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AD-A209 993

NRL Memorandum Report 6481

Scaling Laws for the Spectrum of Interchange Instabilities in the High Latitude Ionosphere

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June 15, 1989

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SECURITY CLASSIFICATION OF THIS PAGE

REPORT DOCUMENTATION PAGE				Form Approved OMB No 0704 0188	
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			5 MONITORING ORGANIZATION REPORT NUMBER(S)		
4 PERFORMING ORGANIZATION REPORT NUMBER(S) NRL Memorandum Report 6481			7a NAME OF MONITORING ORGANIZATION		
6a NAME OF PERFORMING ORGANIZATION Naval Research Laboratory		6b OFFICE SYMBOL (if applicable) Code 4780		7b ADDRESS (City, State, and ZIP Code)	
6c ADDRESS (City, State, and ZIP Code) Washington, DC 20375-5000			9 PROCUREMENT INSTRUMENT IDENTIFICATION NUMBER		
8a NAME OF FUNDING/SPONSORING ORGANIZATION DNA/ONR		8b OFFICE SYMBOL (if applicable)		10 SOURCE OF FUNDING NUMBERS	
8c ADDRESS (City, State, and ZIP Code) DNA, Washington, DC 20305 ONR, Arlington, VA 22203			PROGRAM ELEMENT NO MIPR 88-5 26		PROJECT NO RB/RC RR033-02- 44
			TASK NO 00166		WORK UNIT ACCESSION NO
11 TITLE (Include Security Classification) Scaling Laws for the Spectrum of Interchange Instabilities in the High Latitude Ionosphere					
12 PERSONAL AUTHOR(S) Keskinen, M.J.					
13a TYPE OF REPORT Interim		13b TIME COVERED FROM _____ TO _____		14 DATE OF REPORT (Year, Month, Day) 1989 June 15	
15 PAGE COUNT 28					
16 SUPPLEMENTARY NOTATION					
17 COSATI CODES			18 SUBJECT TERMS (Continue on reverse if necessary and identify by block number)		
FIELD	GROUP	SUB-GROUP	Interchange instabilities, Power spectra, Scaling laws.		
19 ABSTRACT (Continue on reverse if necessary and identify by block number) Constraints on and scaling laws associated with the power spectrum of interchange ($E \times B$) instabilities in the inertial regime in the high latitude ionosphere are derived using conservation laws implied by the fundamental nonlinear plasma fluid equations. Applications to large scale convecting high latitude ionospheric plasma density structures, patches, and enhancements are made. <i>Key words:</i>					
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, JUN 86

Previous editions are obsolete

S/N 0102-LF-014-6603

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SCALING LAWS FOR THE SPECTRUM OF INTERCHANGE INSTABILITIES IN THE HIGH LATITUDE IONOSPHERE

1. INTRODUCTION

In the last several years a considerable amount of both experimental [Basu et al., 1984; Basu et al., 1988; Gurnett et al., 1984; Kintner et al., 1987; Weimer et al., 1985; Vickrey et al., 1980; Vickrey et al., 1986; Cerisier et al., 1985; Baker et al., 1986; Curtis et al., 1982] and theoretical [Mitchell et al., 1985; Lysak and Carlson 1981; Lotko et al., 1987; Chaturvedi and Huba, 1987; Keskinen et al., 1988] research has been directed to the origin, modeling, and general interpretation of plasma turbulence and structure in the high latitude ionosphere and magnetosphere. It is now known [for recent reviews, see Kintner and Seyler, 1985; Temerin and Kintner, 1988; and Tsunoda, 1988] that the high latitude ionosphere and magnetosphere can be characterized as a highly turbulent and structured plasma containing density, electric field, and magnetic field fluctuations with scale sizes ranging from approximately hundreds of kilometers to centimeters.

