REPORT NO.
LAPS-33
SAI 77-983-LJ
December, 1977

LABORATORY FOR APPLIED PLASMA STUDIES

TEMPERATURE GRADIENT AND ELECTRON GYRORADIUS EFFECTS ON LOWER HYBRID DRIFT-DRIFT CYCLOTRON INSTABILITIES

by

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Temperature Gradient and Electron Gyroradius Effects on Lower Hybrid Drift - Drift Cyclotron Instabilities *

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ABSTRACT

Electron and ion temperature gradients as well as finite electron gyroradius ($k\rho_e \gtrsim 1$) have been retained to analyze the stability of drift waves from the lower hybrid down to the ion cyclotron frequency. The instability conditions indicate the importance of these modes over this entire frequency range in a variety of applications.

*This work was supported under Contract #EY-76-C-03-1018 for the Department of Energy.

Due to the large number of practical situations in which the lower hybrid drift instability is expected to play an important role in the dynamics of plasma experiments, an extensive amount of theoretical work has been done to understand its properties in detail. Early calculations[1,2] considered the frequency range $\omega_{ci} < \omega < \omega_{ce}$, and thus treated the electrons as strongly magnetized but the ions as unmagnetized on the instability time scale. Recently Freidberg and Gerwin[3] allowed for magnetized ions to treat lower frequency and lower drift speed regions where $\omega \leqslant \omega_{ci}$. Keeping only density gradients to drive drift waves they demonstrated that as the drift speed decreases (density gradient scale length increases) the dispersion relation for lower hybrid modes goes over into that for drift cyclotron modes indicating the associated importance of understanding these latter modes in environments where lower hybrid instabilities evolve steep plasma and magnetic profiles into broad profiles. They then gave a relation between the diamagnetic and ion thermal speeds necessary for linear instability at $\omega \simeq \omega_{ci}$. Catto and Aamodt[4] have recently looked at drift cyclotron instabilities including ion thermal gradients, and arbitrary loss cone fill. For no loss cone, which is the case of interest to us here, they find that ion gradients can have an influence on stability. In particular they find that profiles in which the density gradient and ion thermal gradients are oppositely directed are always unstable. This prompts them to speculate that drift cyclotron modes may be at work, for example, in tokamak gas puffing experiments to admit gas introduced at the outside into the plasma core. Their analysis focuses on wavelengths much larger than an electron gyroradius, $k^2 \rho_e^2 << 1$. Because of the growing variety of applications (pinches, liner implosions of plasma, tokamak, etc.) for the lower hybrid drift-drift cyclotron modes, and the sensitivity of numerical modeling studies to linear instability conditions, we present here results which extend the treatments of Refs. 3 and 4 by allowing for both electron and ion temperature gradients as well as density gradients, retaining finite electron gyroradius to treat the important case $k \rho_e \sim 1$ and thus unstable systems in which $T_e \gtrsim T_i$. We also retain the effects of gradient B drifts to lowest order. Principal results are (1) the instability condition in the drift cyclotron limit when $\omega \simeq \ell \omega_{ci}$, ℓ an integer, is identical in form to that in the lower hybrid limit, $|\omega| >> \omega_{ci}$. (2) Electron temperature gradients are an important factor in the instability condition for $k \rho_e \gtrsim 1$ and, therefore influence the stability of many systems with $T_e \gtrsim T_i$. When the electron and ion thermal gradients are of the same size these modes are always unstable. (3) At least some wavelengths of the drift cyclotron instability can usually be expected to be unstable whenever it is possible to maintain $\omega = \omega_{ci}$ during the course of an experiment.

We consider systems in which curvature is weak or absent (the effect of curvature on these modes has been treated previously[5,6]) and ignore the localizing effects of magnetic shear which could be important in some applications (the shear problem has been treated for the lower hybrid modes[7] but not as yet for the drift cyclotron modes). We also treat only electrostatic perturbations and thus lose finite beta information.

