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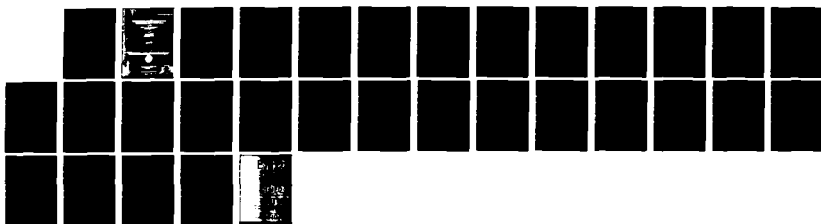
SHORT WAVELENGTH STABILIZATION OF THE GRADIENT DRIFT
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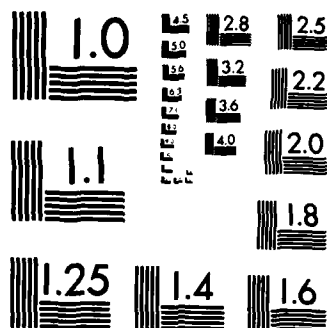
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SHORT WAVELENGTH STABILIZATION OF THE GRADIENT DRIFT INSTABILITY DUE TO VELOCITY SHEAR

I. INTRODUCTION

The gradient drift instability (also known as the $\mathbf{E} \times \mathbf{B}$ instability or the $\mathbf{E} \times \mathbf{B}$ gradient drift instability) has been of interest to space and plasma physicists for two decades. The instability was initially studied by Simon (1963) and Hoh (1963) to explain turbulence observed in laboratory experiments. Since then the dominant interest in the instability has been in regard to natural ionospheric irregularities [e.g., equatorial electrojet irregularities (Farley, 1979; Fejer and Kelley, 1980) and high latitude irregularities (Keskinen and Ossakow, 1982)] and to man-made ionospheric irregularities [e.g., plasma cloud releases (Linson and Workman, 1970)]. The instability can arise in a weakly collisional plasma which contains a density gradient that is orthogonal to the current flow generated by an ambient electric field perpendicular to the magnetic field. Depending upon the ratios ν_e/Ω_e and ν_i/Ω_i , where $\nu_{e(i)}$ is the electron (ion)-neutral collision frequency and $\Omega_{e(i)}$ is the electron (ion) cyclotron frequency, an ambient electric field can produce two types of current in a weakly collisional plasma. For the case of $\nu_e/\Omega_e \ll 1$ and $\nu_i/\Omega_i \ll 1$, a Pedersen current is produced by the ions in the direction of the electric field; when $\nu_e/\Omega_e \ll 1$ and $\nu_i/\Omega_i \gg 1$, a Hall current is produced by the electrons in the direction orthogonal to the electric field. The Pedersen current driven instability is relevant to F region irregularities, while the Hall current driven instability is relevant to E region irregularities. This letter will discuss the influence of velocity shear on the gradient drift instability driven by the Hall current.

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The motivation for this paper is two-fold. First, studies of the Rayleigh-Taylor instability (Guzdar et al., 1982) and the Pedersen current driven gradient drift (or $\mathbf{E} \times \mathbf{B}$) instability (Perkins and Doles, 1975; Huba et al., 1983) indicate that velocity shear can have an important effect on these instabilities. Namely, the stabilization (or suppression) of the short wavelength modes which causes a preferential excitation of a longer wavelength mode. Thus, it is anticipated that velocity shear will have a similar effect on the Hall current driven gradient drift instability, and, in fact, we show that this is the case. Second, recent studies of long wavelength turbulence in the equatorial electrojet (Kudeki et al., 1982; Pfaff et al., 1982) indicate that a long wavelength mode with $\lambda \sim$ few km is dominant during type I conditions. Thus, inclusion of velocity shear in the plasma configuration of the electrojet may account for this observation.

