Stability of electrostatic modes in a levitated dipole

J. Kesner
Plasma Fusion Center, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139

(Received 4 September 1996; accepted 13 November 1996)

Plasma confined in a magnetic dipole is stabilized by the expansion of the magnetic flux. The stability of low beta electrostatic modes in a magnetic dipole field is examined when the distribution function is to lowest order Maxwellian. It is shown that for sufficiently gentle density and temperature gradients the configuration would be expected to be stable to magnetohydrodynamic interchange, as well as to dissipative trapped ion and collisionless trapped particle modes. These results are applicable to any magnetic configuration for which the curvature drift frequency exceeds the diamagnetic drift frequency. © *1997 American Institute of Physics.* [S1070-664X(97)03602-1]

I. INTRODUCTION

The dipole magnetic field is the simplest and most common magnetic field configuration in the universe. It is the magnetic far field of a single, circular current loop, and it represents the dominant structure of the middle magnetospheres of magnetized planets and neutron stars. The use of a dipole magnetic field generated by a levitated ring to confine a hot plasma for fusion power generation was first considered by Hasegawa.[1] As a confinement configuration for magnetic fusion, a dipole possesses uniquely good properties. The coil set is simple and axisymmetric. Operation is inherently steady state and the large flux expansion is expected to simplify the divertor design. *Vis-à-vis* a tokamak, there is no need for current drive and no disruptions. It is expected to have good magnetohydrodynamic (MHD) properties, including plasma pressures that can locally exceed the magnetic pressure, i.e., $\beta > 1$, and excellent confinement properties. By levitating the dipole magnet in order to prevent end losses in a conceptual reactor, studies have supported the possibility of a dipole based fusion reactor.[2–4] In this paper we will focus on the stability of drift modes that are thought to degrade confinement in fusion grade plasmas.

For a plasma confined in a levitated dipole the pressure falls off (moving away from the internal coil) in a region of ''bad'' curvature. In this situation it is well known that stability can be obtained due to the so-called compressibility, and there is a critical value of the pressure gradient that can be confined stably. The dipole reactor concept is based on the idea of generating pressure profiles near marginal stability for low-frequency MHD fluctuations. From ideal MHD, the marginal stability of interchange modes results when the pressure profile satisfies the adiabaticity condition, $\delta(pV^{\gamma})=0$, with p the plasma pressure, V is the flux tube volume, and $\gamma=5/3$. We derive the equivalent condition from the drift kinetic equation. In the derivation of the dispersion relation for low frequency interchange modes from the drift kinetic equation the stabilizing term derives from the square of the curvature drift frequency [Eq. (8)] and there is no need to make the assumption of an equation of state. This derivation therefore derives from first principles and clarifies the origin of compressibility in MHD.

The ability to confine plasma at high beta makes the dipole configuration particularly well suited as an advanced fuel reactor. Ignition in advanced fuel plasmas such as D^3He requires particularly good confinement properties. Since the magnetic field is entirely in the poloidal plane there are no particle drifts off the flux tubes (which in a tokamak result in a ''neoclassical'' degradation of confinement). In this paper we show that plasma confined in a levitated dipole may be expected to be free of drift wave turbulence and therefore a dipole based reactor may be expected to attain classical confinement.

Hasegawa has pointed out[1,2] that when the plasma is sufficiently collisionless, the equilibrium distribution function may be described by $F_0=F_0(\mu,J,\psi)$, with μ the first invariant, $\mu=v_\perp^2/2B$, J the second invariant, $J=\oint ds\ v_\parallel$, and ψ the flux invariant. For fluctuations in the range of the curvature drift frequency, flux is not conserved and a collisionless plasma can approach the state $\partial F/\partial\psi\rightarrow 0$. Furthermore, when $\partial F/\partial\psi=0$ the plasma can be shown to be stable to drift frequency fluctuations. This condition leads to dipole pressure profiles that scale with radius as $r^{-20/3}$, similar to energetic particle pressure profiles observed in the planetary magnetospheres.[5,6]

In a conceptual reactor, confinement must be maintained on a collisional time scale. Therefore, we would expect the distribution function to be, to lowest order, Maxwellian, i.e., $F_0(\mu,J)\rightarrow F_0(\epsilon,\psi)$ and therefore $\partial F/\partial\psi\neq 0$. In this article we assume to lowest order a Maxwellian distribution function for both ions and electrons and we derive the condition for marginal stability to MHD interchange modes, as well as collisionless and dissipative trapped ion modes. We find that each of these collective modes becomes stable when the density and temperature gradients are sufficiently gentle. Therefore, a plasma confined in a levitated dipole field may be expected to be particularly stable to collective modes and, for sufficiently gentle gradients, may exhibit classical confinement.