The equilibrium distribution is[5]

$$f_{o\sigma} = f_{M\sigma}\left[1 + \left(\epsilon'_\sigma + \alpha_\sigma \delta_\sigma v^2\right)\left(x + \frac{v_y}{\omega_{c\sigma}}\right)\right] \tag{1}$$

where

$$f_{M\sigma} = n_o \left(\alpha_\sigma/\pi\right)^{\frac{3}{2}} e^{-\alpha_\sigma v^2}$$

$$\alpha_\sigma = m_\sigma/2T_\sigma$$

$$\delta_\sigma = \frac{1}{T_\sigma}\frac{dT_\sigma}{dx} \tag{2}$$

$$\epsilon'_\sigma = \frac{1}{n}\frac{dn}{dx} - \frac{3}{2}\delta_\sigma$$

and $\sigma = e, i$ refers to electrons, ions. We do not include electric fields explicitly since these can be Doppler shifted away for the frequencies of interest here. The perturbed distribution function obtained from the Vlasov equation linearized about the equilibrium of Eq. (1) for $k_\parallel = 0$ in a slab geometry is the familiar[5]

$$\delta f_\sigma = - \frac{q_\sigma f_{M\sigma}}{T_\sigma} \delta\varphi \left\{ 1 - \left[\omega - \frac{k v_\sigma^2}{\omega_{c\sigma}} \left(\epsilon'_\sigma + \alpha_\sigma \delta_\sigma v^2 \right) \right] \sum_{\ell m} e^{i(\ell - m)(\pi/2 - \theta)} \right.$$

$$\left. \cdot \frac{J_\ell\left(\frac{k v_\perp}{\omega_{c\sigma}}\right) J_m\left(\frac{k v_\perp}{\omega_{c\sigma}}\right)}{\left(\omega - \ell\omega_{c\sigma} - k \frac{\epsilon v_\perp^2}{2\omega_{c\sigma}}\right)} \right\} \tag{3}$$

where

$$\epsilon = \frac{1}{B} \frac{dB}{dx} , \qquad \omega_{c\sigma} = q_\sigma B / m_\sigma c$$

The relevant quantities are the electron and ion density perturbations δn_σ . These are found to be

$$\delta n_\sigma = - \frac{q_\sigma n_0}{T_\sigma} \delta\varphi \left\{ 1 - \sum_\ell \frac{e^{-b_\sigma} I_\ell(b_\sigma)}{(\omega/\omega_{c\sigma} - \ell)} \left[\omega - \omega_\sigma^* - \omega_\sigma^* \eta_{T\sigma} b_\sigma \left(\frac{I'_\ell(b_\sigma)}{I_\ell(b_\sigma)} - 1 \right) \right. \right.$$

$$\left. \left. + \omega_\sigma^* \eta_B \left(1 + b_\sigma \left(\frac{I'_\ell(b_\sigma)}{I_\ell(b_\sigma)} - 1 \right) \right) \right] \right\} \tag{4}$$

where

$$\omega_\sigma^* = k \frac{v_\sigma^2}{\omega_{c\sigma}} \frac{1}{n} \frac{dn}{dx} ,$$

$$v_\sigma^2 = T_\sigma / m_\sigma \ ,$$

$$\eta_{T\sigma} = \frac{1}{T_\sigma} \frac{d\,T_\sigma}{dx} \bigg/ \frac{1}{n} \frac{dn}{dx} \ ,$$

$$\eta_B = \frac{1}{B} \frac{dB}{dx} \bigg/ \frac{1}{n} \frac{dn}{dx} \ ,$$

$$b = k^2 \rho_\sigma^2 \ ,$$

$$\rho_\sigma^2 = v_\sigma^2 / \omega_{c\sigma}^2$$

We have kept only the lowest order term in ϵ from Eq. (3) assuming $\omega - \ell\,\omega_{c\sigma} >> k\,\epsilon\,v_\sigma^2 / 2\,\omega_{c\sigma}$. This is well known to be a low beta approximation.[5] For the ions, we are interested in $k^2 \rho_i^2 >> 1$, and for the electrons we are interested in $\omega << \omega_{ce}$ and thus retain only $\ell = 0$ for the electrons. Therefore, for the ions