II. THEORY

The plasma configuration and slab geometry used in the analysis are shown in Fig. 1. The ambient magnetic and electric fields are in the z and $-x$ directions, respectively (i.e., $\underline{B} = B_0 \hat{e}_z$ and $\underline{E} = -E_0(x) \hat{e}_x$). The density and electric field are assumed to be inhomogeneous in the x direction. The collision frequencies are such that $\nu_e \ll \Omega_e$ and $\nu_i \gg \Omega_i$ where ν_α is the α specie-neutral collision frequency (assumed to be constant) and $\Omega_\alpha = |e_\alpha| B_0 / m_\alpha c$ is the cyclotron frequency of the α species. The electrons can be considered to be strongly magnetized, while the ions are unmagnetized. Thus, the electrons experience a sheared $\underline{E} \times \underline{B}$ drift in the y direction ($\underline{V}_d(x) = -cE_0(x)/B_0 \hat{e}_y$) while the ions do not. This differential motion between the electrons and ions produces a Hall current.

We assume that perturbed quantities vary as $\delta p \sim \delta p(x) \exp[i(k_y y - \omega t)]$ where k_y is the wave number in the y direction and $\omega = \omega_r + i\gamma$ so that $\gamma > 0$ implies wave growth. We neglect perturbations along the magnetic field ($\underline{k} \cdot \underline{B}_0 = 0$) so that only the two dimensional mode structure in the xy plane is obtained. We assume $\omega \ll \Omega_e$ so that electron inertia terms can be neglected. An important feature of our analysis is that a nonlocal theory is developed. That is, the mode structure of the potential in the x direction (the direction in which the density and electron drift velocity are assumed to vary) is determined by a differential equation rather than an algebraic equation obtained by Fourier analysis. This is crucial to the analysis since previous studies of interchange instabilities using a sheared velocity flow equilibrium [Perkins and Doles, 1975; Guzdar et al., 1981; Huba et al., 1983] have shown that a nonlocal theory is necessary to demonstrate the stabilizing influence of velocity shear.

