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A Laboratory Investigation of the High-Frequency Farley-Buneman Instability

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A laboratory investigation of the high-frequency Farley-Buneman instability is described. The instability was studied theoretically by Lee et al. (1971) and is predicted to occur in the low E region of the ionosphere when the E/B drift velocity of the electrons relative to the ions is several times C_s, the ion acoustic speed. In our experiments an increase of the electric field well above the Lee et al. "threshold" merely enhances the general power level of the fluctuations but does not affect appreciably their spectral shape. The observed frequency spectra fall off very rapidly with increasing frequency, with a spectral shape of the type $P(f) \propto f^{-3.5}$. This result agrees well with recent ionospheric observations and suggests that a recently proposed mechanism of anomalous wave electron heating in the lower E region should not be linked to the presence of short-wavelength (10–30 cm) irregularities.

1. INTRODUCTION

In the ~100- to ~115-km height range of the ionosphere, $\nu_i > \omega_{ci}$, whereas $\nu_e \ll \omega_{ce}$, where $\nu_{i(e)}$ is the ion (electron) collision frequency with the neutral gas molecules of the atmosphere and $\omega_{ci}(\omega_{ce})$ is the ion (electron) gyrofrequency. If an electric field, **E**, is present, normal to the geomagnetic field, **B**, the electrons drift with a velocity $\mathbf{v}_D \simeq \mathbf{E} \times \mathbf{B}/B^2$, while the ions are essentially "held back" by collisions with the neutral gas background. When the velocity of the electrons relative to the ions exceeds the local ion acoustic speed, C_s, *Farley* [1963] and *Buneman* [1963] have shown that longitudinal plasma waves grow, with $K_{\parallel}^2 \ll K_{\perp}^2$ and a phase velocity perpendicular to **B** essentially equal to the electron drift velocity. Here $K_{\parallel}$ and $K_{\perp}$ are the components of the propagation vector, **K**, parallel and perpendicular to **B**, respectively.

The Farley-Buneman instability is responsible for the so-called "type 1" irregularities observed by radar backscatter in the equatorial and in the auroral electrojet [e.g., *Bowles et al.*, 1963; *Balsley*, 1969; *Chesnut*, 1969]. Rocket measurements of the instability in the polar cap E region have been reported, for example, by *Olesen et al.* [1976] and *Bahnsen et al.* [1978]. Laboratory experiments on the instability were performed by *Saito et al.* [1964], *D'Angelo et al.* [1974], *John and Saxena* [1975], *Alport et al.* [1981], and *Mikkelsen et al.* [1981].

In 1971, Lee, Kennel, and Kindel [*Lee et al.*, 1971] extended the calculations of *Buneman* [1963] and *Farley* [1963] to modes of higher frequencies and shorter wavelengths, by retaining the term $K^2\lambda_{De}^2$ (λ_{De} is the electron Debye length) in the dispersion relation and considering electron drift velocities, relative to the ions, several times larger than the ion thermal speed. They found that, the dispersion relation being now density dependent, there is an electron density threshold for excitation of higher-frequency modes. For a given density, the fastest growing mode shifts to higher frequency as v_D increases. For v_D equal to three times the ion thermal speed, the fastest mode has a real frequency $\omega_r \simeq 0.7(\omega_{ci}\omega_{ce})^{1/2}$. This means, typically, wavelengths in the lower E region ionosphere of 10–20 cm, or less, rather than the several meters of the ordinary low-frequency Farley-Buneman instability.

To the best of our knowledge, only one rather detailed study of this high-frequency instability, by means of rocket-borne instrumentation, has been performed in the high-latitude E region [*Kelley and Mozer*, 1973]. We shall comment on their results in the last section of this paper.

Related to the high-frequency Farley-Buneman instability is the question of anomalous electron heating in the polar E region, as reported, for example, by *Schlegel and St. Maurice* [1981]. With the incoherent scatter facility at Chatanika, Alaska, they measured electron temperatures of up to 1200°K near the 110-km altitude in the polar E region, in the presence of strong electric fields. They argued that no classical heating process can account for these high temperatures, and they suggested that the temperature enhancements are caused by the unstable plasma waves produced in the disturbed polar E region by electric fields $E \gtrsim$ 50–60 mV/m, through the instability discussed by *Lee et al.* [1971]. An extensive theoretical analysis of the anomalous electron heating is given in a paper by *St. Maurice et al.* [1981].

In this paper we describe an experiment performed to study in the laboratory the high-frequency Farley-Buneman instability. The experimental setup is similar to that employed by *Alport et al.* [1981] in their study of the low-frequency Farley-Buneman instability, except that in the present experiment radial electric fields of up to ~5 V/cm (rather than the ~0.5 V/cm of *Alport et al.* [1981]) are utilized, following *Lee et al.* [1971], to shift the fastest growing mode to frequencies of the order of the lower hybrid frequency.

