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A REVIEW OF RECENT RESULTS ON SPREAD F THEORY. (U)  
JAN 79 S L OSSAKOW  
NRL-MR-3909

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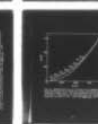
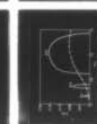
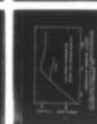
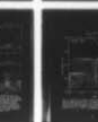
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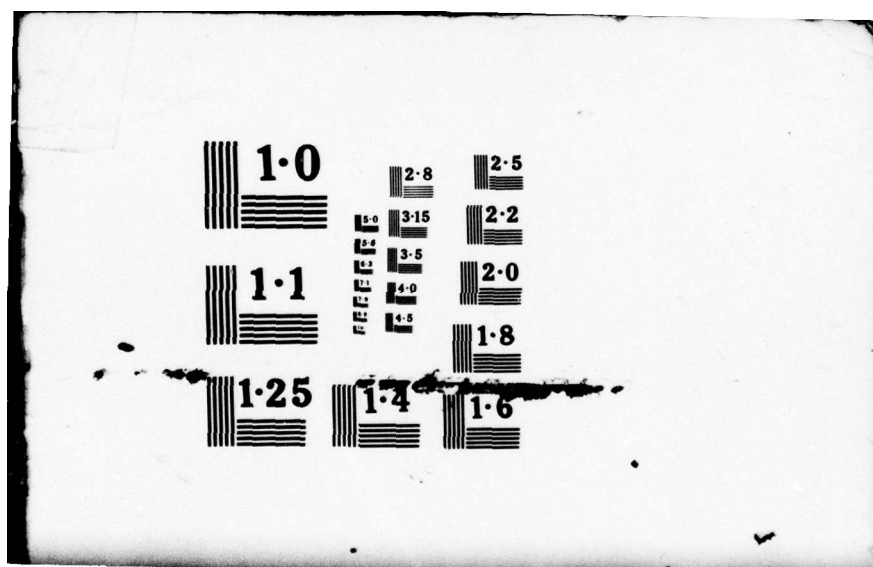
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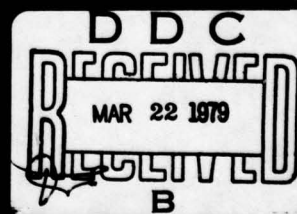
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REPORT DOCUMENTATION PAGE		READ INSTRUCTIONS BEFORE COMPLETING FORM	
1. REPORT NUMBER NRL Memorandum Report 3909	14 NRL-MR-3909	GOVT ACCESSION NO.	3. REPORT'S CATALOG NUMBER 9
4. TITLE (and Subtitle) A REVIEW OF RECENT RESULTS ON SPREAD F THEORY		5. TYPE OF REPORT & PERIOD COVERED Interim report on a continuing NRL problem	6. PERFORMING ORG. REPORT NUMBER
7. AUTHOR(s) S. L. Ossakow		8. CONTRACT OR GRANT NUMBER(s)	
9. PERFORMING ORGANIZATION NAME AND ADDRESS Naval Research Laboratory Washington, D.C. 20375		10. PROGRAM ELEMENT, PROJECT, TASK AREA & WORK UNIT NUMBERS NRL Problem H02-42D and A03-16B DNA Subtask I25AAXYX960	
11. CONTROLLING OFFICE NAME AND ADDRESS Defense Nuclear Agency, Washington, D.C. 20305 and Office of Naval Research, Arlington, Virginia 22217	119	12. REPORT DATE January 1979	13. NUMBER OF PAGES 52
14. MONITORING AGENCY NAME & ADDRESS (if different from Controlling Office)	12 57	15. SECURITY CLASS. (of this report) UNCLASSIFIED	16. DECLASSIFICATION/DOWNGRADING SCHEDULE
16. DISTRIBUTION STATEMENT (of this Report) Approved for public release; distribution unlimited. 18 SBIE		DDC RECEIVED MAR 22 1979 B	
17. DISTRIBUTION STATEMENT (of the abstract entered in Block 20, if different from Report) 19 AD-E 000 273			
18. SUPPLEMENTARY NOTES This Research was sponsored partially by the Defense Nuclear Agency under Subtask I25AAXYX960, work unit code 17 and work unit title Scintillation Description, and partially by the Office of Naval Research. (Continues)			
19. KEY WORDS (Continue on reverse side if necessary and identify by block number) Equatorial Spread F Theory Numerical simulations Review			
20. ABSTRACT (Continue on reverse side if necessary and identify by block number) Although Ionospheric Spread F was discovered some four decades ago, yet only in the past few years has significant progress been made in the theoretical explanation of such phenomena. In particular, considerable effort has been expended to explain equatorial Spread F and the attendant satellite signal propagation scintillation phenomena. The present review dwells mainly in this low latitude area. The various linear plasma instabilities thought to initiate equatorial Spread F will be discussed. Recent theoretical and numerical simulation studies of the nonlinear evolution of the collisional (Continues)			

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S/N 0102-014-6601

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## 18. Supplementary Notes (Continued)

This paper is based on the presentation made at the XIX<sup>th</sup> URSI General Assembly, Helsinki, Finland, August 1978.

## 20. Abstract (Continued)

Rayleigh-Taylor instability in equatorial Spread F ~~will be~~ reviewed. <sup>45</sup> Also, <sup>9</sup> analytical studies of rising equatorial Spread F bubbles in the collisional and collisionless Rayleigh-Taylor regime ~~will be~~ <sup>are</sup> discussed, as well as the nonlinear saturation of instabilities in these two regimes. Current theories on very small scale ~~(<sup>15</sup> 10 meters)~~ size irregularities observed by radar backscatter during equatorial Spread F and their relation to the larger wavelength scintillation causing irregularities ~~will be~~ <sup>are</sup> discussed. Application of turbulence theory to equatorial Spread F phenomena ~~will be~~ <sup>are</sup> reviewed. Remaining problems to be dealt with at equatorial latitudes ~~will be~~ summarized. <sup>13</sup>

(approximately  $\leq$  or  $=$  10 meters)

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## **A REVIEW OF RECENT RESULTS ON SPREAD F THEORY**

### **I. Introduction**

Spread F, as exhibited by diffuse echoes on ionograms, was discovered some forty years ago by Booker and Wells (1938). Until recently, there has been only much statistical data concerning Spread F. However, in the last few years with advances in radar backscatter measurements, in situ measurements, and theoretical and numerical simulation techniques a clearer picture of the fundamental plasma instability mechanisms causing equatorial Spread F phenomena has been evolving. It should be emphasized that this paper will deal only with equatorial Spread F theory and, in addition, dwell only on those theories using plasma mechanisms as a basis. For experimental results, the works of Farley et al. (1970), Dyson et al. (1974), Kelley et al. (1976), Woodman and La Hoz (1976), Morse et al. (1977), and McClure et al. (1977) should serve as good references for the interested reader.

In speaking about Spread F one should at least show its basic manifestation on an ionogram (see Fig. 1). After all, the terminology Spread F emerged from the results of ionosonde traces such as those exhibited in Fig. 1. The multiple traces are caused by magnetic field aligned irregularities. If there were no irregularities, a single crisp trace rising to the right would be exhibited on the ionogram. At this point we want to remember the basic equatorial geometry. Figure 2 exhibits the basic equatorial nighttime ionospheric F region geometry, i.e., the geometry under which equatorial Spread F occurs.

Note: Manuscript submitted November 16, 1978.



$N(y)$  represents the background electron density as a function of altitude ( $y$ ). Gravity,  $g$ , points down, the ambient magnetic field,  $B$ , is horizontal (pointing north) and  $k$  represents a horizontal perturbation (in the westward direction). The maximum in the electron density profile is the F peak. The underside of the profile steepens at night due to chemical recombination effects and electrodynamic forces. The E region has been severely reduced by chemical recombination and plays a negligible role. To a plasma physicist this geometry is a classical flute mode geometry and one might expect this equatorial ionospheric geometry to be unstable to a variety of plasma instabilities. However, until recently, one of the basic difficulties was getting the unstable irregularities to the topside (i.e., above the F peak) when they were initiated on the bottomside. One must remember that the experimental evidence has exhibited both top and bottomside irregularities.

At this point a general brief review will be given of equatorial Spread F (ESF) theories. First, we will discuss the linear theories. Dungey (1956) was the first to suggest that ESF was initiated on the bottomside by a Rayleigh-Taylor instability. Dagg (1957) suggested that the ESF phenomena was due to E to F region coupling, i.e., irregularities in the E region coupled up to the F region. In 1959, Martyn (1959) was the first to suggest that ESF was a manifestation of the  $E \times B$  gradient drift instability. Calvert (1963) proposed that the downward motion of the neutral atmosphere at night was responsible for ESF. This mechanism is essentially equivalent to the  $E \times B$

instability because of the relative motion between ions and neutrals in determining the instability. All of these previously invoked linear instability mechanisms could only explain the formation of bottomside irregularities. The collisional Rayleigh-Taylor instability with field line averaging was also proposed (Balsley et al., 1972; Haerendel, 1974) as a linear instability mechanism. By averaging (integrating) the density along the magnetic field, the total electron content profile becomes steeper on the bottomside and its peak is raised in altitude with respect to the local electron density peak. This would allow the linear mechanism to operate to slightly higher altitudes ( $\sim 100$  km greater), but still would not explain the existence of irregularities above this "new peak". Hudson and Kennel (1975) pointed out the importance of the collisional drift mode in ESF in the wavelength regime 30m - 100m. This mode could be excited on both the top and bottomside but still would not explain the longer wavelengths. In their paper, finite larmor radius (FLR) corrections were also applied to the collisionless and collisional Rayleigh-Taylor instability.

Several nonlinear theories have been invoked to explain the different ESF observations. For example, Hudson et al. (1973) suggested that the very smallest scale ( $\leq 10$ m) irregularities (e.g., those seen by radar coherent backscatter) were due to a two step process. In this prescription a longer wavelength instability sets up the driving conditions for the shorter wavelengths to become unstable. This is similar in spirit to the successful two step



theory (Sudan et al., 1973) proposed for Type II equatorial E region electrojet irregularities. Haerendel (1974) suggested that the range of wavelengths (many kilometers down to centimeters) exhibited by ESF phenomena was due to a multi-step process. This scenario is as follows: (i) the collisional Rayleigh-Taylor (R-T) instability with horizontal wavevectors is driven by gravity and the background, zero order electron density gradient scale length on the bottomside; then (ii) the  $\underline{E} \times \underline{B}$  gradient drift instability with vertical wavevectors arises due to the horizontal density, large amplitude variations set up by the collisional R-T instability; then (iii) the inertia (collisionless) dominated R-T instability arises; and finally (iv) kinetic drift waves grow upon these irregularities after they reach large amplitude. Chaturvedi and Kaw (1976) interpreted the  $k^{-2}$  measured power spectrum of the ESF plasma density irregularities in terms of a two step theory. In this theory longer wavelength R-T modes couple to kinetic collisional drift waves in such a manner that the mode coupling results in the observed  $k^{-2}$  spectrum.

