Astron. Astrophys. 94, 100–108 (1981)

ASTRONOMY AND ASTROPHYSICS

Coronal Evolution and Solar Type I Radio Bursts: An Ion-acoustic Wave Model

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Received July 23, 1979; accepted June 26, 1980

Summary. We propose a model for type I burst emission that can accommodate both the main burst observations and an origin for the continuum. We suppose that ion-acoustic waves are generated in the burst source by a current that is related to the coronal magnetic evolution, in particular to magnetic nonequilibrium caused by photospheric changes (e.g. emerging magnetic field) in active regions. Radio emission arises from coalescence of ion-acoustic and plasma waves. Contrary to other plasma wave models, emission at the harmonic of the plasma frequency is below the present detection threshold ($\lesssim 0.1\%$). The ion-acoustic wave density, having a high saturation value, determines the optical depth, which reaches unity within a few meters. The brightness temperature is thus entirely given by the level of Langmuir waves. These waves may be produced by trapped non-thermal electrons from previous burst sources. The same population also provides sufficient plasma waves for the type I continuum, which may arise from interactions with low-frequency waves present in the corona during times of type I activity.

Key words: solar radio noise storms – type I emission – ion-acoustic waves – current instability

1. Introduction

Shortly after his radio detection of the Sun, Hey (1946) identified a long enduring radio emission connected with the appearance of large sunspots on the solar disk. Now we recognize two principal noise storm components: the short narrowband type I bursts and the broadband type I continuum. There is no generally accepted explanation for this radio emission. Several theories have been summarized by Elgarøy (1977). The most widely cited cyclotron model by Mangeney and Veltri (1976) requires a beam of spiraling electrons for each burst and a suitable spectrum of magnetic turbulence. The theory may also account for the associated type III storm phenomenon at decametric wavelengths (Aubier et al., 1978). None of the published theories explain the cause of the energetic phenomenon, such as a beam of suprathermal electrons (Mangeney and Veltri, 1976; Takakura, 1963, and Sy, 1973), a strong MHD wave (Trakhtengert, 1966; Vereshkov, 1974) or a strong shock (Zaitsev and Fomichev, 1973).

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* Supported by the US National Science Foundation, Atmospheric Research Section and by the Radio Astronomy Group, ETH, Zürich

The type I bursts are definitely not related to flares. They occur for days at rates as high as one per second. Their association with the growth of complex active regions suggests that they are intimately connected with the evolution of the solar corona, and perhaps with basic processes in stellar atmospheres. In particular, they may be the signatures of many small steps in coronal evolution whose cumulative effect is the "gradual" evolution of the corona.

We propose a new model for type I bursts. It includes a plausible energizing mechanism associated with coronal evolution, namely the generation of intense current-driven ion-acoustic waves in a small portion of the corona. The ion-acoustic waves cause anomalous electrical resistance and produce radio emission at the plasma frequency by combining with Langmuir plasma waves. We argue in Sects. 3 and 4 that the conditions for such a situation are plausible. In Sect. 5 the emission mechanism is investigated in detail. We note that Melrose (1979) is considering similar processes in connection with noise storms.

The electrons heated in sites of anomalous resistance escape into the ambient corona. There they may become trapped, with two important consequences. First, the associated Langmuir waves yield the observed radio continuum. Second, the suprathermal electrons at the sites of anomalous resistivity lead to the Langmuir waves that cause the type I radio emission while the resistivity lasts. If the electrons are not trapped, storm-type III bursts may result. We discuss these matters in Sect. 6. A brief summary appears in Sect. 7.

This model should only be considered as a first step toward a more complete understanding of a complex situation. We make assumptions concerning relevant coronal *mhd* and plasma physics that still need to be justified or modified. These constitute possible crystallization points for future work. Most importantly, however, this model suggests that type I radio bursts may constitute an important tracer of the physics that leads to coronal evolution, and perhaps the bursts may indicate the sites where coronal evolution actually takes place. If so, it points to an exceedingly restless corona, even at non-flaring times (Wentzel, 1980).

2. Observational Considerations

Radio observations of type I emission have been reviewed extensively by Kundu (1965) and Elgarøy (1977). Here we summarize the most relevant facts we use in the theory.