Several different source mechanisms [Kintner and Seyler, 1985; Temerin and Kintner, 1988; Tsunoda, 1988] have been proposed to account for high latitude magnetospheric and ionospheric plasma turbulence, e.g., electric fields, particle precipitation, and plasma instabilities. The relative importance of each mechanism depends on several factors, e.g., geomagnetic latitude, seasonal effects, and scale size perpendicular to the geomagnetic field. Keskinen and Ossakow [1982, 1983] have proposed the $\underline{E} \times \underline{B}$ drift instability to be a source of high latitude ionospheric plasma density and electric field fluctuations for scale sizes from tens of kilometers to meters. This interchange-like instability, a plasma analogue of the well-known Rayleigh-Taylor instability which results when a heavy fluid is supported against gravity by a lighter fluid, is driven by both a plasma density gradient and an electric field perpendicular to the geomagnetic field. The electric field will result in a relative velocity between convecting high latitude ionospheric plasma ions and the neutral thermosphere. The $\underline{E} \times \underline{B}$ drift instability has been applied to the small scale dynamics, stability and evolution of convecting large scale ionospheric density structures recently observed in the high latitude ionosphere [Vickrey et al., 1980; Muldrew and Vickrey 1982; Buchau et al., 1983, 1985; Weber et al., 1984, 1986; Foster and Doupnik, 1984; de la Beaujardiere et al., 1985]. Several studies [Sojka and Schunk, 1986;

Schunk and Sojka, 1987; Anderson et al., 1987] have modeled the large scale, global features of these plasma density enhancements and structures. Vickrey and Kelley [1982] have considered the effects of a conducting E layer on the evolution of convecting large scale ionospheric density structures. At smaller scales, Huba et al [1983] have studied the linear theory of the $\underline{E} \times \underline{B}$ instability with an inhomogeneous electric field. Mitchell et al. [1985] have investigated the nonlinear evolution of the $\underline{E} \times \underline{B}$ instability in the high latitude ionosphere with magnetospheric coupling. They find that the primary effect of the magnetosphere is to incorporate inertia into the development of the $\underline{E} \times \underline{B}$ instability. In addition they find that plasma interchange-like instabilities in the high latitude ionosphere develop in a fundamentally different manner when magnetospheric coupling is included as opposed to the case when magnetospheric coupling is absent. Chaturvedi and Huba [1987] have added three-dimensional effects to the linear theory of the $\underline{E} \times \underline{B}$ instability in the high latitude ionosphere. It is now generally believed [Temerin and Kintner, 1988; Tsunoda, 1988] that plasma interchange-like instabilities are a major source of density and electric field fluctuations in the high latitude ionosphere in the scale size regime of a few tens of kilometers to tens of meters perpendicular to the geomagnetic field.

However, the spectrum of density or electric field fluctuations associated with the nonlinear evolution of the $\underline{E} \times \underline{B}$ instability in the coupled magnetosphere-ionosphere system has not been analytically investigated in detail. In this study we derive constraints on and scaling laws associated with the nonlinear spectrum of the $\underline{E} \times \underline{B}$ interchange instability as it may occur in the high latitude ionosphere-magnetosphere coupled system. These constraints are derived by analysis of the fundamental nonlinear equations governing the $\underline{E} \times \underline{B}$ instability in Sec. 2. We discuss and summarize our results in Sec. 3.

2. MODEL EQUATIONS AND ANALYSIS

We consider the high latitude ionospheric plasma to be a weakly ionized low $\beta = 8\pi n_e(T_e + T_i)/B_0^2$ plasma where n_e is the electron plasma density, $T_e(T_i)$ the electron (ion) temperature, and B_0 the geomagnetic field. For our coordinate system we take the ambient density gradient to be in the x-direction, i.e., $n = n_0(x)$, the ambient electric field to be in the y-

direction, i.e., $\underline{E} = E_0 \hat{y}$, and the geomagnetic field to be in the z-direction $\underline{B} = B_0 \hat{z}$. The equations used in this analysis are continuity, momentum, and charge conservation:

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

$$0 = -\frac{e}{m_e} (\underline{E} + c^{-1} \underline{V}_e \times \underline{B}) - (T_e/m_e)(\nabla n/n) - \nu_{ei}(\underline{V}_e - \underline{V}_i) \quad (2)$$

$$\left(\frac{\partial}{\partial t} + \underline{V} \cdot \nabla \right) \underline{V}_i = \frac{e}{m_i} (\underline{E} + c^{-1} \underline{V}_i \times \underline{B}) - (T_i/m_i)(\nabla n/n) - \nu_{ie}(\underline{V}_i - \underline{V}_e) - \nu_{in} \underline{V}_i \quad (3)$$

$$\nabla \cdot \underline{J} = \nabla \cdot [n(\underline{V}_i - \underline{V}_e)] = 0 \quad (4)$$