A major breakthrough was made by Scannapieco and Ossakow (1976) who performed a nonlinear numerical simulation of the collisional R-T instability for ESF geometry. The simulation results showed that the collisional R-T instability generated irregularities and bubbles (plasma density depletions) on the bottomside of the F region which subsequently rose beyond the F peak by nonlinear polarization induced  $\underline{E} \times \underline{B}$  forces. This was the first theoretical result to explain how long wavelength irregularities could appear on both the bottomside and



topside of the F region. The bubble phenomena was in accord with the recent observations (Kelley et al., 1976; McClure et al., 1977, Woodman and La Hoz, 1976) of plasma density depletions. An analytical nonlinear mode-mode coupling theory for the coherent development of the collisional R-T instability was performed by Chaturvedi and Ossakow (1977). This theory suggested that vertical modes would be dominant and result in a  $k^{-2}$  power spectrum. Hudson (1978) extended the previous results to the collisionless R-T regime and reached similar conclusions. Analytical models for the rise of collisional and collisionless R-T ESF bubbles, in analogy with fluid bubbles, was presented by Ott (1978). At the same time, Ossakow and Chaturvedi (1978) presented analytical models for the rise of collisional R-T ESF bubbles within the context of the electrical analogy with barium clouds.

Costa and Kelley (1978a,b) suggested that coherent steepened structures and not turbulences would give a  $k^{-2}$  power spectrum. Moreover, these sharp gradients could cause small scale sizes ( $\sim 20m$ ) by collisionless low frequency (much less than the ion gyrofrequency,  $\Omega_i$ ) kinetic drift waves via a two step process. Their analysis was a linear one carried out on a nonlinear state, i.e., one achieves the steepened gradients by nonlinear processes and then one performs linear theory on this state. Kelley and Ott (1978) suggested that the ESF bubbles, in the collisionless R-T regime, generate a wake with vortices. They then applied two dimensional fluid turbulence theory to the model. This resulted in the development of turbulence at

shorter and longer wavelengths than the bubble size. This in turn led to a prediction of  $k^{-1}$  for the power spectrum (which does not appear to be in agreement with existing experimental observations) in the range  $L_S^{-1} < k < L_D^{-1}$ , where  $L_S$  is the stirring (bubble) size and  $L_D$  is a dissipation length cutoff. In a continuation of the numerical simulation work, Ossakow et al. (1978) showed a more rapid ESF development and higher bubble rise velocities resulting from sharper bottom-side background electron density gradients and higher altitudes of the F peak. In Huba et al. (1978) very small scale (wavelengths  $\sim 1\text{m}$  and  $36\text{cm}$ ) irregularities are reported. A two step process, utilizing high frequency ( $\geq \Omega_i$ ) kinetic drift cyclotron or lower hybrid drift instabilities, is invoked to explain them. Linear theory for these instabilities, in the ESF environment, was performed on the nonlinear state with encouraging results.

The above introductory remarks glaringly point out that much work and significant progress in the theoretical area of equatorial Spread F has been accomplished in the past few years (indeed just look at the number of publications during 1978 alone). Notwithstanding the recent successes, much work still needs to be done. Indeed, the theoretical and numerical simulation efforts in ESF are continuing along a hot and heavy path. Section II of this paper presents outlines of some of the theoretical efforts briefly mentioned in the preceding paragraphs. Given the length limitations, it would be exceedingly difficult to outline all of the theoretical works mentioned or even to present all of the details of a few works.

Section II, hopefully, will wetten the reader's appetite to read the referenced works. Section III presents a summary concerning ESF theory.

## II. Theory

In this section we present some representative theoretical and numerical simulation works with the appropriate references.

a. General. The basic plasma fluid equations applicable to the equatorial Spread F ionosphere are as follows:

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

$$-T_e \nabla n_e - en_e (-\nabla \phi + \frac{\underline{V}_e \times \underline{B}_0}{c}) = 0 \quad (2)$$

$$\begin{aligned} m_i n_i \left( \frac{\partial}{\partial t} + \underline{V}_i \cdot \nabla \right) \underline{V}_i &= -T_i \nabla n_i \\ &+ en_i (-\nabla \phi + \frac{\underline{V}_i \times \underline{B}_0}{c}) + m_i n_i \underline{g} - m_i n_i \nu_{in} \underline{V}_i \end{aligned} \quad (3)$$

$$\nabla \cdot \underline{J} = 0 \quad (4)$$

$$\underline{J} = ne (\underline{V}_i - \underline{V}_e) \quad (5)$$

In the above equations the subscript  $\alpha$  denotes species ( $e$  is electron,  $i$  is ion),  $n$  is density,  $\underline{V}$  is velocity,  $P$  is the production,  $\nu_R$  is the chemical recombination rate,  $T$  is temperature,  $\nabla$  is the gradient



operator,  $\underline{B}_0$  is the ambient magnetic field (taken to be uniform),  $e$  is the electronic charge,  $c$  is the speed of light,  $m$  is mass,  $g$  is gravity,  $\nu_{in}$  is the ion-neutral collision frequency,  $\underline{J}$  is current, and the electrostatic approximation has been made where  $\underline{E} = -\nabla\phi$ .

Equation (1) is the continuity equation, (2) and (3) are the electron and ion momentum equations, respectively, (4) is the divergence of the current and (5) is the current equation. What we have in mind is to apply the set of equations (1) - (5) to the two dimensions perpendicular to  $\underline{B}_0$  at the geomagnetic equator, making various approximations.

Assuming a harmonic perturbation dependence of the form  $\exp-i(\underline{k} \cdot \underline{x}_\perp - \omega t)$ , where  $\perp$  denotes perpendicular to  $\underline{B}_0$  ( $\underline{k}$  is horizontal) and linearizing equations (1) - (5) we obtain for the linear growth rate

$$\gamma = \left[ -\nu_{in} + (\nu_{in}^2 - 4g \cdot \frac{\nabla n_0}{n_0})^{1/2} \right] / 2 - \nu_R \quad (6)$$

$$\omega \equiv \omega_r + i\gamma$$

which reduces to

$$\gamma = \begin{cases} \frac{g}{\nu_{in} L} - \nu_R, & \nu_{in}^2 \gg 4g/L \\ \frac{(g)^{1/2}}{L} - \nu_R, & \nu_{in}^2 \ll 4g/L \end{cases} \quad (7a)$$

$$(7b)$$

$$|L| = \left| \frac{\nabla n_0}{n_0} \right|$$

in the collision dominated and inertia dominated regimes, for  $kL \gg 1$ . Equations (7a) and (7b) represent the R-T instability, including recombination damping, in the respective regimes. (It should be noted that this only represents instability on the bottom-side of the F region where the first term in (7a) and (7b) is positive.)

b. 2D Computer Simulation Results. In this section we will outline the two dimensional ( $\perp \underline{B}_0$ ) computer simulation results (Scannapieco and Ossakow, 1976; Ossakow et al., 1978). The simulations follow the nonlinear evolution of the collisional R-T instability; consequently, in eqn. (3) inertial terms (i.e., the left hand side of the eqn.) are neglected. Furthermore, one takes for the F region  $v_{in}/\Omega_i \ll 1$  ( $\Omega_i = eB_0/m_i c$ ) and the quasi-neutrality assumption is made, i.e.,  $n_e \approx n_i \approx n$ . Equations (1) - (5) then become

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot (n_\alpha \underline{V}_\alpha) = -v_R (n_\alpha - n_{\alpha 0}) \quad (8)$$

$$\underline{V}_e = \frac{c}{B_0} \underline{E} \times \hat{z} \quad (9)$$

$$\underline{V}_i = \left( \frac{\underline{g}}{\Omega_i} + \frac{c}{B_0} \underline{E} \right) \times \hat{z} + \frac{v_{in}}{\Omega_i} \left( \frac{\underline{g}}{\Omega_i} + \frac{c}{B_0} \underline{E} \right) \quad (10)$$

$$\nabla \cdot \underline{J} = 0, \quad \underline{J} = ne(\underline{V}_i - \underline{V}_e) \quad (11)$$

where one has taken  $T_e = T_i = 0$  for simplicity (see Ossakow et al., 1978),  $\underline{B}_0 = B_0 \hat{z}$  and the subscript o refers to equilibrium quantities.

Making the electrostatic assumption

$$\underline{E} = - \nabla \phi \quad (12)$$

and breaking the potential into an equilibrium and a perturbed quantity,

$$\phi = \phi_0 + \gamma \quad (13)$$

eqns. (8) - (11) become

$$\frac{\partial n}{\partial t} - \frac{c}{B_0} (\nabla \gamma \times \hat{z}) \cdot \nabla n = - v_R (n - n_0) \quad (14)$$

$$\nabla \cdot (v_{in} n \nabla \gamma) = \frac{B_0}{c} (g \times \hat{z}) \cdot \nabla n \quad (15)$$

where eqns. (14) and (15) are taken to be two dimensional ( $\perp B_0$ ). Linearizing eqns. (14) and (15), taking a horizontal perturbation results, as in eqn. (7a) but in a more illustrative form, in the linear growth rate

$$\gamma = - \frac{g}{v_{in}} \cdot \frac{\nabla n_0}{n_0} - v_R$$

This clearly shows that only the bottomside of the F region where  $\nabla n_0$  is positive can be linearly unstable (and only if the first term  $> v_R$ ).

Equations (14) and (15) were solved numerically using a vertical mesh spacing of  $\Delta y = 2$  km and total y extent of 200km, and an east-west, horizontal mesh spacing  $\Delta x = 200$ m and a total horizontal extent of 8km. Realistic profiles of  $v_{in}$  and  $v_R$  as a function of



altitude were utilized. The system was initialized with a perturbation of a few percent in the horizontal ( $x$ ), east-west direction with a wavelength  $\sim 3\text{km}$  and the evolution in time of (14) and (15) was followed for different background electron density profiles. Figure 3 shows the results for a background electron density,  $n_0$ , profile with an F peak at 354km and a minimum bottomside background electron density gradient scale length,  $L \sim 10\text{km}$ . In this case at  $t = 4000$  sec a bubble (plasma density depletion) is clearly forming and beginning to rise in the central portion of the mesh (note  $n \equiv n_0 + n_1$ ). The isodensity contours are such that the maximum absolute value of the enhancement or depletion is in the center and the contours decrease (in absolute percentage) as one goes toward the outer contours. At  $t = 4000$  sec the maximum depletion within the bubble is 54% and the maximum enhancement over the mesh is 84%. At  $t = 8000$  sec we notice that the bubble has reached the altitude of the F peak, with the innermost contour of the rising bubble representing a 41% depletion. At  $t = 10^4$  sec the main bubble is clearly through the F peak with an innermost depletion contour of 41%. However, in the ionosphere below the bubble near  $x = 0$  there is a 71% depletion contour, similarly in the wings near  $|x| = 4\text{km}$ . The innermost enhancement contour, at this time, represents a 236% enhancement with a maximum inside this contour of 294%. Note that the top of the main bubble is at an altitude of 375km while the bottom trail of the bubble is at an altitude of 270km. Between  $t = 8000$  and  $10^4$  sec the bubble has risen  $\sim 24\text{km}$  which represents a rise velocity  $\sim 12\text{m/sec}$ .