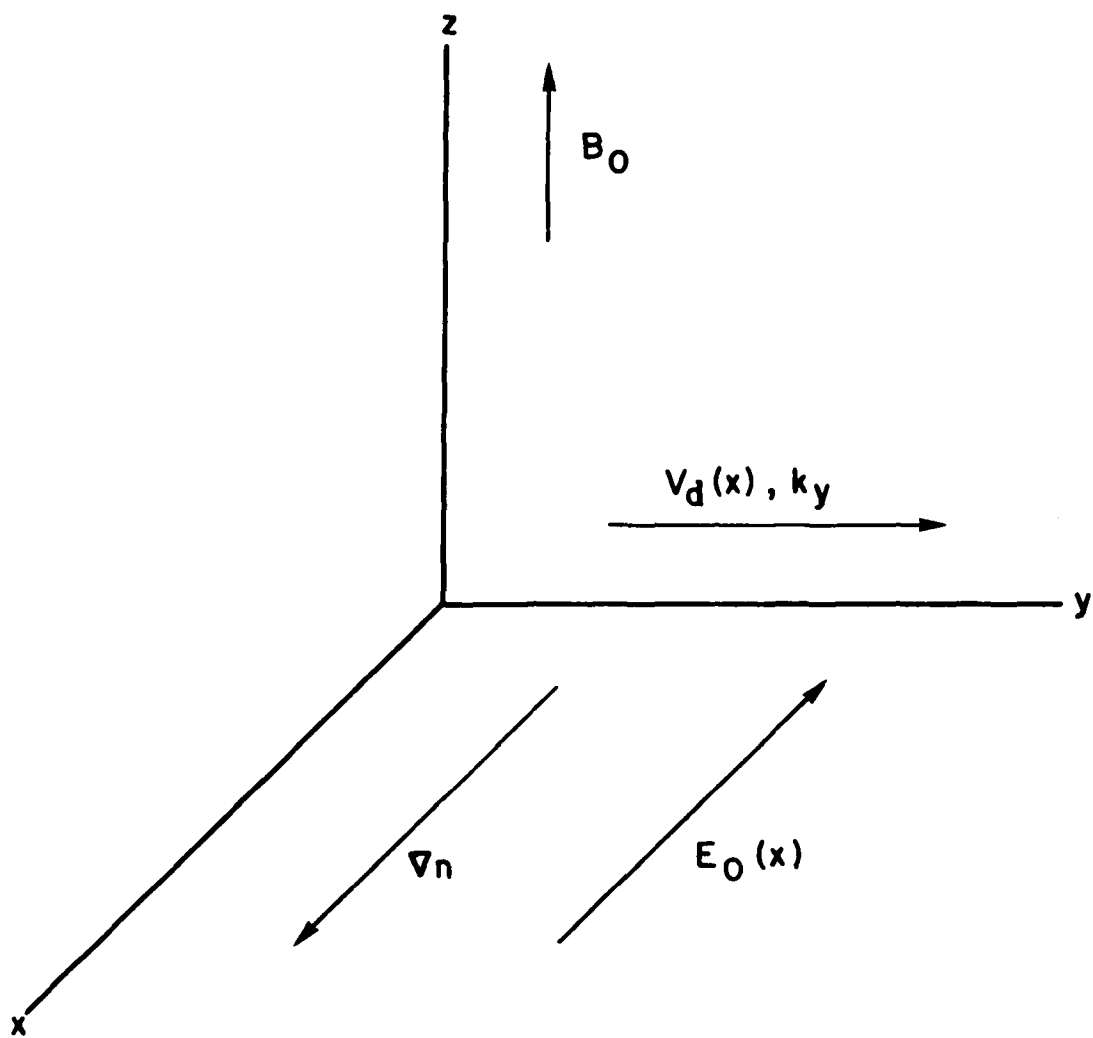


Fig. 1 Plasma configuration and slab geometry.

The fundamental equations used in the analysis are continuity and momentum transfer in the frame of reference of the neutrals:

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot (n_\alpha \mathbf{v}_\alpha) = 0 \quad (1)$$

$$0 = -\frac{e}{m_e} \left(\mathbf{E} + \frac{1}{c} \mathbf{v}_e \times \mathbf{B}_0 \right) - \frac{T_e}{nm_e} \nabla n - \nu_e \mathbf{v}_e \quad (2)$$

$$\frac{d\mathbf{v}_i}{dt} = \frac{e}{m_i} \mathbf{E} - \frac{T_i}{nm_i} \nabla n - \nu_i \mathbf{v}_i \quad (3)$$

where α denotes particle species (e: electrons; i: ions) and other variables have their usual meanings. From Eqs. (2) and (3) we find that the equilibrium drifts of the electrons and ions, to first order in ν_e/Ω_e and Ω_i/ν_i , are given by

$$\mathbf{v}_e = \frac{\nu_e}{\Omega_e} (V_d - v_{de}) \hat{e}_x + (V_d - v_{de}) \hat{e}_y \quad (4)$$

and

$$\mathbf{v}_i = -\frac{\Omega_i}{\nu_i} (V_d - v_{di}) \hat{e}_x \quad (5)$$

where $V_d = -cE_0/B_0$ is the $\mathbf{E} \times \mathbf{B}$ velocity and $v_{d\alpha} = (cT_\alpha/e_\alpha B_0) \partial \ln n / \partial x$ is the diamagnetic drift velocity of species α . An equilibrium relationship between n and V_d can be obtained by substituting Eqs. (4) and (5) into Eq. (1) and setting $\partial n_\alpha / \partial t = 0$. Assuming $V_d \gg v_{di}$, the equilibrium relation between n and V_d is

$$n(x) V_d(x) = \text{constant} \quad (6)$$

where we have neglected the x variation of the collision frequencies. We hasten to add that Eq. (6), although a self-consistent relationship within the context of Eqs. (1) - (3), does not necessarily model the equatorial electrojet or auroral plasmas.

We linearize Eqs. (1)-(3) by assuming $n_\alpha = n_\alpha + \delta n_\alpha$, $\underline{v}_\alpha = \underline{v}_\alpha + \delta \underline{v}_\alpha$ and $\underline{E} = \underline{E}_0 - \nabla \phi$ where ϕ is the perturbed electrostatic potential. Since we assume $v_e \ll \Omega_e$, $v_i \gg \Omega_i$, and $V_d \gg v_{d\alpha}$, we take the equilibrium drifts to be, to lowest order,