Here, n_α is the electron ($\alpha=e$) or ion ($\alpha=i$) density, $\underline{V}_\alpha$ is the electron or ion velocity, $\underline{E}$ is the electric field, e is the electron charge, m_α is the electron or ion mass, c is the speed of light, T_α is the electron or ion temperature, $\nu_{ei}(\nu_{ie})$ is the electron-ion (ion-electron) collision frequency, and ν_{in} is the ion-neutral collision frequency.

We now perturb Eq. (1)-(4) about an equilibrium and let $n_\alpha = n_0 + \delta n$, $\underline{E} = \underline{E}_0 - \nabla \delta \phi$, and $\underline{V}_\alpha = \underline{V}_{0\alpha} + \delta \underline{V}$ with δn , $\delta \underline{E}$, $\delta \underline{V} \propto \exp [i(k_x x + k_y y) - \omega t]$ where $(k_x^2 + k_y^2)^{1/2} L > 1$ and $L^{-1} = n_0^{-1}(\partial n_0 / \partial x)$. After transforming to a reference frame moving with velocity $\underline{V}_0 = (c \underline{E}_0 \times \underline{B}) / B^2 = (c \underline{E}_0 / B) \hat{x}$, Eqs. (2) and (3) yield.

$$\underline{V}_e = \underline{V}_{e0} + \delta \underline{V}_e \quad (5)$$

$$\underline{V}_i = \underline{V}_{i0} + \delta \underline{V}_i \quad (6)$$

with

$$\underline{V}_{e0} = - (v_e^2 / \Omega_e) (\partial \ln n_0 / \partial x) \hat{y}$$

$$\delta \underline{V}_e = - \frac{c}{B} \nabla \delta \tilde{\phi} \times \hat{z} - \frac{v_{ei}}{\Omega_e} \frac{v_i^2}{\Omega_i} \left(1 + \frac{T_e}{T_i} \right) \nabla \frac{\delta n}{n_0}$$

$$\underline{V}_{i0} = \left[\left(v_i^2 / \Omega_i \right) (\partial \ln n_0 / \partial x) + (v_i / \Omega_i) (c E_0 / B) \right] \hat{y}$$

$$\begin{aligned} \delta \underline{V}_i = & - \frac{c}{B} \nabla \delta \psi \times \hat{z} - \frac{v_i}{\Omega_i} \frac{c}{B} \nabla \delta \psi - \frac{v_{ie}}{\Omega_i} \frac{v_i^2}{\Omega_i} \left(1 + \frac{T_e}{T_i} \right) \nabla \frac{\delta n}{n_0} \\ & - \frac{c}{B} \frac{1}{\Omega_i} \left(\frac{\partial}{\partial t} + \underline{V}_{i0} \cdot \nabla \right) \nabla \delta \psi \end{aligned}$$

where $\delta \tilde{\phi} = \delta \phi - (T_e/e)(\delta n/n_0)$, $\delta \psi = \delta \phi + (T_i/e)(\delta n/n_0)$, $\Omega_e(\Omega_i)$ is the electron (ion) gyrofrequency, and $v_e(v_i)$ is the electron (ion) thermal velocity. Mitchell et al [1985] and Keskinen et al [1988] have shown that electric fields generated by the $\underline{E} \times \underline{B}$ instability in the high latitude ionosphere can map far into the magnetosphere on time scales comparable to the $\underline{E} \times \underline{B}$ instability growth time. As a result, a magnetic-field-line integrated model of the $\underline{E} \times \underline{B}$ instability is appropriate for application to the high latitude, near-earth space plasma [Mitchell et al., 1985]. Using Eqs. (1) and (4) one can then write [Mitchell et al., 1985],

$$\frac{\partial \Sigma_p}{\partial t} + \nabla \cdot (\Sigma_p \underline{V}) = D \nabla^2 \Sigma_p \quad (7)$$