Also note that the bubble is  $\sim 1$  km wide. Figure 4 depicts contours of constant induced potential  $\phi_1$  at  $t = 10^4$  sec. This shows that the more isolated part of the high altitude bubble depicted in Fig. 3 is acted on by an induced electric field which points from west to east and is dipolar in nature. This causes the bubble to rise with a  $(-c/B_0) \nabla \phi_1 \times \hat{z}$  velocity. However, the lower portion of the mesh is acted on by an induced electric field which points from east to west. This field is much weaker than the induced field acting on the isolated portion of the central bubble. The lower altitude electric field causes the enhancements and depletions to move downward. Thus, the lower altitude portion of the central bubble becomes captured by the enhancements.

Figure 5 presents for comparison a case in which the background electron density profile (shape) was kept the same, but the entire profile moved up in altitude so that the F peak was at 434 km. All other parameters are the same as in Figure 3, except  $v_{in}$  and  $v_R$  are taken for the altitude range 332 km to 532 km (those used in this simulation). One can immediately note the more rapid time evolution of the Spread F process with respect to that presented in Fig. 3. At  $t = 700$  sec, a rising bubble with an innermost contour of 79% and a maximum depletion inside this contour of 84% were noted. At  $t = 1000$  sec the bubble has reached the peak and at this time the innermost depletion contour is 85%. The long trail associated with the high altitude bubble has a 100 km extension to lower altitudes. At  $t = 1400$  sec the top of the main bubble is at an altitude  $\sim 500$  km.

and has a long trail connecting to an altitude of 357km. There is a maximum 70% depletion within the innermost contour of the high altitude bubble. Between  $t = 10^3$  and 1400 sec the top part of the bubble rose  $\sim 65$ km and this represents a rise velocity  $\sim 160$ m/sec. Potential contour results for this simulation show similar patterns to those exhibited in Fig. 4. Naturally, the induced electric fields causing the bubble to rise in the present case are stronger. This spread in bubble rise velocities has been observed by AE satellite data (McClure et al., 1977).

Other numerical simulations in this series have been performed and the paper by Ossakow et al. (1978) should be consulted for more details. The basic conclusions reached from these simulations are as follows: (i) the collisional R-T instability causes linear growth on the bottomside of the equatorial Spread F region; (ii) plasma density depletions (bubbles) steepen on their top and nonlinearly rise to the topside ionosphere, beyond the F peak, by polarization (induced)  $\mathbf{E} \times \mathbf{B}$  forces; and (iii) high altitude of the F peak, small bottomside background electron density gradient scale lengths, and large percentage depletions yield large vertical bubble rise velocities, with the first two conditions favoring collisional R-T linear growth (instability). In addition, large spatial bubbles with similar rise velocities to those presented here, but with almost 100% depletions, have been produced by numerical simulations (Zalesak et al., 1978). In these cases the horizontal mesh covered 200km in extent (intermesh spacing  $\Delta x = 5$ km) and a long wavelength ( $\sim 75$ km) initial perturbation was



used. Bubbles with horizontal dimensions  $\sim 50\text{km}$  resulted. The NRL group has also added neutral wind effects to these simulations and found that an eastward neutral wind results in a westward motion of the bubbles in addition to its rise. This is also in agreement with observations of bubble motion (see McClure et al., 1977).

c. Analytical 2D Coherent Mode Coupling Results. A two dimensional nonlinear quasi-final state of the collisional R-T instability was investigated by Chaturvedi and Ossakow (1977) using analytical means and considering coherent mode coupling as the saturation mechanism. Equations (14) and (15) were utilized with  $n = n_0 + \tilde{n}$ , etc. This yields the following coupled nonlinear equations

$$\frac{\partial \tilde{n}}{\partial t} - \frac{c}{B_0} \nabla \phi \times \hat{z} \cdot \nabla n_0 = - \nu_R \tilde{n} + \frac{c}{B_0} \nabla \phi \times \hat{z} \cdot \nabla \tilde{n} \quad (16)$$

$$\frac{g_{xz}}{\Omega_i} \cdot \nabla \tilde{n} - \frac{c}{B_0 \Omega_i} \left[ n v_{in} \nabla^2 \phi + \nabla \phi \cdot \nabla (\tilde{n} v_{in}) \right] = 0 \quad (17)$$

where the second term on the LHS of (16) represents growth, the first term on the RHS represents damping, and the second term on the RHS is the nonlinear term. Comparing the nonlinear term in (16) with the last term in (17) results in

$$\frac{c}{B_0} \nabla \phi \times \hat{z} \cdot \nabla \tilde{n} : \frac{c}{B_0 \Omega_i} \nabla \phi \cdot \nabla (\tilde{n} v_{in}) \approx \frac{\Omega_i}{v_{in}} \gg 1$$

Therefore, eqn. (17) is treated linearly and the nonlinearity is retained in (16). A perturbation of the form

$$\frac{\tilde{n}}{n_0} = A_{1,1} \sin(k_y y - \omega t) \cos k_x x + A_{2,0} \sin 2k_x x \quad (18)$$

is chosen. In this analysis  $x$  is vertical (altitude) and  $y$  horizontal (east-west). This is the way it appears in the reference. For the convenience of the reader, in referring to the original work, we have kept the coordinate system of the reference.

The coupled mode equations for the amplitudes  $A_{1,1}$  and  $A_{2,0}$  become

$$\frac{\partial A_{1,1}}{\partial t} = \gamma_{1,1} A_{1,1} - 2\alpha A_{1,1} A_{2,0} \quad (19)$$

$$\frac{\partial A_{2,0}}{\partial t} = \gamma_{2,0} A_{2,0} + \frac{\alpha}{2} A_{1,1}^2 \quad (20)$$

where the coupling coefficient  $\alpha = k_x k_y^2 g/k^2 v_{in}$ , and  $\omega \equiv \omega_r + i\gamma$ . It should be noted that the linear growth rate  $\gamma_{2,0}$  is negative, i.e., the purely vertical mode  $A_{2,0}$  is linearly damped. Also  $\gamma_{1,1}$  is positive and represents linear growth of the mode  $A_{1,1}$ . In the saturated steady state one has  $\partial A_{1,1}/\partial t = \partial A_{2,0}/\partial t = 0$  and from (19) and (20) this results in

$$A_{2,0} = \frac{\gamma_{1,1}}{2} \approx \frac{1}{2 k_x L} \quad (21)$$

$$A_{1,1} \approx \left( \frac{k^2}{k_y^2} \frac{v_R v_{in}}{gL k_x^2} \right)^{1/2} \approx \left( \frac{-2\gamma_{2,0}}{\alpha} A_{2,0} \right)^{1/2} \quad (22)$$

For typical values of the parameters,  $A_{2,0} \gg A_{1,1}$  and shows  $A_{2,0} \propto k_x^{-1}$  (i.e., the power spectrum would  $\propto k_x^{-2}$ ). This represents a coherent nonlinear evolution where the linearly damped mode  $A_{2,0}$  is generated nonlinearly by the linearly unstable mode  $A_{1,1}$ , by a harmonic generation ( $A_{1,1}$  contains  $k_x$  and  $A_{2,0}$  contains  $2k_x$ ). More detailed information can be found in Chaturvedi and Ossakow (1977).

d. Analytical Models for ESF Bubbles. First we will discuss collisional R-T bubbles (Ossakow and Chaturvedi, 1978) in which the electrical analogy with plasma density enhancement (e.g., barium clouds) was utilized and results obtained for general 2D ( $\underline{B}_0$ ) bubble shapes. Here eqns. (14) and (15) are utilized with the further simplifying assumptions of neglecting recombination chemistry ( $v_R = 0$ ) and the explicit altitude dependence of the ion-neutral collision frequency,  $v_{in}$ . The following set of equations result.

$$\frac{\partial n}{\partial t} - \frac{c}{B_0} (\nabla \phi \times \hat{z}) \cdot \nabla n = 0 \quad (23)$$

$$\nabla \cdot (n \nabla \phi) = \underline{E}^* \cdot \nabla n \quad (24)$$

$$\underline{E}^* = \underline{E}_0 + \frac{\Omega_i}{v_{in}} \frac{m_i}{e} \underline{g} \times \hat{z} \quad (25)$$

where an ambient horizontal electric field  $\underline{E}_0$  has been included to show more generality. Equation (24) can be thought of, in general, as a potential equation for a dielectric immersed in the applied electric field,  $\underline{E}^*$ . The plasma density depletion (bubble) is analogous to the case of a cavity immersed in a dielectric with a



uniform electric field ( $\underline{E^*}$ ). In general (23) and (24) do not admit of a two dimensional ( $\perp B_0$ ), steady state solution. However, for the case of a constant density inside the depletion and a constant density outside the depletion, two dimensional steady state solutions can be obtained.

For a constant density depletion with an elliptical shape one has

$$n(x,y) = n_0 - n_D(x,y)$$

$$= n_0 \left[ 1 - \frac{\delta n}{n_0} F(x,y) \right] \quad (26)$$

$$n_D = \delta n H \left[ 1 - \left( \frac{x}{a} \right)^2 - \left( \frac{y}{b} \right)^2 \right] \quad (27)$$

$$H(x) = \begin{cases} 1, & x > 0 \\ 0, & x < 0 \end{cases}$$

where H is the Heaviside function (note the geometry is taken such that x is vertical and y horizontal). For this elliptical shape, the solution to (24), neglecting  $\underline{E_0}$ , is

$$\frac{\partial \tilde{\phi}}{\partial y} = - \frac{\Omega_i}{v_{in}} \frac{m_i g}{e} \left( \frac{a}{b + a(1 - \delta n/n_0)} \frac{\delta n/n_0}{1} \right) \quad (28)$$

Using  $-(c/B_0) \nabla \tilde{\phi} \times \hat{z}$  this further results in the nonlinear vertical bubble rise velocity,  $v_B$ , given by

$$v_B = \frac{g}{v_{in}} \left( \frac{a}{b + a(1 - \delta n/n_0)} \frac{\delta n/n_0}{1} \right) \quad (29)$$

Limiting cases of (29) for sheet, cylindrical and slab bubble

geometries are given by

$$\frac{v_B}{g/v_{in}} = \begin{cases} (\delta n/n_o) (1 - \delta n/n_o)^{-1}, & b \ll a & (30a) \\ (\delta n/n_o) (2 - \delta n/n_o)^{-1}, & b = a & (30b) \\ 0 & , b \ll a & (30c) \end{cases}$$

For typical values of  $v_{in}$  as a function of altitude, Fig. 6 exhibits rise velocities given by (30b) for various percentage depletions. Table 1 shows the rise velocities, in units of  $g/v_{in}$ , for various shapes. (Note that the linear case comes from linearizing (23) and (24)). All of the geometry results can be expressed in a concise formula

$$v_B = \frac{g}{v_{in}} f\left(\frac{\delta n}{n_o}\right) \quad (31)$$

where  $f(\delta n/n_o)$  is an increasing function of the percentage depletion  $\delta n/n_o$ . Basically, the results predict that high altitudes and/or large percentage depletions yield high vertical rise velocities for the bubbles (in agreement with experimental observations).

