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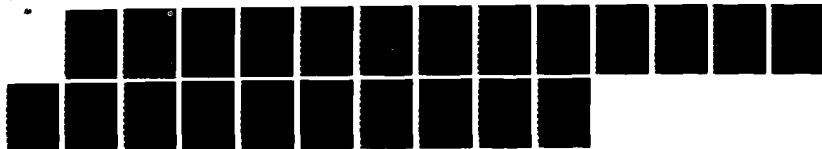
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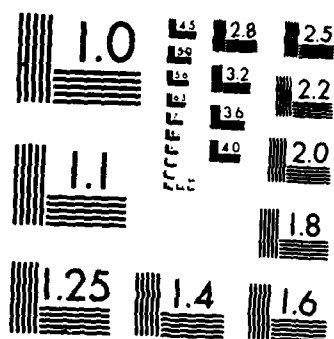
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On the Stability of a Stratified Shear Layer

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ON THE STABILITY OF A STRATIFIED SHEAR LAYER

I. INTRODUCTION

The Kelvin-Helmholtz instability is a widely studied instability driven by velocity shear in neutral fluids as well as in ionized gases¹⁻⁶. In addition to various examples in neutral fluids, the situation of a shear layer under the action of gravity is also encountered in space (e.g., equatorial spread F, the plasmopause) and laboratory (e.g., laser fusion experiments) plasma phenomena. Recent observations show that large velocity shears exist at the equatorward edge of the diffuse auroral boundary⁷. Furthermore, the plasmopause boundary is conjectured to be unstable to a ballooning type shear flow instability⁸. In the case of the equatorial ionosphere, the plasma is confined by a uniform magnetic field and the flow velocity perpendicular to the magnetic field is sheared as a function of the altitude. Gravity is directed opposite to the density gradient so that the system is prone to gravity driven interchange as well as to the shear driven Kelvin-Helmholtz instability. Based on this geometry and physics, linear stability analysis has been performed on a collisional plasma in the equatorial F region of the ionosphere, and a general mode structure equation has been derived^{4,5} which reduces to the general Rayleigh equation dealt with in detail by Drazin².

In the neutral fluid literature (for example, Drazin²) and by Vinas and Madden⁸, the Boussinesq approximation is generally assumed in performing the analysis of the general Rayleigh equation; this assumption

amounts to ignoring all the density gradient terms except the density gradient term that contributes to the buoyancy. In this paper, we relax the Boussinesq approximation and show that the new stability boundary is different than that obtained by Drazin²: the stability boundary is dependent on the density gradient, leading to more stringent restrictions on the Richardson number, and the stability boundary is determined for modes with phase velocity half that of the peak background flow velocity. The following analysis is based on a smooth velocity profile and an exponential density gradient.

II. THEORY

The geometry of the plasma and field configuration used in the analysis is as follows: the magnetic field is uniform and in the z direction ($\underline{B} = B \hat{z}$), the plasma is inhomogeneous along the x direction, $n = n_0(x)$, gravity is acting along the negative x direction $\underline{g} = -g \hat{x}$, and the flow velocity along the y direction is sheared in the x direction, $\underline{V} = -V_0(x) \hat{y}$.

We consider low frequency fluctuations ($\omega \ll \omega_{ce}$ where ω_{ce} is the electron cyclotron frequency) so that electron inertia is ignored. Two dimensional perturbations are considered. The perturbed quantities vary as $\phi = \phi(x) \exp[-i(k_y y - \omega t)]$, where $\omega = \omega_r + i\gamma$. The equation that describes the perturbed electrostatic potential is given as⁴

$$\frac{\partial^2 \phi}{\partial x^2} + p(X) \frac{\partial \phi}{\partial x} + q(X) \phi = 0, \quad (1)$$