Section 2 of the paper describes the experimental arrangement, while section 3 contains the experimental results. In section 4 these results are discussed, also in relation to the important question of the anomalous electron heating in the lower E region and to other ionospheric observations.

2. EXPERIMENTAL SETUP

We used, for the present experiment, the same plasma device employed by *Alport et al.* [1981] in their study of the low-frequency Farley-Buneman instabilitiy and of the gradient drift instability. The geometry of this experiment is cylindrical and does not, strictly speaking, reproduce the Cartesian geometry associated with the ionospheric situation. In particular, in a proper theoretical treatment, the usual term associated with a Doppler shift in the ionosphere, $(\omega - \mathbf{k}\cdot\mathbf{V})\mathbf{v}$, should be replaced by a term taking the rotational effects into account. However, the source of free energy (perpendicular currents) is the same, and the laboratory plasma does indeed show the

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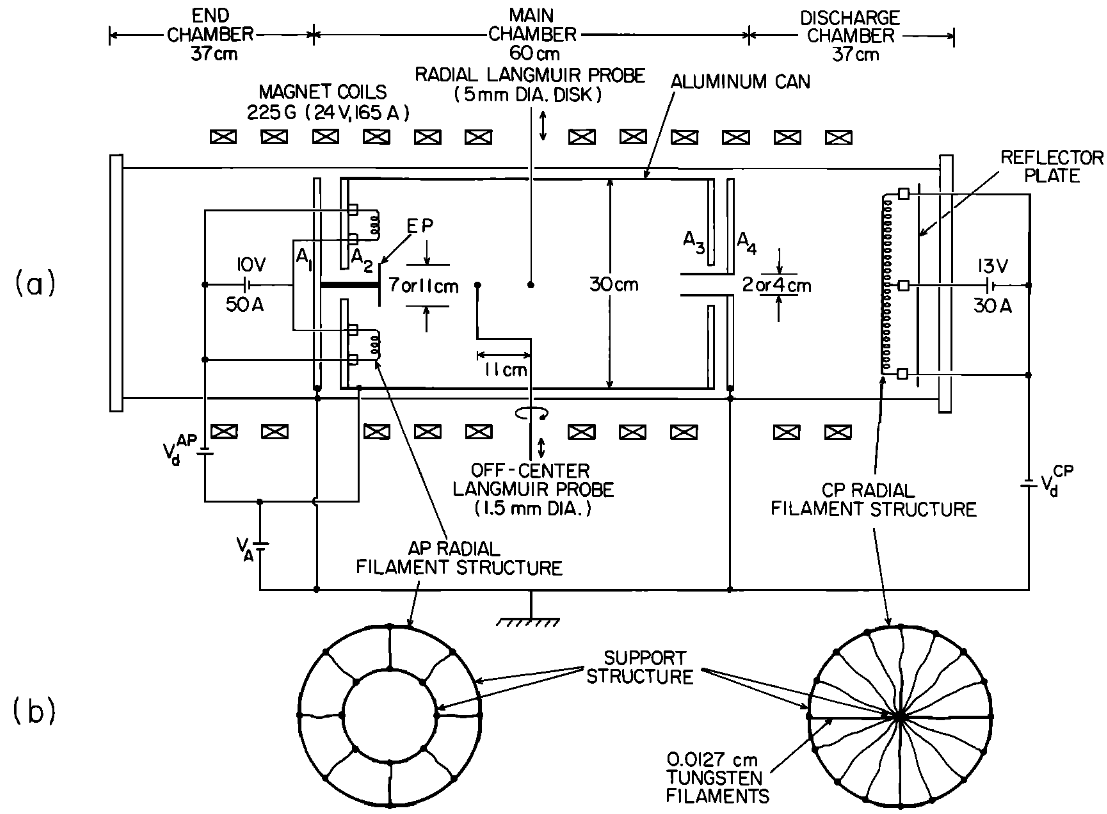

Fig. 1. Schematic diagram of the experimental setup. (*a*) Top view. (*b*) Central plasma and annular plasma filament structures.

production of irregularities of the Farley-Buneman and gradient drift type. As with the *Alport et al.* [1981] setup, the neutral gas (argon) pressure was still of the order of 10^{-3} torr, while the axial magnetic field was $\sim$225 G. This insures that the two conditions $\omega_{ci} < \nu_i$ and $\omega_{ce} \gg \nu_e$ (i.e., collisional ions and collisionless electrons) are satisfied, as required for the development of both the low-frequency and the high-frequency Farley-Buneman instability. However, with a radial E field of $\sim$0.5 V/cm, as in the *Alport et al.* [1981] experiment, the azimuthal drift speed of the electrons, relative to the nearly stationary ions, is only $\sim 2.2 \times 10^5$ cm/s, i.e., only slightly in excess of the ion acoustic speed, $C_s \sim 2 \times 10^5$ cm/s.