Collisional and collisionless (inertia dominated) two dimensional cylindrical R-T bubbles have been studied analytically by Ott (1978). This study is based on the analogy with fluid dynamic flows and brings forth some of the work done on two-dimensional fluids. This study begins with the basic equations (1) - (3), considers two dimensions ( $\perp B_o$ ), sets  $n_e \approx n_i \approx n$ , and makes the assumption that

$$\Omega_i \gg \frac{\partial}{\partial t}, \quad \underline{v}_i \cdot \nabla, \quad v_{in} \quad (31)$$

To lowest order, using (31), one obtains from (3) a lowest order ion

velocity (with  $\hat{z} = \underline{B}_0 / |\underline{B}_0|$ ),

$$\underline{v}_i^{(0)} = \frac{\hat{z} \times \nabla \phi}{B_0} + \frac{\hat{z} \times \nabla p_i}{neB_0} + \frac{m_i \underline{g} \times \underline{B}_0}{eB_0^2} \quad (32)$$

(Note that mks units are being used here to coincide with the units used by Ott (1978)). Quasi-neutrality, i.e.,  $\nabla \cdot \underline{J} = 0$ , with the assumption of two dimensionality implies that  $\underline{J}$  can be specified in terms of a single scalar potential function,  $\psi$ , such that

$$\underline{J} = -\hat{z} \times \nabla \psi \quad (33)$$

Using eqn. (2) for the electron velocity, and a next order ion velocity equation obtained by putting (32) into (3), the following ion velocity equation is obtained

$$nm_i \left[ \frac{\partial \underline{v}_i^{(0)}}{\partial t} + \underline{v}_i^{(0)} \cdot \nabla \underline{v}_i^{(0)} \right] = -\nabla \tilde{p} - nm_i \underline{v}_{in-i}^{(0)} + nm_i \underline{g} \quad (34)$$

with  $\tilde{p} = p_e + p_i + B \psi$ . Using (32) and the assumption of either isothermal or adiabatic ions one has

$$\nabla \cdot \underline{v}_i^{(0)} = 0 \quad (35)$$

Using  $\underline{v}_i^{(0)}$  in the ion continuity equation (1) with  $v_R = P = 0$ , eqn. (1) becomes

$$\frac{\partial n}{\partial t} + \underline{v}_i^{(0)} \cdot \nabla n = 0. \quad (36)$$



Equations (34) - (36) form a complete set of equations sufficient to determine the unknown quantities  $\underline{v}_i^{(0)}$ ,  $n$ , and  $\tilde{p}$ . In the limit of  $v_{in} \rightarrow 0$ , eqns. (34) - (36) are identical to those of an ideal incompressible fluid. At this point the philosophy taken is that there is much to be learned concerning bubbles in the ESF ionosphere from the extensive studies of bubbles in fluids. One then uses a stream function  $\xi$  such that

$$\underline{v}_i^{(0)} \equiv \hat{z} \times \nabla \xi \quad (37)$$

After a series of manipulations one can obtain an equation for the stream function. In the collisional and collisionless case one finds that for certain values of the bubble rise velocity  $\xi$  will satisfy the equation, for a cylindrical shape (see Ott, 1978 for more details).

The results of this study by Ott (1978) predicts the bubbles to be cylindrical (circular cap at top) in two dimensions. The bubble rise velocity in the collision-dominated regime, for a 100% depletion, is given by

$$v_B = \frac{g}{v_{in}} \quad (38)$$

which is altitude dependent. This result is the same as that predicted by (30b) for  $\delta n/n_0 = 1$ , i.e., a 100% depletion. In the inertia (collisionless) dominated regime, for a 100% depletion, the bubble rise velocity is given by

$$v_B \approx \frac{1}{2} (Rg)^{1/2} \quad (39)$$

where R is the radius of curvature at the top of the bubble so that (39) is size dependent.

e. Analytical Work on the Very Small Scale ( $\leq 10m$ )

Irregularities. The work in this section represents essentially multilinear calculations using kinetic theory resulting in plasma kinetic drift wave modes. It is multilinear because it depends on say a two-step process whereby linear theory is performed on the non-linear state. The driving density gradients, in these calculations, are thought to arise through a primary instability, driven by the zero order background ionospheric equatorial F region electron density gradient, achieving a large amplitude state. The zero order ionospheric electron density gradient is of larger scale size than the primary instability electron density gradient scale size, which would be of the order of the instability wavelength. Because the calculations are kinetic, they employ particle distribution functions. Kinetic drift waves have been investigated for laboratory plasma fusion conditions for over twenty years, so a well developed formalism exists.

Before proceeding with specific calculations let us present some general concepts regarding kinetic drift waves which will be useful for both types of calculations presented in this section. The basic geometry is such that

$$\underline{B} = B_0 \hat{z} \quad (40a)$$

$$n_0 = n_0(x) \quad (40b)$$

(similar to the zero order background equatorial ionospheric geometry) where here  $n_o(x)$  arises due to the primary instability. The orthogonal coordinate system is completed by the ion and electron diamagnetic velocities being along the  $y$  axis,

$$\underline{V}_d = (V_{di} - V_{de}) \hat{y} \quad (41a)$$

$$V_{di} = (V_i^2 / 2\Omega_i) (\partial \ln n_o / \partial x) \quad (41b)$$

$$V_{de} = - (V_e^2 / 2\Omega_e) (\partial \ln n_o / \partial x) \quad (41c)$$

where  $V_{i,e} = (2T_{i,e}/m_{i,e})^{1/2}$  and the larmor radius is defined by  $r_\alpha = V_\alpha / \Omega_\alpha$ . The linear analysis is then performed with this in mind.

First the low frequency ( $\omega \ll \Omega_i$ ) collisionless drift wave calculations of Costa and Kelley (1978a,b), for  $kr_i \sim 1$  ( $r_i$  is the ion gyroradius), will be presented. These calculations were meant to provide a basis within which to try and account for the 3 meter radar backscatter observations of Jicamarca (see Woodman and La Hoz, 1976). A linear kinetic dispersion relation is derived using perturbations of the form  $\exp i [\underline{k} \cdot \underline{x} - \omega t]$ , where  $\underline{k} = k_\parallel \hat{z} + k_\perp \hat{y}$  and  $\omega = \omega_r + i\gamma$ . Furthermore, assumptions are made such that  $\omega \ll \Omega_i$ ,  $V_i \ll |\omega/k_\parallel| \ll V_e$  and  $\omega > \nu_{in}, \nu_{ie}, \nu_{ii}$ , which are the ion-neutral, ion-electron, and ion-ion collision frequencies, respectively. This analysis depends on having a  $k_\parallel$ , a component of the wavevector along the ambient geomagnetic field. Figure 7 shows some growth rate results from these calculations. The growth rate is in units of the ion thermal velocity



divided by the electron density gradient scale length ( $L = n_o (dn_o/dx)^{-1}$ ). Figure 8 depicts some measured inverse electron density gradient scale lengths during ESF and maximum growth rates as a function of these inverse scale lengths. The basic results of these calculations show maximum growth for  $k_{\perp} r_i \approx 1.5$ , i.e.,  $\lambda_{\perp} \approx 20m$  (for typical ESF parameters) with growth rates  $\geq 1 \text{ sec}^{-1}$ . For more detailed analysis concerning these low frequency drift waves applied to ESF see Costa and Kelley (1978a,b).

Now we present the high frequency drift wave analysis of Huba et al. (1978). In this reference radar backscatter observed irregularities with wavelengths of 1 meter and 36 cm at Kwajalein during ESF conditions are shown. In an effort to explain these very short wavelength irregularities high frequency ( $\omega \geq \Omega_i$ ) drift waves were analyzed for ESF conditions. These waves are the so-called drift cyclotron (DC) or lower hybrid drift (LHD) instabilities with maximum growth rates for  $k_{\perp} r_e \sim 1$  (and  $k_{\parallel} = 0$ ). No  $k_{\parallel}$  is required for these instabilities. The parameter determining which instability operates in a collisional plasma is

$$C_f = (v_{ii}/\Omega_i) (k_{\perp} r_i)^2 \quad (42)$$

Utilizing the linear dispersion relation for high frequency drift waves for  $C_f \ll 1$ , instability results for

$$L/r_i < (1/2\ell) (m_i/m_e)^{1/2} \quad (43)$$

where  $\ell$  is the harmonic number ( $\omega_r \approx \ell \Omega_i$ ). For the  $O^+$  ESF ionospheric plasma this requires the electron density gradient scale length

$L < 340\text{m}$ , which is satisfied (see Fig. 8). Growth rates for these instabilities are given by

$$\gamma \approx (m_e/m_i)^{1/2} \ell \Omega_i \quad (44)$$

However, the condition  $C_f \ll 1$  implies that  $(k_{\perp} r_i)^2 n \ll 2 \times 10^7$  and for  $k_{\perp} r_e \sim 1$  this means  $n \ll 10^3 \text{ cm}^{-3}$ , which is quite restrictive. Longer wavelengths, i.e., smaller values of  $k_{\perp} r_e$  would raise the density restriction somewhat, but still be restrictive for ESF conditions.

For  $C_f \geq 1$ , the lower hybrid drift instability is operative and there is no threshold condition on  $L$ . Basically, the collisions which increase  $C_f$  destroy the ion gyroresonances needed for the DC instability to operate. The real and imaginary part of the frequency for the LHD instability are given by

$$\frac{\omega_r}{\Omega_i} \sim \frac{r_i}{L} \left( \frac{m_i}{m_e} \right)^{1/2}, \quad \frac{\gamma}{\Omega_i} \sim \left( \frac{r_i}{L} \right)^2 \left( \frac{m_i}{m_e} \right)^{1/2} \quad (45)$$

In the collisionless limit it should be noted that the DC instability transforms into the LHD instability for high enough ion diamagnetic velocities such that

$$L/r_i \lesssim (m_i/m_e)^{1/2}$$

which for ESF conditions implies that  $L \lesssim 30\text{m}$ . Figures 9 and 10 show some typical results from the analysis of the high frequency drift wave linear dispersion relation.

