AD-A151 414 EFFECTS OF FINITE CURRENT CHANNEL WIDTH ON THE CURRENT CONVECTIVE INSTABILITY(U) NAVAL RESEARCH LAB WASHINGTON DC P SATYANARAYANA ET AL. 05 MAR 85 NRL-MR-5494 1/1

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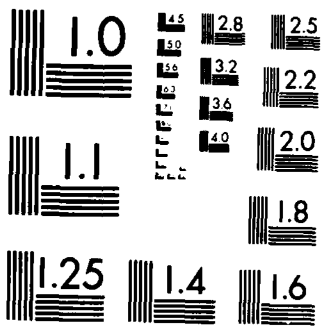

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NRL Memorandum Report 5494

AD-A151 414

Effects of Finite Current Channel Width on the Current Convective Instability

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March 5, 1985

This research was partially sponsored by the Office of Naval Research and the Defense Nuclear Agency under Subtask S99QMXBC, work unit 00102 and work unit title "Plasma Structure Evolution."

NAVAL RESEARCH LABORATORY
Washington, D.C.

REPORT DOCUMENTATION PAGE

1a. REPORT SECURITY CLASSIFICATION UNCLASSIFIED	1b. RESTRICTIVE MARKINGS
2a. SECURITY CLASSIFICATION AUTHORITY	3. DISTRIBUTION / AVAILABILITY OF REPORT Approved for public release; distribution unlimited.
2b. DECLASSIFICATION / DOWNGRADING SCHEDULE	
4. PERFORMING ORGANIZATION REPORT NUMBER(S) NRL Memorandum Report 5494	5. MONITORING ORGANIZATION REPORT NUMBER(S)
6a. NAME OF PERFORMING ORGANIZATION Naval Research Laboratory 6b. OFFICE SYMBOL (If applicable) Code 4780	7a. NAME OF MONITORING ORGANIZATION
6c. ADDRESS (City, State, and ZIP Code) Washington, DC 20375-5000	7b. ADDRESS (City, State, and ZIP Code)
8a. NAME OF FUNDING / SPONSORING ORGANIZATION ONR and DNA 8b. OFFICE SYMBOL (If applicable)	9. PROCUREMENT INSTRUMENT IDENTIFICATION NUMBER
8c. ADDRESS (City, State, and ZIP Code) Arlington, VA 22217 Washington, DC 20305	10. SOURCE OF FUNDING NUMBERS PROGRAM ELEMENT NO (See page ii) / PROJECT NO / TASK NO / WORK UNIT ACCESSION NO

11. TITLE (Include Security Classification)

Effects of Finite Current Channel Width on the Current Convective Instability

12. PERSONAL AUTHOR(S)

Satyanarayana, P.,* Chen, J.* and Keskinen, M.J.

13a. TYPE OF REPORT	13b. TIME COVERED	14. DATE OF REPORT (Year, Month, Day)	15. PAGE COUNT
Interim	FROM 10/83 TO 10/84	1985 March 5	41

16. SUPPLEMENTARY NOTATION

*Science Applications, Inc., McLean, VA 22102 (Continues)

17. COSATI CODES			18. SUBJECT TERMS (Continue on reverse if necessary and identify by block number)
FIELD	GROUP	SUB-GROUP	Plasma instability; Fluid instability; High latitude ionosphere; Finite current channel width

19. ABSTRACT (Continue on reverse if necessary and identify by block number)

The effects of the finite current channel width on the current convective instability are studied both analytically and numerically. First using a sharp-boundary field-aligned current distribution which has a finite width along the plasma density gradient, the dispersion relation is obtained analytically. It is found that, for the long wavelength modes ($k_y d \ll 1$) where the nonlocal effects are most prominent, the growth rate γ is proportional to $|\nabla_d d/L^2|$ in the collisional limit $(\nu_{in} \gg |\omega|)$, (Continues)

20. DISTRIBUTION / AVAILABILITY OF ABSTRACT ☒ UNCLASSIFIED/UNLIMITED ☐ SAME AS RPT ☐ DTIC USERS	21. ABSTRACT SECURITY CLASSIFICATION UNCLASSIFIED	
22a. NAME OF RESPONSIBLE INDIVIDUAL J. D. Huba	22b. TELEPHONE (Include Area Code) (202) 767-3630	22c. OFFICE SYMBOL Code 4780

DD FORM 1473, 84 MAR 83 APR edition may be used until exhausted. All other editions are obsolete

SECURITY CLASSIFICATION OF THIS PAGE

10. SOURCE OF FUNDING NUMBERS

PROGRAM ELEMENT NO.	PROJECT NO.	TASK NO.	WORK UNIT ACCESSION NO.
62715H		RR033-02-44	DN880-024
61153N			DN230-291

16. SUPPLEMENTARY NOTATION (Continued)

This research was partially sponsored by the Office of Naval Research and the Defense Nuclear Agency under Subtask S99QMXBC, work unit 00102 and work unit title "Plasma Structure Evolution."

19. ABSTRACT (Continued)

where d is the half-width of the current channel, L is the plasma density gradient scale length, $\overline{V}_d$ is the field-aligned current velocity, ν_{in} is the ion-neutral collision frequency and ω is the perturbation frequency. For the long wavelength modes in the inertial limit ($\nu_{in} << |\omega|$), the growth rate γ scales as $|\overline{V}_d d/L^2|^\alpha$, where α = 1/2 (2/3) for $k_z^2/k_y^2\, \Omega_e/\nu_{ei}$ much less than (greater than) $|\omega|/\Omega_i$ and $k_z(k_y)$ is the wavenumber parallel (perpendicular) to the magnetic field, $\Omega_e(\Omega_i)$ is the electron (ion) gyrofrequency, and ν_{ei} is the electron-ion collision frequency. Numerical results are also presented for a diffuse-boundary current velocity distribution. Applications to the high latitude ionosphere are discussed.

CONTENTS

EFFECTS OF FINITE CURRENT CHANNEL WIDTH ON THE CURRENT CONVECTIVE INSTABILITY

I. INTRODUCTION

From a variety of experimental observations [see recent reviews by Fejer and Kelley, 1980, Vickrey and Kelley, 1983, and Hanuise et al., 1981], it is now known that the high latitude ionosphere, from the auroral zone into the polar cap, is a highly structured and nonequilibrium medium containing plasma fluctuations and structures with scale sizes ranging from hundreds of kilometers to meters. Several processes, e.g., particle precipitation, plasma instabilities, and neutral fluid dynamics, have been proposed to account for high latitude ionospheric plasma structure [Keskinen and Ossakow, 1983a]. The plasma instabilities that have received the most attention thus far are the ExB instability [Simon, 1963; Linson and Workman, 1970; Keskinen and Ossakow, 1982, 1983b; Huba et al; 1983, Keskinen, 1984] and the current convective instability [Lehnert, 1958; Kadomtsev and Nedospasov, 1960; Ossakow and Chaturvedi, 1979; Huba and Ossakow, 1980; Keskinen et al., 1980; Chaturvedi and Ossakow, 1981, 1983; Satyanarayana and Ossakow, 1984; Gary, 1984; Huba, 1984]. The current convective instability results from the coupling of a magnetic field-aligned current and a plasma density gradient transverse to the magnetic field in the presence of electron-neutral collisions — a configuration likely to occur in the auroral zone ionosphere [Vickrey et al., 1980]. All of the previous studies of the current convective instability have assumed that the width of the current transverse to the magnetic field is infinite. However, recent observational results [see, for example, Bythrow et al., (1984)] indicate that the transverse dimension of the field-aligned current distributions may be comparable to or less than the density gradient scale length. Since the instability requires both a density gradient and a field-aligned current, the properties of the instability may be modified if one uses a more realistic geometry with a current channel of finite width. Indeed, physical intuition would lead one to expect reduction in the growth rate and modification in the stability boundary for narrow current channels. In this analysis we remove the infinite-width current simplification and study the effects of finite current channel width on the current convective instability. The analytical and numerical

Manuscript approved October 4, 1984.

results will show that, in the long wavelength limit, $k_y d \ll 1$, where the nonlocal results are most prominent, the growth rate depends on the global quantities such as $(\overline{V}_d d/L^2)$, where $\overline{V}_d$ is the current velocity, d is the half-width of the current channel and L is the density gradient scale length. These intrinsically nonlocal effects may be of interest for interpretation of the observational data. We discuss the nonlocal results and apply them to the high latitude ionosphere.