where $p(X)$ and $q(X)$ are given by

$$p(X) = \frac{\partial \ln n_0}{\partial x}, \quad (2)$$

$$q(X) = -\hat{k}^2 + \frac{1}{\Omega} \left[\frac{\partial^2 U}{\partial X^2} + \frac{\partial \ln n_0}{\partial X} \frac{\partial U}{\partial X} + \frac{(gL_v/\bar{V}_y^2)}{\Omega} \frac{\partial \ln n_0}{\partial X} \right], \quad (3)$$

where $V_0 = \bar{V}_y U(X)$, $c = \omega/k_y \bar{V}_y$, $\Omega = c - U(x)$, $X = x/L_v$, and $\hat{k} = k_y L_v$. Here L_v and L_n are the gradient scale lengths of the velocity and density, respectively. Equation (1) is a general equation that describes Rayleigh-Taylor and Kelvin-Helmholtz instabilities for arbitrary density profiles and drift velocities. This equation is identical to that obtained for counter-streaming neutral fluids in a gravitational field².

In this paper we consider an exponential profile for the density such that $n_0(x) = n_0 \exp(\beta X)$ where $\beta = L_v/L_n$ and represent the drift velocity by a smooth profile $U(X) = \tanh(X)$. We let β be arbitrary and thus generalize the problem analyzed by Drazin². This is the main objective of this paper. The coefficients $p(X)$ and $q(X)$ can be rewritten as

$$p(X) = \beta \quad (4)$$

$$q(X) = -\hat{k}^2 + \frac{2 \operatorname{sech}^2 X (-2 \tanh X + \beta)}{(c - U)} + \frac{J}{(c - U)^2} \quad (5)$$

where $J = (g/L_n)/(\bar{V}_y^2/L_v^2)$. Using the Boussinesq approximation Drazin² set $p = 0$ and ignored β in the second term in q ; this leads to the condition $c = 0$. We relax these assumptions in the following analysis.

Using the transformation $\phi = \psi \exp[\int dX p(X)/2]$ we write (1) as

$$\frac{\partial^2 \phi}{\partial X^2} - Q(X)\phi = 0 \quad (6)$$

where $Q(X) = q - p'/2 - p^2/4$. Equations (4) and (5) together with (6) yield

$$\frac{\partial^2 \phi}{\partial X^2} + \left[-\hat{k}^2 - \frac{\beta^2}{4} + \left(\frac{\operatorname{sech}^2 X (-2 \tanh X + \beta)}{(c - U)} + \frac{J}{(c - U)^2} \right) \right] \phi = 0. \quad (7)$$

With U as the independent variable, (7) can be written as

$$(1 - U^2) \frac{\partial^2 \phi}{\partial U^2} - 2U \frac{\partial \phi}{\partial U} + \left[\frac{-(\hat{k}^2 + \beta^2/4)}{(1 - U^2)} + 2 + \frac{2c - \beta}{(U - c)} + \frac{J}{(U - c)^2(1 - U^2)} \right] \phi = 0. \quad (8)$$

Equation (8) is a second order differential equation with four singular points at $U = \infty$, c , $+1$, and -1 . The new terms in this equation are $-(\beta^2/4)/(1 - U^2)$ and $(2c - \beta)/(U - c)$; the former, obtained by retaining the first derivative term in (1), introduces a lower cut-off in the wavenumber, and the latter, obtained by retaining the cross term involving the density inhomogeneity and velocity inhomogeneity, allows c to take the value $\beta/2$.

We assume a solution to (8) of the form

$$\phi = (U + 1)^{\mu_+} (U - 1)^{\mu_-} (U - c)^\lambda, \quad (9)$$

where the exponents $\mu_\pm$ and λ are determined by

$$2\mu_\pm = \left[\hat{k}^2 + \beta^2/4 - \frac{J}{(1 \pm c)^2} \right]^{1/2} \quad (10)$$

$$(\lambda^2 - \lambda)(1 - c^2) = - \frac{J}{(1 - c^2)}. \quad (11)$$

Note that for $c = 0$ and $\beta = 0$ the exponents (10) and (11) reduce to those given by Drazin². These exponents have to satisfy the condition

$$(\mu_+ + \mu_- + \lambda)(\mu_+ + \mu_- + \lambda + 1) - 2 = 0, \quad (12)$$

as can be seen by matching the coefficients of the constant terms. From (12) we note that $\mu_+ + \mu_- + \lambda = 1$ or -2 . The latter condition yields a convergent solution with the constraint $2c - \beta = 0$ or $c = \beta/2$.