The results of this analysis predicts that the lower hybrid drift instability is dominant for most typical ESF ionospheric parameters. Also, maximum growth of the instability occurs for  $k_{\perp} r_e \sim 1$  ( $\lambda_{\perp} \sim 21$  cm), although good growth rates can occur for  $\lambda_{\perp} \sim 1$  m. Finally, from this instability, large growth rates ( $\gamma \ll \Omega_i$ ) resulting in growth times,  $\tau = \gamma^{-1}$ , less than a second can occur. For more details of this work see Huba et al. (1978).

### III. Summary

Although much progress has been made in the theoretical efforts directed toward the equatorial Spread F ionosphere, especially in the past three years, more has to be done. Also the burden cannot be placed on the theoretician alone. Correlative measurements have to be made prior to and during ESF conditions. It is not sufficient to make a single measurement with one instrument and then expect a complete theoretical description of ESF. One needs to know the state of the ionosphere with respect to driving parameters such as background electron density profiles (bottomside electron density gradient scale lengths and height of the F peak), d.c. electric fields, neutral winds, and ionic mass composition in order to build a predictive model. In addition, in order to compare results from the predictive model, measurements of the in situ fluctuating component of the electric field and plasma density have to be made, as well as radar backscatter measurements of the very small scale irregularities ( $\leq 10$ m) and ground measurements of satellite signal propagation amplitude and phase scintillations.



Some of the remaining theoretical problems could be listed as follows. (1) What are the effects of changing initial conditions, including ion inertia, and including neutral winds in the numerical simulations? (2) An analytical description of many bubbles rising. (3) How do bubbles decay and what role does diffusion, etc. play? (4) How does a turbulent development occur? (5) What are the effects of other regions of the ionosphere (e.g., E region) on ESF? (6) A more complete study of collisional effects on drift waves is needed. What determines when the small scale irregularities should occur (a more quantitative description)? (7) A determination of the nonlinear saturation of small scale irregularities (instabilities) is needed. (8) What are the effects of  $k_{\parallel}$ ? Indeed points (6) and (7) are tied to the more general question of what is the relation between the very small scale ( $\leq 10m$ ) irregularities (which radar backscatter observes) and the longer wavelength fluid type (e.g., R-T) irregularities (which are primarily responsible for Spread F seen on ionograms)? I am sure that some more questions and points could be raised. However, the above list should keep the theoreticians busy for a reasonable time.

### Acknowledgements

The author is grateful to the dedication of his colleagues in their efforts to provide an understanding of equatorial Spread F. In particular, I am deeply indebted to Pradeep Chaturvedi, Joe Huba, Mike Keskinen, Ed McDonald, Tony Scannapieco, and Steve Zalesak whose interest and work in ionospheric irregularities is the sine qua non without which this article could not have been written. In addition, this work was supported by the Defense Nuclear Agency and the Office of Naval Research.

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Table I - Bubble rise velocity (in units of  $g/v_{in}$ ),  $V'_B$ , as a function of fractional depleted density,  $\delta n/n_o$ , for various bubble shapes

$V'_B \backslash \frac{\delta n}{n_o}$	.25	.5	.75	.9	1
Linear	.25	.5	.75	.9	1
Sheet	.33	1	3	9	$\infty$
Cylindrical	.14	.33	.6	.82	1
Elliptical (5:1)	.26	.71	1.67	3	5
Elliptical (10:1)	.29	.83	2.14	4.5	10

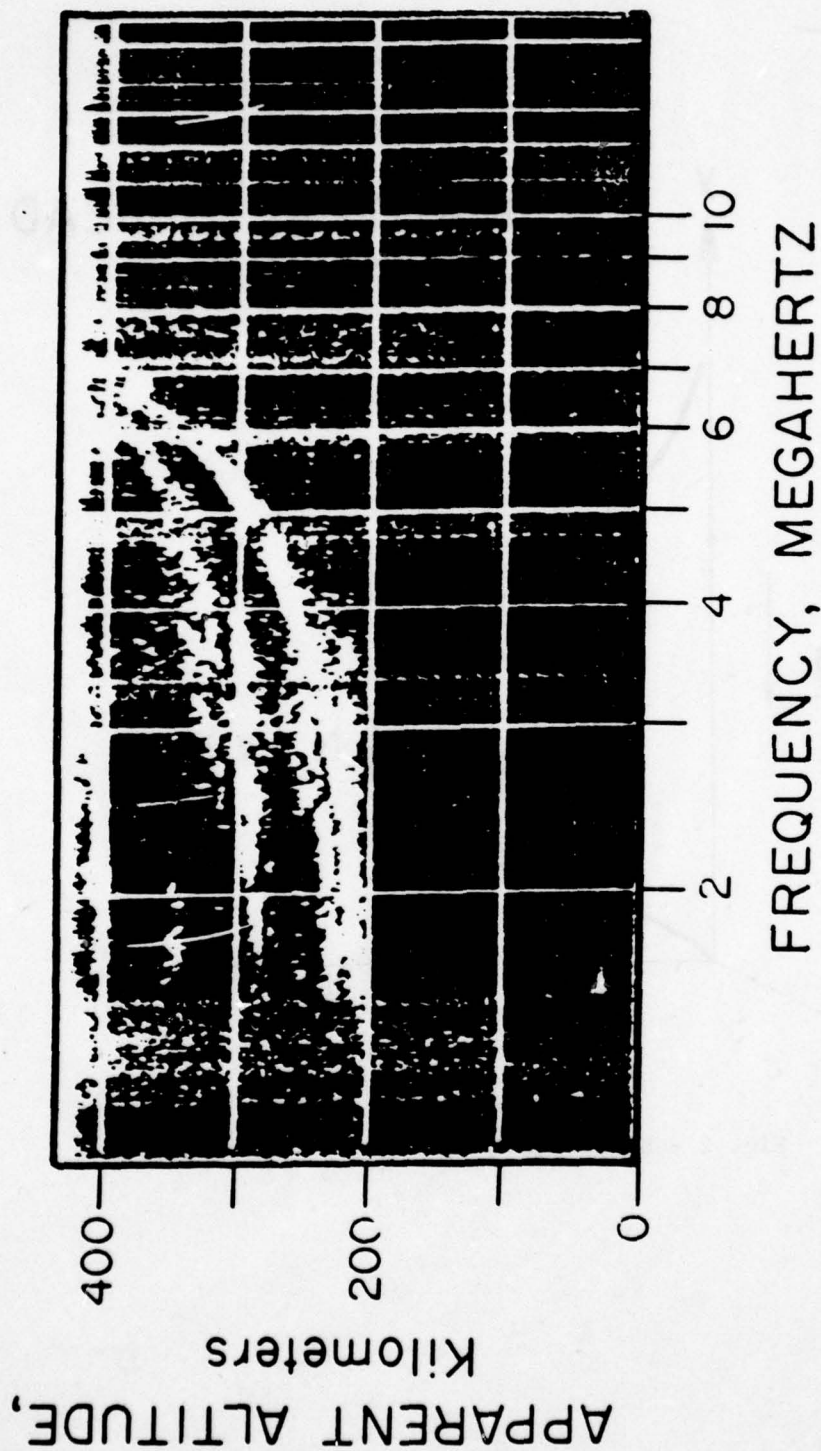


Fig. 1 - Ionogram recorded at Natal, Brazil (equatorial station), on November 18, 1973 at 2122-2123 UT. This figure was taken from Hudson and Kennel (1975)