The organization of the paper is as follows. In Section II we derive the general dispersion relation describing the mode structure of the current convective instability with a finite current channel width. In Section III we present a theoretical analysis of the mode structure equation while in Section IV we give numerical results. Finally, in Section V we summarize our findings and discuss applications to the high latitude ionosphere.

II. MODE STRUCTURE EQUATION

In this paper we use a slab model shown in Fig. (1). The plasma is assumed inhomogeneous along the north-south direction $(\hat{\underline{x}})$ and is assumed to support equilibrium currents along the magnetic field. The magnetic field is uniform and vertical $(\hat{\underline{z}})$ with the field-aligned currents taken to be of the form $\underline{J} = J_o(x)\ \hat{\underline{z}}$. The primary objective of this paper is to confine the field-aligned currents to a finite region along the direction of the density gradient $(\hat{x})$ and to study the effects of finite current width on the current convective instability. We neglect weak altitude dependent density gradients as well as east-west $(\hat{\underline{y}})$ gradients and transverse equilibrium electric fields. The temperature effects responsible for diffusive damping are ignored as is magnetic field shear. Electron inertia is neglected since current convective instability is a low frequency instability $(\omega \ll \Omega_e)$. Furthermore, we confine ourselves to the F-region of the ionosphere and ignore the electron-neutral collision frequency (ν_{en}) compared to the electron-ion collision frequency (ν_{ei}). Due to the variation of the field-aligned currents and the background density along x, we perform a nonlocal analysis and derive a mode structure equation for the perturbed potential. This procedure allows us to study short wavelength and long wavelength modes. The basic two-fluid equations describing the electron and ion dynamics in the rest frame of the neutrals are

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot (n_\alpha \underline{V}_\alpha) = 0 \ , \qquad (1)$$

$$m_\alpha n_\alpha \left(\frac{\partial}{\partial t} + \underline{V}_\alpha \cdot \nabla\right) \underline{V}_\alpha = q_\alpha n_\alpha \left(\underline{E} + \frac{\underline{V}_\alpha \times \underline{B}}{c}\right) - m_\alpha n_\alpha \nu_{\alpha n} \underline{V}_\alpha - \underline{R}_\alpha \ , \qquad (2)$$

where α denotes species (e for electrons; i for ions), n_α is the plasma density, $\underline{V}$ is the fluid velocity, m_α is the mass of the species, q_α is the charge of the species, c is the speed of light, $\underline{E}$ is the electric field, $\nu_{\alpha n}$ is the plasma particle collision frequency with neutrals, and $\underline{R}_\alpha$ is the friction force given by

$$R_i = m_i n_i \nu_{ie} (\underline{V}_i - \underline{V}_e) \ , \qquad (3)$$

$$R_e = m_e n_e \nu_{ei} (\underline{V}_e - V_i). \qquad (4)$$

The equilibrium velocities are given by

$$V^0_{ez} = - \frac{e \ E_{0z}}{m_e \nu_{ei}} \qquad (5)$$

$$V^0_{iz} = 0 \qquad (6)$$

In obtaining these equations, terms proportional to $\nu_{en}/\Omega_e << 1$ have been neglected. By setting $\underline{E}_{0\perp} = 0$, it can be easily shown that $\underline{V}_{e\perp} = \underline{V}_{i\perp} = 0$. We choose the friction force to be $R_e = - R_i$ so that momentum conservation demands $\nu_{ie} m_i = \nu_{ei} m_e$ and thus $V^0_{iz} = 0$. The net drift velocity $V_d \equiv V^0_{ez} - V^0_{iz}$ then becomes $V_d = V^0_{ez}$. More generally, however, neutral particles also participate in momentum conservation, and V^0_{iz} need not be zero. The above simplification implies that the electrons are the primary current carriers.

In order to derive the nonlocal mode structure equation, we use a perturbation of the form $\hat{f} \sim f(x) \exp[-i(\omega t - k_y y - k_z z)]$, where $\omega = \omega_r + i\gamma$, $\gamma > 0$ for growth and k_y and k_z are the mode numbers in y and z directions. We then linearize Eqs. (1) and (2), and finally we subtract the electron continuity equation from the ion continuity equation [Eq. (1)] and use quasi-neutrality condition to derive the mode structure equation.

We consider only electrostatic perturbations and do not perturb the magnetic field [Ossakow and Chaturvedi, 1979]. By perturbing Eq. (2) for both electrons and ions, we obtain the perturbed velocities in the z direction

$$\hat{V}_{ez} = -\frac{e\,\hat{E}_z}{m_e \nu_{ei}} \tag{7}$$

$$\hat{V}_{iz} = 0 \tag{8}$$

The perpendicular perturbed velocities are given as

$$\hat{V}_{e\perp} = \frac{c}{B}\left[\hat{\underline{E}}_\perp \times \hat{\underline{z}} - (\nu_{ei}/\Omega_e)\hat{E}_\perp\right] \tag{9}$$

$$\hat{\underline{V}}_{i\perp} = \frac{c}{B}\left[\hat{\underline{E}}_\perp \times \hat{\underline{z}} - i\,\frac{(\omega + i\nu_{in})}{\Omega_i}\,\hat{\underline{E}}_\perp\right] \tag{10}$$

We substitute Eqs. (7)-(10) in Eq. (1) and use $\underline{E} = -\nabla\phi$ in the electrostatic limit. Then we subtract the electron continuity equation from the ion continuity equation and impose quasineutrality to obtain

$$\frac{\partial^2\hat{\phi}}{\partial x^2} + p(x)\,\frac{\partial\hat{\phi}}{\partial x} + q(x)\hat{\phi} = 0, \tag{11}$$

where

$$p(x) = n_o'/n_o \tag{12}$$

$$q(x) = -k_y^2 - \frac{i\,k_z^2\,(\Omega_e/\nu_{ei})\,\Omega_i}{(\omega + i\nu_{in})}$$

$$+ \frac{k_z\,V_d(x)}{(\omega - k_z V_d(x))}\,\frac{\Omega_i}{(\omega + i\nu_{in})}\left[k_y(n_o'/n_o) - i(\Omega_e/\nu_{ei})\,k_z^2\right] \tag{13}$$

This equation agrees with the nonlocal mode structure equation derived earlier by Satyanarayana and Ossakow (1984) and Huba (1984).

The effects of finite current width are contained in the third term in Eq. (13). We solve Eq. (11) exactly for two specific profiles of the current velocity $V_d(x)$. In the next section, we obtain the analytic dispersion relation for a "waterbag" distribution.

$$V_d(x) = \begin{cases} \bar{V}_d, & |x| \leq d \\ 0\ , & |x| > d \end{cases} \tag{14}$$

where 2d is the current channel width. We choose an exponential plasma density profile $n_o(x) = \bar{n}_o \exp(x/L)$ so that $n_o'/n_o = 1/L$. In section IV we present the numerical results for both the waterbag current model and the diffuse Gaussian current model.

III. DISPERSION RELATION

The dispersion relation is obtained by matching the logarithmic derivatives of the wavefunction inside and outside the waterbag at the boundaries, $x = \pm d$, and requiring that the solution be evanescent for $x \to \pm \infty$. First, we will cast Eq. (11) into a Schrödinger type equation by using the transformation

$$\hat{\phi} = \psi(x) \exp\left[-\frac{1}{2} \int p(x)dx\right]. \tag{15}$$

Then, we have

$$\frac{\partial^2 \psi}{\partial x^2} + Q^2(x)\ \psi = 0 \tag{16}$$

where

$$Q^2(x) = q(x) - p^2/4 - p'/2 \tag{17}$$

and $p' = \partial p/\partial x$. For an exponential density profile, $n_o = \bar{n}_o \exp(x/L)$, $p = 1/L$ and $p' = 0$. Thus, for the waterbag profile $V_d(x)$ given by Eq. (14) and for the exponential density profile, Eq. (16) can be written as

$$\frac{\partial^2 \psi_{out}}{\partial x^2} - Q_o^2\ \psi_{out} = 0 \qquad |x| > d \tag{18}$$

$$\frac{\partial^2 \psi_{in}}{\partial x^2} + Q_i^2\ \psi_{in} = 0 \qquad |x| \leq d \tag{19}$$

with

$$Q_i^2 = -[k_y^2 + 1/(4L^2)] - \frac{i\,k_z^2(\Omega_e/\nu_{ei})\Omega_i}{(\omega + i\nu_{in})}$$