$$\nabla \cdot \left[\Sigma_p \nabla \tilde{\phi} + C_m \left(\frac{\partial}{\partial t} + \underline{V} \cdot \nabla \right) \nabla \tilde{\phi} + \frac{T_i}{e} \nabla \Sigma_p \right] = 0 \quad (8)$$

$$\underline{V} = - \frac{c}{B} \nabla \tilde{\phi} \times \hat{z} \quad (9)$$

with

$$\Sigma_p = \int dz \frac{ne c v_i}{B \Omega_i} = \int dz \sigma_p$$

$$C_m = \frac{1}{4\pi} \int dz (c^2/V_A^2)$$

with $V_A = B/(4\pi m_i)^{1/2}$ the Alfvén speed, $D = (v_{ei} v_i^2 / \Omega_e \Omega_i) [1 + T_e/T_i]$, and $\underline{E} = -\nabla \tilde{\phi}$ with $\tilde{\phi} = \phi - (T_e/e) \ln n$. Within the context of this model, Eq. (8) then defines an effective frequency $\nu = \Sigma_p / C_m$. For fluctuations with characteristic frequency $\partial/\partial t \sim \gamma$, then ν/γ is essentially the flux-tube-integrated ratio of the Pedersen current in the ionosphere and polarization current in the magnetosphere. The quantity $\nu/\gamma < 1$ implying magnetospheric control for small Pedersen conductivities, large inertial capacitances, and large fluctuation growth rates. On the other hand, $\nu/\gamma > 1$ implying ionospheric control for large Pedersen conductances, small inertial capacitances, and small fluctuation growth rates. In order to solve Eq. (7)-(8) one must specify σ_p and V_A along a specified flux tube. We make the assumption that the magnetosphere, to lowest order, is uniform and incompressible. As a result, a continuity equation for the plasma density

in the magnetosphere is not necessary. We further take the Pedersen current $\Sigma_p \propto \omega N$ and the inertial capacitance $C_m \propto N$ with $N = \int n \, dz$ the integrated plasma density. We make the assumption that the magnetosphere is uniform and incompressible. Eqs. (7)-(9) can be further simplified giving:

$$\frac{\partial N}{\partial t} + \nabla \cdot N \underline{V} = D \nabla^2 N \quad (10)$$

$$\nabla \cdot \left[\omega N \nabla \tilde{\phi} + n \left(\frac{\partial}{\partial t} + \underline{V} \cdot \nabla \right) \nabla \tilde{\phi} + (T_i/e) \omega \nabla N \right] = 0 \quad (11)$$

$$\underline{V} = - \frac{c}{B} \nabla \tilde{\phi} \times \hat{z} \quad (12)$$

Eq. (10)-(12) are now a closed system in N and $\tilde{\phi}$.

Eq. (10)-(11) have been used, in the limit $T_i, T_e = 0$, to study the nonlinear evolution of the interchange $\underline{E} \times \underline{B}$ instability in the high latitude ionosphere by Mitchell et al. [1985]. They find that the $\underline{E} \times \underline{B}$ instability develops in a fundamentally different manner depending on whether the collisional ($\omega/\gamma > 1$) or inertial ($\omega/\gamma < 1$) limit is taken. Here γ is the linear growth rate of the $\underline{E} \times \underline{B}$ instability. We now proceed to derive scaling laws associated with spectrum of the interchange $\underline{E} \times \underline{B}$ instability in the inertial limit. Expanding $N = n_0(x) + \delta n(x, y)$ and $\tilde{\phi} = \phi_0 + \delta \phi(x, y)$ we find

$$\frac{\partial \delta n}{\partial t} + \frac{c}{B} \hat{z} \times \nabla \delta \phi \cdot \nabla n_0 + \frac{c}{B} \hat{z} \times \nabla \delta \phi \cdot \nabla \delta n = D \nabla^2 \delta n \quad (13)$$