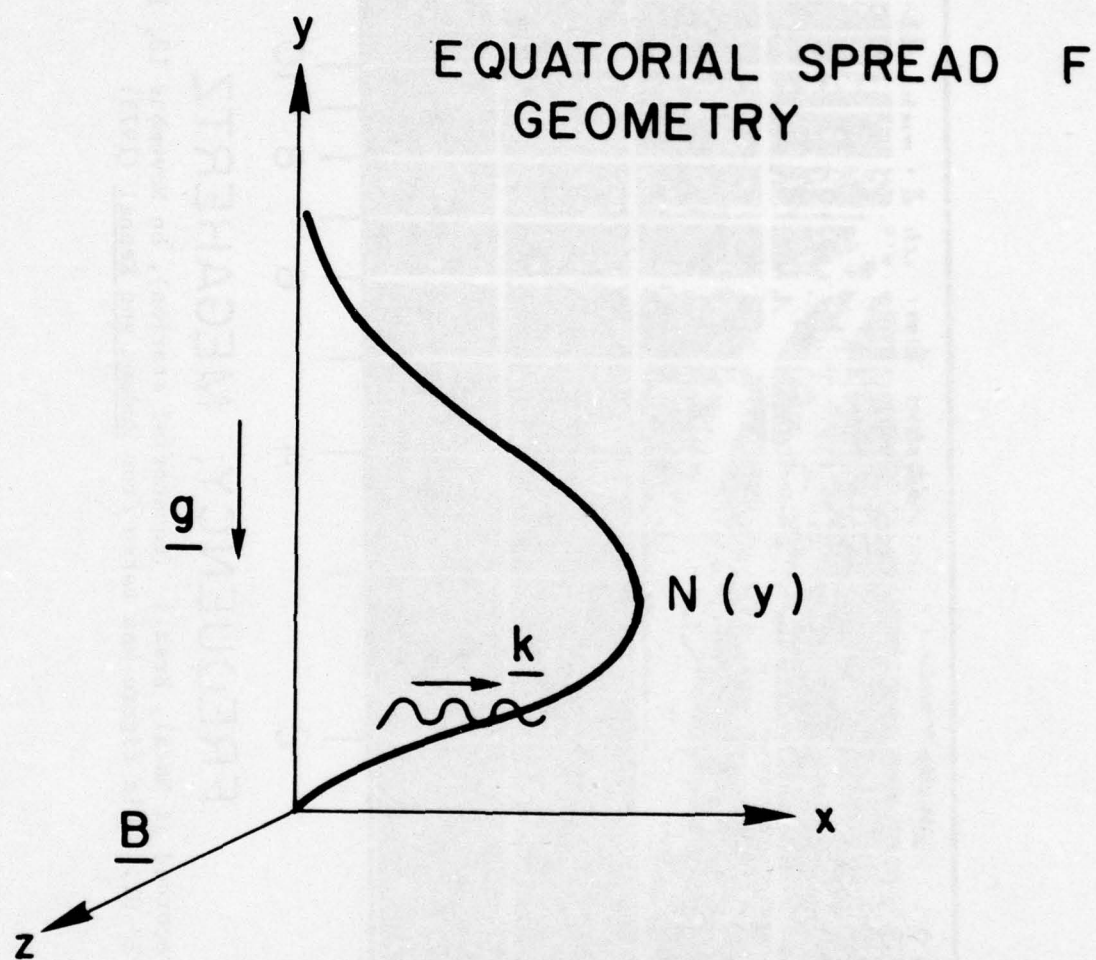


Fig. 2 - Basic equatorial Spread F geometry



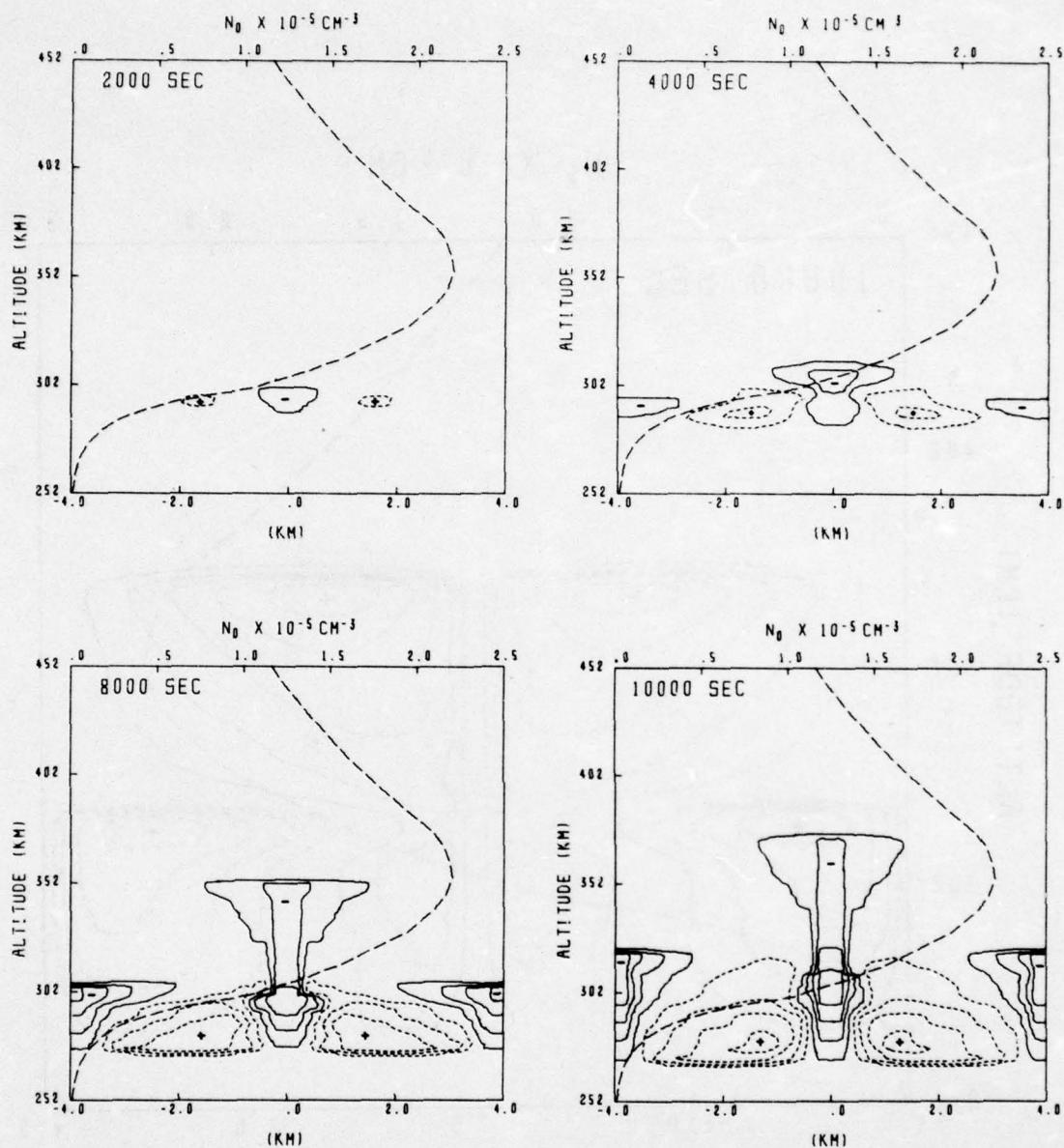


Fig. 3 - Contour plots of constant  $n_1/n_0$  for the simulation with an F peak at 354km at  $t = 2000, 4000, 8000$ , and  $10,000$  sec. The small dashed contours with a plus sign inside and the solid contours with a minus sign inside indicate enhancement and depletions over the ambient electron number density. The large dashed curve depicts the ambient electron number density (values on upper horizontal axis),  $n_0$ , as a function of altitude. The vertical y axis represents altitude, the lower horizontal x axis is east-west range, and the ambient magnetic field is along the z axis, out of the figure. Taken from Ossakow et al. (1978)

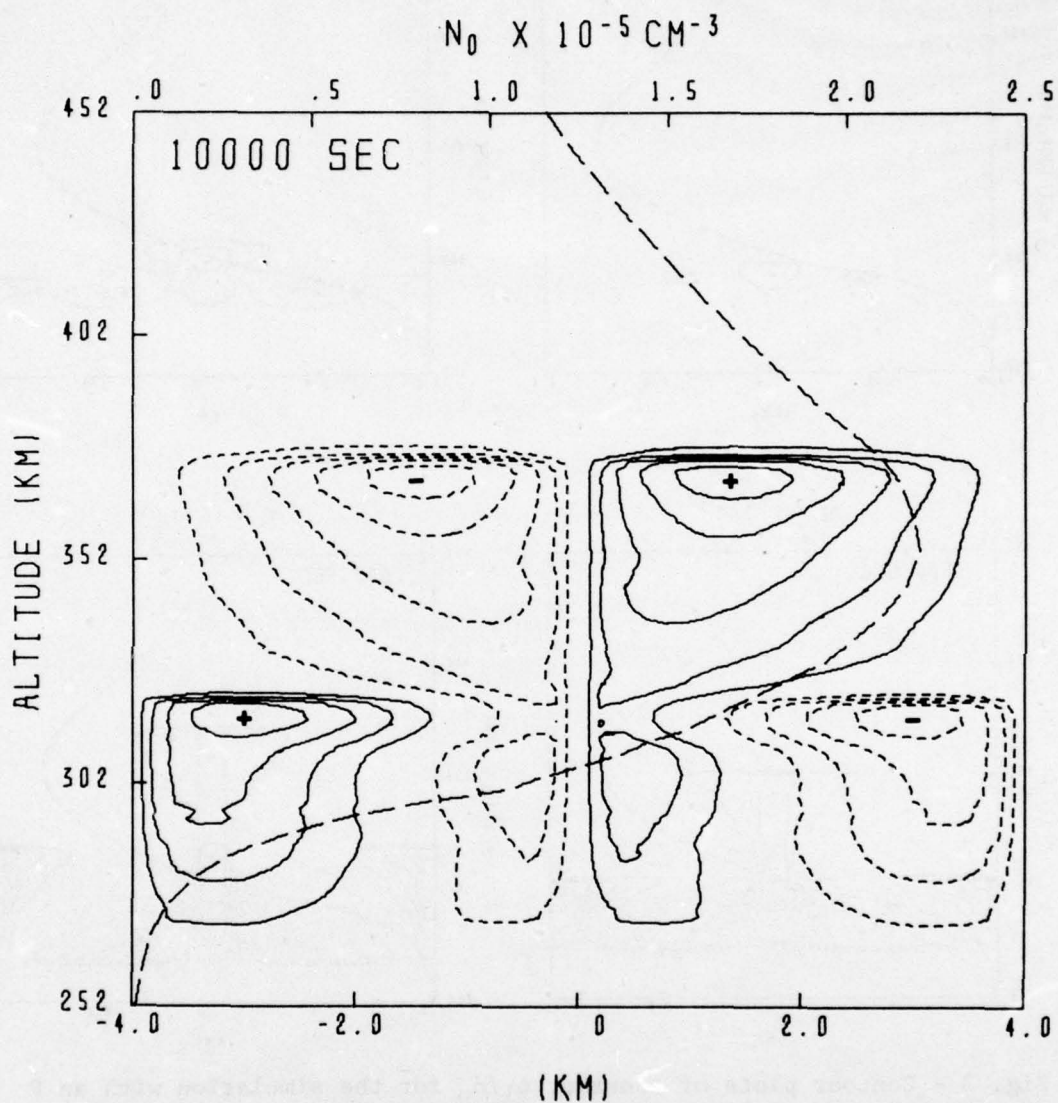


Fig. 4 - Contours of constant induced (polarization) potential,  $\phi_1$ , over the mesh (see Fig. 3) at  $t = 10,000$  sec. Plus and minus denote positive and negative values, with values decreasing in magnitude as one goes from the innermost to the outermost contours. The large dashed curve is  $n_0$ . Taken from Ossakow et al. (1978)

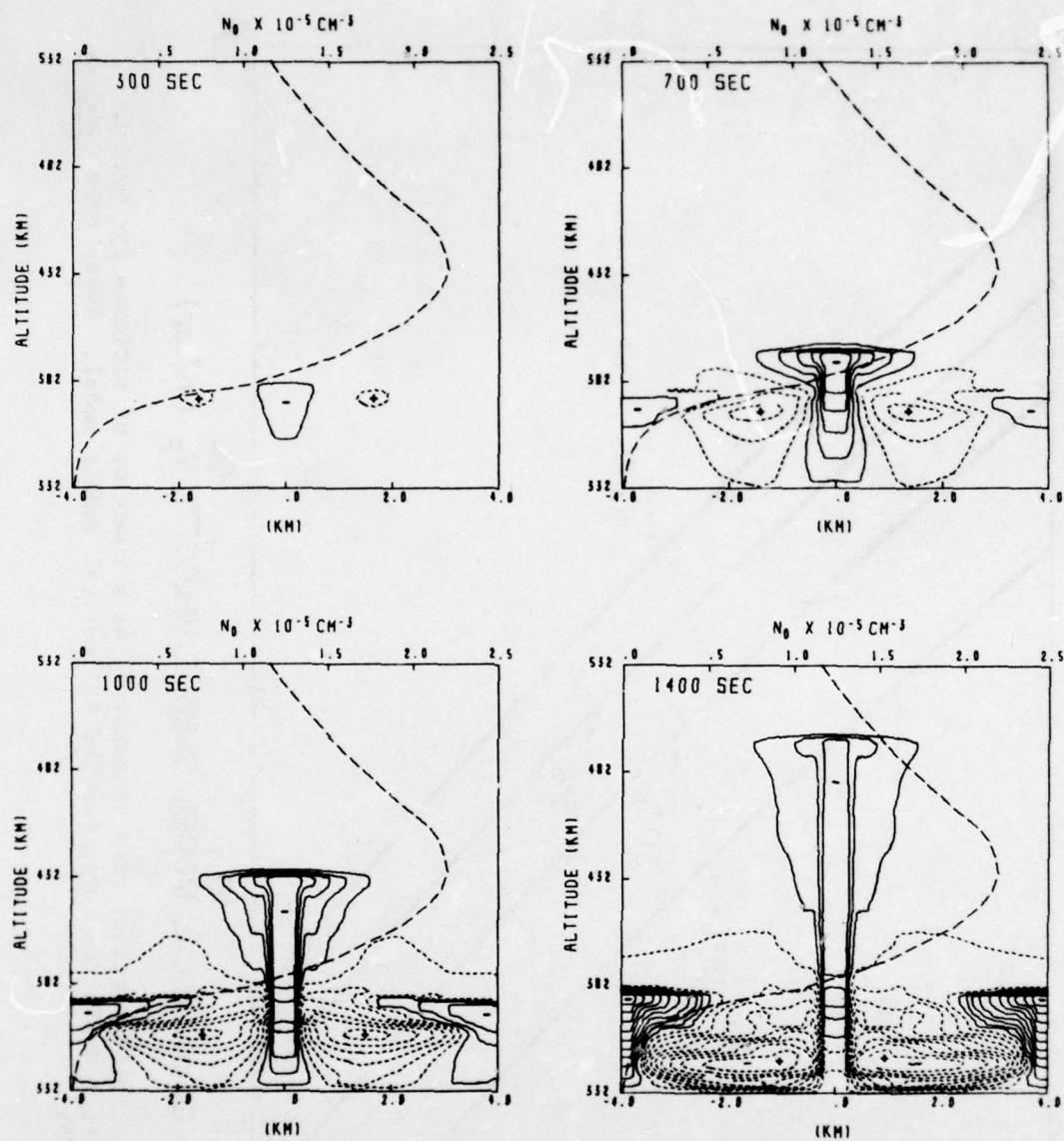


Fig. 5 - Contour plots of constant  $n_1/n_0$  for the F peak at 434km simulation at  $t = 300, 700, 1000$ , and  $1400$  sec. All other nomenclature is same as Fig. 3. Taken from Ossakow et al. (1978)