These results are applicable to any magnetic configuration for which the curvature drift frequency is comparable to the diamagnetic drift frequency, such as a low aspect ratio tokamak.

II. ELECTROSTATIC, TRAPPED PARTICLE MODES

To derive the stability criterion for electrostatic modes we consider a fluctuating electrostatic potential, ϕ and ignore any equilibrium electrostatic potential. From Faraday's law it is possible for a perturbation to leave the magnetic field undisturbed if $E=-\nabla\phi$, which is consistent with $\beta\ll 1$. If ϕ varies along a field line, there will be a finite $E_\parallel$ (a situation not possible in ideal MHD theory).

We analyze the stability of such a perturbation under the assumptions that the wave frequency ω is less than the cyclotron frequency Ω_c and that the ion Larmor radius ρ_i is shorter than the perpendicular wavelength $\lambda=2\pi/k_\perp$ which is, in turn, short compared to a parallel wavelength, $2\pi/k_\parallel$. The appropriate equation for the distribution function $\tilde{f}$ is[7,8]

$$\tilde{f}=q\phi F_{0\epsilon}=J_0(k_\perp\rho)h \tag{1}$$

and h satisfies

$$(\omega-\omega_d+iv_\parallel \mathbf{b}\cdot\nabla')h=-(\omega-\omega_*)q\phi F_{0\epsilon}J_0(k_\perp\rho)+iC(h). \tag{2}$$

$F_0(\epsilon,\psi)$ is the equilibrium distribution function and ∇' is the gradient at constant ϵ and μ

$$F_{0\epsilon}\equiv\frac{\partial F_0}{\partial\epsilon}, \tag{3a}$$

$$\omega_*=\frac{\mathbf{b}\times\mathbf{k}_\perp\cdot\nabla' F_0}{m\Omega_c F_{0\epsilon}}, \tag{3b}$$

$$\omega_d=\mathbf{k}_\perp\cdot\mathbf{b}\times\frac{(mv_\parallel^2\mathbf{b}\cdot\nabla\mathbf{b}+\mu\nabla B)}{m\Omega_c}, \tag{3c}$$

$$\mathbf{B}=\nabla\psi\times\nabla\theta, \tag{3d}$$

$$\mathbf{b}=\mathbf{B}/|\mathbf{B}|. \tag{3e}$$

$J_0(k_\perp\rho)$ is the Bessel function of the first kind and θ is the azimuthal angle.

We consider a perturbation whose growth time is long compared to a particle bounce time and obtain the result that h is a constant along a field line $h=h_0(\epsilon,\mu,\psi)$. We can then determine the constant by taking the time average of Eq. (3),

$$h_0=\frac{-(\omega-\omega_*)q\overline{\phi F_{0\epsilon}J_0}}{(\omega-\overline{\omega_d}+i\nu_q)}. \tag{4}$$

The overbar indicates a time average

$$\overline{\phi}=\frac{1}{\tau_B}\int\frac{dl}{|v_\parallel|}\phi, \tag{5a}$$

$$\tau_B=\int\frac{dl}{|v_\parallel|}. \tag{5b}$$

For simplicity the collision operator has been replaced by a Krook model in Eq. (4), i.e., $C(h)\rightarrow-\nu_j h$ with ν_j the appropriate collision frequency.