$$\underline{v}_e = V_d(x) \hat{e}_y = - c E_0(x) / B_0 \hat{e}_y \quad (7)$$

$$\underline{v}_i = 0. \quad (8)$$

The perturbed velocities are given by

$$\begin{aligned} \delta \underline{v}_e = & \left[- \frac{ik_y}{\Omega_e} \left(\frac{e}{m_e} \delta \phi - v_e^2 \frac{\delta n_e}{n} \right) + \frac{v_e}{\Omega_e} \frac{ik_x}{\Omega_e} \left(\frac{e}{m_e} \phi - v_e^2 \frac{\delta n_e}{n} \right) \right] \hat{e}_x \\ & + \left[\frac{ik_x}{\Omega_e} \left(\frac{e}{m_e} \delta \phi - v_e^2 \frac{\delta n_e}{n} \right) + \frac{v_e}{\Omega_e} \frac{ik_y}{\Omega_e} \left(\frac{e}{m_e} \phi - v_e^2 \frac{\delta n_e}{n} \right) \right] \hat{e}_y \end{aligned} \quad (9)$$

$$\delta \underline{v}_i = - \frac{ik_y}{v_i} \left(\frac{e}{m_i} \phi + v_i^2 \frac{\delta n_i}{n} \right) \quad (10)$$

where $v_\alpha = (T_\alpha / m_\alpha)^{1/2}$ is the thermal speed of species α and $\underline{k} = k_x \hat{e}_x + k_y \hat{e}_y = -i \partial / \partial x \hat{e}_x + k_y \hat{e}_y$, i.e., k_x is an operator. In writing Eqs. (9) and (10) we have assumed $k_y^2 L_n^2 \gg 1$ and $k_x^2 L_n^2 \gg 1$ where $L_n = (\partial \ln n / \partial x)^{-1}$, so that we are only considering unstable modes which have wavelengths small compared to the scale length of the density inhomogeneity.

We substitute Eqs. (9) and (10) into Eq. (1) and obtain

$$[-i(\omega - k_y v_d) + k^2 \rho_e^2 v_e + i k_y \rho_e v_e / L_n] \frac{\delta n_e}{n} =$$

$$\left[i k_y \frac{1}{L_n} + k^2 \frac{v_e}{\Omega_e} \right] \frac{e}{m_e} \frac{1}{\Omega_e} \phi \quad (11)$$

and

$$[\omega(\omega + i v_i) - k^2 v_i^2] \frac{\delta n_i}{n} = \frac{e}{m_i} k^2 \phi. \quad (12)$$

where $\rho_e = v_e / \Omega_e$ and $L_n = (\partial \ln n / \partial x)^{-1}$. Assuming quasi-neutrality ($\delta n_e = \delta n_i$) we find that Eqs. (11) and (12) reduce to

$$i \frac{v_e}{\Omega_e} c_s^2 k^4 \phi + [(\omega - k_y v_d) \Omega_i - i \omega(\omega + i v_i) \frac{v_e}{\Omega_e} - \frac{k_y}{L_n} c_s^2] k^2 \phi + \frac{i k_y}{L_n} v_i \omega \phi = 0 \quad (13)$$

Finally, we assume $k_x^2 \ll k_y^2$ and expand k^4 about the small parameter k_x^2/k_y^2 , retaining terms to second order in k_x . We then make the identification $k_x^2 = -\partial^2/\partial x^2$ and arrive at the mode equation

$$\frac{\partial^2 \phi}{\partial x^2} - k_y^2 Q(x) \phi = 0 \quad (14)$$

where $Q(x) = q(x)/p(x)$ and

$$p(x) = \omega - k_y v_d(x) - k_y^2 \rho_i^2 [(k_y L_n)^{-1} - 2i \tilde{v}_e] \Omega_i - i \omega(\omega / \Omega_i + i \tilde{v}_i) \tilde{v}_e \quad (15)$$

$$q(x) = \omega - k_y v_d(x) + [\omega(\omega / \Omega_i + i \tilde{v}_i) - k_y^2 \rho_i^2 \Omega_i] [(k_y L_n)^{-1} - i \tilde{v}_e] \quad (16)$$

where $\rho_i = c_s / \Omega_i$, $\tilde{v}_i = v_i / \Omega_i$ and $\tilde{v}_e = v_e / \Omega_e$.

III. RESULTS

In order to understand the influence of velocity shear on the gradient drift instability we contrast two cases. In the first case we assume $E_0 = \text{constant}$ (i.e., $V_d = \text{constant}$; no shear) and in the second case we consider $E_0 \neq \text{constant}$ (i.e., $V_d = V_d(x)$; shear).

