b_i}{4\varepsilon}\right), \tag{3}$$

where $\varepsilon \equiv r_n/L_c$ is the inverse aspect ratio, $\beta_c^{MHD} \simeq \varepsilon$ and ω is approximately given by

$$\omega = \frac{\omega_{*i}}{2} \pm i\omega_A \left(\beta/\beta_c^{FLR} - 1\right)^{1/2} . \tag{4}$$

When $\beta > \beta_c^{FLR}$ there is one stable and one unstable mode which connect to two discrete stable modes when $\beta < \beta_c^{FLR}$. As β is further reduced to β_c^{MHD}, the real parts of the frequencies of these modes approach zero and ω_{*i} respectively. Finally, when $\beta < \beta_c^{MHD}$, these discrete stable modes are replaced by a stable continuum, since then $\omega(\omega - \omega_{*i}) > 0$ (c.f. $\omega^2_{MHD} > 0$ in the ideal case [1].

Now in the absence of FLR effects it is known that unstable high n resistive ballooning modes are present for $\beta < \beta_c^{MHD}$ [6]. It is the purpose of this note to examine the effect of resistivity and perpendicular ion viscosity on the stable discrete modes present in the range $\beta_c^{MHD} < \beta < \beta_c^{FLR}$.

In Sec. II we derive the required eigenvalue equation from a kinetic description of ions and electrons in the appropriate collisional regime. For the electrons we solve the kinetic equation in the weakly collisional fluid limit $\nu_e > \omega$, but $\omega\nu_e \sim k_\parallel^2 V_{Te}^2$, where ν_j is the collision frequency, V_{Tj} the thermal velocity of species j and $k_\parallel$ is the wave number parallel to the magnetic field. For the ions we consider $k_\perp$ sufficiently large so that $\omega \sim \omega_* > k_\parallel V_{Ti}$ and, to include ion viscosity, $\omega \sim \nu_i b_i$. Substitution of the electron and ion responses in the quasi-neutrality and Maxwell's equations provides an eigenvalue equation. Solutions of this equation (both qualitative and numerical) for a large aspect-ratio equilibrium with circular surfaces are discussed in Sec. III.

II. THE EQUATIONS FOR RESISTIVE MODES

In this section we introduce the equilibrium, the ballooning representation, the gyro-kinetic equation, from which we derive useful moment equations, and Maxwell's equations.

A. Equilibrium.

We adopt the axisymmetric orthogonal flux coordinates ψ, ζ, χ with Jacobian J, where ψ is the poloidal flux within a magnetic surface, ζ the toroidal angle, and χ a poloidal angle variable. In these coordinates the magnetic field is expressible as [1]

$$\underset{\sim}{B} = -\underset{\sim}{\nabla}\psi \times \underset{\sim}{\nabla}\zeta + I(\psi)\,\underset{\sim}{\nabla}\zeta \quad , \tag{5}$$

where $I(\psi)$ is the prescribed toroidal field function, and the gradient operator may be written as

$$\underset{\sim}{\nabla} = \underset{\sim}{e}_\psi\, RB_\chi \frac{\partial}{\partial\psi} + \underset{\sim}{e}_\zeta \frac{1}{R} \frac{\partial}{\partial\zeta} + \underset{\sim}{e}_\chi \frac{1}{JB_\chi} \frac{\partial}{\partial\chi} \tag{6}$$

with R being the major radius. The safety factor then becomes

$$q = \frac{1}{2\pi} \oint \nu\, d\chi \tag{7}$$

with $\nu = IJ/R^2$.

B. Ballooning Representation.

Perturbed quantities are represented in the form [2]

$$\phi(\psi,\zeta,\chi,) = \sum_{p=-\infty}^{\infty} \bar{\phi}(\psi,\zeta,\chi - 2\pi p) \tag{8}$$

which is automatically periodic in χ. Therefore, for high-n perturbations, one can use an eikonal form to represent modes with short wavelength perpendicular to the magnetic field but long wavelength along it, without a conflict with periodicity in the presence of shear. Thus we write