For a study of the high-frequency Farley-Buneman instability, radial electric fields several times larger are required. For this reason the device had to be partially modified, to allow us to apply radial electric fields of several volts per centimeter. It was also desirable that the radial field be varied while keeping the plasma density approximately constant.

A schematic diagram of the experimental setup is shown in Figure 1. A comparison of this figure with Figure 4 of *Alport et al.* [1981] will make clear which modifications to the original device were needed to obtain the larger radial electric field. We have added a cylindrical aluminum can, 30 cm in diameter, which is electrically connected to anode rings A_2 and A_3, and an additional set of filaments (annular plasma (AP) radial filament structure) mounted on anode ring A_2. The anode end plate (EP) and ring A_4 are still connected to the vacuum chamber, which is grounded. Plasma and primary electrons from the discharge chamber (right side) stream through the aperture in anode ring A_4, thus producing a plasma (with a diameter determined mainly by the aperture in A_4) which is terminated in the main chamber on the (grounded) end plate attached to A_1.

Typically, the main discharge, or center plasma (CP), is operated with a background argon pressure $p \simeq 10^{-3}$ torr, a discharge current $I_d^{CP} \simeq 1$–4 A, a discharge voltage $V_d^{CP} \simeq 50$ V, and a magnetic field $B = 225$ G in the center of the main chamber. The axial variation of the magnetic field is $\sim$15% over the entire length of the plasma column. The plasma density in the center plasma is, generally, 1–5 $\times 10^{10}$ cm^{-3}. The annular plasma is produced by a discharge between the AP filaments and anode rings A_2, A_3, and the 30-cm-diameter aluminum can. This discharge is operated, typically, with $I_d^{AP} \simeq 0.01$–0.03 A and $V_d^{AP} \simeq 50$ V. Thus only a very tenuous plasma is independently produced in the outer, annular discharge region. The potential of the annular plasma is controlled by means of the outer anode bias V_A. All power supplies for producing and biasing the annular plasma are independent of those for the core plasma.

The operation of our device in its present form is somewhat analogous to that of a standard double-plasma (DP) device. In a DP device [*Taylor et al.*, 1972], a "driver" plasma and a "target" plasma are produced in a common vacuum chamber and are separated by a negatively biased grid which largely prevents the electrons of one plasma from intermixing with those of the order. In our present setup, a role similar to that of the grid of a DP device is played by the axial magnetic field, which reduces the mobility of all plasma particles and of the primary ionizing electrons. However, a substantial ($\lesssim 1$ A) radial current flows generally between the annular plasma and the central plasma.

A radial (inward) electric field is produced when the AP anode structure (A_2, A_3, and can) is biased to some positive potential V_A, generally less than $\sim$20 V. With V_A applied, the space potential of the annular plasma rises to a value $\gtrsim V_A$. With A_2 and A_4 grounded, the central plasma potential in-

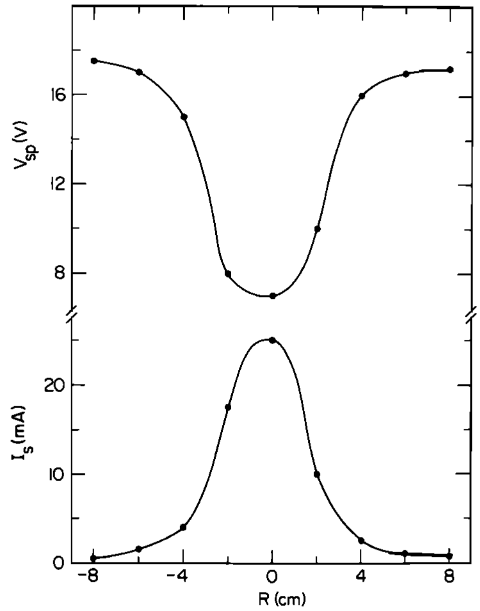

Fig. 2. Radial profiles of the space potential, V_{sp} (volts), and of the electron saturation current, I_s (milliamperes) ($I_d^{CP} = 1.75$ A, $I_d^{AP} = 0.025$ A, $V_d^{CP} = V_d^{AP} = 50$ V, $V_A = 15.2$ V, plasma density in the range 10^9–10^{10} cm^{-3}).

creases somewhat with increasing V_A, but only by a small fraction of V_A. Typical results for the radial profiles of Langmuir probe electron saturation current, I_s, and space potential, V_{sp}, across the central plasma and a large portion of the annular plasma, are shown in Figure 2 ($I_d^{CP} = 1.75$ A, $I_d^{AP} = 0.025$ A, $V_d^{AP} = V_d^{CP} = 50$ V, $V_A = 15$ V). In the region 0 cm < R < +4 cm, the potential profile provides an average radial electric field of ~2 V/cm.