III. FAST GROWING MODES

To explore modes that grow on the MHD time scale we assume that $\omega>\overline{\omega_d}>\nu_j$ and expand the denominator of Eq. (4) to obtain for the perturbed particle distribution function

$$\tilde{f}=q\phi F_{0\epsilon}-q\overline{\phi F_{0\epsilon}}J_0^2+\left[\frac{\mathbf{b}\times\mathbf{k}_\perp\cdot\nabla' F_0}{m\Omega_c\omega}-\frac{\overline{\omega_d}}{\omega}F_{0\epsilon}+\frac{\mathbf{b}\times\mathbf{k}_\perp\cdot\nabla' F_0}{m\Omega_c\omega}\frac{\overline{\omega_d}}{\omega}-\frac{\overline{\omega_d^2}}{\omega^2}F_{0\epsilon}\right]q\overline{\phi J_0^2}. \tag{6}$$

We determine the eigenfrequency ω by requiring that the mode be quasineutral. We expand J_0^2 as $J_0^2\propto 1-(k_\perp v_\perp)^2/2\Omega_c^2$ in the first term, but neglect the $k_\perp^2\rho_i^2$ correction in the second term. With these assumptions the quasineutrality condition becomes

$$O=\sum_q q^2\int d^3v\left\{F_{0\epsilon}(\phi-\overline{\phi})+\overline{\phi\frac{k_\perp^2 v_\perp^2}{2\Omega_c^2}F_{0\epsilon}}+(\omega_*\overline{\omega_d}-\overline{\omega_d^2})\frac{F_{0\epsilon}}{\omega^2}\overline{\phi}\right\}. \tag{7}$$

The terms proportional to $1/\omega$ have canceled in the sum over species. Consider first the case when the eigenmode is flute-like, i.e., $\overline{\phi}\approx\phi$, $E_\parallel=0$. In this limit we can solve for ω^2 to obtain

$$\omega^2=\frac{\Sigma_q q^2\int d^3v\;F_{0\epsilon}\overline{\omega_d}(\overline{\omega_d}-\omega_*)}{\Sigma_q q^2\int d^3v\;F_{0\epsilon}k_\perp^2v_\perp^2/2\Omega_c^2}. \tag{8}$$

In our convention $\omega_*>0$ and for a dipole there is ''bad'' curvature, i.e., $\overline{\omega_d}>0$. Flute-like (interchange) modes are nevertheless stable when $\int d^3v\;F_0(\overline{\omega_d^2}-|\omega_*\overline{\omega_d}|)>0$. Defining $\tilde{\omega}_*=\int F_0\omega_* d^3v$, $\tilde{\omega}_d^2=\int F_0\overline{\omega_d^2}\,d^3v$ and approximating $\int F_0\omega_*\overline{\omega_d}\,d^3v\approx(1+\eta)\tilde{\omega}_*\tilde{\omega}_d$ with $\eta=\nabla\ln(T)/\nabla\ln(n)$, we see that stability requires

$$\tilde{\omega}_d>(1+\eta)\tilde{\omega}_*, \tag{9}$$

i.e., $r_*>R_c$ with $r_*\sim p/\nabla p$ and R_c is the radius of curvature.

Under some circumstances non-flute-like ''trapped particle'' modes can grow on the MHD time scale.[9] These fast growing modes can be investigated by constructing a quadratic form. We multiply Eq. (7) by ϕ^*/B and perform a flux tube ($\int dl/B$) integration. Writing

$$\int d^3v=\frac{2\pi B}{m^2}\int\frac{d\epsilon\,d\mu}{v_\parallel}$$

interchanging the order of integration, and solving for ω^2 gives

$$-\omega^2=\frac{\Sigma_q q^2(2\pi/m^2)\int d\epsilon\,d\mu\;\tau_B\overline{\phi}^2F_{0\epsilon}(\overline{\omega_d}^2-\omega_*\overline{\omega_d})}{\Sigma_q q^2(2\pi/m^2)\int d\epsilon\,d\mu\;\tau_B(-F_{0\epsilon})\{(\overline{\phi^2}-\overline{\phi}^2)+\overline{\phi}^2k_\perp^2\rho_i^2\}}. \tag{10}$$

The quadratic form, Eq. (10) is variational with respect to ϕ and since the denominator is positive definite, if we can find a trial function for ϕ such that $\omega^2<0$ the true eigenfunction will give an even larger growth rate. The numerator of Eq. (10) can be positive for regions where the magnetic field curvature is concave with respect to the plasma if $\omega_* > \overline{\omega_d}$. When the drive is localized the eigenfunction tends to concentrate in these regions. The denominator, however, is small for trial functions which are spread out. The actual eigenfunction is determined by the balance between concentrating in regions of high curvature to make the numerator larger and spreading as much as possible to make the denominator smaller. For a mode that localizes on the outer midplane of the torus the deeply trapped particles feel the full strength of the mode while the passing particles feel a reduced ''average'' fluctuation.