1. Noise storms are associated with strong magnetic fields and complex sunspot structure, thus with large gradients in magnetic fields, that is, on scales of the order of 10^3–10^4 km. Noise storms often appear in the growing phase of a spot group and sometimes

are clearly associated with changes in the photospheric field (Sakurai, 1975). Theoretically this implies changes in the electrical currents imposed on the corona from below and/or strong induced electric fields.

2. The continuum extends from about 400 MHz well into decameter wavelengths and has a total duration of a few hours to several days, commensurate with the time scales of photospheric and coronal evolution. Superimposed are individual type I bursts at frequencies above 40 MHz and type III storm bursts below. Bursts may occur every few seconds and, given that the radiation is somewhat directive (Steinberg et al., 1974) and many bursts may be missed, they may occur at the rate of once per second or more.

3. The position of type I bursts scatters around that of the continuum source. The centroid of burst emission is usually within the observed source region of the continuum (Daigne et al., 1970). Thus a given source location is associated with the same frequency for both continuum and bursts. Bursts and continuum are equally and usually strongly polarized. Thus we adopt fundamental radiation near the plasma frequency for both.

4. The burst radiation has a high brightness temperature. This observation is used below to eliminate enhanced Bremsstrahlung from suprathermal isotropic electrons as an explanation for the bursts. How high may the brightness temperature T_b become? One observes $T_b < 10^{10}$ K when using the observed angular dimension of 3′. However, this is limited by coronal scattering.

Temporal structure is observed down to 0.1 s. Suppose first that the waves within the source travel at the speed of light, c. Then the temporal duration implies a size of $3\ 10^4$ km or 40″. If the source has two such dimensions, T_b without coronal scattering would be $2\ 10^{11}$ K. We feel that this is a minimum estimate and note that T_b corresponds to 10 MeV.

Our model makes the radio waves within the source move at the velocity $n_t c$, the index of refraction $n_t \approx 0.2$. The source diameter is reduced by a factor n_t from 40″, and the emission area by a factor n_t^2. However, the intensity leaving the corona is higher than at the source by a factor $1/n_t^2$. Thus the deduced T_b at the source is independent of n_t at the source.

The brightness temperature may in fact exceed the minimum of $2\ 10^{11}$ K. Our theory can accommodate values up to roughly $T_b = 10^{13}$ K. Therefore, it can accommodate a source some 50 times smaller in area than the above $(6000\ \mathrm{km})^2$, or it can accommodate absorption by a factor of 50 as the radiation escapes from the plasma level. Such absorption may be expected for plasma-frequency radiation at a few 100 MHz if the gas density decreases outwards in a smooth manner. However, strong density-dependent absorption would introduce a strong frequency-dependence into the burst intensity, which is not noted. Indeed, if the emission is from a site of somewhat enhanced density, as seems likely theoretically (see below), absorption is not important.

5. The bandwidth of a type I burst is only about 2% of the observed frequency. If the radiation is at the plasma frequency, the bandwidth corresponds to a range in density varying by 4% or a range in height of the order of 7000 km (10 times Baumbach-Allen density model for emission of 100 MHz). This is only an upper limit. It might be much less if the emission at one location has a finite bandwidth. Indeed, our model invokes a burst site shaped more like a crinkled sheet, some 6000 km across in two dimensions and of the order of 0.1 km thick.

6. We emphasize the lack of recognizable harmonic emission from type I bursts. Since harmonic emission from isotropic Langmuir waves easily outshines the fundamental, the waves must be rather weak. The high T_b then requires an intense level of low-frequency waves.

A numerical estimate on the maximum T_b consistent with the lack of harmonic emission appears in Sect. 4. (We discount the possibility of intense unidirectional Langmuir waves. See Sect. 4.)

7. Often, type I bursts cluster in "drifting chains", which generally (but not always) drift to lower frequency. Taking observations by de Groot et al. (1976) and assuming fundamental plasma emission we deduce an average velocity of the exciter of 160 km/s. At lower frequencies the drift rates and derived speeds are considerably smaller (Elgarøy and Ugland, 1970; de la Noe, 1978). The exciter velocity has nothing to do with the signal velocity internal to the burst source discussed in point 4) above. For instance, if field lines emerge or are compressed or twisted at the foot of a loop, and if triggering of the burst instability requires a finite amplitude disturbance, the site of this amplitude may migrate at some fraction of the Alfvén speed.