A. No shear ($V_d = \text{constant}$)

We assume $E_0 = \text{constant}$ so that $V_d = \text{constant}$, and use local theory previous authors have considered (e.g., Rogister and D'Angelo, 1970; Si et al., 1973), i.e., Eq. (6) is violated. The local approximation assumes $k_x^2 L_n^2 \gg 1$ and $k_y^2 L_n^2 \gg 1$, i.e., the wavelengths of the unstable modes are small compared to the density inhomogeneity. The dispersion equation (Eq. (13)) can be written as

$$\omega(1 + \tilde{v}_e \tilde{v}_i) - k_y V_d + i(\omega \frac{k_y}{k} \frac{\tilde{v}_i}{k L_n} + \frac{\tilde{v}_e}{\Omega_i} (\omega^2 - k^2 c_s^2)) = 0 \quad (17)$$

where we have assumed $\tilde{v}_e \gg 1/k L_n$ for simplicity. Assuming $|\gamma| \ll |\omega_r|$, we find the solution of Eq. (17) to be

$$\omega_r = k_y V_d / (1 + \tilde{v}_i \tilde{v}_e) \quad (18)$$

$$\gamma = - \frac{\tilde{v}_e \tilde{v}_i}{1 + \tilde{v}_e \tilde{v}_i} \left[\frac{\omega_r}{\tilde{v}_e} \frac{k_y}{k} \frac{1}{k L_n} - \frac{1}{\tilde{v}_i \Omega_i} (\omega^2 - k^2 c_s^2) \right] \quad (19)$$

Equations (18) and (19) agree with earlier results (Rogister and D'Angelo, 1970; Sudan et al., 1973) and describe both the Farley-Buneman and gradient drift instabilities. The Farley-Buneman instability can occur when $\omega > k c_s$ or $V_d > c_s(1 + \tilde{v}_i \tilde{v}_e)$. On the other hand, the gradient drift instability is excited when

$$\frac{k_y (cE_0/B)}{1 + \tilde{v}_i \tilde{v}_e} \frac{1}{\tilde{v}_e} \frac{k_y}{k} \frac{1}{kL_n} > 0 \quad (20)$$

for $k\rho_i \ll 1$. When $V_d < c_s(1 + \tilde{v}_i \tilde{v}_e)$ and $k\rho_i \sim 1$, the gradient drift modes are stabilized because of electron-neutral collision (i.e., diffusion damping). The important condition represented by Eq. (20) is that the electric field and density gradient be in the same direction for instability.

We plot the growth rate γ_m/Ω_i vs. $k_y\rho_i$ (and k_yL) in Fig. 2 (curves A and B) for the case of no shear. Equation (13) is solved numerically for the following parameters: $V_d/c_s = -1.5$ (curve A) and $V_d/c_s = -0.2$ (curve B), $\rho_i/L = 10^{-3}$, $L_n/L = 1.0$, $\nu_e/\Omega_e = 10^{-2}$, and $\nu_i/\Omega_i = 25.0$. Here, γ_m is the growth rate maximized with respect to k_x/k_y . We note the following. First, $\gamma_m \propto k_y$ for $k_y\rho_i \lesssim 0.04$ (or $k_yL \lesssim 40$) in both curves A and B. Second, $\gamma_m \sim \text{constant}$ for $0.04 \lesssim k_y\rho_i \lesssim 0.40$ (or $40 \lesssim k_yL \lesssim 400$). And finally, γ_m increases sharply for $k_y\rho_i > 1$ when $V_d/c_s = -1.5$ which is due to the Farley-Buneman instability (curve A); γ_m becomes negative ($\gamma_m < 0$) for $k_y\rho_i \sim 0.65$ when $V_d/c_s = -0.2$ which is due to the collisional electron damping of the gradient drift instability (curve B).