$$+ \frac{k_z \bar{V}_d}{(\omega - k_z\bar{V}_d)} \frac{\Omega_i}{(\omega + i\nu_{in})} [k_y/L - i\,k_z^2 (\Omega_e/\nu_{ei})] \quad (20)$$

and

$$Q_o^2 = [k_y^2 + 1/(4L^2)] + \frac{i\,k_z^2 (\Omega_e/\nu_{ei})\Omega_i}{(\omega + i\nu_{in})} \quad (21)$$

where primes indicate derivatives with respect to x. The current convective instability is a nearly purely growing mode with $\omega = \omega_r + i\gamma \sim i\gamma$. This will be shown to be true a posteriori in the next section. For $\omega_r < \gamma$, the sign convention applicable in Eqs. (18) and (19) is such that $Q_o^2 > 0$ and $Q_i^2 > 0$. In Eq. (20) the sign of $\bar{V}_d/L$ (< 0) is chosen so that $Q_i^2 > 0$ and the growth rate is positive for positive k_y. Furthermore, in the following we assume that $|k_z\bar{V}_d| \ll |\omega|$ in the denominator and $k_z^2\ (\Omega_e/\nu_{ei}) < k_y/L$. First, we rewrite Q_i^2 and Q_o^2 as follows:

$$\hat{Q}_o^2 = \frac{\hat{k}^2(\omega + i\bar{\nu})}{(\omega + i\nu_{in})} \quad (22)$$

$$\hat{Q}_i^2 = -\frac{\hat{k}^2(\omega - \omega_+)(\omega - \omega_-)}{\omega(\omega + i\nu_{in})} \quad (23)$$

where

$$\bar{\nu} = \nu_{in} + (\hat{k}_z^2/\hat{k}^2)(\Omega_e/\nu_{ei})\Omega_i \quad (24)$$

$$\omega_o^2 = \omega_+\omega_- = (\hat{k}_z\hat{k}_y/\hat{k}^2)|\bar{V}_d/L|\Omega_i \quad (25)$$

$$\omega_\pm = -(i\bar{\nu}/2)[1 \pm [1 + 4(\omega_o^2/\bar{\nu}^2)]^{1/2}] \quad (26)$$

$$\hat{k}^2 = \hat{k}_y^2 + 1/4 \quad (27)$$

and $\hat{k}_y$, $\hat{k}_z$, and $\hat{d}$ are the normalized quantities k_yL, k_zL, and d/L respectively. We note that $\omega_{\pm}$ are the two roots of the local dispersion relation $(Q_i = 0)$. Since the potential, $\hat{Q}_i^2(x)$, in Eq. (16) is symmetric about $x = 0$, we can restrict our analysis to eigenfunctions of definite parity.

In the region $|x| < d$ we have

$$\psi_{in} = A \cos (\hat{Q}_i x/L) \qquad \text{for even parity} \tag{28}$$

$$\psi_{in} = B \sin (\hat{Q}_i x/L) \qquad \text{for odd parity} \tag{29}$$

where A and B are constants. In the region $|x| > d$, we have an exponentially decreasing solution for either case

$$\psi_{out} = C\, e^{-\hat{Q}_o |\frac{x}{L}|} \tag{30}$$

where C is a constant. The matching condition in the case where $V_{iz}^0 = 0$ is that the logarithmic derivative be continuous at $x = \pm d$

$$(\psi_{in}'/\psi_{in})_{d_-} = (\psi_{out}'/\psi_{out})_{d_+}$$

where $d_{\pm} = \lim_{\varepsilon \to o} (d \pm \varepsilon)$. This condition yields the dispersion relation

$$\hat{Q}_i \tan (\hat{Q}_i \hat{d}) = \hat{Q}_o \qquad \text{even modes} \tag{31a}$$

$$\hat{Q}_i \operatorname{cotan} (\hat{Q}_i \hat{d}) = -\hat{Q}_o \qquad \text{odd modes} \tag{31b}$$

The effect of finite current channel width are contained in the argument of the tangent (cotangent) function for the even modes (odd modes). We analyze the dispersion relation for the even modes by first inverting Eq. (31a) to obtain

$$\hat{Q}_i = \frac{1}{\hat{d}} [\operatorname{Arc} \tan (\hat{Q}_o/\hat{Q}_i)] \tag{32}$$

We can rewrite Eq. (32) as

$$\hat{Q}_i = \frac{1}{\hat{d}} \{ \tan^{-1}(\hat{Q}_o/\hat{Q}_i) + m\pi \} \tag{33}$$

where $\tan^{-1}(\hat{Q}_o/\hat{Q}_i)$ is the principal value of Arc tan $(\hat{Q}_o/\hat{Q}_i)$ and m is an integer. Equation (33) is the fundamental nonlocal dispersion relation. We note that when d = 0, the dispersion relation $Q_o = 0$ from Eq. (31a) does not give instability since the currents driving the instability vanish.

We can compare Eq. (33) and Eq. (19) to understand the local and nonlocal limits. We recall that the local limit means that the current channel width is infinite so that Eq. (19) describes the mode structure throughout the space. In the local theory, Eq. (19) is also Fourier analyzed with respect to x to obtain

$$Q_i^2 = k_x^2 \tag{34}$$

and if we assume $k_y^2 \gg k_x^2$ we regain the earlier results of Ossakow and Chaturvedi (1979), $Q_i \sim 0$, corresponding to $\omega \sim \omega_+$. Comparing Eqs. (33) and (34) we see that in the nonlocal analysis, the right hand side of Eq. (33) provides the inverse scale length, L_x^{-1} . The role played by the current channel, then, is to form a wavepacket of width of the order of L_x localized within the current channel.

In the limit $|\hat{Q}_o/\hat{Q}_i| \gg 1$, $\tan^{-1}(Q_o/Q_i)$ approaches $\pi/2$. This happens in the local limit when $\omega \simeq \omega_\pm$ as can be seen by dividing Eq. (22) by Eq. (23) to obtain

$$\frac{\hat{Q}_o^2}{\hat{Q}_i^2} = - \frac{\omega(\omega + i\bar{\nu})}{(\omega - \omega_+)(\omega - \omega_-)}$$

As we approach the local limit we have from Eq. (33)

$$\hat{Q}_i = (m + 1/2)\pi/\hat{d} \tag{35}$$

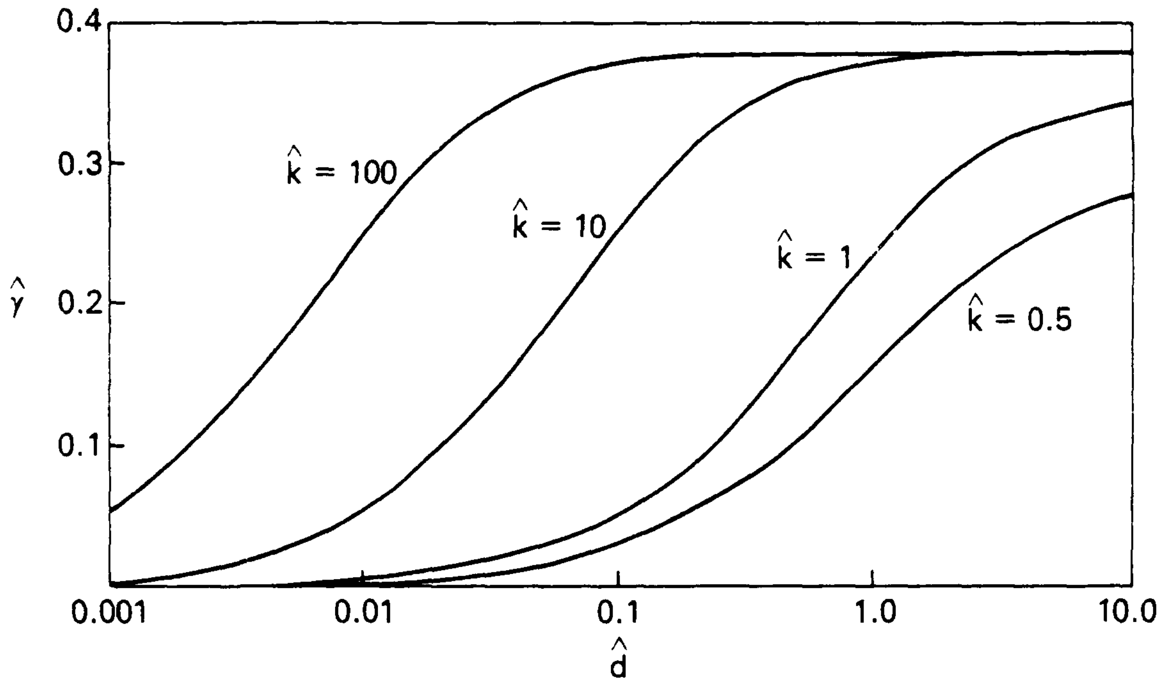

Figure 4

Plot of growth rate $\hat{\gamma}$ versus wavenumber $\hat{k}$ in the collisional limit for diffuse boundary case. The profiles used for the current and the density are given in Eqs. (53) and (54), respectively. The parameters used are the same as in Fig. 3.