3. MHD Considerations

Our model is based on the occurrence of anomalous electrical resistivity. We discuss this in Sect. 4 but treat here the major problem associated with most attempts to explain the magnetic evolution of the solar corona: if magnetic reconnection occurs in the corona, it requires electrical currents and matching field gradients or field twists that are highly localized, typically on scales of the order of a km or less (Wentzel, 1978). This statement is independent of whether or not the magnetic evolution also causes heating of the corona.

We argue that emerging magnetic flux and photospheric changes in existing fields *inevitably* lead to *field reconnection* and we identify type I bursts with some of these necessarily small sites of reconnection. The arguments are: (i) the imposition or change of currents entering or leaving the corona can lead to a force-free equilibrium only if the distant footpoints of the field lines accommodate these currents. This is unlikely. Coronal disequilibrium results. Mathematically, the equations of magnetostatic equilibrium are overspecified (Parker, 1972). (ii) The disequilibrium manifests itself in field changes traveling about, until they become singular. For example, emerging flux punches through and separates old field lines until their own tangling prohibits further separation. We shall call such processes "knot-tying" (Wentzel, 1980) but the image should not be taken too literally. The field can simplify fast enough only when the current becomes so localized and strong that instabilities and anomalous dissipation result. Indeed our model implies a rate of reconnection comparable to that needed for coronal evolution (Wentzel, 1980).

The picture we propose is sketched in Fig. 1. The volume of emission is a thin sheet of high current density. Although many weak coronal currents may travel parallel to the field, the driving motions for the instability require a finite component $j_\perp$, of $\boldsymbol{j}$ normal to $\boldsymbol{B}$. We shall adopt $j_\perp = 0.1\, j$. The current sheet may well be crinkled since the driving process is inherently inhomogeneous, but all the plasma phenomena occur on such short spatial scales that the sheet may be considered flat for this purpose.

We expect that the electrons are heated rapidly. If their pressure were to become comparable to the magnetic pressure, they would tend to widen the current channel and shut off the ion-acoustic instability. However, before that happens they tend to escape along the field lines: with $j_\perp = 0.1\, j$, a sheet 0.1 km thick extends only 1 km along the field lines. This we use in Sect. 6 as the source of suprathermal coronal electrons. For the source itself it means that the electron heating is limited. We adopt $T_e = 2\ 10^7$ K in the following. We note that an electron at this

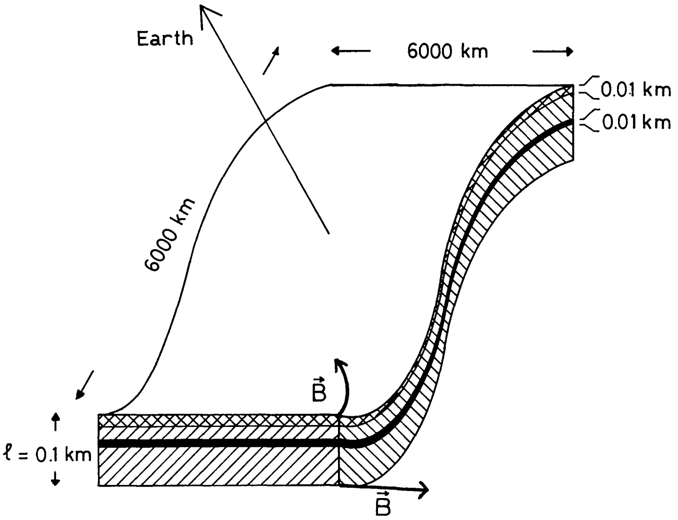

Fig. 1. Schematic view of proposed type I burst source. The current is unstable in the black region, from where the ion-acoustic waves spread into the shaded volume causing an enhanced Langmuir wave level. Fundamental emission by coalescence of an ion-acoustic wave and a Langmuir wave becomes optically thick in the cross-hatched region

temperature crosses half of the source size (3000 km) in 0.2 s, comparable to the observed time scales.

The energy released within the type I burst source is not sufficient to heat the corona (see Sect. 6).