Multiplying Eq. (13) by δn and integrating over all x and y with $d^2x = dx dy$ we find

$$\begin{aligned}
& \frac{1}{2} \frac{\partial}{\partial t} \int d^2x (\delta n)^2 + \frac{c}{B} \int d^2x \delta n \hat{z} \times \nabla \delta \phi \cdot \nabla n_0 \\
& + \frac{c}{B} \int d^2x \delta n \hat{z} \times \nabla \delta \phi \cdot \nabla \delta n - D \int d^2x \delta n \nabla^2 \delta n = 0
\end{aligned} \tag{14}$$

In Eq. (14), the second term on the left-hand side, resembling a source term, is proportional to the initial density gradient while the fourth term, resembling a sink, is proportional to the cross-field diffusion. The third term, which is nonlinear and represents a mode-mode coupling among the fluctuations, vanishes upon averaging over d^2x since

$$\begin{aligned}
\int d^2x \delta n \hat{z} \times \nabla \delta \phi \cdot \nabla \delta n &= \frac{1}{2} \int d^2x \hat{z} \times \nabla \delta \phi \cdot \nabla (\delta n)^2 \\
&= \frac{1}{2} \int d^2x \nabla \cdot [(\delta n)^2 \hat{z} \times \nabla \delta \phi] = 0
\end{aligned}$$

since $(\delta n)^2 \rightarrow 0$ as $x, y \rightarrow \infty$ for a finite sized ionosphere where $\nabla \cdot [(\delta n)^2 \hat{z} \times \nabla \delta \phi] = \hat{z} \times \nabla \delta \phi \cdot \nabla (\delta n)^2$. We Fourier expand δn and $\delta \phi$ as follows:

$$\begin{pmatrix} \delta n(x, y) \\ \delta \phi(x, y) \end{pmatrix} = \int d^2k \begin{pmatrix} \delta n_{\underline{k}} \\ \delta \phi_{\underline{k}} \end{pmatrix} \exp i \underline{k} \cdot \underline{x}$$

with $\underline{k} = k_x \hat{x} + k_y \hat{y}$. As a result, Eq. (14) can be written

$$\frac{1}{2} \frac{\partial}{\partial t} \int d^2k I_{\underline{k}} + i \frac{c}{B} \int d^2k \hat{z} \times \underline{k} \cdot \nabla n_0 \delta n_{-\underline{k}} \delta \phi_{\underline{k}}$$

$$+D \int d^2k \, k^2 I_{\underline{k}} = 0 \quad (15)$$

with $I_{\underline{k}} \equiv |\delta n_{\underline{k}}|^2$. Eq. (11) gives

$$\delta \phi_{\underline{k}} = \left[n_0 k^2 (i\omega - \nu) \right]^{-1} \left[\nu (i\underline{k} \cdot \underline{E}_0 + (T_i/e)k^2) \right] \delta n_{\underline{k}} \quad (16)$$

which can be inserted into Eq. (15). We note that Eq. (15) using Eq. (16) can be analyzed either in the collisional ($\nu > \omega$) limit, as was done by Keskinen and Ossakow [1981], or the inertial ($\nu < \omega$) limit. In the inertial limit we have

$$\frac{1}{2} \frac{\partial}{\partial t} \int d^2k I_{\underline{k}} = \int d^2k \gamma_{\underline{k}} I_{\underline{k}} \quad (17)$$

with

$$\gamma_{\underline{k}} = \left[\left(\nu c / B k^2 n_0 \right) \left(\underline{k} \cdot \underline{E}_0 \hat{z} \times \underline{k} \cdot \nabla n_0 \right) \right]^{1/2} - D k^2 = (\nu c E_0 \cos \theta / B L)^{1/2} - D k^2$$

with $L^{-1} = (1/n_0)(\partial n_0 / \partial x)$ and θ is the angle defined by $\underline{k}$ and $\underline{E}_0$.