For a dipole we have seen from Eq. (9) that interchange stability requires that $\tilde{\omega}_d > \omega_*$. Since for a dipole, ω_d is largest at the outer midplane, we see that when a levitated dipole is stable to interchange modes it will also be stable to fast growing trapped particle modes.

IV. DRIFT FREQUENCY MODES

Lower frequency modes can be destabilized by the resonance at $\omega=\overline{\omega_d}$ [Eq. (4)]. To evaluate the stability of the resonant modes we begin with Eq. (4) and apply quasineutrality

$$\sum_q q^2 \int d^3v\ F_{0\epsilon}\phi = \sum_q q^2 \int d^3v\ J_0^2 F_{0\epsilon} \overline{\phi} \left[\frac{\omega - \mathbf{b}\times\mathbf{k}_\perp \cdot \boldsymbol{\nabla}' F_0/(m\Omega_c F_{0\epsilon})}{\omega - \overline{\omega_d} + i\nu_q} \right]. \tag{11}$$

Multiplying Eq. (11) by ϕ^* and taking a flux tube average yields

$$(1+T_e/T_i) \int d^3v\ F_0 \overline{\phi^2} = \sum_q \int d^3v\ J_0^2 F_{0\epsilon} \overline{\phi}^2 \left[\frac{\omega - \mathbf{b}\times\mathbf{k}_\perp \cdot \boldsymbol{\nabla}' F_0/(m\Omega_c F_{0\epsilon})}{\omega - \overline{\omega_d} + i\nu_q} \right]. \tag{12}$$

The bounce-averaged quantities, $\overline{\phi^2}$ and $\overline{\phi}^2$ are in general functions of the pitch angle. For a flute-like mode $\overline{\phi^2} = \overline{\phi}^2$ and we will show below that such modes are not unstable. A trapped particle driven mode is localized to the outer midplane so that trapped particles experience the mode more strongly than passing particles and $\overline{\phi^2} > \overline{\phi}^2$. To obtain an approximate dispersion relation for the collisionless trapped particle modes we extract an average value of the bounce average quantities, $\overline{\phi}$ and $\overline{\phi}^2$ and rewrite Eq. (12) as

$$\frac{(1+T_e/T_i)}{f_t} = \sum_q \int d^3v \times J_0^2 F_{0\epsilon} \left[\frac{\omega - \mathbf{b}\times\mathbf{k}_\perp \cdot \boldsymbol{\nabla}' F_0/(m\Omega_c F_{0\epsilon})}{\omega - \overline{\omega_d} + i\nu_q} \right]. \tag{13}$$

For constant $\phi(\phi=\overline{\phi}=(\overline{\phi^2})^{1/2})$, $f_t=1$. For a mode that is localized near the outer midplane or which changes sign from the outside to the inside the trapped particles feel the mode more strongly ($\phi_{\text{tapped}} \sim \phi_{\max}$) than the passing particles ($\phi_{\text{pass}} \sim 0$) and it can be seen that f_t is related to the trapped particle fraction.

For a dipole the curvature drift is fairly uniform along the field line and we can approximate $\overline{\omega_d} \to (\epsilon/T)^* \hat{\omega}_d$ in Eq. (13) to obtain

$$\left\langle \frac{\omega - \omega_{*n}[1+\eta_e(\epsilon/T_e - (3/2))]}{\omega - \hat{\omega}_d(\epsilon/T_e) + i\nu_e} \right\rangle + \left\langle \frac{\omega\tau + \omega_{*n}[1+\eta_i(\epsilon/T_i - (3/2))]}{\omega + (\hat{\omega}_d/\tau)(\epsilon/T_i) + i\nu_i} \right\rangle = (1+\tau)/f_t, \tag{14}$$