B. Shear ($V_d \neq \text{constant}$)

We assume $E_0 \neq \text{constant}$ so that $V_d = V_d(x)$ and a sheared velocity flow exists in the plasma. As in the case of the $\mathbf{E} \times \mathbf{B}$ instability [Perkins and Doles, 1975; Huba et al., 1983] local theory is inadequate to properly describe the linear properties of the mode. We solve Eq. (14) numerically using a finite difference scheme [i.e., the Numerov method (Gladd and Horton, 1973)] to obtain eigenvalues and eigenfunctions. The boundary conditions used are

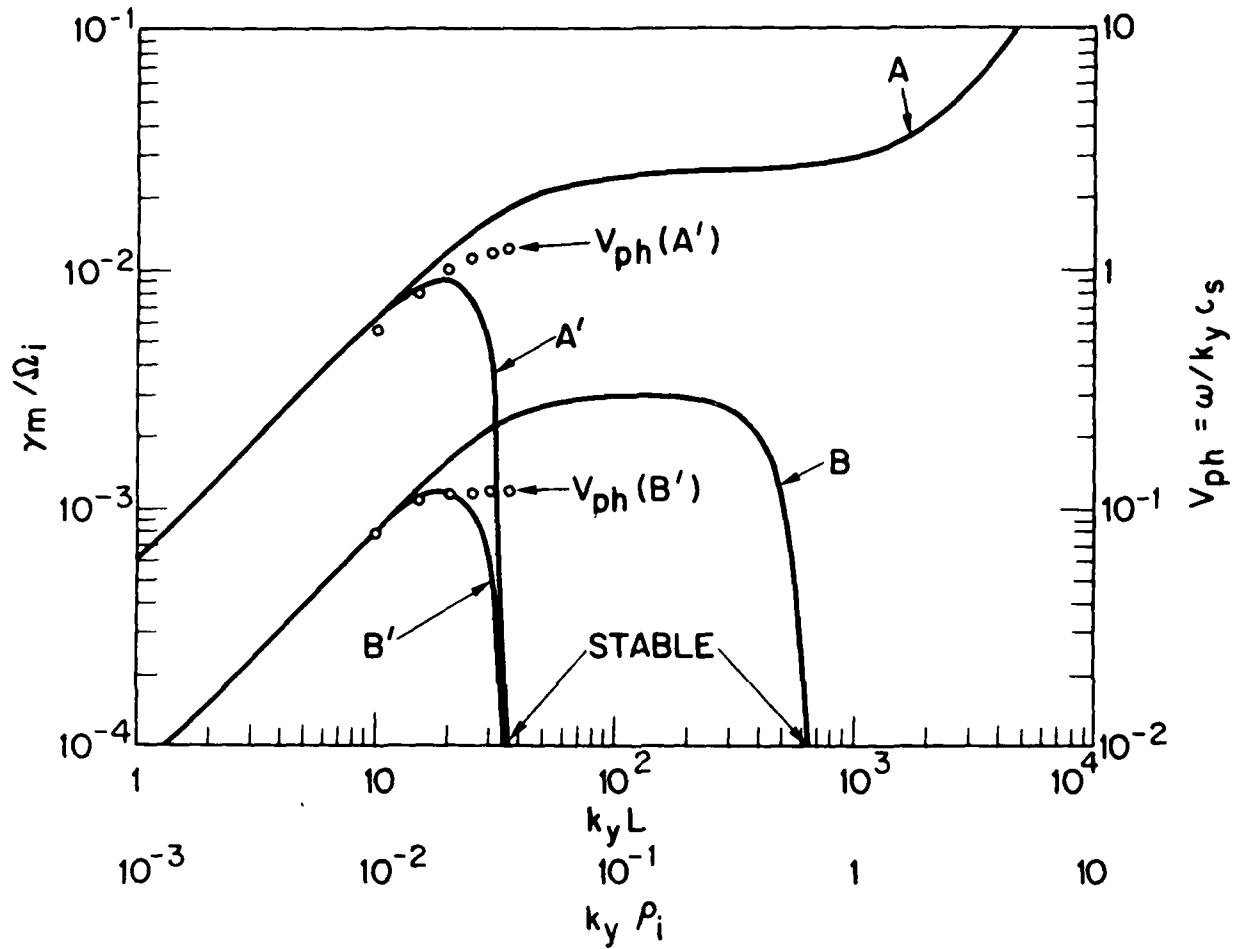


Fig. 2 Plot of the maximum growth rate vs. wavenumber [γ_m / Ω_i vs. $k_y L (k_y \rho_i)$] for the cases of no shear (A and B) and shear (A' and B'); and of the phase velocity vs. wavenumber [$V_{ph} = \omega / k_y c_s$ vs. $k_y L (k_y \rho_i)$], denoted by circles, for the sheared case. See the text for a description of the parameters.

$$\phi(x) = \frac{1}{Q^{1/4}(x)} \exp\left[\pm i \int^x dy Q^{1/2}(y)\right] \text{ for } |x| \rightarrow \infty \quad (21)$$

with the sign of the WKB solution chosen such that a damped solution is obtained in the limit $|x| \rightarrow \infty$.