Eq. (17) can be written in polar coordinates ($k^2 = k_x^2 + k_y^2$, $\theta = \tan^{-1} k_x/k_y$) assuming steady state conditions in the following manner:

$$\int_{k_{\min}}^{k_c} dk \, k \int_0^{2\pi} d\theta \, \gamma_{\underline{k},g} I(k, \theta) = \int_{k_c}^{k_{\max}} dk \, k \int_0^{2\pi} d\theta \, \gamma_{\underline{k},d} I(k, \theta) \quad (18)$$

where $\gamma_{\underline{k},g} = (\nu c \cos \theta c E_0 / B L)^{1/2}$ is the growth rate and $\gamma_{\underline{k},d} = -D k^2$ is the damping rate. As a result, the spectral power generated in the unstable range of wave numbers between $k_{\min}$ and k_c , the critical wave number, is balanced by the power dissipated in the stable range of wave numbers between k_c and $k_{\max}$ under steady state conditions.

In the following we assume a general form for the density spectrum $I(k, \theta)$ and use Eq. (18) to find constraints on the parameters used to specify $I(k, \theta)$. To our knowledge, the steady state spectrum $I(k, \theta)$ of the $\underline{E} \times \underline{B}$ instability in the high latitude ionosphere, in the inertial regime, has not been studied in detail. Keskinen and Ossakow [1982, 1983] have studied aspects of the spectrum in the collisional regime. However, Hassam et al. [1986] have investigated the nonlinear spectrum of the gravitational Rayleigh-Taylor instability, a closely related instability to the $\underline{E} \times \underline{B}$ instability, in both the inertial and collisional regimes. Hassam et al. [1986] show that the spectrum in the inertial regime is more isotropic than in the collisional regime. We take, consistent with the computations of Hassam et al. [1986], the spectrum to be of the form

$$I(k, \theta) = I_0 f(\theta) (1 + (k/k_0)^2)^{-(n+1)/2}$$

with $f(\theta) = \cos^m \theta + \eta \sin^m \theta$, $\eta < 1$, and $m > 0$. Substituting $I(k, \theta)$ into Eq. (18) we find

$$\left(\frac{k_0}{k_{\max}}\right)^{n-1} = \frac{3-n}{n-1} \frac{\Gamma\left(\frac{m+2}{2}\right)\Gamma\left(\frac{m+2}{2}\right)}{\Gamma\left(\frac{m+1}{2}\right)\Gamma\left(\frac{m+3}{2}\right)} \frac{(vcE_0/BL)^{1/2}}{Dk_{\max}^2} \quad (19)$$

where $\Gamma(q)$ is the gamma function defined by

$$\Gamma(q) = \int_0^{\infty} dy y^{q-1} e^{-y}$$

and we have assumed $n > 1$. From Eq. (19) we see that $1 < n < 3$. We note that the scaling in the inertial limit for k_0 as given by Eq. (19) is different from the collisional limit [Keskinen and Ossakow, 1981] where $(k_0/k_{\max}) \propto (cE_0/BL D k_{\max}^2)^{1/n-1}$. In order to verify the spectral scaling laws outlined in this paper, density power spectra in or near convecting patches and plasma enhancements in the high latitude ionosphere must be measured. The spectral index, outer scale k_0 , and inner scale $k_{\max}$ wavenumbers would then need to be computed. The scaling of $k_0/k_{\max}$ with

the ambient electric field E_0 in the neutral frame, say, could then be investigated. Taking $n = 2$, $m = 2$, $(v_{BL}/cE_0) = 0.2$ [Mitchell et al., 1985], $cE_0/B = 500$ m/s, $D = 1$ m²/s, $L = 10$ km and $2\pi/k_{\max} = 10$ m [R.T. Tsunoda, 1989] we find from Eq. (19) $2\pi/k_0 \approx 600$ m which is consistent with recent observations [R.T. Tsunoda, 1989].