The ion-acoustic instability generates an anomalous resistance. The dissipation reduces j. To avoid j falling below the threshold for the ion-acoustic instability, j must be continuously driven from the outside. With an effective collision frequency $\nu_{\rm eff}=\beta_s\omega_p =10^{-2}\omega_p$ (Papadopoulos, 1977 and Sect. 4 for β_s) the conductivity is given by $4\pi\sigma=\omega_p^2/c^2\nu_{\rm eff}=10^{-19}\omega_p$, and the current dissipation time $\tau=4\pi l^2\sigma$ where $4\pi j=B/l$. With $j=nv_D e/c$ and $n=10^8\ {\rm cm}^{-3}$, $B=3\ G$ (Alfvén speed $650\ {\rm km\ s}^{-1}$) and $v_D=300\ {\rm km\ s}^{-1}$, $l=0.03$ km and $\tau \doteq 10^{-3}$ s. The rate of driving, $l/\tau=30\ {\rm km\ s}^{-1}$, is reasonable. The duration of the burst is the duration of the driving motions.

The initial current sheet must be narrower than l, say 0.01 km, because a current thickness l would imply a field reversal and neutral sheet, which is unlikely.

The distribution of the ion-acoustic waves is wider than the current thickness, since the waves travel some distance before damping. This distribution gives the extent of electron heating and, if suitable Langmuir waves are present, radio emission. We adopt 0.1 km for the thickness of the hot, radiating volume, which corresponds to the product of damping time and group velocity of ion-acoustic waves.

4. Energy Density in Waves

a) Plasma Waves

Plasma radiation is common in the laboratory. It is generally associated with nonthermal electrons. For example, Hutchinson et al. (1978) consider a thermal level of ion-acoustic waves plus a level of Langmuir waves made by Bremsstrahlung from suprathermal electrons. However, for type I bursts the minimum T_b corresponds to 10 MeV (Sect. 2). Relativistic electrons (Papadopoulos, 1969) appear unlikely. Complex electron streams with energies in the range 10–100 keV have been found in interplanetary space by Lin and Anderson (1967) and Lin (1970); these originate in active regions, usually with intense noise storm radio emission. Therefore, the type I bursts are probably not caused by the enhanced Bremsstrahlung of relativistic electrons.

The source for the Langmuir waves is not likely to be runaway electrons, for two reasons. First, if the current density is at the ion-acoustic threshold, then $E/E_D\approx v_i/v_e<10^{-2}$, where E_D is the Dreicer field with the anomalous collision frequency that also appears in the resistivity. Thus $\boldsymbol{E}$ produces very few nonthermal particles. (Strong heating of the ions would quench the ion-acoustic instability.) Second, the Langmuir waves of runaway electrons would be in the same direction as the ion-acoustic waves, namely both opposite to $\boldsymbol{j}$ and $\boldsymbol{E}$, and the waves would not combine into radiation.

We shall argue that the faster of the thermal electrons escape from the type I burst site, later to be trapped and used to produce Langmuir waves in other burst sites (or returning to the same site). Without such trapped electrons, the site of ion-acoustic waves would remain unobservable.

The same phenomenon that causes the rapid dissipation, namely the collective behavior of the ions, also enhances the Bremsstrahlung emission of Langmuir waves by the fast electrons. This emission requires that the fast electrons transit the ion waves at effective rates comparable to the plasma frequency, i.e. $(k\lambda_D)v_f/v_e>1$, which is marginally satisfied for our values. Thus a nonthermal level of Langmuir waves may well be in equilibrium with the fast electrons despite the ion waves. Finally, we should recognize that the level of ion-acoustic waves implies large deformations in the velocity distributions. The differences in potential energy between peaks and troughs of the waves are roughly $0.3\,\beta_s^{1/2}$ times the energy of the hot electrons. This will certainly make the velocity distributions of both the "thermal" and the "hot" electrons non-Maxwellian and increase the emission of Langmuir waves.

Nonrelativistic electrons that are isotropic can explain T_b only to $T_b\simeq 10^{10}$ K ($\approx mc^2$, Robinson, 1977). This may be adequate for the continuum (Sect. 6) but not for the bursts. *Anisotropic* electrons are necessary, presumably with some marginally stable velocity distribution. We expect such a distribution in the presence of the greatly enhanced effective collision frequency due to the ion-acoustic waves that we invoke (Benz and Kuijpers, 1976). The required intense level of plasma waves appears at least plausible.