At different values of V_A we obtain similar radial potential profiles, with a monotonic increase of the radial electric field with increasing V_A. This is illustrated in Figure 3, where the difference in space potential, ΔV_{sp}, between $R = 5$ cm and $R = 0$ cm is plotted versus V_A. Figure 3 was constructed by averaging, for any given V_A, the ΔV_{sp}'s obtained under several somewhat different CP and AP discharge conditions. With our present setup, evidently we do not observe the saturation of the radial field at $E_r \simeq 0.5$ V/cm reported by *Alport et al.* [1981] (see their Figure 7).

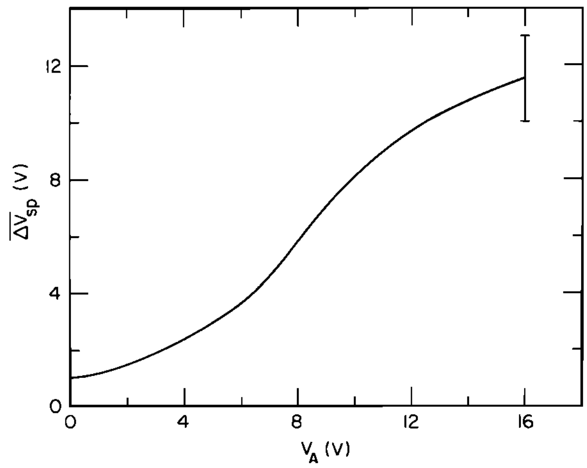

Fig. 3. The average potential difference, $\overline{\Delta V}_{sp}$ (volts), between $R = +5$ cm and $R = 0$ cm, versus V_A (volts).

Experiments were performed with an A_4 electrode opening of either 2 cm or 4 cm, thereby producing plasma columns of different diameters. On some occasions, a metal cylinder, of either 2 cm or 4 cm in diameter and 4.5 cm in length, was also attached to A_4, as shown in Figure 1. The terminating plate connected to A_1 was either 7 cm or 11 cm in diameter. Several combinations of these various electrodes were used in the experiments to be described in section 3, to ensure that the results obtained were largely independent of any particular electrode configuration. We also used Langmuir probes of different sizes, from a minimum diameter of ~0.15 cm to a maximum diameter of ~0.5 cm.

To obtain the spectral distribution of the plasma potential fluctuations, $\tilde{\phi}$, the signal from a floating Langmuir probe was fed into a Hewlett Packard model 8557A spectrum analyzer. Possible effects of the Langmuir probe cable capacitance on the spectral shape of the observed fluctuations were repeatedly checked by either varying the length of the cable or by adding in parallel with the cable a capacitance comparable to that of the cable itself. It was concluded that, in the frequency range investigated, cable capacitance had a negligible effect on the spectral shape of the fluctuations.

3. EXPERIMENTAL RESULTS

From the linear theory of *Lee et al.* [1971] and under the general experimental conditions described in section 2, we expected that the frequency spectra of electrostatic fluctuations measured for $E/B \approx$ several times C_s would each exhibit a broad peak at $f \simeq (0.5$–$0.7)f_{LH}$ (f_{LH} is the lower hybrid frequency), namely, around a frequency $f \sim 1500$ kHz.

Figure 4 shows the measured power spectrum ($\propto \tilde{\phi}^2$) of the potential fluctuations ($\tilde{\phi}$) between ~100 kHz and ~2.0 MHz for $p \simeq 1 \times 10^{-3}$ torr, $B \simeq 225$ G and $V_A = 0$, 5, 10, and 16 V. The hole diameter of the A_4 electrode in this case was 2 cm, and the spectra were measured at a fixed radial position of ~3 cm where E_r is greatest. Typically we find fluctuation levels $\Delta n/n = e\tilde{\phi}/kT_e \sim 20$–30%. Evidently, for all four values of V_A, the power spectra fall very rapidly with increasing frequency, in the same general manner reported by, for example, *D'Angelo et al.* [1975] and *Alport et al.* [1981]. Apparently, at