The gravitational Rayleigh-Taylor instability, thought to be responsible for rising bubble and plumelike phenomena in the equatorial ionosphere, [Ossakow, 1981; Kelley and McClure, 1981] is analogous to the interchange $\mathbf{E} \times \mathbf{B}$ instability. The Rayleigh-Taylor instability can be described by equations identical to Eq. (1)-(4) except that the driving electric field $\mathbf{E}_0 = E_0 \hat{\mathbf{y}}$ is replaced by gravity $\mathbf{g} = -g \hat{\mathbf{x}}$. The formulae for k_0 in Eq. (19) which are applicable to interchange instability due to the gravitational Rayleigh-Taylor instability may then be computed simply by taking $E_0 \rightarrow Bg/cv_{in}$. Recently, experimental observations gathered during the PLUMEX rocket campaign [Rino et al, 1981; Kelley et al., 1982] have been further analyzed [R.C. Livingston, private communication]. These data indicate that the power spectrum of density fluctuations in equatorial Spread-F usually contains a break at a wavelength λ_B which is altitude dependent. The power spectrum is steeper (shallower) for wavelengths greater (less) than the break wavelength. Typical power spectra $P_k \propto k^{-n}$ have spectral indices $n = 1.5 - 3$ and are computed in the altitude interval of approximately 275-500 km. Furthermore, the data indicate [R.C. Livingston private communication] that λ_B is directly proportional to altitude, i.e., the break wavelength λ_B is large (small) at high (low) altitudes with $\lambda_B \approx 0.2 - 1$ km in the altitude interval $z \approx 275-500$ km. In addition, it was found [R. C. Livingston, private communication] that the altitude dependence given by v_{in} , the ion-neutral collision frequency, provides a good fit to the data, i.e., $\lambda_B \propto v_{in}^{-1}$ or $k_B \propto v_{in}$ with $k_B = 2\pi/\lambda_B$. Using Eq. (19) and the previous results of Keskinen and Ossakow [1981] we can then write for the gravitational Rayleigh-Taylor instability, $(k_0/k_{\max})^{n-1} \propto (g/L)^{1/2}/Dk_{\max}^2$ in the inertial limit, i.e., $v_{in} < (g/L)^{1/2}$ while $(k_0/k_{\max})^{n-1} \propto g/v_{in}LDk_{\max}^2$ in the collisional limit $v_{in} > (g/L)^{1/2}$. These scalings can be rewritten as $(k_0/k_{\max})^{n-1} \propto \gamma_G/\gamma_D(k_{\max})$ where γ_G is the growth rate and $\gamma_D(k_{\max})$ is the damping rate for mode k evaluated at $k = k_{\max}$. As a result $k_0^I/k_0^C \propto \gamma_G^I/\gamma_G^C$ where $k_0^I(k_0^C)$ is the spectral break wavenumber in the inertial (collisional) limit and $\gamma_G^I = (g/L)^{1/2}$ is the

growth rate in the inertial limit with $\gamma_G^C = g/v_{in}L$ the growth rate in the collisional limit. Since $\gamma_G^I < \gamma_G^C$, typically, $k_O^I/k_O^C < 1$, i.e., the spectral break wavelength $\lambda_O \propto 2\pi/k_O$ is larger in the inertial regime (high altitudes) than in the collisional regime (low altitudes) in agreement with the observations of Livingston et al. Furthermore, the scaling of the spectral break wavenumber with altitude $k_O^I/k_O^C \propto \gamma_G^I/\gamma_G^C \propto v_{in}$ is also consistent with the PLUMEX data [R.C. Livingston, private communication].

3. SUMMARY

We have derived constraints and scaling relationships for the power spectrum of the inertial $\underline{E} \times \underline{B}$ drift instability in the weakly nonlinear regime in the high latitude ionosphere by analyzing conservation laws implied by the fundamental plasma fluid equations. By assuming $I(k, \theta) \propto f(\theta)(1 + (k/k_O)^2)^{-(n+1)/2}$ where $I(k, \theta)$ is the two-dimensional density power spectrum of the $\underline{E} \times \underline{B}$ instability, $\cos \theta = \underline{k} \cdot \underline{E}_O$, $\underline{E}_O$ is the ambient electric field, $f(\theta) = \cos^m \theta + \eta \sin^m \theta$, m, η constants, and k_O is the outer scale we find that $1 < n < 3$. Furthermore, we derive the scaling of the ratio of outer (k_O) and inner (k_{max}) scales $k_O/k_{max} \propto (v_e E_O / B L D^2 k_{max}^4)^{1/2(n-1)}$ with k_{max} the inner scale. Using typical values of v, E_O, L, D , and k_{max} we find values for k_O and n which are consistent with observations.

ACKNOWLEDGMENTS

We thank S.L. Ossakow for useful discussions. This was supported by DNA and ONR.

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