Let us derive an upper limit on β_L, the energy density in Langmuir waves compared to nT_e, assuming that the Langmuir waves are distributed isotropically. The power emitted at the harmonic is about

$$P\approx 12\beta_L^2 nT_e\omega_p(k_L^m\lambda_D)^{-3}(v_e/c)^5\ {\rm erg\ cm^{-3}\ s^{-1}}, \quad (1)$$

where $\lambda_D=$ Debye length, $v_e^2\equiv T_e/m$, and k_L^m is the maximum wavenumber assuming a broadband distribution. Let us choose $n=10^8\ {\rm cm}^{-3}$, $T_e=2\ 10^7$ K ($3\ 10^{-9}$ in energy units), and $k_L^m\lambda_D=0.3$. Then $P=5\ 10^4\ \beta_L^2\ {\rm erg\ cm^{-3}\ s^{-1}}$. At values of interest, the emission is optically thin, so we must specify an emission volume. We choose the value $(6000\ {\rm km})^2\times 0.1\ {\rm km}=10^{21.5}\ {\rm cm}^3$. When the radiation is distributed over an area of $10^{27}\ {\rm cm}^2$ at 1 A.U. and over a frequency range of 3 MHz, the flux received is $4\ 10^{-8}\ \beta_L^2$ erg s^{-1} Hz^{-1} cm^{-2}. The observed upper limit is about 10^{-20} erg s^{-1} Hz^{-1} cm^{-2} (Jaeggi and Benz, in preparation). We obtain $\beta_L<5\ 10^{-7}$. If the Langmuir waves were the most intense waves, the emission at the fundamental in a volume as thin as 0.1 km would be negligible. In contrast, as we show below, with intense (low-frequency ion-acoustic) waves the optical depth can reach unity over a distance of merely 0.1 km. The thermal level of Langmuir waves under the same circumstances (integrated up to the same k_L^m is about $\beta_L = 10^{-1.5}\, k_L^{m3}/n = 0.5\; 10^{-12}$ [see Eq. (23)]. Our upper limit exceeds the thermal value by a factor of 10^6. Therefore, if the fundamental is optically thick, one may explain a fundamental $T_b = 10^6\, T_e = 2\; 10^{13}$ K without exceeding the limit on the harmonic emission.

b) Low-frequency Waves

Several kinds of low-frequency waves could be considered. We propose ion-acoustic waves. The essential feature of the ion-acoustic instability is that it saturates at a relatively high level. This has two vital consequences for the theory of the type I bursts, as we shall see in Sect. 5. First, since the fundamental emission is proportional to β_s, the high β_s causes intense emission. Therefore, the observed high fundamental brightness temperature is compatible with a very thin source of radiation, as is necessary for the current-driven instability. Second, high β_s permits the intensity of Langmuir waves to be relatively weak, consistent with the lack of harmonic radio emission. This would not be the case for low-frequency waves that saturate at a low level, for then the Langmuir waves would have to be intense to produce enough fundamental in a small volume, and the the harmonic would exceed the fundamental emission.

5. Wave-wave Interactions and Radio Emission

The condition for efficient energy transfer by $s + L \to t$ wave coalescence is

$$\omega_s \pm \omega_L = \omega_t \tag{2}$$

$$\boldsymbol{k}_s \pm \boldsymbol{k}_L = \boldsymbol{k}_t . \tag{3}$$

We assume here that weak-turbulence theory is applicable despite the large β_s, at least to demonstrate that high optical depth occurs even in a very thin source.

We shall proceed assuming that some simple spectrum of Langmuir waves is produced. Of interest, of course, is that part of the spectrum with $\boldsymbol{k}_L \simeq \boldsymbol{k}_s$. Landau damping requires $k_L \lambda_D \lesssim 0.3$. We also require that the emission be in the o-mode to explain the total polarization, which amounts to

$$(k_r \lambda_D)^2 < \tfrac{1}{3} Y \cos\theta_t \tag{4}$$

from the dispersion relations of Langmuir and x-modes (see below). Y is the ratio of electron gyrofrequency to plasma frequency and θ_t is the propagation angle of the transverse wave relative to the magnetic field. We will adopt $\sin^2\theta_t = 0.1$ in numerical estimates. A typical $Y = 0.1$ then yields the limit $k_L \lambda_D < 0.2$. (If free-free absorption is taken into account, the emission need not be totally in the o-mode and Eq. (4) can be relaxed.)