Using the definitions for μ_+ and λ from (10) and (11) in the condition

$$\mu_+ + \mu_- + \lambda = 1, \quad (13)$$

we obtain

$$J = \hat{k}^2(1 - \beta^2/4 - \hat{k}^2). \quad (14)$$

From (14) the maximum value of J is given as $J_{\max} = (1 - \beta^2/4)^2/4$, which occurs at $\hat{k}_{\max} = (1/\sqrt{2})(1 - \beta^2/4)^{1/2}$. So the new stability condition is $J > (1 - \beta^2/4)^2/4$.

From (14) we see that for $\beta = 1$ the critical Richardson number is $9/64$, which means that for $J > 9/64$ the flow is stable; the cut off wavenumber is $\hat{k} = \sqrt{3}/2$. In comparison, for a Boussinesq fluid the critical J is $1/4$ and the cutoff $\hat{k}$ is 1 . Finally, the stability boundary is not neutral as was the case with a Boussinesq fluid, but the waves have a phase velocity that is half that of the peak background flow velocity, i.e., $c = 1/2$ as opposed to $c = 0$ for a Boussinesq fluid. Finally, for a fluid in a gravity free field ($J = 0$) the Kelvin-Helmholtz unstable domain is given by $0 < \hat{k} < (1 - \beta^2)^{1/2}$. These conclusions can be clearly seen in Fig. 1, where we plot J versus $\hat{k}$. The figure shows the stability boundary for various values of β . It is worth noting that for $\beta > 2$, that is if the density gradient scale length is half the velocity gradient scale length or less, the Kelvin-Helmholtz instability is stable and the system is unstable to the gravity driven interchange, i.e., only when $J < 0$.

We apply the results to the plasmopause boundary based on the observations of velocity shears observed by Kelley⁷ and compare our results with those of Vinas and Madden⁸. The velocity shear V_0/L_v is estimated to be 0.17 Hz; when mapped to the plasmopause region⁷ this yields a local Richardson number of ~ 0.18 [see Ref. 8]. Vinas and Madden⁸ show

that $\beta = 1.5$, which when substituted into (15) yields the critical Richardson number $J < 0.047$ for instability. Thus, our calculations suggest that such strong shears may not drive shear flow instability if steep density gradients exist at the same time.