The density profile used in the analysis is given by

$$n(x) = n_0 \frac{1 + \epsilon \tanh(x/L)}{1 - \epsilon} \quad (22)$$

where $0 < \epsilon < 1$, L characterizes the width of the boundary layer, and $n_0 = n(x = -\infty)$. We assume $\epsilon = 0.8$ for the results presented so that

$$(L_n^{-1})_{\max} = \left(\frac{\partial \ln n}{\partial x}\right)_{\max} = L^{-1} \text{ at } x/L = -0.55 \quad (23)$$

The equilibrium velocity profile chosen to satisfy Eq. (6), i.e.,

$$V_d(x) = V_{d0} n_0 / n(x) \quad (24)$$

where $V_{d0} = V_d(x = -\infty)$.

In Fig. 2 we plot γ_m / Ω_i vs. $k_y L$ (or $k_y \rho_i$) based upon Eqs. (14), (22) and (23) for the following parameters: $V_{d0}/c_s = -7.5$ (curve A') and $V_{d0}/c_s = -1.0$ (curve B'), $\rho_i/L = 10^{-3}$, $v_e/\Omega_e = 10^{-2}$, and $v_i/\Omega_i = 25.0$. Here, γ_m is the growth rate maximized with respect to mode number, which is analogous to k_x . We point out that the mode structure of the unstable waves (i.e., the eigenfunction) is not localized about the position of maximum density gradient ($x = -0.55 L$) but rather at $x = 0$. The local drift velocity at $x = 0$ is approximately $V_{d\ell} \approx V_{d0}/5$ so that $V_{d\ell} \approx -1.5$ (curve A') and $V_{d\ell} \approx -0.2$ (curve B'). Thus, the local plasma conditions

are roughly the same between curves A and A', and between B and B'. The only difference being that the primed curves contain the effects of velocity shear. In the region $k_y L \lesssim 10$ (or $k_y \rho_i \lesssim 10^{-2}$) it is clear that velocity shear has no effect on the gradient drift instability since curves A and A', and B and B' coincide. However, for $k_y L > 10$ (or $k_y \rho_i > 10^{-2}$) the primed and unprimed curves diverge. We find that the gradient drift instability achieves maximum growth at $k_y L \approx 20$ and is stable ($\gamma < 0$) at $k_y L \approx 35$. The value of $k_y L$ for which the mode is "shear stabilized" is more than an order of magnitude smaller than the diffusive cutoff associated with curve B ($k_y L \sim 650$). Thus, velocity shear has the effect of stabilizing the short wavelength gradient drift modes. Note that the position of maximum growth is independent of the drift velocity V_d . Thus, the wavelength of the most unstable modes is dictated by L , not V_d .

We also plot the phase velocity of the waves as a function of k_y in Fig. 2. Specifically, we show $V_{ph} = \omega/k_y c_s$ vs. $k_y L$ for the sheared drift cases (A' and B'); we denote V_{ph} by small circles in Fig. 2. For $k_y L < 20$ we find that $V_{ph} \propto k_y$, while for $k_y L > 20$ that $V_{ph} \sim \text{constant}$. Unlike the position of maximum growth, the phase velocity of the waves is a function of the drift velocity V_d ; the larger the drift velocity, the larger the phase velocity of the waves. For the parameters considered in Fig. 2 we note that $\omega/k_y c_s \approx 0.1$ for $V_{d0}/c_s = -1.0$ (B') and $\omega/k_y c_s \approx 1.0$ for $V_{d0}/c_s = -7.5$ (A').