$$\delta n_i \simeq - \frac{e\,n_o\,\delta\varphi}{T_i} \left\{ 1 - \sum_\ell \frac{1}{\sqrt{2\pi}\,k\,v_i} \frac{1}{\left(\omega/\omega_{ci} - \ell\right)} \left[\omega - \omega_i^* \left(1 - \frac{\eta_i}{2} \frac{\eta_B}{2} \right) \right] \right\} \qquad (5a)$$

Using the identity

$$\sum_\ell \frac{1}{\omega/\omega_{ci} - \ell} = \pi \cot\left(\pi\,\omega/\omega_{ci}\right)$$

it is useful to write δn_i alternatively as

$$\delta n_i \simeq - \frac{e\, n_o\, \delta\varphi}{T_i} \left\{ 1 - \sqrt{\frac{\pi}{2}}\, \frac{1}{k\, v_i} \cot\left(\frac{\pi\,\omega}{\omega_{ci}}\right) \left[\omega - \omega_i^* \left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right)\right]\right\} \tag{5b}$$

For the electrons we obtain the well known result[1, 2]

$$\delta n_e \simeq \frac{e\, n_o\, \delta\varphi}{T_e} \left[1 - e^{-b_e} I_o(b_e) + \frac{k}{\omega} V_\Delta\right] \tag{6}$$

$$k\, V_\Delta = \omega_e^* I_o(b_e)\, e^{-b_e} \left\{1 - \eta_B \left[1 - b_e \left(1 - \frac{I_1(b_e)}{I_o(b_e)}\right)\right] - \eta_{Te}\, b_e \left(1 - \frac{I_1(b_e)}{I_o(b_e)}\right)\right\}$$

The dispersion relation is obtained as usual from Poissons equation

$$1 + \chi_e + \chi_i = 0$$

$$\chi_\sigma = - \frac{4\pi\, e\, \delta n_\sigma}{k^2\, \delta\varphi} \tag{7}$$

We now consider separately two limits in which Eq. (7) simplifies to yield analytic information.

A. Lower Hybrid Drift Limit $\left(|\omega| >> \omega_{ci}\,,\ \ \gamma > \omega_{ci}\right)$

For drift speeds sufficiently large that the growth rate is greater than the ion gyrofrequency, $\cot\left(\pi\,\omega/\omega_{ci}\right) \to -i$ and

$$\delta n_i \simeq - \frac{e\, n_o\, \delta\varphi}{T_i} \left\{1 + i\sqrt{\pi/2}\, \frac{1}{k\, v_i} \left[\omega - \omega_i^* \left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right)\right]\right\} \tag{8}$$

This limit corresponds to unmagnetized ions on the instability time scale. Equation (8) is appropriate to drift speeds below the ion thermal speed, which is the case of interest to us here. The case of drift speeds greater than the ion thermal speed was originally discussed by Krall and Liewer[1] for application to the implosion phase of theta pinches. Davidson and Gladd[2] first showed that instability persists even when the drift speed falls below the ion thermal speed indicating the importance of the modes in the broader sheath, post implosion phase of pinch experiments. The necessary and sufficient conditions for the persistence of a lower hybrid drift instability obtained from Eqs. (6)-(8) are

$$\left\{1 - \eta_B\left[1 - b_e\left(1 - \frac{I_1(b_e)}{I_o(b_e)}\right)\right] - \eta_{Te}\, b_e\left(1 - \frac{I_1(b_e)}{I_o(b_e)}\right)\right\}\left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right) > 0 \tag{9a}$$