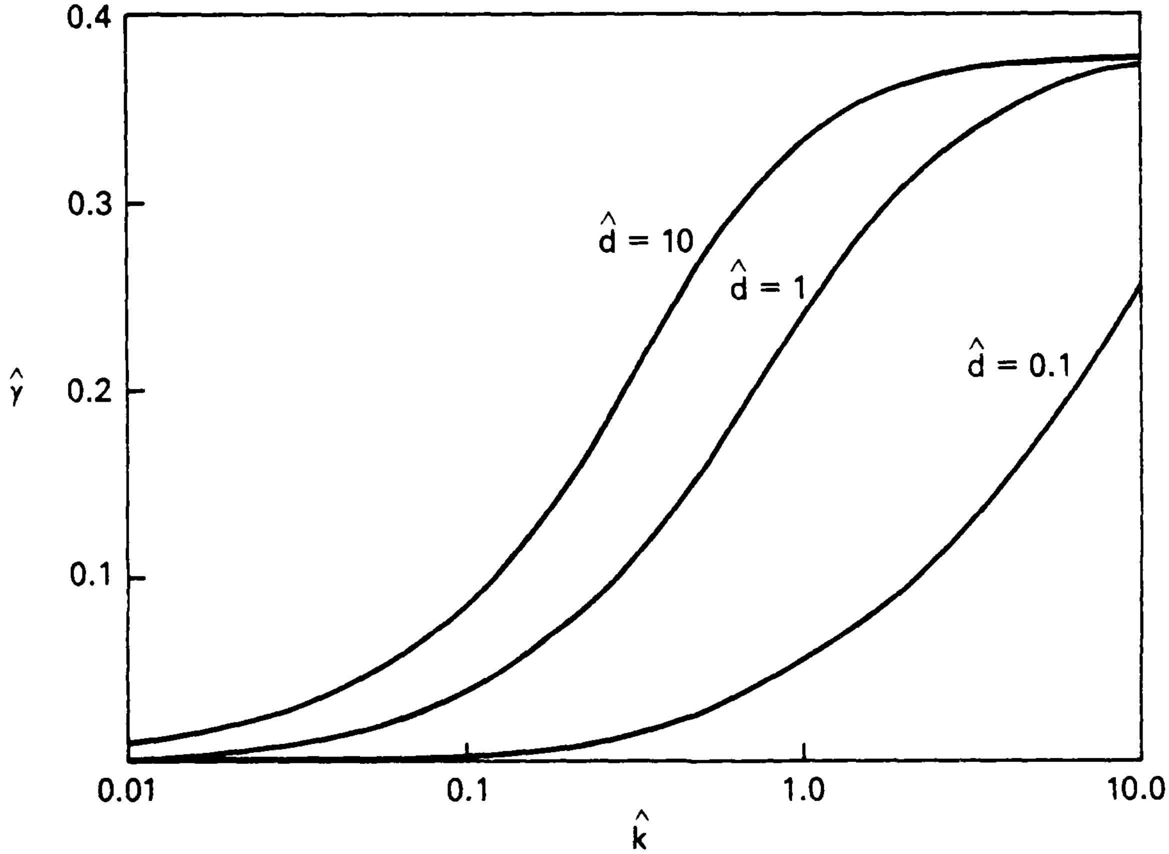

Figure 3

Plot of growth rate $\hat{\gamma}$ versus $\hat{k}$ for $\hat{d}$ = 0.1, 1.0, and 10.0 for sharp boundary case. The parameters chosen are $\nu_{in}/|\bar{V}_d/L| = 1$, $\Omega_i/|\bar{V}_d/L| = 10^4$, and $\nu_{ei}/|\bar{V}_d/L| = 10^{-4}$.

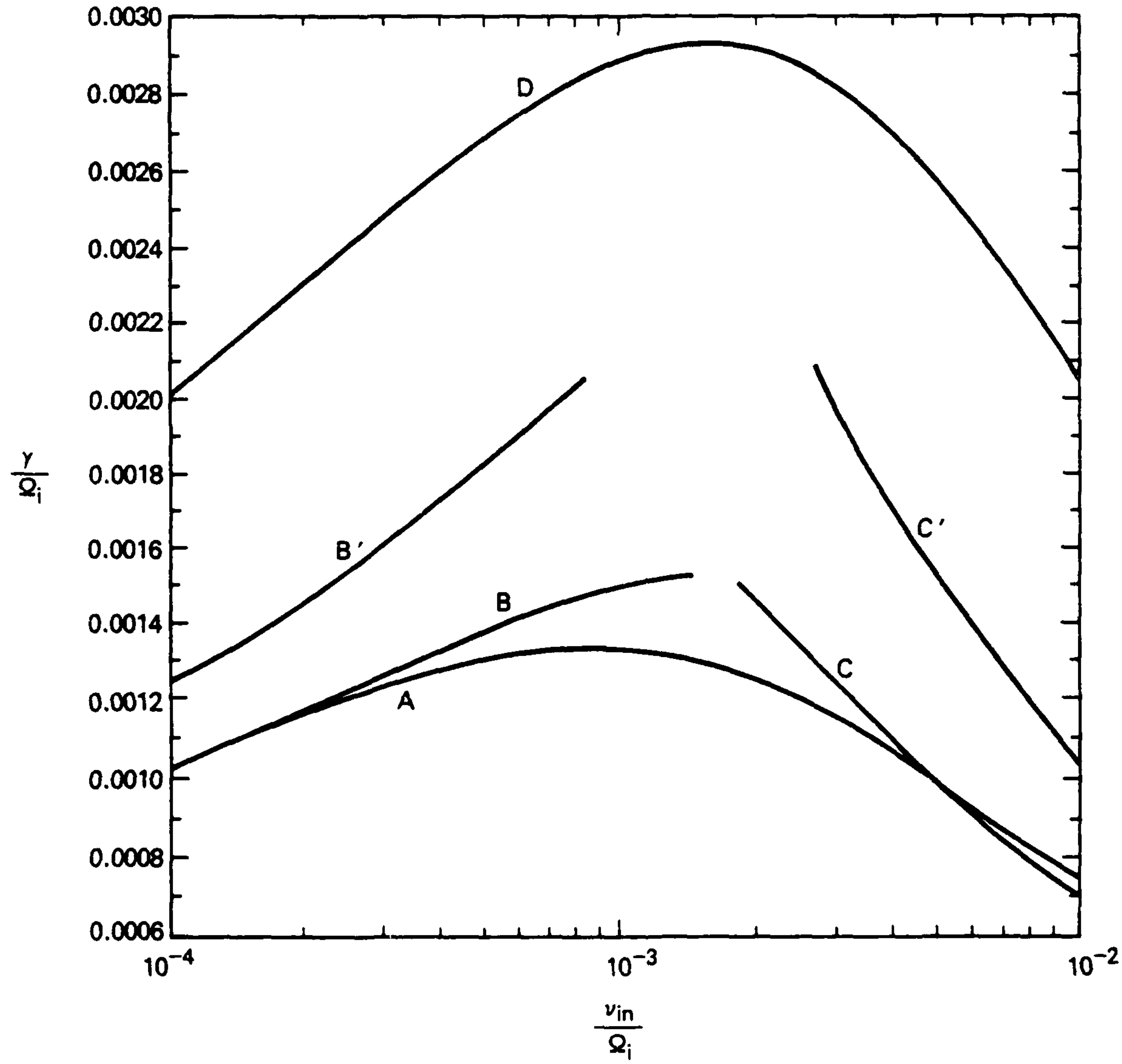

Figure 2

Plot of growth rate (γ/Ω_i) versus the ion-neutral collision frequency (ν_{in}/Ω_i) for $k_y d = 0.3$, $\hat{V}_d/L = 5.0$, and $\nu_{ei}/\Omega_e = 10^{-4}$. Curve A shows the exact numerical solutions of Eq. (33). Curves B and C represent Eqs. (46) and (42) which are solutions in the inertial and collisional limits for long wavelength modes. Curves B′ and C′ represent the simplified analytical expressions, Eqs. (52) and (43), in the same limits. Curve D is the local solution, ω_-, given by Eq. (38).