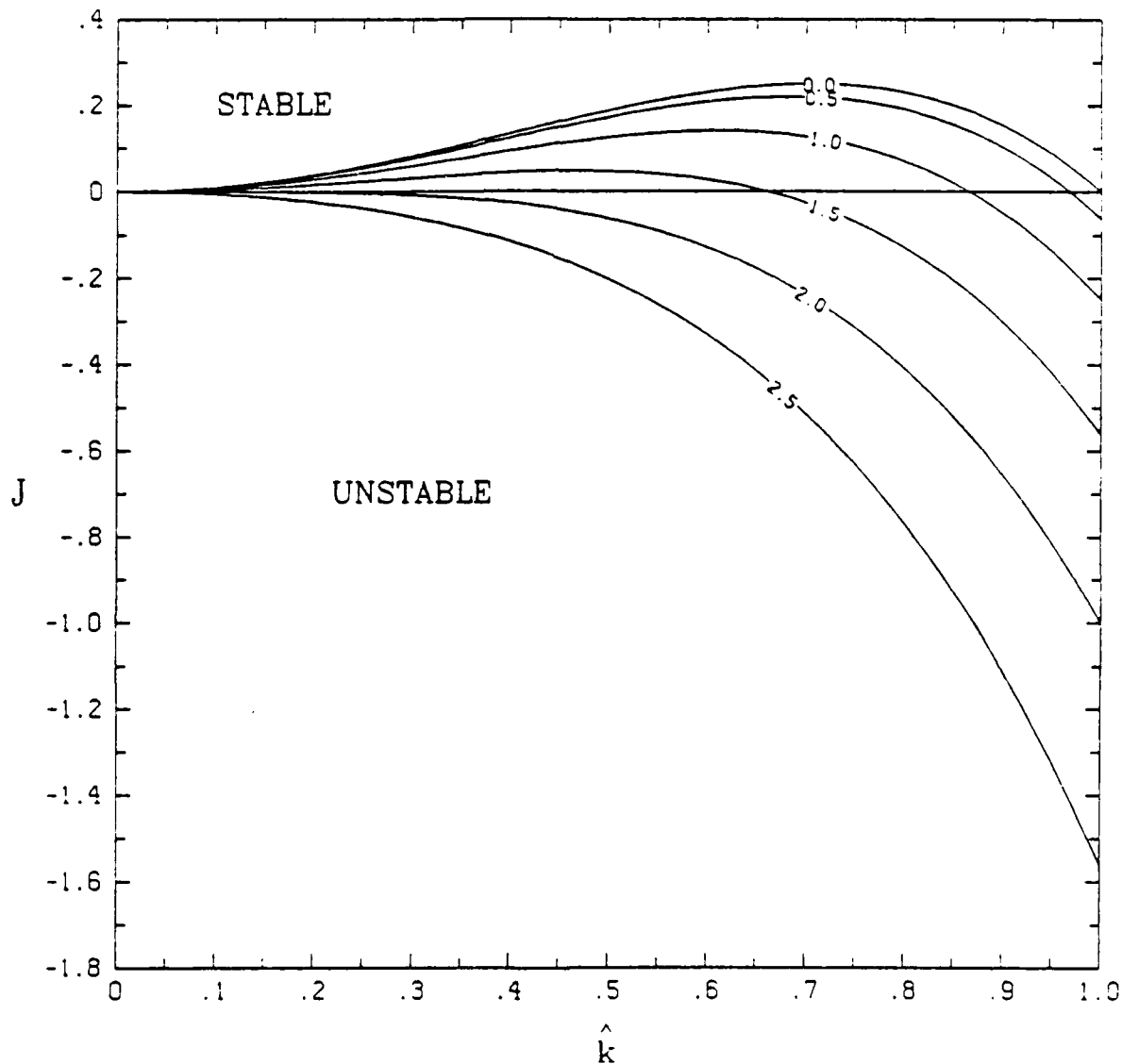


Figure 1. The stability boundary J versus $\hat{k}$ for $\beta = 0.0, 0.5, 1.0, 1.5, 2.0$, and 2.5 . The plasma is unstable (stable) in regions below (above) each curve. The case $\beta = 0$ corresponds to the Boussinesq fluid treated by Drazin. Note that for $\beta \geq 2$, $J < 0$ is required for instability.

III. SUMMARY

In summary, we have shown that the mode equation for a stratified shear layer under the action of gravity can be solved for stability boundary without the need to impose the Boussinesq approximation. For $L_n = L_v$, $\beta = 1$, the critical Richardson number is $9/64$ as opposed to $1/4$ and the unstable wavenumber domain is smaller with a cut off for $J = 0$ at $\hat{k} = \sqrt{3}/2$ rather than at $\hat{k} = 1$. In addition, the waves do not have a zero phase velocity but a phase velocity that is half the peak background flow velocity. These changes are due to the density gradient terms that were ignored using the Boussinesq approximation. In addition, we have shown that plasmopause boundary may not be unstable to shear flow ballooning instability if strong density gradients exist in conjunction with strong shears at the diffuse auroral boundary at ionospheric heights.

ACKNOWLEDGMENTS

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