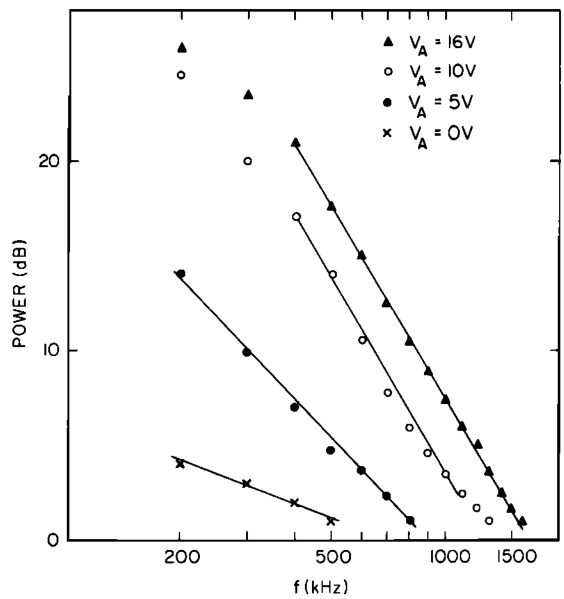

Fig. 4. Power spectra of the potential fluctuations for $V_A = 0$, 5, 10, and 16 V ($p = 1 \times 10^{-3}$ torr, $B = 225$ G, $R = +3$ cm).

the largest V_A's investigated, the effect of increasing V_A and, thereby, the radial electric field is mainly to enhance the general power level of the fluctuations, without appreciably affecting the spectral shape. The same results are obtained when the neutral gas pressure is increased to $\sim 3 \times 10^{-3}$ torr.

The spectra shown in Figure 4 have been plotted on a logarithmic frequency scale with the power given in decibels to emphasize a dependence of the type $P(f) = Af^{-\alpha}$. This assumed dependence gives the power in decibels as, P(dB) $=$ const $- \alpha \log f$. For $V_A = 5$ V this gives a spectral index $\alpha \approx 2.1$, whereas for $V_A = 10$ V and $V_A = 16$ V, $\alpha \approx 3.3$–3.5. A spectral form of the type $P(f) \propto f^{-3.5}$ was also reported for the standard low-frequency Farley-Buneman instability by *D'Angelo et al.* [1975], *John and Saxena* [1975], and *Alport et al.* [1981].

The general shape of the power spectra shown in Figure 4 was consistently observed, in the present experiment, in all spectra measured under conditions we regard as suitable to the excitation of the high-frequency instability of *Lee et al.* [1971]. These spectra do not seem to be a peculiarity of one special experimental arrangement. Not only were such spectra measured with the various combinations of electrode A_4 hole size and end plate EP size; they were also observed (although in an experiment not designed specifically to test the *Lee et al.*

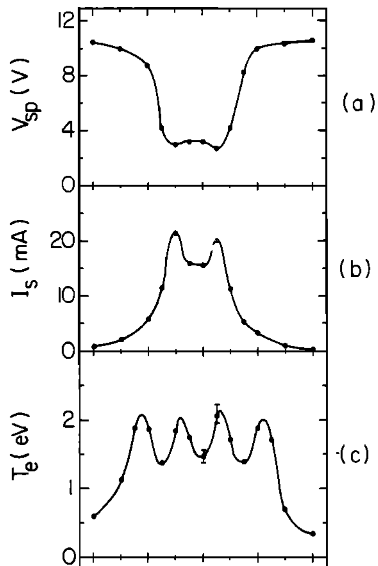

Fig. 5. Radial profiles of (*a*) space potential V_{sp}, (*b*) electron saturation current I_s to a Langmuir probe, (*c*) electron temperature T_e, and (*d*) power in the potential fluctuations at $f = 200$, 400, 1000, and 1400 kHz ($I_d^{CP} = 4$ A, $I_d^{AP} = 10$ mA, $V_d^{CP} = V_d^{AP} = 50$ V, $V_A = 8$ V, $p = 1 \times 10^{-3}$ torr, $B = 225$ G). The hole size was 4 cm with a 4.5-cm-long metal cylinder connected to A_4.

[1971] calculations) by *John and Saxena* [1975], in a plasma produced by an RF discharge, and with radial electric fields such that $E/B \simeq$ several times C_s.