where $\langle A\rangle \equiv (2/\pi^{1/2}T_j^{3/2})\int_0^\infty d\epsilon\ \epsilon^{1/2} \exp(-\epsilon/T_j)A$, $\tau \equiv T_e/T_i$, $\omega_{*n} \equiv -\omega_{*ni}\tau = -k_\theta \rho_i v_i \tau/2r_n$, $v_i=(2T_i/m_i)^{1/2}$, $\rho_i = v_i/\Omega_i$, $\Omega_i = eB/m_i c$, $k_\theta = m/r$, $r_n = -(d \ln n/dr)^{-1}$, $\omega_d \equiv (\omega_d)_e = -(\omega_d)_i \tau = k_\theta \rho_i v_i \tau/2R_C > 0$, $\eta_j \equiv d \ln T_j/d \ln n_j$, and $\omega_{*n} \gtreqless 0$ for $dn/dr \lesseqgtr 0$.

For $T_e = T_i$ Eq. (14) simplifies

$$\left\langle \frac{\omega - \omega_{*n}[1+\eta(\epsilon/T - (3/2))]}{\omega - \hat{\omega}_d(\epsilon/T) + i\nu_e} \right\rangle + \left\langle \frac{\omega + \omega_{*n}[1+\eta(\epsilon/T - (3/2))]}{\omega + \hat{\omega}_d(\epsilon/T) + i\nu_i} \right\rangle = 2/f_t. \tag{15}$$

V. COLLISIONLESS RESONANT MODE

Equation (14) was analyzed by Tagger *et al.*[10] and by Tang *et al.*[11] for the tokamak case, i.e., when $\hat{\omega}_* > \hat{\omega}_d$. We have seen that for a dipole MHD stability requires $\tilde{\omega}_* < \tilde{\omega}_d$.

Consider first the collisionless mode ($\nu_e \sim \nu_i \sim 0$). If we define $\Omega = \omega/\hat{\omega}_d$, $\Omega_* = \omega_{*n}/\hat{\omega}_d$ and $\Omega_{*T} = \eta\Omega_*$ Eq. (15) becomes

$$\left\langle \frac{\Omega + \Omega_{*T}(3/2 - 1/\eta - \epsilon/T)}{\Omega - \epsilon/T} \right\rangle + \left\langle \frac{\Omega - \Omega_{*T}(3/2 - 1/\eta - \epsilon/T)}{\Omega + \epsilon/T} \right\rangle = \frac{2}{f_t}. \tag{16}$$

Noting that

$$\left\langle \frac{\Omega + A}{\Omega + \epsilon/T} \right\rangle = 1 + \left\langle \frac{A - \epsilon/T}{\Omega + \epsilon/T} \right\rangle$$

Eq. (16) can be written as

$$\left\langle \frac{\Omega_{*T}(3/2 - 1/\eta)/(1-\Omega_{*T}) + \epsilon/T}{\Omega - \epsilon/T} \right\rangle - \left\langle \frac{\Omega_{*T}(3/2 - 1/\eta)/(1-\Omega_{*T}) + \epsilon/T}{\Omega + \epsilon/T} \right\rangle = \frac{2/f_t - 2}{1 - \Omega_{*T}}. \tag{17}$$

Equation (17) has complex solutions, $\Omega=\Omega_r+\gamma$. Consider marginal stability, $\Omega=\Omega_r=\Omega_0$, with $\Omega_0>0$. The singularity in Eq. (17) is removed when

$$\Omega_0=-\frac{\Omega_{*T}(3/2-1/\eta)}{1-\Omega_{*T}}. \tag{18}$$

Substituting into Eq. (17) we obtain an equation for Ω_0

$$\left\langle\frac{\Omega_0-\epsilon/T}{\Omega_0+\epsilon/T}\right\rangle=1+\frac{(2/f_t-2)}{(1-\Omega_{*T})}. \tag{19}$$

Since $\langle(\Omega_0-\epsilon/T)/(\Omega_0+\epsilon/T)\rangle<1$, Eq. (19) does not have a solution when $f_t=1$, which indicates that a flute-like mode would not be unstable. For a marginally stable solution to exist the right hand side of Eq. (19) must be <1 and therefore $\Omega_{*T}>1$. From Eq. (18) we observe that when $\Omega_{*T}>1$, marginal stability is obtained for $\eta>2/3$. A comparison with the solutions to the dispersion relation obtained in Ref. 10 indicates that this is a necessary condition for the existence of unstable solutions. Therefore, when $\eta>2/3$, the condition $\Omega_{*T}<1$ is a sufficient condition for stability. From the solutions that appear in Ref. 10 we observe that typically $0.5<\Omega_0<2$. Rewriting the stability condition we obtain

$$\hat{\omega}_d>\hat{\omega}_{*T} \tag{20}$$

with $\hat{\omega}_{*T}=\eta\hat{\omega}_*$. Normally $\eta>1$ and this condition is less restrictive than the interchange condition: $|\hat{\omega}_d|>\hat{\omega}_*(1+\eta)$.