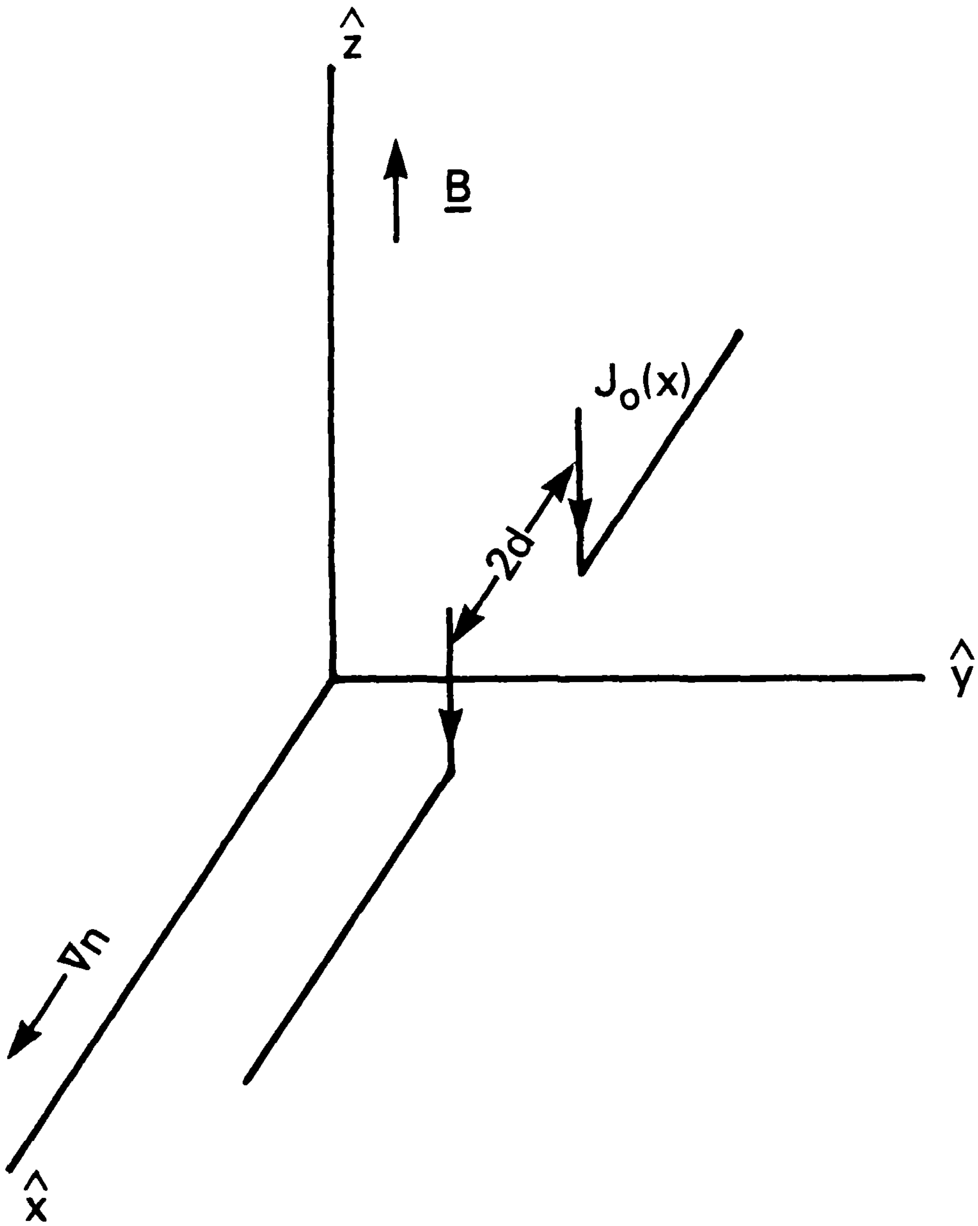

Figure 1

Schematic of the geometry used displaying the finite current channel of width 2d along $\hat{x}$.

ACKNOWLEDGMENTS

We would like to thank J. Huba for the critical reading and suggestions for improvement of the manuscript and G. Ganguli, Y.C. Lee, P. Chaturvedi, and S.L. Ossakow for constructive discussions. This work has been supported by the Office of Naval Research and the Defense Nuclear Agency.

that both the long and short wavelength modes driven by the current convective instability should have longer growth times than previously estimated. Other data sets have also shown very narrow high latitude F-region field-aligned currents. From the results of this paper, we expect the nonlocal nature to have an important impact.

Recently Bythrow et al. (1984) using HILAT satellite data have measured very large downward currents (~ 100 $\mu A/m^2$), presumably carried by cold upward drifting electrons, in the high latitude topside F-region ionosphere. The dimensions of these currents perpendicular to the magnetic field is of order the transverse scale size (~ 3 km) of precipitation induced density structures. Turbulent horizontal ion drifts were also measured in conjunction with the large currents and transverse density gradients. Their data shows that $d \sim 1.9$ km, $L = 6$ km, and $\overline{V}_d = 30$ km/sec. In the case of ionospheric F region, $\nu_{ei}/\Omega_e \sim 10^{-4}$ and $\nu_n/\Omega_i \sim 10^{-3}$. For these parameters Eq. (33) yields the growth rates as $\gamma \sim 0.0028\ \Omega_i$, $0.0026\Omega_i$, and $0.00146_i\Omega$, for $\lambda = 1$, 2, and 10 km, respectively. Whereas, for $d/L \gg 1$, the local growth rate is $\sim 0.0031\ \Omega_i$. For narrower current channel and large density gradients or longer wavelength modes the growth rate could be substantially lower.

Finally, several aspects of the theory of the effects of finite current channel width on the current convective instability have not been discussed here. For very large currents, the self-consistent role of magnetic shear must be included [Huba and Ossakow, 1980]. An approximate value of the scale length for magnetic shear is $L_s \approx cB/4\pi J$ with c, the speed of light, B the magnetic field, and J the current density. For $J \approx 10^{-6} A/m^2$ and $B \sim 0.5$ G we find $L_s \sim 10^4$ km. For $d \ll L_s$ where d is the current channel width we do not expect magnetic shear effects to play an important role. In addition, we have not considered self-consistent velocity shear and its effects on the current convective instability with a finite current channel width, the finite parallel extent of the auroral ionosphere, electromagnetic effects, or nonlinear aspects. We are presently investigating several of these aforementioned effects.

total integrated currents in order to predict the behaviour of the plasma at any point. Second, because there is no plane wave in the x direction, an observer passing through the current channel in the x direction does not see the usual Doppler shift $k_x V_o$ that would be predicted by a local theory, where V_o is the velocity of the observer. This point has also been noted by Ganguli et al. (1984) in connection with the current driven ion cyclotron instability. The two perpendicular directions, x and y, are no longer equivalent.

The wavefunction as given by Eqs. (21) and (30) shows that it is localized around the point $x = 0$ with a width $L_x \approx 1/Q_i$ given by Eq. (33). The fastest growing mode ($m = 0$) has a single peak around $x = 0$. Note that the width of the wave packet is directly proportional to d, which implies that in the presence of wider channels the wavefunction is spread out, whereas in the presence of narrow channels the wavefunction is sharply peaked.

An important result is that in the long wavelength limit, the growth rate scales as some power of the product of the magnitude of the drift velocity and the current channel width $(\overline{V}_d d/L^2)^\alpha$, $\alpha < 1$. This suggests that the growth rate would be the same whether a weaker current is distributed over a larger channel or a stronger current is concentrated in a narrow channel. With regard to the auroral ionosphere, if the density gradient scale length is around 20 kms, then we see from Fig. 4, that the growth rate of smaller scale modes, say $\lambda \sim 1$ km corresponding to $k_y L \sim 100$, is reduced from $0.3\ \overline{V}_d/L$ to $0.15\ \overline{V}_d/L$ by a current channel of width $2d = 40$ km. The above growth rate corresponds to $0.75\ 10^{-3}\ s^{-1}$ if $\overline{V}_d = 50$ m/s (corresponding to currents of order $\mu A/m^2$). Since $\nu_{in} \sim \overline{V}_d/L$, these parameters correspond to the high latitude F region. We also point out that the short wavelength modes ($k_y L < 1$) are not affected by the current channels of 40 kms width.