In Figure 5 we show a set of measurements (A_4 hole size of 4 cm, $p \simeq 10^{-3}$ torr, $B \simeq 225$ G, and $V_A = 8$ V) of radial profiles of (Figure 5*a*) space potential V_{sp}, (Figure 5*b*) electron saturation current I_s to a Langmuir probe, (Figure 5*c*) electron temperature T_e measured by the probe, and (Figure 5*d*) power in the potential fluctuations of frequencies $f = 200$, 400, 1000, and 1400 kHz. From Figure 5 we notice the following: (1) the largest fluctuation amplitude is observed, at all four frequencies, in the annular spatial region where the radial electric field is largest, and (2) although the electron temperature at the outer edges of the plasma column is as low as $\sim$0.4–0.5 eV (consistent with rapid cooling of the electrons by collisions with neutral gas molecules at pressures $p \simeq 10^{-3}$ torr), the electron temperature within the column is of the order of 1.4–2 eV and exhibits four well-defined peaks. These peaks in T_e do not occur at the same radial positions where E_r or the fluctuation amplitudes, $\tilde{\phi}$, are largest. Instead, the minima of T_e at $R = \pm 3$ cm are observed at or near the locations where both E_r and $\tilde{\phi}$ are largest. The most straightforward explanation of this effect seems to be that the electron temperature enhancements observed in Figure 5 are not due to electron wave heating nor to resistive heating by the radial Pedersen currents, which must predominantly flow radially inward wherever the radial **E** field is largest, but rather to currents flowing along **B**, just inside and just outside the annulus of largest E_r, to complete the current paths.

This view is strengthened by similar measurements, performed under somewhat different experimental conditions, and shown in Figure 6. The four T_e peaks are still evident here, although with different magnitudes, and the minima of T_e at $R = \pm 3$ cm again occur at the locations of largest E_r and largest fluctuation amplitudes.

The results presented in this section will be discussed in section 4, together with the relevant information available from ionospheric observations of the high-frequency Farley-Buneman instability and of the electron heating by large **E** fields in the lower E region.

4. DISCUSSION AND CONCLUSIONS

As pointed out in section 3, a shift of γ_{max} (the maximum growth rate of the Farley-Buneman instability) to higher and higher frequencies as E_r/B is continuously increased to several times C_s, is not apparent from our experimental results. This shift of γ_{max} is predicted by the linear theory of *Lee et al.* [1971], and we should have expected it to give rise, in the nonlinear evolution of the instability, to spectra dominated by a broad peak around $0.7 f_{LH} \approx 1.3$–1.7 MHz, this frequency range on f_{LH} reflecting the density variation in the region of strongest radial electric field. It seems to us quite surprising, if the *Lee et al.* [1971] instability is indeed excited as expected, that the instability should give rise in the observed turbulent spectra to no recognizable signature or clue of its presence at the linear stage.

What is the reason for this puzzling result? We have considered the possible effect of the finite ion drift which may be present in our experiment. At a pressure of $\sim 1 \times 10^{-3}$ torr, the mean free path for ion-atom collisions is of the order of a few centimeters. For $V_A = 10$ V, the ions would be accelerated radially inward reaching energies up to 2 eV. However, for $V_A = 16$ V, we observed no significant changes in the spectrum although E_r is larger and presumably the ion drift would be

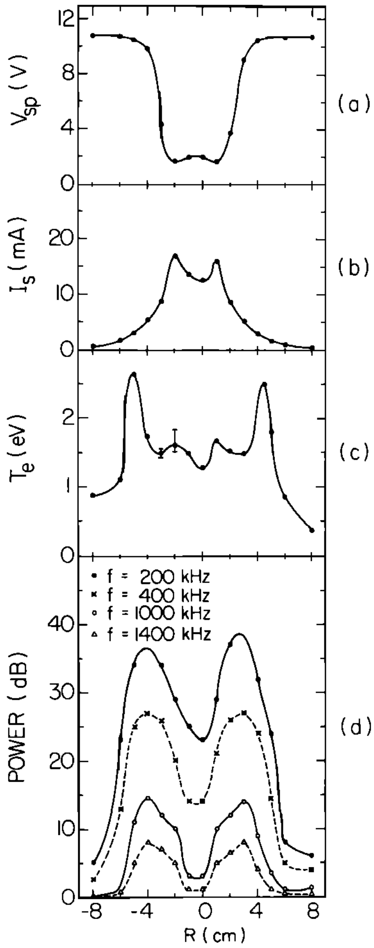

Fig. 6. As in Figure 5 with a 4-cm hole size, but without the metal cylinder connected to A_4.

more significant. Also, we observed no significant differences in the spectrum when the pressure was raised to $\sim 3 \times 10^{-3}$ torr. We believe, furthermore, that the size of the Langmuir probes used to detect the fluctuations and the cables connecting our probes to the spectrum analyzer were appropriate for the expected wavelengths and frequencies. We tried different experimental arrangements all with the same results as far as the spectral shape was concerned. Also, as already noted in sections 2 and 3, we used different combinations of electrodes A_4 and A_1 to make sure that the observed spectra were not a peculiarity of one particular device configuration.

In addition, *John and Saxena* [1975], working with an entirely different device and radial **E** fields large enough to give $E/B \sim$ several times C_s, found also spectra very similar to ours. This strengthens our experimental results, adding to our confidence that they are not device dependent.