$$\bar{\phi} = \hat{\phi}(\psi,\chi) \exp(inS) \tag{9}$$

with

$$S = \zeta - \int^{\chi} \nu d\chi + \int^{\psi} k(\psi) d\psi, \tag{10}$$

so that $\underset{\sim}{n} \cdot \underset{\sim}{\nabla} S = 0$ where $\underset{\sim}{n} = \underset{\sim}{B}/B$. $\hat{\phi}$ is a slowly varying function of ψ and χ which satisfies the same equation as ϕ but is no longer a periodic function of χ and must be suitably well behaved as $|\chi| \to \infty$. We shall frequently use the notation $\underset{\sim}{k}_\perp = n\underset{\sim}{\nabla} S$. Thus $\underset{\sim}{k}_\perp = k_\psi \underset{\sim}{e}_\psi + k_b \underset{\sim}{n} \times \underset{\sim}{e}_\psi$ with

$$k_\psi = - nRB_\chi \left[\int^{\chi} d\chi \frac{\partial \nu}{\partial \psi} - k(\psi) \right] , \tag{11}$$

and

$$k_b = nB/RB_\chi . \tag{12}$$

C. Gyrokinetic Equation.

As a result of the small ratio of their Larmor radii to mascroscopic scales the plasma ions and electrons of a Maxwellian plasma have perturbed distribution functions of the form [5]

$$\hat{f} = \frac{-e\hat{\phi}}{T} F_M + \hat{h}(\varepsilon,\mu,\chi)\exp(iL) \quad , \tag{13}$$

where ϕ is the electrostatic potential, F_M is a Maxwellian,

$$F_M = n_o \left(\frac{m}{2\pi T}\right)^{3/2} \exp\left(\frac{-m\varepsilon}{T}\right) \quad , \tag{14}$$

$\varepsilon = V^2/2$ is the energy and $\mu = V_\perp^2/2B$ is the magnetic moment per unit mass, η is the gyro-phase angle so that $\underset{\sim}{V} = V_\parallel \underset{\sim}{n} + V_\perp (\cos\eta \underset{\sim}{e}_\psi + \sin\eta\, \underset{\sim}{n}\times\underset{\sim}{e}_\psi)$,

$$d^3v = B d\mu\, d\varepsilon\, d\eta/|V_\parallel| \quad ,$$

and

$$L = (V_\perp/\Omega)\,(k_\psi \sin\eta - k_b \cos\eta) \tag{15}$$

with Ω being the gyrofrequency.

The function $\hat{h}$ is a solution of the gyro-kinetic equation [5,7]

$$V_\parallel \underset{\sim}{n}\cdot\underset{\sim}{\nabla}\hat{h} - i\hat{h}(\omega-\omega_D) - \langle e^{-iL} C(\hat{h}e^{iL})\rangle_\eta \tag{16}$$

$$= \frac{-ie}{T} F_M(\omega-\omega_*^T)\left[J_o(\alpha)\left(\hat{\phi} - \frac{V_\parallel}{c}\hat{A}_\parallel\right) + J_1(\alpha)\frac{V_\perp}{k_\perp c}\,\delta\hat{B}_\parallel\right] \quad ,$$

where C is the usual Fokker-Planck collisional operator, $\langle\ \rangle_\eta$ represents a gyrophase average, and

$$\omega_D = \frac{m\mu B}{T}\omega_B + \frac{mV_\parallel^2}{T}\omega_\kappa \tag{17}$$

with $\omega_B = (T/m\Omega B)\ \underset{\sim}{n}\times\underset{\sim}{\nabla}B\cdot\underset{\sim}{k}_\perp$ and $\omega_\kappa = (T/m\Omega B)\ \underset{\sim}{n}\times(\underset{\sim}{n}\cdot\underset{\sim}{\nabla}n)\cdot\underset{\sim}{k}_\perp$. J_0 and J_1 are Bessel functions of argument $\alpha = k_\perp V_\perp/\Omega$, $\omega^{*T} = \omega^*\left[1 + \eta\left(m\varepsilon/T - 3/2\right)\right]$ with $\omega_* = (ncT/e)d(\ln n_0)/d\psi$ and $\eta = d(\ln T)/d(\ln n_0)$. $\hat{A}_\parallel$ and $\delta\hat{B}_\parallel$ are the components of the vector potential and the perturbed magnetic field along the equilibrium field. Here we use the Coulomb gauge div $\hat{\underset{\sim}{A}} = 0$.

We emphasize that this equation contains secular terms in χ through ω_D and α arising from shear and that we seek non-periodic solutions defined on the interval $-\infty < \chi < \infty$ which are suitably well behaved at infinity. In this work we shall consider the long wavelengths limit $b \ll 1$ where $b = k_\perp^2 T/m\Omega^2$ so that in this equation the collision term and Bessel functions can be expanded in L and α.