and

$$\tau\, e^{-b_e} I_o(b_e)\, \frac{\left\{1 - \eta_B\left[1 - b_e\left(1 - \frac{I_1(b_e)}{I_o(b_e)}\right)\right] - \eta_{Te}\, b_e\left(1 - \frac{I_1(b_e)}{I_o(b_e)}\right)\right\}}{\left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right)} < \tag{9b}$$

$$1 + \tau - e^{-b_e} I_o(b_e) + k^2 \lambda_e^2$$

where

$$\tau \equiv T_e/T_i \quad , \qquad \lambda_e^2 = v_e^2/\omega_{pe}^2 = \text{electron Debye length.}$$

We will examine these conditions more closely further on.

B. Drift Cyclotron Limit $\left(\omega \simeq \ell\,\omega_{ci}\;,\;\gamma < \omega_{ci}\right)$

In the weak drift limit when the growth rate falls below the ion gyrofrequency, we consider modes with $\omega \simeq \ell\,\omega_{ci}$ and retain only the resonant term in the Bessel series summation of Eq. (5a). In this case Eqs. (5a)-(7) yield a quadratic equation for ω,

$$A\,\omega^2 + B\,\omega + C = 0 \tag{10}$$

$$A \equiv 1 - \frac{\tau}{\sqrt{2\pi}\,k\rho_i}\;\frac{1}{1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2}$$

$$B \equiv \left[kV_\Delta + \frac{\tau\omega_i^*\left(1 - \frac{\eta_{T_i}}{2} - \frac{\eta_B}{2}\right)}{\sqrt{2\pi}\,k\rho_i}\right]\frac{1}{1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2} - \ell\,\omega_{ci}$$

$$C \equiv - \frac{k\,V_\Delta\,\ell\,\omega_{ci}}{1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2}$$

with the solutions

$$\omega \simeq \frac{1}{2A}\left(\ell\omega_{ci} - \frac{kV_\Delta + \frac{\tau\omega_i^*}{\sqrt{2\pi}\,k\rho_i}\left(1 - \frac{\eta_{T_i}}{2} - \frac{\eta_B}{2}\right)}{1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2}\right)$$

$$\pm \frac{1}{2A}\left\{\left(\ell\omega_{ci} + \frac{kV_\Delta + \frac{\tau\omega_i^*}{\sqrt{2\pi}\,k\rho_i}\left(1 - \frac{\eta_{T_i}}{2} - \frac{\eta_B}{2}\right)}{1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2}\right)^2 \right. \tag{11}$$

$$- 2\sqrt{\frac{2}{\pi}}\frac{1}{k\rho_i}\,\frac{\tau\ell\omega_{ci}}{\left(1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2\right)^2}\left[kV_\Delta + \omega_i^*\left(1 - \frac{\eta_{T_i}}{2} - \frac{\eta_B}{2}\right)\right.$$

$$\left.\left. \cdot \left(1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2\right)\right]\right\}^{\frac{1}{2}}$$

For $\eta_{T\sigma} = \eta_B = 0$, $k^2 \rho_e^2 << 1$, the drift cyclotron instability discussed in Ref. 5 is recovered. Examining Eq. (11) we see that instability first of all requires

$$\ell\omega_{ci} \simeq \frac{-kV_\Delta}{1 + \tau - e^{-b_e} I_o(b_e) + k^2\lambda_e^2} \tag{12}$$

which insures that $\omega \simeq \ell\, \omega_{ci}$ as assumed. When Eq. (12) is satisfied, we find further that the conditions of Eqs. (9a), (9b) are also necessary to give an instability; i.e., Eqs. (9a, b) and (12) taken together are the necessary and sufficient conditions for instability. Hence, the statement in the introduction that the mathematical condition for drift cyclotron instability at $\omega = \ell\, \omega_{ci}$ is identical to that for lower hybrid drift instability.