It could be argued that the cylindrical plasma geometry of both our experiment and that of John and Saxena's is not suitable to the excitation of the high-frequency instability. This argument, however, would seem very specious indeed, since we have previously reported, under similar geometrical configuration, the excitation of the low-frequency, long-wavelength Farley-Buneman instability. It would appear that, if at all, a cylindrical geometry of some given radius R for the plasma would be more prejudicial to the excitation of wavelengths $\gtrsim R$ than of wavelengths $\lesssim R$. This argument becomes even more cogent when possible end effects are considered.

Since we do not know of any obvious flaw in the linear calculations of *Lee et al.* [1971], we must somehow come to the conclusion that, irrespective of the conditions of linear excitation of the Farley-Buneman instability (namely, γ_{max} largest at low frequency or at high frequency) the turbulent spectrum evolves generally in such a way as to contain most of the energy at the longest wavelengths allowed by the size of the system and that a $P(K) \propto K^{-3.5}$ spectral shape is, somehow, an almost "universal" shape. This type of argument would be in line with the ideas prevalent on two-dimensional magnetohydrodynamic turbulence, as presented, for example, in the work of *Knorr* [1974].

Is this (tentative) conclusion warranted by what we know about the excitation of the high-frequency instability in the high-latitude ionospheric E region? Here the situation seems also quite puzzling. *Olesen et al.* [1976] and *Bahnsen et al.* [1978] reported observations, by rocket-borne instruments, of the Farley-Buneman instability in the polar cap E region. The measured DC electric field was $\sim$70–80 mV/m, large enough to provide electron drift velocities of several times the ion acoustic speed. The high-frequency Farley-Buneman instability should, thus, have been easily excited in their case. They observed in the ELF region (18–180 Hz) very regular wave forms (see, for example, Figure 4 of *Bahnsen et al.* [1978]), i.e., "a nearly sinusoidal wave form, with a frequency close to 100 Hz." This fact already suggests a spectral shape of the fluctuations in which the wave amplitude drops very rapidly with increasing frequency, since, for example, a nearly flat spectrum would hardly be reconcilable with "a nearly sinusoidal wave form." The VLF sensor provided a signal amplitude at 7 kHz of $\sim 10^{-7}$ V m^{-1} Hz$^{-1/2}$, at 100- to 110-km altitude. Since the corresponding ELF amplitude is given as $\sim$5 mV/m, we may compare the two directly if we assume, conservatively, that the line width around 100 Hz was $\sim$50 Hz. This gives for the ELF amplitude a figure of $\sim 7 \times 10^{-4}$ V m^{-1} Hz$^{-1/2}$. Thus the amplitude at $\sim$7 kHz is $\sim 10^{-4}$ times smaller than the amplitude at $\sim$100 Hz, which indicates a spectral shape with very rapidly decreasing power to the higher frequencies.

On the other hand, *Kelley and Mozer* [1973], in an earlier paper, have reported the observation of the high-frequency Farley-Buneman instability in the auroral E region, when DC electric fields were present of $\sim$70–80 mV/m. In particular, a sonogram of the AC electric field showed electrical noise extending from zero frequency all the way to $\sim$6.5 kHz, the lower hybrid frequency. AC electric field spectra, in the altitude range 96 km to 113 km, were essentially flat from $\sim$20 Hz to $\sim$3 kHz (see Figure 6 of *Kelley and Mozer* [1973]). There are, however, certain features in the Kelley and Mozer report which appear rather puzzling. For example, the broadband (50–1000 Hz) electric field was detected by two independent probe detectors of different antenna lengths (2.36 m and 4.48 m). The voltages detected by the two instruments were proportional to the antenna length, between altitudes of $\sim$95 km and $\sim$120 km, indicating, as the authors remark, that the detectors responded primarily to waves with wavelengths comparable to or larger than several meters. This conclusion seems to be somewhat at odds with a spectrum which is essentially flat between $\sim$20 Hz and $\sim$3 kHz.

With regard to specific calculations involving the nonlinear evolution of the Farley-Buneman and gradient drift instabilities in the ionosphere, we note that *Sudan* [1983] has predicted a spectrum of the form $K^{-5/3}$. Sudan also predicted that when high enough frequencies are directly excited by the linear mechanism, the slope of the spectrum should become less pronounced, without changing its sign. Within the precision allowed by our experimental setup, we did not find such a change in slope from our measurements.