D. Electron Solution.

Using the symbol $k_\parallel = \underset{\sim}{n}\cdot\underset{\sim}{\nabla}\phi/\phi$, we solve the electron gyro-kinetic equation in the limit $\omega/k_\parallel V_{Te} \sim k_\parallel V_{Te}/\nu_e \ll 1$, $b_e \ll 1$. To simplify notation, we now suppress electron species suffices.

Thus, in leading order (and writing ϕ for $\hat{\phi}$, etc.), we have

$$C(h_0) = 0 \qquad , \tag{18}$$

which implies a Maxwellian solution

$$h_0 = \frac{n'}{n_0} F_M \qquad , \tag{19}$$

where n' is the non-adiabatic density perturbation. For simplicity, we shall ignore equilibrium temperature gradients and the resulting temperature perturbations.

In next order, we have

$$C(h_1) = V_\parallel \underset{\sim}{n}\cdot\nabla h_o - i(\omega-\omega_*)\left(\frac{eF_M}{Tc}\right) V_\parallel A_\parallel \tag{20}$$

and, finally,

$$C(h_2) = V_\parallel \underset{\sim}{n}\cdot\underset{\sim}{\nabla} h_1 - i(\omega-\omega_D)h_o + i(\omega-\omega_*)\frac{eF_M}{T}\left(\phi + \frac{V_\perp^2}{2\Omega c}\delta B_\parallel\right) \quad . \tag{21}$$

The parallel current arises from the solution of Eq. (20) for h_1, using Eq. (19) for h_o. This resembles the problem solved by Spitzer [8] and we find

$$j_\parallel = \sigma\left[\frac{T_e}{en_o}\underset{\sim}{n}\cdot\underset{\sim}{\nabla} n_e' + i(\omega-\omega_{*e})\frac{A_\parallel}{c}\right] \quad , \tag{22}$$

where σ is the Spitzer conductivity.

E. Ion Solution.

In the ion case we consider $\omega \sim \omega_* \gg k_\parallel V_{Ti}$ (imposing a lower limit on the toroidal mode number n) and $b_i < 1$, but $\nu_i b_i \sim \omega$, which introduces ion - ion collisions as a collisional viscosity.

Thus, in leading order, we note that ion - ion collisions dominate the collision term and, hence, obtain for the ions

$$C(h_o) = 0 \quad , \tag{23}$$

i.e.,

$$h_o = \frac{n'}{n_o} F_M \tag{24}$$

where ion species' suffices are again suppressed. In next order, we have

$$-i(\omega-\omega_D)h_o - C(h_1) + \langle C(h_o\frac{L^2}{2})\rangle_\eta - \langle LC(h_oL)\rangle_\eta \tag{25}$$

$$= -i(\omega-\omega_*)\frac{eF_M}{T}\left[\phi - \frac{V_\parallel A_\parallel}{c} + \frac{V_\perp^2}{2\Omega c}\delta B_\parallel\right] \quad .$$

Taking the density moment, we obtain

$$(\omega-\omega_\kappa-\omega_B)n' = (\omega-\omega_*)\frac{en_o}{T}\left(\phi + \frac{T^2}{m\Omega c}\delta B_\parallel\right) \quad , \tag{26}$$

where in Eq. (25) density and momentum conservation in ion - ion collisions annihilate the collision term.

F. Parallel Current Moment Equation.

We form the moment $\sum e_j \int d^3Ve^{iL}$ of the gyrokinetic equation, Eq.(16), to obtain

$$\underset{\sim}{B}\cdot\underset{\sim}{\nabla}\frac{j_\parallel}{B} = i\omega\sum\frac{e^2}{T}\int d^3VF_M\{[1-(1-\frac{\omega_*}{\omega})J_o^2(\alpha)]\phi - (1-\frac{\omega_*}{\omega})J_o(\alpha)J_1(\alpha)\frac{V_\perp}{k_\perp c}\delta B_\parallel\} \tag{27}$$

$$- i\sum e\int d^3V\omega_Dhe^{iL} + \sum e\int d^3Ve^{iL}\langle e^{-iL}C(he^{iL})\rangle_\eta + \sum e\int d^3VV_\parallel h\underset{\sim}{n}\cdot\underset{\sim}{\nabla}e^{iL}$$