Vickrey et al. (1980) have computed the growth rates from the current convective instability to explain small scale plasma density fluctuations near high latitude F region large scale plasma enhancements. However, their data indicate that the transverse dimensions of the currents driving the instability are comparable to the perpendicular (to the geomagnetic field) scale sizes of the plasma enhancements. Our results suggest that their growth rates are overestimates. From our analysis we would predict

V. DISCUSSION AND SUMMARY

We have presented an analysis of the effects of finite current channel width on the current convective instability. The current convective instability results from the coupling of a magnetic field aligned current and a density gradient perpendicular to the magnetic field in the presence of electron-neutral collisions. For a current with perpendicular scale size 2d and a plasma density gradient scale length L, our analytical results indicate a monotonic decrease in the growth rate γ of the current convective instability for $d/L < 1$ both in the collisional $(\nu_{in} >> \omega)$ and inertial limits $(\nu_{in} << \omega)$. For $d/L << 1$, we find $\gamma \propto \overline{V}_d d/L^2$ in the collisional limit while $\gamma \propto (\hat{V}_d d/L^2)^\alpha$, $\alpha < 1$ in the inertial limit. In general, the growth rate of the current convective instability is reduced by the finite width of the current channel from that of the local case in which the current distribution is infinite in extent.

It is significant to note that, in the long wavelength limit $k_y d << 1$ where the nonlocal effects are most prominent, the dependence on $V_d(x)$ enters the nonlocal results in the sharp-boundary case in the form of $(\overline{V}_d d)$, the "area" under the function $V_d(x)$ in the current channel. It is clear that this quantity is proportional to the total current in the channel per unit length in the y-direction (divided by the density of the current carriers and the electric charge). This point is to be contrasted with the local dispersion relation (34) in which the growth rate depends on the current velocity itself, a local quantity. In addition, the nonlocal growth rate depends on the fractional plasma density gradient across the current channel (i.e. d/L). The growth rate is reduced because in a growth period $(\sim (V_d/L)^{-1})$ the fluid element does not sample the entire gradient $(1/L)$ but a smaller fraction of the density gradient $(f/L\ ,\ f = d/L)$. Furthermore, a single value of γ corresponds to the entire wavepacket which occupies the current channel. These nonlocal features are well-known and general to all intrinsically nonlocal perturbations such as the tearing mode (see, for example, Chen and Palmadesso, 1984). With respect to perturbations that have valid local limits, certain differences in the interpretation of theoretical and observational results must be kept in mind as one goes from the local to nonlocal regime. First, with finite current channels, observations must provide global quantities such as the

50%. Furthermore, we notice that the longer wavelength modes are more easily stabilized. That is, a wider channel ($\hat{d} \gtrsim 10$) significantly reduces the growth rate of modes with wavelength of $4\pi L$ ($\hat{k} = 0.5$), whereas channel width of $\hat{d} = 0.1$ reduces the growth rate of modes with wavelengths of $0.2\pi L(\hat{k} = 10)$ by the same amount. From figures 3 and 4, we can conclude that the growth rate of the current convective instability is reduced if the current channel width is such that $kd \lesssim 1$. Similar results were noticed in the context of collisionless current driven ion-cyclotron waves by Bakshi et al. (1983). We note that we find the real part of the frequency to be much smaller than the growth rate, $\hat{\omega}_r \sim 10^{-4} << \hat{\gamma}$.

We have also numerically solved equation (11) for a more realistic diffuse profile of density and current. We choose a density profile given by

$$n(x) = n_o \frac{1 + \varepsilon \tanh(x/L)}{1 - \varepsilon} \tag{53}$$

so that

$$\frac{n'}{n} = \frac{1}{L} \frac{\varepsilon \operatorname{sech}^2(x/L)}{1 + \varepsilon \tanh(x/L)}$$

We take $\varepsilon = 0.8$ so that n'/n is a maximum at $x/L \equiv x_o = -0.55$. We find then that $(n'/n)_{max} = 1/L$. In addition we consider a smooth profile for the current velocity $V_d(x)\hat{z}$ with

$$V_d(x) = \bar{V}_d \exp\left[-(x-x_o)^2/d^2\right] \tag{54}$$

We take the same parameters as in the sharp boundary case. In Fig. 5 we plot $\hat{\gamma}$ versus k_yL for several values of d/L. Similar nonlocal behavior is seen for the growth rate $\hat{\gamma}$. In Fig. 6, we plot $\hat{\gamma}$ versus d/L. The results shown in these figures are similar to those obtained by using a simple waterbag model (Figs. 3 and 4) leading us to believe that the simple waterbag model adequately models the essential physics, namely that the finite current channel width has a stabilizing influence on the current convective instability.

is that for $d \ll \lambda$ the growth rate scales as $(V_d d/L^2)^\alpha$, where $\alpha < 1$. This result shows that in the long wavelength limit the growth rate depends on the total current in the region where the wave packet is localized rather than on V_d as is true in the local case.

In solving equation (33), we consider the fastest growing mode ($m = 0$) and use the following parameter values: $\nu_{in}/|\overline{V}_d/L| = 1$, $\Omega_i/|\overline{V}_d/L| = 10^4$, $\nu_{ei}/\Omega_e = 10^{-4}$ and $k_z/k_y = (\nu_{ei}\nu_{in}/\Omega_e\Omega_i)$ (Chaturvedi and Ossakow, 1981; Satyanarayana and Ossakow, 1984). In Fig. 3, we plot the normalized growth rate $\hat{\gamma} \equiv \gamma/|\overline{V}_d/L|$ as a function of $\hat{k} = k_y L$ for various values of the width of the current channel normalized to density gradient scale length $\hat{d} \equiv d/L = 0.1$, 1.0, and 10.0. First, we see the nonlocal behaviour of the instability where the growth rate is reduced for $\hat{k} \equiv k_y L < 1$ and slowly increases and saturates to the local limit at $\hat{k} \gg 1$ (Huba, 1984). Second, the larger the value of $\hat{d}$, the smaller is the $\hat{k}$ value for which the local limit $\hat{\gamma}_L$ is achieved; for $\hat{d} = 10$, $\hat{\gamma}$ approaches $\hat{\gamma}_L$ for $\hat{k} \gtrsim 2$ whereas for $\hat{d} = 0.1$, $\hat{\gamma}$ approaches $\hat{\gamma}_L$ for $\hat{k} \sim 100$(not shown in figure). The local limit, $\hat{\gamma}_L$, is given by the imaginary part of Eq. (38) normalized to $|\overline{V}_d/L|$. Thus we can conclude that the current channel width has be to small ($\hat{d} \lesssim 1$) to substantially reduce the growth rate of the short wavelength modes ($\hat{k} > 1$).

In Figure 4 we plot $\hat{\gamma}$ as a function of $\hat{d}$ for $\hat{k} = 0.5$, 1.0, 10.0, and 100.0. As we discussed earlier, we see that as $\hat{d} \to \infty$ we recover the conventional current convective growth rate ($\hat{\gamma}_L$ given in Eq. (38)). The second point is that, as $\hat{d}$ is decreased, the growth rate decreases, eventually going to zero for $\hat{d} \to 0$. The quenching of the instability should be expected, because as $\hat{d} \to 0$ the currents that are driving the current convective instability become highly localized and thus do not sample the entire density gradient. Since the instability is driven by the density gradient in the presence of field aligned currents, as the current channel width (d) is reduced in relation to the density gradient scale length (L), the growth rate of the instability is reduced. The instability is most sensitive to the finite width of the current channel when the current channel width, d, is comparable to or less than the gradient scale length, L. In fact, when the wavelength is comparable to the current channel width, $\lambda \sim d$ implying $kd \sim 2\pi$, the growth rate drops by at least

In this limit where $|Q_o/Q_i| \ll 1$, $\nu_{in} \ll |\omega|$ and $k_{z1}^2 \nu_{in} \ll |\omega|$, the growth rate depends on the square root of the product of $\overline{V}_d$ and d. An important conclusion is that in both limits discussed above the growth rate scales as $(\overline{V}_d d/L^2)^\alpha$, $\alpha < 1$. We note that the growth rates given in Eqs. (42)-(46) can be shown to be less than the local growth rate, $|\omega_{\pm}|$, as assumed earlier. In the next section we solve numerically the full nonlocal dispersion relation [Eq. (33)] in order to study the growth rate over a larger range of $k_y L$ and d/L.