We now examine the instability condition more carefully in several limits to make contact with earlier work and to highlight new results:

(a) $b_e = k^2 \rho_e^2 << 1$

In the long wavelength limit, $b_e << 1$, Eq. (12) gives

$$\frac{L_n}{\rho_i} = \frac{\tau\, k\, \rho_i}{\ell} \frac{\left(1 - \eta_B\right)}{\tau + b_e\left(1 + \omega_{ce}^2/\omega_{pe}^2\right)} \tag{13a}$$

where we choose to write the result as a ratio of density gradient scale length, $L_n = [(1/n)(dn/dx)]^{-1}$, to ion gyroradius.

Equations (9a), (9b) give

$$\left(1 - \eta_B\right)\left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right) > 0 \tag{13b}$$

$$b_e > \frac{\tau}{2} \frac{\left(\eta_{Ti} - \eta_B\right)}{\left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right)} \frac{1}{\left(1 + \frac{\omega_{ce}^2}{\omega_{pe}^2}\right)} \tag{13c}$$

To make contact with earlier work, we disregard η_B at this point. Equation (13a) reduces to that of Freidberg and Gerwin[3] for $\eta_{Ti} = 0$ and $\omega_{ce}^2/\omega_{pe}^2 \to 0$. Eqs. (13b, c) are always satisfied in this case. However, for $\eta_{Ti}/2 > 1$, Eq. (13b) cannot be satisfied and all modes are stable for $b_e << 1$. If $\eta_{Ti}/2 < 1$, only those wavenumbers that satisfy Eq. (13c) are unstable. For example Eq. (13a) is maximum for fixed τ, ρ_i at $b_e = \tau$ in the high density limit yielding the maximum sheath thickness $L_n/\rho_i = (1/2)\sqrt{m_i/m_e}$ obtained by Freidberg and Gerwin.[3] However, Eq. (13c) requires that $\eta_{Ti} < 1$ for instability; if $\eta_{Ti} > 1$ there would be some decrease in L_n relative to ρ_i. Further, we note that $\eta_{Ti} < 0$ satisfies Eqs. (13b, c) and thus always leads to instability, as has recently been pointed out by Aamodt and Catto, provided Eq. (13a) is satisfied. We note that Eq. (13c) restricts the instability in this limit, $b_e << 1$, to systems with $T_e/T_i << 1$, unless $\omega_{ce}^2 >> \omega_{pe}^2$, or $\eta_{Ti} \lesssim 0$.

(b) $b_e = k^2 \rho_e^2 >> 1$

In the short wavelength limit, $b_e >> 1$, Eq. (12) gives

$$\frac{L_n}{\rho_i} = \sqrt{\frac{\tau m_i}{2\pi m_e}} \; \frac{\left(1 - \frac{\eta_{Te}}{2} - \frac{\eta_B}{2}\right)}{\ell\left(1 + \tau + k^2 \lambda_e^2\right)} \tag{14a}$$

Equations (9a, b) give

$$\left(1 - \frac{\eta_{Te}}{2} - \frac{\eta_B}{2}\right)\left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right) > 0 \qquad (14b)$$

$$\frac{\tau}{\sqrt{2\pi}\, k\rho_e} \frac{\left(1 - \frac{\eta_{Te}}{2} - \frac{\eta_B}{2}\right)}{\left(1 - \frac{\eta_{Ti}}{2} - \frac{\eta_B}{2}\right)} < 1 + \tau + k^2 \lambda_e^2 \qquad (14c)$$