where we have used Eq. (8) and charge neutrality. If we consider small b, i.e., ignoring terms in b_e, $b_i\omega_D/\omega$, $b_i\,(\delta B_\parallel/B)/(e\phi/T)$ (see subsection G) and $b_i\,k_\parallel V_{Ti}/\omega_*$, we obtain

$$\underset{\sim}{B}\cdot\underset{\sim}{\nabla}\,\frac{j_\parallel}{B} = \frac{i n_o e^2}{T_i}\,(\omega-\omega_{*i})\;b_i\phi \;+\;(\omega_{*i}-\omega_{*e})\,\frac{T_i\delta B_\parallel}{m_i\Omega_i c}$$

$$-\; i\sum_j e_j\;(\omega_\kappa+\omega_B)_j\;n'_j \;-\; e\nu_i b_i^{\,2}\alpha n'_i \tag{28}$$

with $\alpha = \frac{3}{10}$ and $\nu_i = 4\pi^{1/2} n_o e^4 \ln\lambda/(3M_i^{\,1/2}T_i^{\,3/2})$.

In deriving this result we have expanded the collision term in Eq. (27) to order $b_i^{\,2}$, the first non-vanishing contribution, using appropriate conservation properties, and Eq. (25) for the order b_1 correction to h_o. The resulting integrals have been evaluated using the projection techniques and expansion procedure of Shkarofsky et al. [9]. The ion collision term representing the lowest-order collisional non-ambipolar diffusion flux corresponds to ion perpendicular viscosity.

G. Maxwell's Equations.

We can eliminate $\delta B_\parallel$ by means of the component of Ampere's law perpendicular to the magnetic surface

$$i k_b\,\delta B_\parallel = \frac{4\pi}{c}\,j_\psi \quad , \tag{29}$$

with $j_\psi = \sum e\int d^3v\;v_\perp\,\cos\eta\;h\,\exp(iL)$. Equation (29), along with (19) and (24), yields, in the limit $b \ll 1$,

$$\delta B_\parallel = 4\pi\sum\frac{eT}{m\Omega}\,n' \quad , \tag{30}$$

confirming that $\delta B_\parallel/B \sim \beta\, n'/n_o$. Using the identity $(\beta_i/2)\;(\omega_{*i}-\omega_{*e}) = (\omega_\kappa-\omega_B)$ where $\beta_i = 8\pi n_o T_i/B^2$, we note that the role of $\delta B_\parallel$ is effectively to replace ω_B by ω_κ in Eq. (28).

Finally, we have the parallel component of Ampere's law, i.e.,

$$k_\perp^2 A_\parallel = \frac{4\pi}{c} j_\parallel \tag{31}$$

H. Eigenvalue Equation.

In the limit ω_D/ω , $\beta < 1$, Eq. (26) for n_i' together with quasi-neutrality implies that

$$n_j' = \frac{n_o e_j}{T_j} \phi \left(1 - \frac{\omega_{*j}}{\omega}\right) . \tag{32}$$

Combining Eqs. (22) and (28) - (32) yields the eigenvalue equation

$$\underset{\sim}{B}\cdot\underset{\sim}{\nabla} \left\{\frac{(k_\perp^2/B^2)\underset{\sim}{B}\cdot\underset{\sim}{\nabla}\phi}{1 + [ik_\perp^2 c/4\pi\sigma(\omega-\omega_{*e})]}\right\} = \frac{-4\pi n_o e^2\phi}{T_i}\{\omega(\omega-\omega_{*i})b_i + 2\omega_{\kappa i}(\omega_{*i}-\omega_{*e})$$

$$+ i\alpha(\omega-\omega_{*i})\nu_i b_i^2\} \tag{33}$$

or, employing fluid notation,

$$\underset{\sim}{B}\cdot\underset{\sim}{\nabla}\left[\frac{(|\underset{\sim}{\nabla}S|^2/B^2)\,\underset{\sim}{B}\cdot\underset{\sim}{\nabla}\phi}{1 + i\eta_s|\underset{\sim}{\nabla}S|^2/4\pi(\omega-\omega_{*e})}\right] = -4\pi\phi\,\{\omega\,(\omega-\omega_{*e})\,m_i n_o|\underset{\sim}{\nabla}S|^2B^2 \tag{34}$$

$$+ 2\underset{\sim}{B}\times\underset{\sim}{\kappa}\cdot\underset{\sim}{\nabla}S\,\underset{\sim}{B}\times\underset{\sim}{\nabla}P\cdot\underset{\sim}{\nabla}S + i\mu_v\,(\omega-\omega_{*i})\,|\underset{\sim}{\nabla}S|^4\}/B^4$$

where p is the total pressure, η_s the Spitzer resistivity, $\underset{\sim}{\kappa} = \underset{\sim}{n}\cdot\underset{\sim}{\nabla}\underset{\sim}{n}$ and $\mu_v = \alpha m_i^2 n_o T_i \nu_i/e^2$ is the ion perpendicular viscosity.