IV. NUMERICAL RESULTS

In order to illustrate the properties of Eq. (33) and the scaling properties of the growth rate in the collisional and inertial domains, we have chosen, for convenience, the case where $k_{z1} = 1$. This choice corresponds to $k_z = k_y(\nu_{ei}/\Omega_e)(\nu_{in}/\Omega_i) \ll k_y$. With this choice, the local dispersion relation (equation (38)) exhibits a maximum at $\nu_{in}/\Omega_i = (1/24)^{2/3} (\omega_o^2/\Omega_i^2)^{2/3}$. This is shown in Fig. 2 where we plot γ/Ω_i versus ν_{in}/Ω_i. As a numerical example, we choose $\nu_{ei}/\Omega_e = 10^{-4}$, $|\bar{V}_d/L| = 5$, d/L = 0.3 and $k_y L = 1$ (Bythrow et al., 1984). Curve A shows the exact solution of Eq. (33) obtained by solving it numerically. Curves B and C represent the growth rates given by the analytical expressions in Eqs. (46) and Eq. (42) in the inertial and collisional domains, respectively. These curves show that there is good agreement between the numerical results and the analytical expressions given in Eqs. (42) and (46). Curves B′ and C′ show the simplified analytical expressions given by Eqs. (43) and Eq. (52). Curve D gives the local growth rate, ω_-, given by Eq. (38). The figure shows that for small current channel widths, $d \sim 0.3\ \lambda/2\pi$, where λ is the wavelength of the perturbation, the growth rate is reduced by an order of magnitude. For the particular choice of the parallel wavenumber the growth rate peaks at $\nu_{in}/\Omega_i \sim 2.0\ 10^{-3}$ whereas it is maximized at $\nu_{in}/\Omega_i \sim 5.0\ 10^{-3}$ in the nonlocal limit. In addition, curves B′ and C′ show that the Eqs. (43) and (52), which contain the scaling properties of the growth rate, agree fairly well with the exact numerical results in $\nu_{in} \gg \Omega_i$ and $\nu_{in} \ll \Omega_i$ regimes, respectively. Based on this agreement we can say that the growth rate scales as $(\nu_{in}/\Omega_i)^{1/2}$ in the collisional domain and as $(\nu_{in}/\Omega_i)^{-1/2}$ in the inertial limit. A more important result

Inertial limit

In the inertial limit where $\nu_{in} << |\omega|$, the solutions of Eq. (41) are

$$\omega_< = -(i\nu_</2)\ \{1 \pm [1 + 4\ \omega_o^2\ (1 + \kappa_2)/\nu_<^2\ ]^{1/2}\}/(1 + \kappa_2) \tag{46}$$

where $\kappa_2 = (1/k_y d)$ and $\nu_< = \overline{\nu} + \nu_{in}\hat{k}_{z1}^2\kappa_2^2/2$. To obtain a simple expression for the growth rate, we consider small channel widths, (d << L), where the growth rate is much smaller than the local growth rate. Thus for $\omega << \omega_\pm$ Eq. (40) yields in the inertial limit

$$\omega_+\omega_-\hat{k}\hat{d} = -\ \omega^2\ (1 + i\hat{k}_{z1}^2\nu_{in}/\omega)^{1/2} \tag{47}$$

Equation (47) can be simplified for the case $\hat{k}_{z1}^2\ \nu_{in} >> |\omega|$, i.e., $(\hat{k}_z^2/\hat{k}^2)(\Omega_e/\nu_{ei}) >> |\omega|/\Omega_i$. The growth rate can then be written as

$$\gamma = (\nu_{ei}\Omega_i/\Omega_e)^{1/3}(\hat{k}_y\hat{d}\ )^{2/3}(|\bar{V}_d|/L)^{2/3} \tag{48}$$

From Chaturvedi and Ossakow, (1981) we have, in the local inertial limit the maximum growth rate

$$\gamma_L = (\nu_{ei}\Omega_i/\Omega_e)^{1/3}\ (|\bar{V}_d|/L)^{2/3} \tag{49}$$

As a result, we have for the nonlocal long wavelength inertial limit,

$$\gamma = (k_y d)^{2/3}\gamma_L \tag{50}$$

Equation (50) shows that the nonlocal growth rate in the inertial limit γ is less than the local growth rate since $k_y d << 1$ in the long wavelength limit. For $k_{z1}^2\nu_{in} << |\omega|$, i.e., $(\hat{k}_z^2/\hat{k}^2)\ (\Omega_e/\nu_{ei}) << |\omega|/\Omega_i$, Eq. (47) yields

$$\gamma = (\omega_+\omega_-\hat{k}\hat{d})^{1/2} \tag{51}$$

By using Eqs. (24) - (27), this expression can be rewritten as

$$\gamma = (\hat{k}_z\hat{k}_y/\hat{k})^{1/2}\ (\Omega_i)^{1/2}\ |\bar{V}_d/L|^{1/2}(d/L)^{1/2} \tag{52}$$

Collisional limit

In this limit where $\nu_{in} >> |\omega|$, Eq.(41) has the following roots:

$$\omega_{>} = -(i\nu_{>}/2) \{1 \pm [1 + 4\,\omega_o^2\,(1 + \kappa_1)/\nu_{>}^2]^{1/2}\}/(1 + \kappa_1) \tag{42}$$

where $\kappa_1 = (1 + \hat{k}_{z1}^2)^{1/2}/k_y d$ and $\nu_{>} = \nu + \nu_{in}\kappa_1$. A simple expression for the growth rate can be obtained by examining the small channel width limit, d << L, where the growth rate is much less than the local growth rate. Thus for $\omega << \omega_{\pm}$, Eq. (40) or (41) yields

$$\omega_+\omega_-\hat{k}\hat{d} = -\,i\omega\nu_{in}\,(1 + \hat{k}_{z1}^2)^{1/2}$$

from which the growth rate $(\omega = \omega_r + i\gamma)$ is given as

$$\gamma = (\hat{k}_z\hat{k}_y/\hat{k})\,(1 + \hat{k}_{z1}^2)^{-1/2}\,(\Omega_i/\nu_{in})\,|\bar{V}_d/L|\,(d/L) \tag{43}$$

It is worthwhile to compare the growth rate in nonlocal long wavelength limit as given by Eq. (43) with the local result as given by Chaturvedi and Ossakow, (1981)

$$\gamma_L = (k_z/k_y)\,(\Omega_i/\nu_{in})\,(\bar{V}_d/L)\,(1 + \hat{k}_{z1}^2)^{-1} \tag{44}$$

We then find the simple result

$$\gamma = k_y d(1 + \hat{k}_{z1}^2)^{1/2}\gamma_L \tag{45}$$

Since $k_y d < 1$ in the nonlocal long wavelength limit, the nonlocal growth rate is less than the local growth rate. Several important conclusions can be drawn from this nonlocal analysis. In this collisional long wavelength limit, (i) the growth rate depends on the product of the magnitude of the drift velocity, $\bar{V}_d$, and the current channel width, d; (ii) the growth rate scales inversely with ν_{in}, the ion-neutral collision frequency; and (iii), the growth rate scales inversely with the gradient scale length, L.