Related to the question of the excitation of the Farley-Buneman instability is the "anomalous" heating of the polar E region observed, among others, by *Schlegel and St. Maurice* [1981]. With the incoherent scatter facility at Chatanika, Alaska, they measured electron temperatures of up to 1200°K near the 110-km altitude in the polar E region when strong DC electric fields were present. They concluded that no classical heating process could account for these high temperatures, and suggested that the temperature enhancements were probably caused by unstable plasma waves of the Farley-Buneman type. This is an interesting suggestion which has been explored in detail in the companion paper by *St. Maurice et al.* [1981]. These original heating calculations, however, depended on the presence of a sufficiently high power at short wavelengths (tens of centimeters) in the Farley-Buneman spectrum. Can the long-wavelength waves be therefore responsible for the apparent electron heating? We are inclined not to think so, from the following argument. *Schlegel and St. Maurice* [1981] estimate that a heat input to the electrons of 1–2 × 10^{-6} erg cm^{-3} s^{-1} is required in order to account for the enhanced electron temperatures. Now, the rate of energy transfer from unstable Farley-Buneman waves to the electrons, Q_w, can be estimated as $Q_w = \varepsilon/\tau$, where ε is the energy density in the waves and τ is a characteristic time for energy transfer. If we write $\varepsilon = 1/2(n_0 m_i v_1{}^2)$ and $v_1 \sim (\Delta n/n_0)C_s$, where n_0 is the unperturbed plasma density, $\Delta n/n_0$ is the fractional density perturbation due to the waves, and C_s is the ion acoustic speed, we find that $\varepsilon \sim 2 \times 10^{-10}$ erg cm^{-3}, when $n_0 \sim 2 \times 10^5$ cm^{-3}, $C_s \sim 300$ m/s, and $\Delta n/n_0 \sim 0.2$. This value of ε is, in fact, likely to be an overestimate, since $(\Delta n/n_0)$'s which are as large as 20%, even at the longest wavelengths, are uncommon.

The Farley-Buneman wave periods, at the long wavelengths, are of the order of 10^{-2} s, and it seems very unlikely that τ should be any smaller than 10^{-2} s, as otherwise no coherent wave [e.g., *Bahnsen et al.*, 1978] would ever be observed. In fact, the best estimate of τ is, probably, $\sim 10^{-1}$ s, i.e., about equal to the linear growth time computed by *Schmidt and Gary* [1973]. We thus obtain $Q_w = \varepsilon/\tau \sim 2 \times 10^{-9}$ erg cm^{-3} s^{-1}, which is smaller than 1–2 × 10^{-6} erg cm^{-3} s^{-1} by about 3 orders of magnitude. It has been suggested, however, that for some reason, in a collisional medium such as we have here, we would need to consider τ to be the shortest time scale of the problem, that is, the electron collision time instead of the linear growth time. In that case, a heat input of 10^{-6} erg cm^{-3} s^{-1} could indeed be reached. We do not believe that the collision time should be used, however, because with an energy transfer time as short as $v_{en}{}^{-1}$, the coherent waves of *Bahnsen et al.* [1978] would hardly seem possible.

If we are wrong about our choice of the characteristic energy transfer time, τ should remain constant for all wavelengths. Then, with a power law of the type $f^{-3.5}$ any wave heating would have to be done by low-frequency waves. But if, on the other hand, we are right about our choice of τ, and if higher-frequency waves are excited, no wave heating at all would, nevertheless, seem possible. This would be in spite of a value of τ which would presumably be much smaller than the $\sim$0.1 s estimated for long wavelengths, making the energy transfer time much smaller. As can readily be seen, if the wave number power spectra of the electrostatic fluctuations were indeed of the type $P(K) \propto K^{-3.5}$, nothing would be gained by invoking a contribution to electron heating from the short-wavelength part of the spectrum. With increasing K, the energy available would decrease much faster than the corresponding increase of $1/\tau$.

It seems to us that a reliable measurement of the power spectrum of the high-frequency Farley-Buneman instability in the lower E region is required, under the condition $E \gtrsim 60$–70 mV/m. At present, the two sets of measurements available, namely, those of *Olesen et al.* [1976] and *Bahnsen et al.* [1978], on the one hand, and of *Kelley and Mozer* [1973] on the other, appear to give contrasting results. Our present laboratory results, as well as those of *John and Saxena* [1975], are not easily reconcilable with those of *Kelley and Mozer* [1973].

After this paper was prepared, we learned of a very recent work by *Pfaff et al.* [1984] on electric field and plasma density measurements in the auroral electrojet. With a DC electric field of 54 mV/m, i.e., large enough for the excitation of the *Lee et al.* [1971] instability, they measured power spectra of the density and electric field fluctuations in which the power decreased rapidly with increasing wave number. They also concluded that there appears to be insufficient power in the short-wavelength (tens of centimeters) fluctuations, so that their amplitude level may not be as high as required by the anomalous heating theory as it is presently formulated [*St. Maurice et al.*, 1981].

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