III. SOLUTION FOR LARGE ASPECT RATIO TOKAMAK WITH CIRCULAR SURFACES

In this section we consider solutions of Eq. (33) for a large aspect ratio tokamak with circular flux surfaces. In this case the equilibrium can be characterised by two parameters: $\alpha = 2Rq^2P/r_nB^2$ representing the pressure gradient and $s = (r/q)\,dq/dr$ representing the shear [10]. It is convenient to introduce the normalised frequency $\Omega = \omega/\omega_{*e}$ and collision frequency $\nu = \nu_e r_n/V_{Te}$ where $\sigma = n_o e^2/m\nu$. Equation (33) becomes a differential equation defined on an infinite range of the poloidal angle θ:

$$\frac{d}{d\theta}\left[\frac{1+s^2\theta^2)}{1+i\nu\, b_{io}^{1/2}\, q^2/\hat{\beta}\,\varepsilon_n(\Omega-1)}\right]\frac{d\phi}{d\theta} + 2\hat{\beta}\,(1+\tau)\,(\cos\theta + s\theta\sin\theta)\,\phi$$

$$+ (\hat{\beta}\, b_{io}/\varepsilon_n\tau)\,\Omega(\Omega+\tau)\,(1+s^2\theta^2)\left[1 + i\,\delta_i\tau\nu b_{io}^{1/2}\right. \tag{35}$$

$$\left.\times\,(1+s^2\theta^2)/\Omega\,\right]\phi = 0 \quad .$$

Here, $\hat{\beta} = 4\pi q^2 n_o T_e/\varepsilon_n B^2$, $b_{io} = m_i T_i n^2 q^2/e^2 B^2 r_n^2$, $\tau = T_i/T_e$, and we use the labeling parameter $\delta_i \sim 1$ to consider ion-ion collisions independently of electron collisions. We have solved this equation for $s = q = \tau = 1$ and $\varepsilon = b_{io} = 0.1$ (consistent with the expansion $\omega_* > k_\parallel V_{Ti}$). With $\nu = 0$ the instability threshold is $\hat{\beta} = 0.1997 \equiv \hat{\beta}_c^{FLR}$. We have considered the variation of Im Ω with ν for the unstable branch when $\hat{\beta} = 0.1995$, i.e., just below $\hat{\beta}_c^{FLR}$, and when $\hat{\beta} = 0.1900$, intermediate in the range $\hat{\beta}_c^{MHD} < \hat{\beta} < \hat{\beta}_c^{FLR}$. The results are shown in Fig. 1 for two choices of δ_i, namely, $\delta_i = 0$, corresponding to ignoring ion-ion collisions, and $\delta_i = 1$, characteristic of their being fully effective.

It should be noted that the growth rate increases with electron-ion collisions. Indeed the behavior can readily be understood qualitatively by perturbation theory. From the general Eq. (34), with $\mu = 0$ for simplicity, one can construct schematically a quadratic form

$$\omega(\omega - \omega_{*i}) + \omega_A^{\,2} \left(\frac{\beta}{\beta_c^{MHD}} - 1 + \frac{i\nu^*}{\omega - \omega_{*e}}\right) = 0 \tag{36}$$

where $\nu^* \propto \nu$. Thus

$$\omega = \frac{\omega_{*i}}{2} \pm i \left[\frac{\beta}{\beta_c^{FLR}} - 1 + \frac{i\nu^*}{\omega - \omega_{*e}} \right]^{1/2} \tag{37}$$

When $\beta = \beta_c^{FLR}$, Im ω has a $\nu^{1/2}$ dependence, whereas when $\beta < \beta_c^{FLR}$ it takes on a linear variation with ν, features characteristic of Fig. 1. It can be seen in Fig. 1 that ion-ion collisions have a small stabilizing effect on this mode.