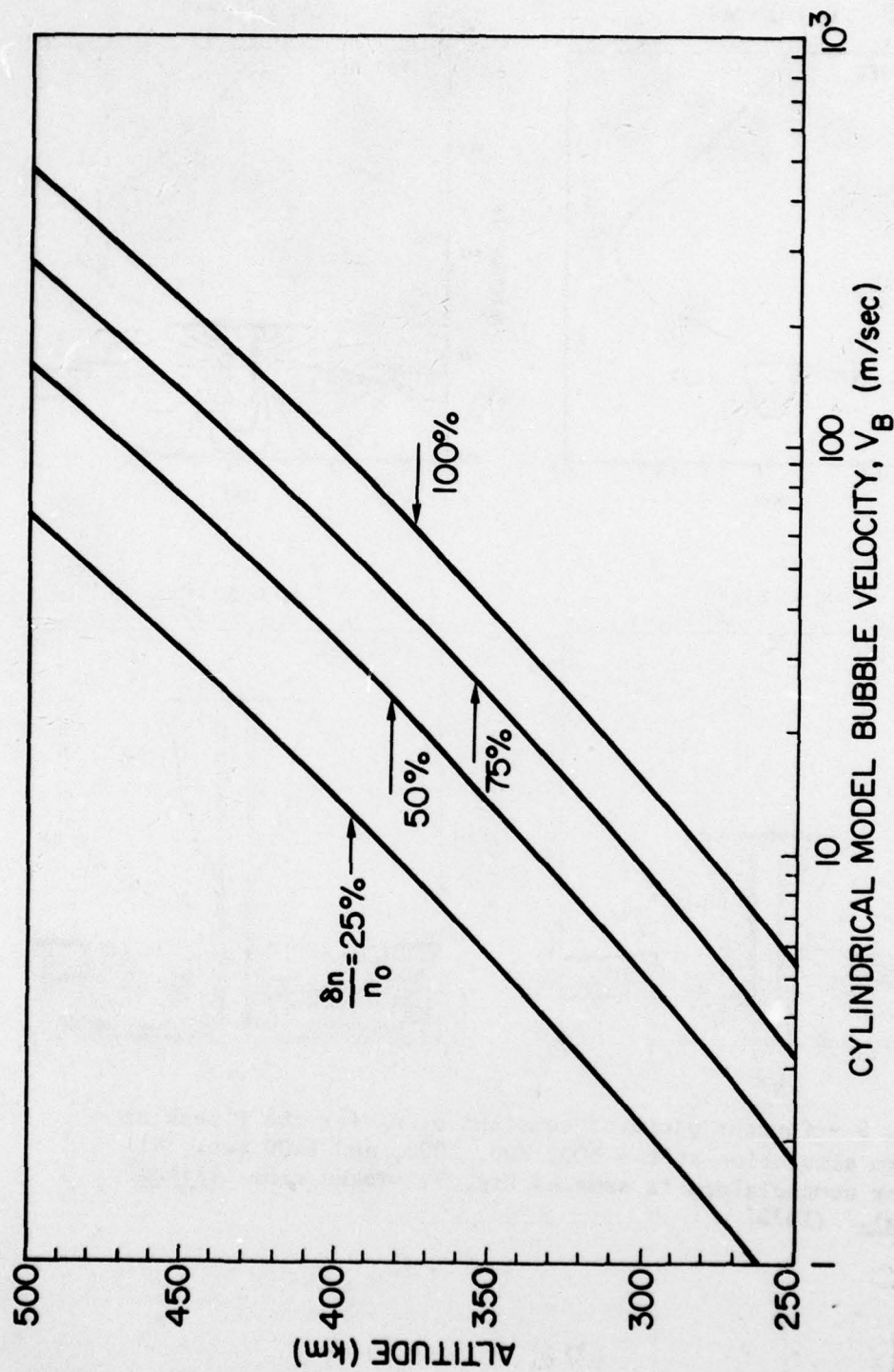


Fig. 6 - Depletion (bubble) vertical rise velocity  $V_B$  as a function of altitude for various values of the percentage depletion  $\delta n/n_0$  for the cylindrical bubble model. Taken from Ossakow and Chaturvedi (1978)

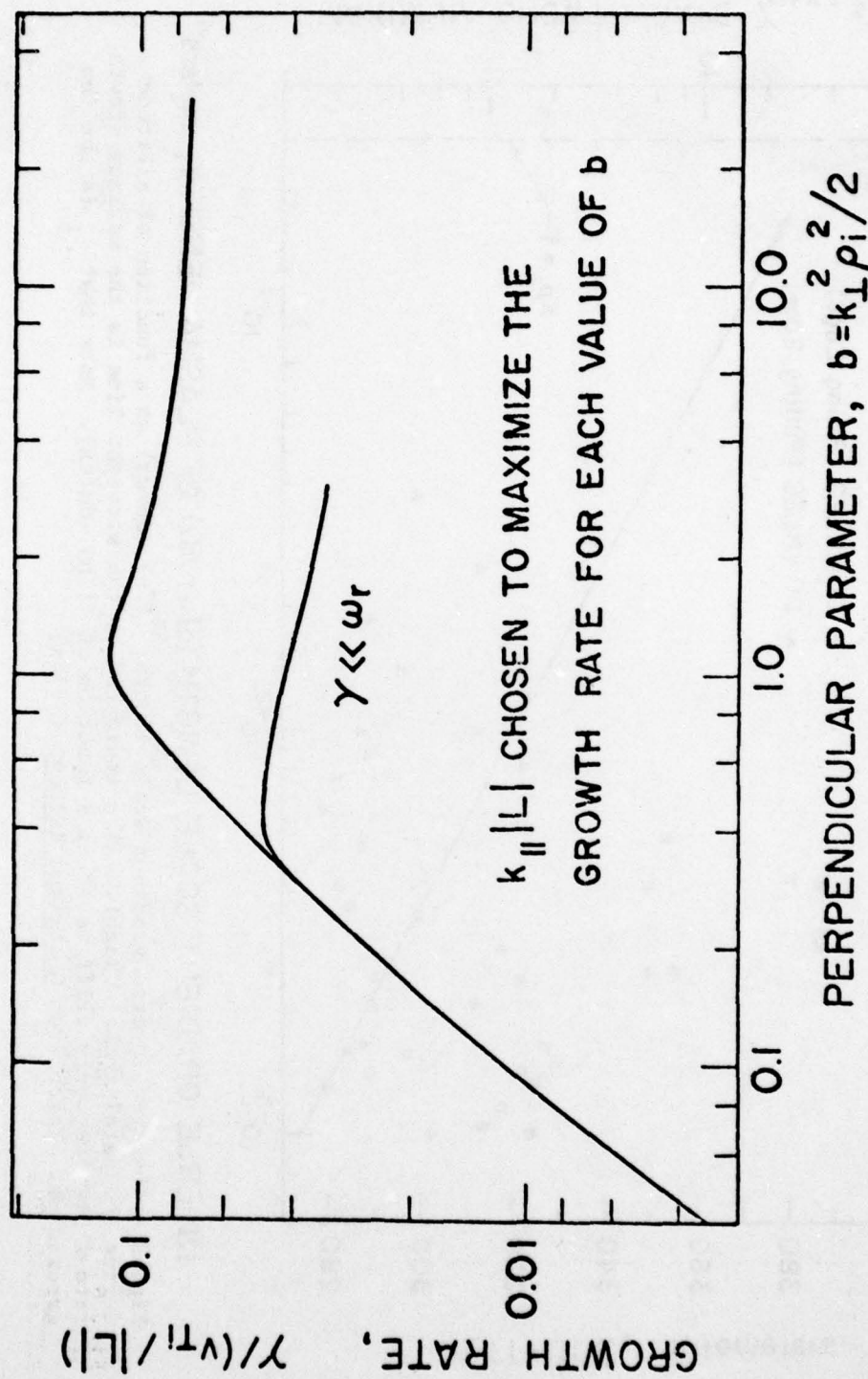
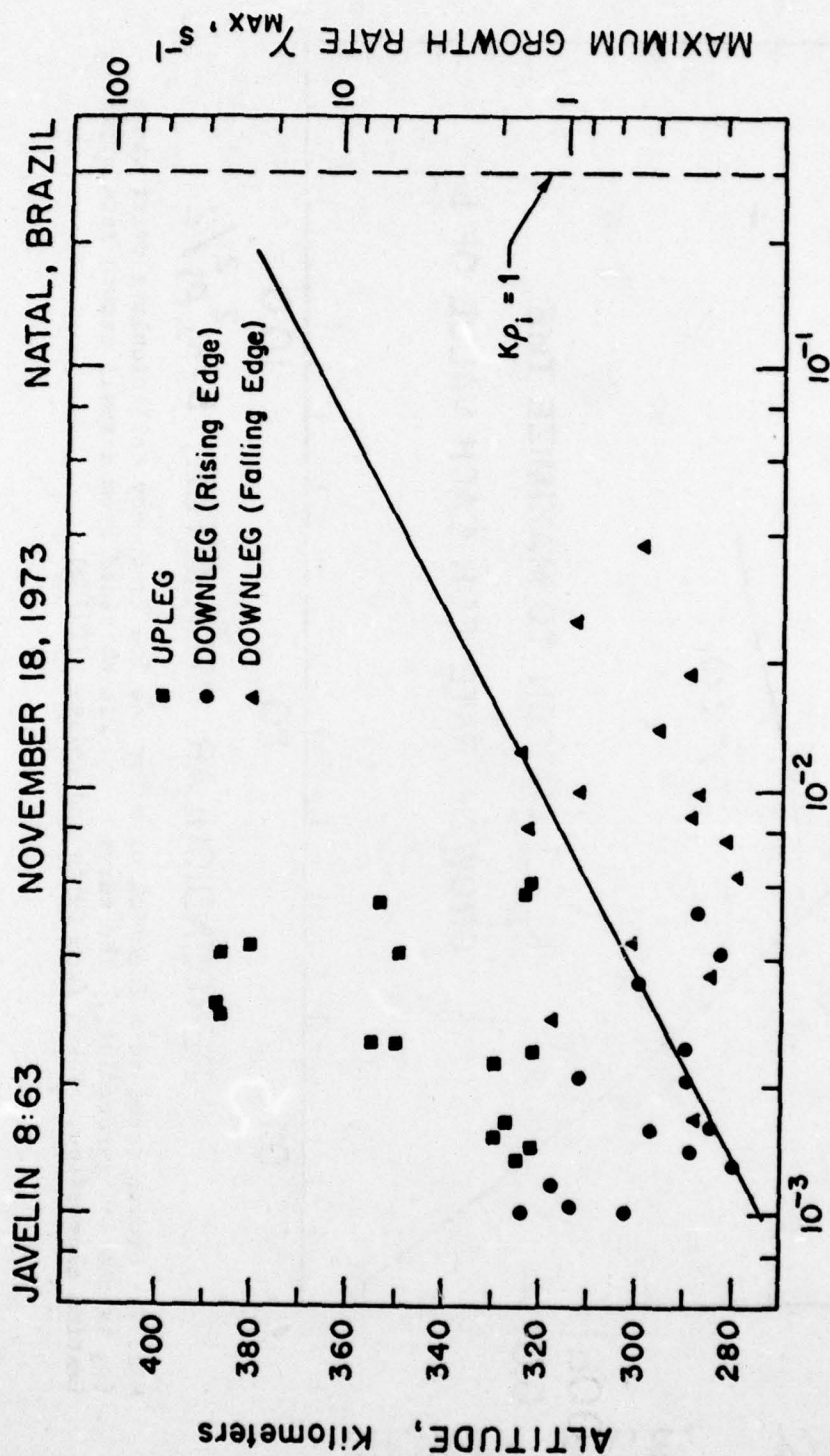