We set $m = 0$ and consider the fastest growing mode. Then, with the definition of $\hat{Q}_i$ from Eq. (23), Eq. (35) becomes

$$- \hat{k}^2 \frac{(1 - \omega_+/\omega)(1 - \omega_-/\omega)}{(1 + i\nu_{in}/\omega)} = \frac{\pi^2}{4\hat{d}^2} \tag{36}$$

Thus, as we approach the local limit, we have $L_x = 2d/\pi$ and L_x plays the role of the wavelength in the x direction. For $\nu_{in}/|\omega| \ll 1$ the new roots are

$$\omega = \frac{-(i\bar{\nu}/2)\,\{1 \pm [1 + 4\,\omega_o^2\,(1 + \pi^2/4\hat{k}^2\hat{d}^2)/\bar{\nu}^2\,]^{1/2}\}}{(1 + \pi^2/4\hat{k}^2\hat{d}^2)} \tag{37}$$

Furthermore, we regain the local growth rates for $k_y d \to \infty$, namely,

$$\omega \equiv \omega_{\pm} = (-i\bar{\nu}/2)\{1 \pm [1 + 4(\omega_o^2/\bar{\nu}^2)]^{1/2}\} \tag{38}$$

In the limit $|\hat{Q}_o/\hat{Q}_i| \ll 1$, the inverse tangent function is approximately equal to $\hat{Q}_o/\hat{Q}_i$. Then Eq. (33) becomes

$$\hat{Q}_o/\hat{Q}_i = \hat{Q}_i\ \hat{d} \tag{39}$$

Using Eqs. (22) and (23) we obtain

$$(\omega - \omega_+)(\omega - \omega_-) = -\,\omega(\omega + i\nu_{in})[1 + i\hat{k}_{z1}^2\nu_{in}/(\omega + i\nu_{in})]^{1/2}/\hat{k}\hat{d} \tag{40}$$

where $\hat{k}_{z1}^2 = (\hat{k}_z^2/\hat{k}^2)(\Omega_e/\nu_{ei})(\Omega_i/\nu_{in})$. The above equation is valid for $|\hat{Q}_o/\hat{Q}_i| \ll 1$ or $|\hat{Q}_i d| \ll 1$ which is essentially the long wavelength limit, $k_y d \ll 1$. Equation (40) can then be solved both in the highly collisional and weakly collisional domains. Using Eqs. (24) and (25), Eq. (40) can be written as

$$\omega^2 + i\bar{\nu}\,\omega + \omega_o^2 \approx \pm\,\omega(\omega + i\nu_{in})[1 + i\,\hat{k}_{z1}^2\nu_{in}/(\omega + i\nu_{in})]^{1/2}/\hat{k}\hat{d} \tag{41}$$

The negative (positive) sign on the right hand side yields a growing (damped) mode. In the following anlaysis we consider the growing modes.

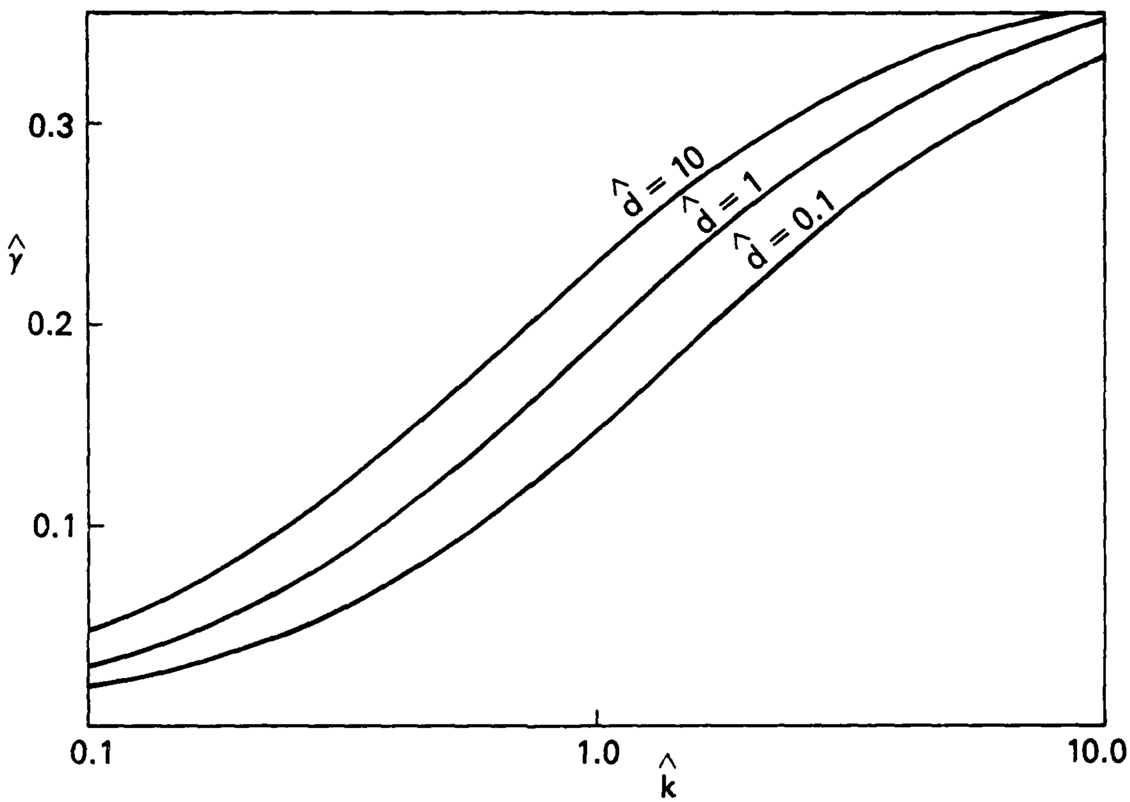

Figure 5

Plot of growth rate $\hat{\gamma}$ versus $\hat{d}$ for $\hat{k}$ = 0.5, 1.0, 10.0 and 100.0, and for the same parameters as in Fig. 3. Note that $\hat{\gamma} \to \hat{\gamma}_L$ (Eq. 38) for $\hat{d} \to \infty$.

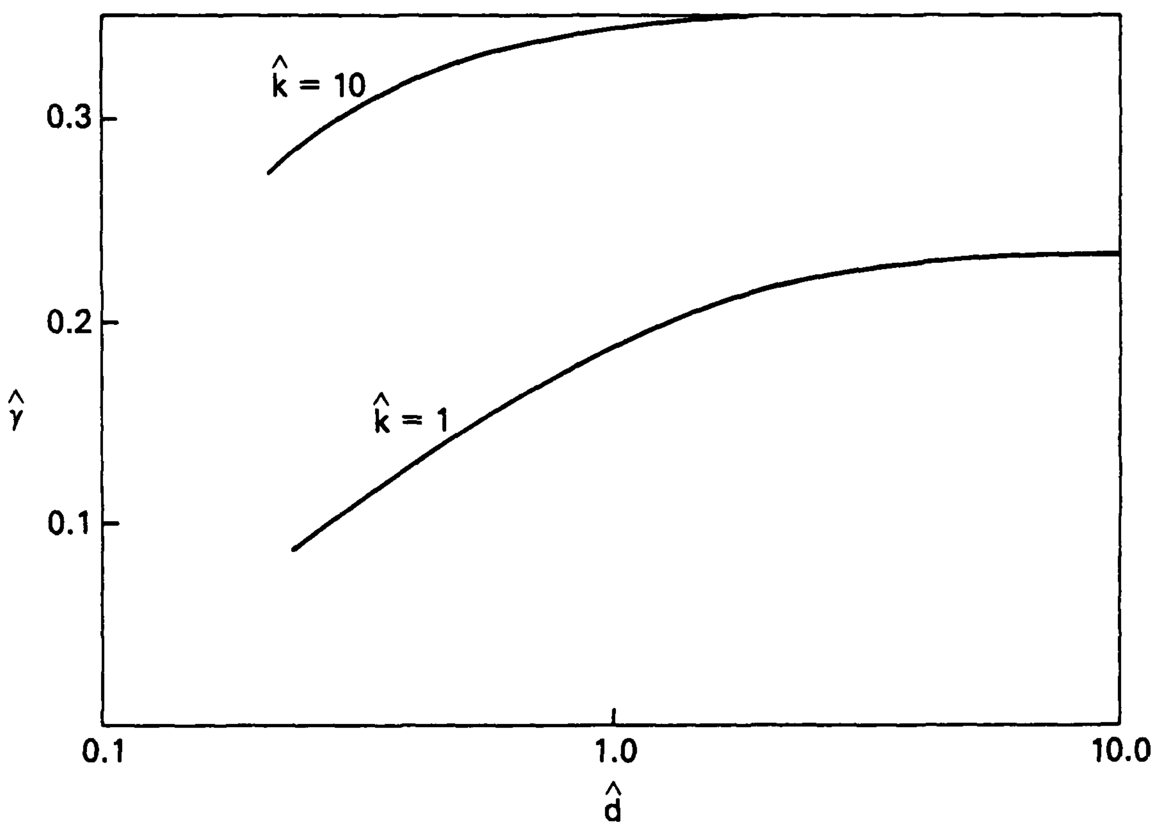

Figure 6

Plot of growth rate $\hat{\gamma}$ versus finite current channel width $\hat{d}$ for diffuse boundary case and for the same parameters as in Fig. 3.

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