IV. Conclusion.

We have derived an equation to investigate the influence of collisions on high-n ballooning modes in the presence of FLR effects. Calculations have been performed for a large aspect ratio tokamak with circular flux surfaces. In the absence of collisional effects there are two modes with real frequencies in the range of β value given by $\beta_c^{MHD} < \beta < \beta_c^{FLR}$, where $\beta_c^{MHD,FLR}$ are the critical β values in ideal MHD theory and in the presence of FLR, respectively. Electron-ion collisions, i.e., resistivity, destabilize one of those branches with a growth rate $\gamma \propto \nu$ when $\beta < \beta_c^{FLR}$, but with $\gamma \propto \nu^{1/2}$ when $\beta \lesssim \beta_c^{FLR}$. Ion-ion collisions, i.e., ion perpendicular viscosity, tend to reduce these growth rates somewhat. Opposite behaviors are observed for the resistively stable branch. Finally, we remark that FLR effects on resistive ballooning modes in the $\beta \lesssim \beta^{MHD}$ regime are currently under investigation and will be reported in a future publication.

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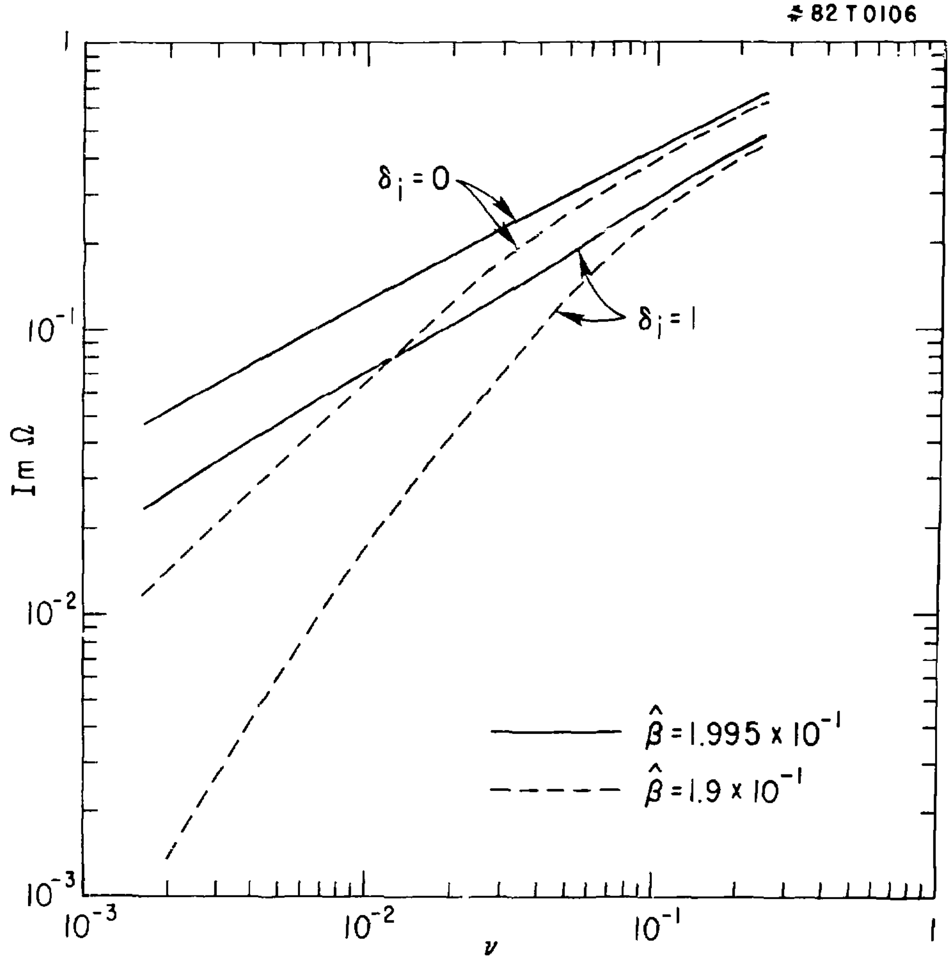
Figure Caption

Fig. 1. Growth rate versus electron-ion Collision frequency ν, for two different values of $\hat{\beta}$. $\delta_i = 1$ and $\delta_i = 0$ correspond, respectively, to calculations with and without ion-ion collisions.



82 T 0106
Im Ω
1
10^-1
10^-2
10^-3
10^-3
10^-2
10^-1
1
ν
δ_i = 0
δ_i = 1
β̂ = 1.995 x 10^-1
β̂ = 1.9 x 10^-1

Fig. 1