Setting $\eta_{T\sigma} = \eta_B = \lambda_e^2 = 0$, Eq. (14c) reduces to the result of Gerwin and Freidberg.[3] We note, however, that the electron temperature gradients are important in determining sheath thicknesses in applications where $\eta_{Te} \sim 1$. Further note that Eq. (14c) is nonrestrictive. Again we see that temperature gradients oppositely directed to density gradients are always unstable at low β according to Eq. (14b), and Eqs. (14a-c) can be satisfied and thus predict instability for systems in which $T_e \sim T_i$ even for $\eta_{Ti} \sim 1$. This reinforces the comments made by Catto and Aamodt[4] regarding the possible applicability of drift cyclotron modes for tokamak in certain cases. In addition to the examples cited by Catto and Aamodt[4] in which these modes could possibly play a role in tokamak experiments with $\eta_{T\sigma} < 0$, we speculate that they could be present during the current penetration phase of the discharge when the temperature profiles are peaked off axis, and if so, might provide a mechanism for anomalous current penetration. The role of magnetic shear on these modes must be understood, however, to provide a basis for their application to tokamak geometry.

(c) $b_e = k^2 \rho_e^2 = 1$

As a final example, we give the instability conditions for $b_e = 1$ to compare with the extreme limits $b_e << 1$, $b_e >> 1$ discussed above. In this case Eq. (12) gives

$$\frac{L_n}{\rho_i} = 0.47 \sqrt{\tau \frac{m_i}{m_e}} \frac{(1 - 0.45\, \eta_B - 0.55\, \eta_{Te})}{\ell \left(\tau + 0.53 + \omega_{ce}^2/\omega_{pe}^2\right)} \qquad (15a)$$

Equations (9a) and (9b) give

$$(1 - 0.45\, \eta_B - 0.55\, \eta_{Te})\ (1 - 0.5\, \eta_B - 0.5\, \eta_{Ti}) > 0 \qquad (15b)$$

$$0.47\, \tau \frac{(1 - 0.45\, \eta_B - 0.55\, \eta_{Te})}{(1 - 0.5\, \eta_B - 0.5\, \eta_{Ti})} < \tau + 0.53 + \omega_{ce}^2/\omega_{pe}^2 \qquad (15c)$$

Note that these results for $b_e = 1$ are similar to those in the limit $b_e >> 1$. In particular Eq. (15b) is very similar to Eq. (14c) and Eq. (15c) is nonrestrictive as is Eq. (14c).

In summary, we have found that the instability conditions for the lower hybrid drift and drift cyclotron instabilities in terms of the various plasma gradients are identical [Eqs. (9a, b)]. The lower hybrid drift limit, $|\omega| >> \omega_{ci}$, corresponds to relatively larger drift speed (larger gradients), e.g., for $k^2 \rho_e^2 << 1$, $\omega_{ce}^2/\omega_{pe}^2 << 1$, the condition $\gamma > \omega_{ci}$ requires[2] $L_n/\rho_i < (\sqrt{\pi}/2)^{\frac{1}{2}} (m_i/m_e)^{\frac{1}{4}}$. The drift cyclotron instability is subject to the constraint $\omega = \ell\,\omega_{ci}$ which is given by Eq. (12), and thus can occur for relatively weaker drifts (milder gradients). Modes with $k^2 \rho_e^2 << 1$ require large ion temperature gradients, $\eta_{Ti} > 2$, for stability and instability in this limit is restricted to $T_e/T_i << 1$ unless $\omega_{ce}^2 >> \omega_{pe}^2$, or $\eta_{Ti} \leqslant 0$ [see Eqs. (13b, c)]. For $k^2 \rho_e^2 \gtrsim 1$ [see Eqs. (14) and (15)], even strong temperature gradients cannot stabilize these modes if the ion and electron temperature gradients are comparable. Temperature gradients opposite to density gradients are unstable for all values of $k^2 \rho_e^2$ and only in special cases are all wavelengths linearly stable, e.g., $\eta_{Ti} > 2$ but $\eta_{Te} < 2$. Thus, the drift waves discussed in this note are potentially important over the whole frequency range from the lower hybrid down to the ion cyclotron frequency in a variety of situations in which finite β, shear and favorable curvature are not completely stabilizing.

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