IV. DISCUSSION

We have presented a nonlocal analysis of the gradient drift instability. The new effect included in this theory is the allowance for an inhomogeneous electric field which produces a sheared drift velocity, i.e., $V_d = V_d(x)$. The major result of this work is that velocity shear can stabilize the short wavelength modes of the instability, and preferentially excite a longer wavelength mode than would be expected from local theory. This result is similar to the short wavelength stabilization of the $\mathbf{E} \times \mathbf{B}$ instability [Perkins and Doles, 1975; Huba et al., 1983]. On the other hand, the long wavelength modes appear to be unaffected by velocity shear. We emphasize that we have not explored the influence of velocity shear on the Farley-Buneman instability; it is possible that γ_m in curve A' of Fig. 2 may become positive for $k_y \rho_i \gtrsim 1$. We also mention that we have only considered linear theory and that a nonlinear cascade of wave energy (Keskinen et al., 1979) could produce wave turbulence in the regime $k_y \rho_i > 20$ even for the velocity sheared cases considered in Fig. 2.

We discuss the possible significance of these results to the equatorial and auroral E regions. Recently, observations of long wavelength irregularities in the equatorial E region have been made (Kudeki et al., 1982; Pfaff et al., 1982). These long wavelength irregularities are characterized by having (1) wavelengths such that $kL \sim 23$ where $L \sim 3-30$ km (i.e., $\lambda \sim$ few kms); (2) wavelengths determined by the gradient scale length, not the drift velocity; (3) phase velocities $V_{ph} \sim 100$ m/sec $< c_s \sim 360$ m/sec; and (4) large amplitudes. Our theory is consistent with the first three observations, but cannot address the fourth since it is only a linear theory. We find that velocity shear stabilizes the short wavelength modes of the gradient drift instability and preferentially

excites a mode such that $k_y L \sim 20$; that the scaling of the dominant wavelength is related to the scale size of the system and not to the drift velocity (i.e., $k_y L \sim 20$ for both curves A' and B'); and that the phase velocity of the dominant wave is such that $V_{ph} \sim \omega/k_y < c_s$ for $V_d < c_s$. These findings suggest that velocity shear may account for the dominant long wavelength irregularities observed in the equatorial electrojet. This conclusion must be regarded as tentative, though, since the density and velocity profiles used in this analysis [Eqs. (22) and (24)], based upon Eq. (6), do not accurately model the equatorial electrojet. However, based upon previous work (Guzdar et al., 1982; Huba et al., 1983) and analysis of Eq. (14) using different plasma profiles, we find that the fundamental conclusion of this paper, viz., velocity shear can stabilize short wavelength modes of the gradient drift instability, is not sensitive to the plasma profiles chosen. Nonetheless, in order to test our hypothesis regarding the equatorial electrojet, an appropriate plasma equilibrium is needed which could perhaps be supplied by experimentalists. Also, the actual parameters used in this analysis (e.g., $V_{d0}/c_s = -7.5$) are not appropriate to electrojet conditions but are chosen to contrast the role of shear to the role of drift velocity.

We suggest that such long wavelength irregularities may exist in the auroral electrojet and could be responsible for the long wavelength structuring of auroral irregularities. Discrete auroras are observed to have two distinct size scales: (i) one is the tens of kilometer width associated with inverted V precipitation, and (ii) the other is the kilometer widths of auroral arc elements which appear to be imbedded in the inverted V precipitation region (e.g., Davis, 1973). The formation of inverted V arcs can be understood in terms of the magnetosphere-ionosphere

coupling (e.g., Kan and Lee, 1980). We suggest here that the gradient drift instability discussed in this paper can lead to the development of the small-scale arc elements within the inverted V precipitation region. During substorms, the ionospheric plasma density in the auroral arc region is greatly enhanced due to the ionization of the atmosphere caused by the precipitation of energetic electrons. In addition, paired electric fields directed toward each other are present in the inverted V region. Therefore, a large plasma density gradient and an electric field gradient may exist in the arc region, which will result in the long wavelength gradient drift instability as obtained here. Take an inverted V region to have a width $L \approx 30$ km. The eigenfunction $\phi(x)$ in (14) for the most unstable mode ($k_y L \approx 20$) is found to have 4-6 peaks. Thus, it is plausible that the large-scale arc will be filamented into 4-6 arc elements as usually observed (Davis, 1978).

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