Fig. 7 - Growth rates as a function of  $b$  for the low frequency collisionless drift waves ( $\rho_i$  is the ion gyroradius). The curve  $\gamma \ll \omega_r$  is obtained from a small growth rate approximation expression. Taken from Costa and Kelley (1978b)



INVERSE GRADIENT SCALE LENGTH  $(1/n) (dn/dx)$  OF PLASMA DENSITY, meters<sup>-1</sup>

Fig. 8 - Observed inverse gradient scale lengths  $(1/n) (dn/dx)$  as a function of altitude during the Natal rocket flight. Also indicated by the straight line is the maximum growth rate of low frequency drift waves as a function of  $(1/n) (dn/dx)$ . Note that  $\rho_i$  is the ion gyroradius. Taken from Costa and Kelley (1978a)



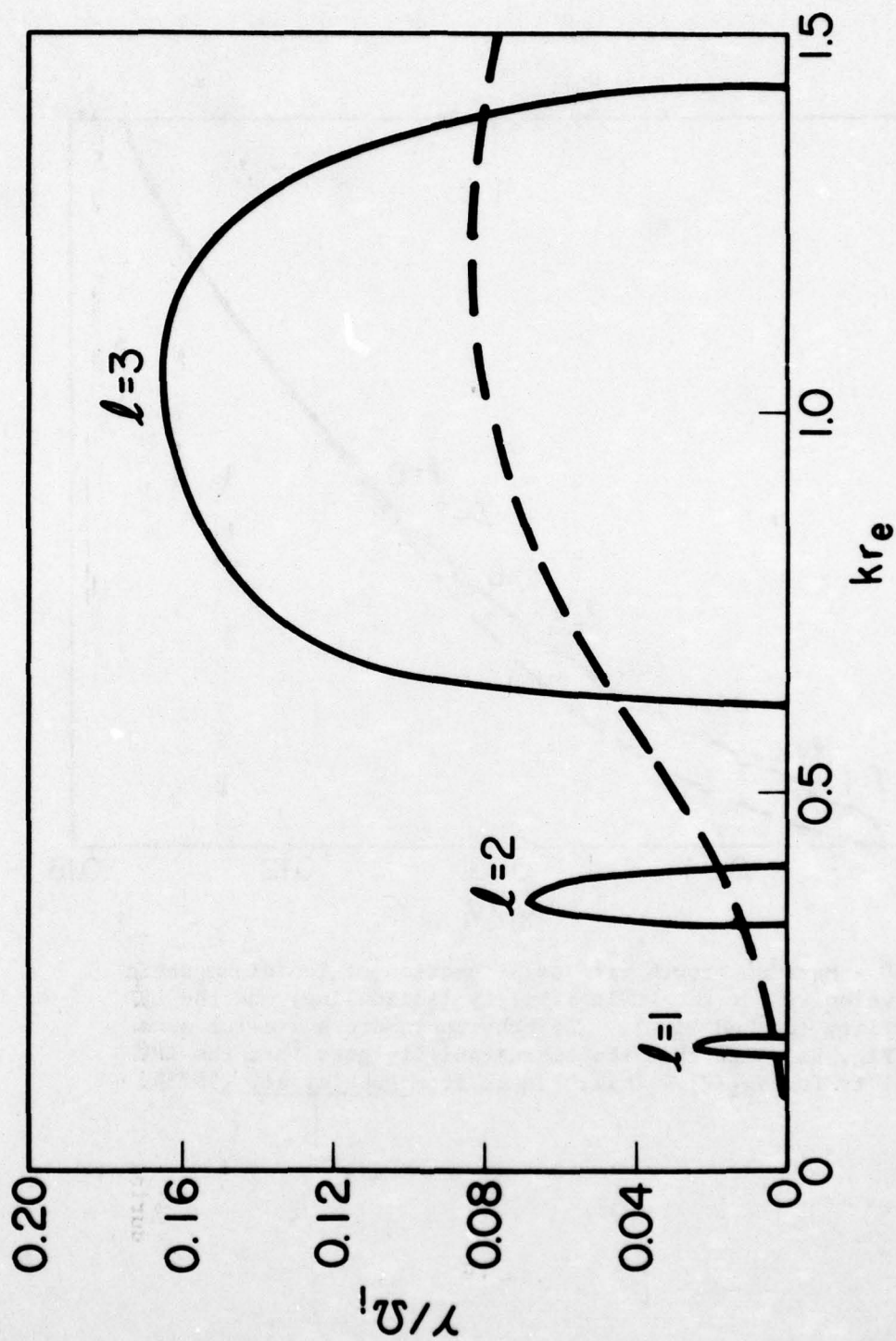


Fig. 9 - Growth rate as a function of perpendicular wavenumber for the DC instability (solid line) and LHD instability (dashed line) for an  $O^+$  plasma with  $T_e = T_i$ ,  $\omega_{pe}/\Omega_e = 10$  (where  $\omega_{pe}$  is the electron plasma frequency), and  $V_{di}/V_i = 0.037$ . Growth also occurs for  $k r_e > 1.5$ , but has not been plotted. Taken from Huba et al. (1978)

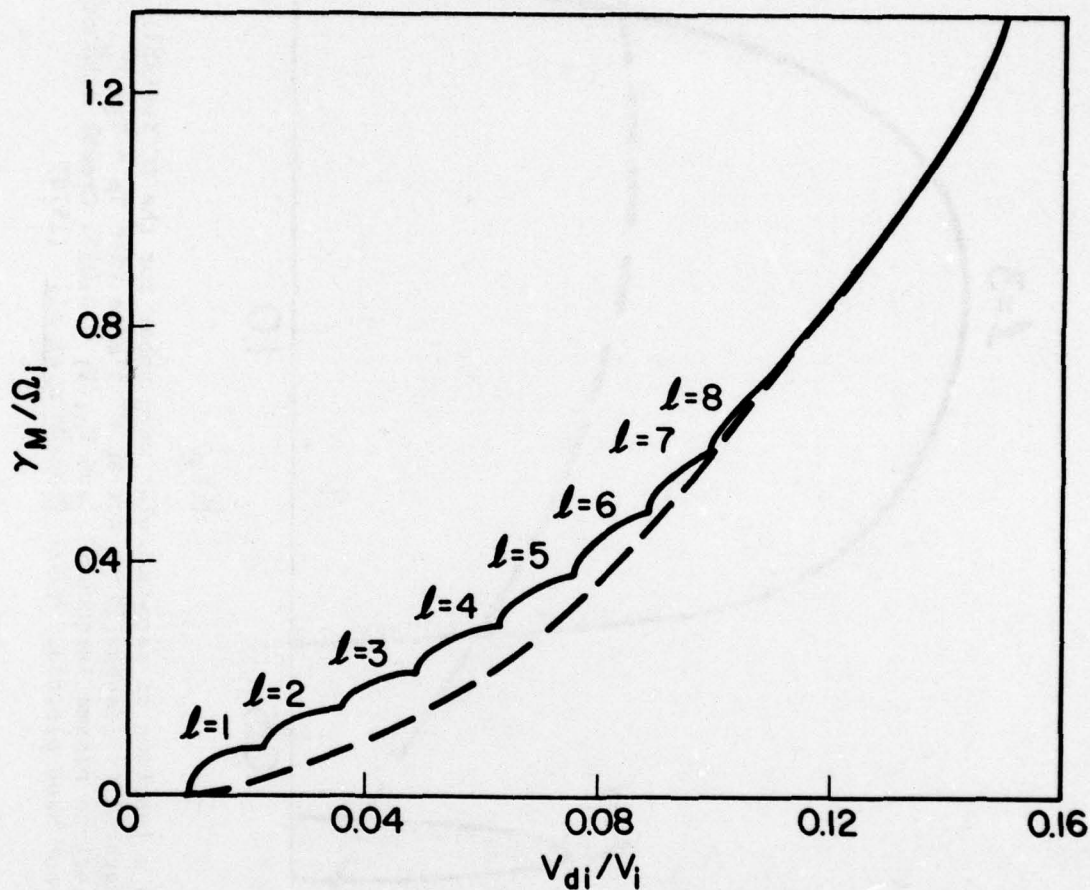


Fig. 10 - Maximum growth rate as a function of ion diamagnetic drift velocity for the DC instability (solid line) and the LHD instability (dashed line). The other parameters are the same as in Fig. 9. Note that the DC instability goes into the LHD instability for  $V_{di}/V_i \gtrsim 0.11$ . Taken from Huba et al. (1978)

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