VI. DISSIPATIVE TRAPPED ION MODE

In the simplest approximation the dissipative trapped ion mode can be derived by assuming collisional electrons and collisionless ions, i.e., by taking the limit $\nu_e\rightarrow\infty$ and $\nu_i\rightarrow 0$. In this limit Eq. (17) would be replaced by

$$\left\langle\frac{\epsilon/T+\Omega_{*T}(3/2-1/\eta)/(1-\Omega_{*T})}{\epsilon/T-\Omega}\right\rangle=\xi, \tag{21}$$

with $\xi=(2/f_t-1)/(\Omega_{*T}-1)$. The marginally stable solution to Eq. (21) is $\Omega=\Omega_0=(3/2-1/\eta)/(\Omega_{*T}-1)$ and $\xi=1$, i.e., $\Omega_{*T}=2/f_t$ (as pointed out in Ref. 11). Notice that this solution of Eq. (21) is possible when $f_t=1$ which corresponds to a flute-like mode. This indicates that the dissipative mode is driven by the difference in the (collisional) electron and the (collisionless) ion response and instability does not require a localization of the mode in the trapped particle region.

If we take $\Omega=\Omega_0+\delta\Omega+i\gamma$ and $\xi=1+\Delta$ with $\Delta>0$ the imaginary part yields a constraint

$$\left\langle\frac{\Omega_0+\epsilon/T}{\Omega^2+\gamma^2}\right\rangle=0, \tag{22}$$

while the real part gives

$$\Delta=-\left\langle\frac{\delta\Omega^2+\gamma^2}{\Omega^2+\gamma^2}\right\rangle,$$

i.e., $\Delta<0$. This indicates that there is no unstable solution when $\xi>1$, i.e., when $\Omega_{*T}<2/f_t$. Thus a sufficient condition for stability is $\hat{\omega}_d>\omega_{*T}/2$. This condition is less restrictive than the stability condition for either the collisionless mode or the interchange mode.

VII. DISCUSSION

The results shown above are generally applicable to a magnetic confinement device that that is stabilized by compressibility, i.e., one that satisfy's the inequality $\widetilde{\omega}_d>\omega_*$. For a tokamak $\omega_*/\overline{\omega_d}\sim A$ with A the aspect ratio, so that the compressibility is usually considered to be a small correction. For a low aspect ratio tokamak, however, the stabilizing compressibility term can become important.

The dipole reactor concept is a radical departure from the tokamak or from other similar toroidal magnetic fusion reactor concepts. The magnetic field lines are closed, the field is poloidal and the flux surfaces are defined by the toroidal drifts. There are no drifts off the flux surfaces and the dipole is not subject to neoclassical effects. In addition the high degree of axisymmetry inherent in the coil set insures the absence of nonaxisymmetry driven ''ripple'' losses.

Hasegawa[2] considered a collisionless plasma confined in a dipole that is characterized by an equilibrium distribution function which is non-Maxwellian. He showed that when the equilibrium distribution function can be characterized by $F_0=F_0(\mu,J)$, i.e., $\partial F_0/\partial\psi=0$, drift frequency modes are not unstable.

We have shown that in the more fusion relevant case, when F_0 is Maxwellian and therefore $\partial F_0/\partial\psi<0$, the dipole may still not be subject to drift frequency fluctuations. Therefore a magnetic dipole based fusion reactor may exhibit classical transport.

ACKNOWLEDGMENTS

The author would like to thank Professor M. Mauel (Columbia University) for many stimulating discussions about dipole confinement and Dr. B. Lane (Plasma Dynamics, Belmont, Mass) for insights about the gyrokinetics.

This research was supported by the U.S. Department of Energy.

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