We shall adopt below a spectrum for Langmuir and ion-acoustic waves near $k_L \lambda_D = 0.1$ with a width Δk which is set to $\Delta k = 2k$ in numerical estimates. We also include a solid angle Ω for the ion-acoustic waves, where $0.1 < \Omega < 1$ steradians is indicated by the experiments cited in the Appendix. Only the Langmuir waves within the opposite cone of directions yield radiation. When we estimate necessary Langmuir wave energies below, we are including only the waves interacting with ion-acoustic waves. The total energy in Langmuir and/or ion-acoustic waves may well be greater.

The relevant dispersion relations are

$$\omega_s^2 = c_s^2 k_s^2 = \omega_p^2 (m/M)\, (k_s \lambda_D)^2 \tag{5}$$

$$\omega_L^2 = \omega_p^2 (1 + 3 k_L^2 \lambda_D^2), \tag{6}$$

and for ω_t close to ω_p (because of observed o-mode polarization)

$$\omega_t^2 = \omega_p^2 (1 - Y \cos\theta_t) + k_t^2 c^2 . \tag{7}$$

Total o-mode polarization requires that the x-mode is evanescent

$$\omega_p^2 < \omega_t^2 < \omega_p^2 (1 + Y\cos\theta_t), \tag{8}$$

which is brought into the form

$$Y\cos\theta_t < (k_t c/\omega_p)^2 < 2\, Y\cos\theta_t . \tag{9}$$

The range in frequency, $1/2\, Y\cos\theta_t \omega_p$, is consistent with the observed average bandwidth of $2\; 10^{-2}\; \omega_p$.

The rate equation for the wave combination $s + l \to t$ is

$$\frac{\partial Nt}{\partial t} = \int w\, [N_s N_L - N_L N_t - N_s N_t]\, \frac{d\boldsymbol{k}_L d\boldsymbol{k}_s}{(2\pi)^6} \tag{10}$$

with $N_t = N_t(\boldsymbol{k}_t)$, etc. being the wave quantum densities. The conversion probability w can be derived from the second order terms in the wave equations,

$$w = \frac{(2\pi)^6 \hbar e^2 M \omega_p \omega_s^3}{8\pi m^3 v_e^4 k_s^2 \omega_t} f(\theta_t, \theta_L, \theta_s)\, \delta(\boldsymbol{k}_L - \boldsymbol{k}_L - \boldsymbol{k}_s)\, \delta(\omega_t - \omega_L - \omega_s). \tag{11}$$

The Planck constant $\hbar$ is reduced by 2π, e is the elementary charge. The function f will later reduce to $F(\theta_t, \theta_L)$ of Melrose and Sy (1972). The resulting radio emission is due to the first term in Eq. (10) which yields an emissivity

$$J = \frac{k_t^2}{d\omega_t/dk_t} \int \hbar\omega_t w N_L N_s \frac{d\boldsymbol{k}_L d\boldsymbol{k}_s}{(2\pi)^9}, \tag{12}$$

and the absorption coefficient follows from the other terms:

$$\mu = \frac{1}{d\omega_t/dk_t} \int w\, [N_L + N_s]\, \frac{d\boldsymbol{k}_L d\boldsymbol{k}_s}{(2\pi)^6}. \tag{13}$$

The probability w is greatly simplified by setting $\omega_L = \omega_t$ and $k_s = k_L$. The integral in Eq. (13) over $\boldsymbol{k}_s$ is then straightforward, so that

$$\mu = \frac{1}{d\omega_t/dk_t} \frac{\hbar e^2}{8\pi m} \frac{c_s}{T_e} \int F(\theta_t, \theta_L) k_L [N_t(\boldsymbol{k}_s = -\boldsymbol{k}_L) + N_L(\boldsymbol{k}_L)]\, \delta(\omega_t - \omega_L)\, d\boldsymbol{k}_L . \tag{14}$$

The group velocity yields the factor $\omega_p (k_t c^2)^{-1}$. As discussed above the relevant ion-acoustic and Langmuir waves fill the same solid angle Ω. We define the angular integral in Eq. (14) to be $F(\theta_t, \Omega)\Omega$. For $\theta_L \ll \theta_t$, $F(\theta_t, \Omega)$ reduces to the function $F(\theta_t, \theta_L = 0)$. According to Melrose and Sy (1972),

$$F(\theta_t, \theta_L \approx 0) = \frac{1}{2} \left\{ \sin^2\theta_t + \left[\tan\theta_t \frac{Y}{1 - \omega_p^2/\omega_t^2} + \sin\theta_t \right]^2 \right\} \tag{15}$$

$$\approx \frac{1}{2} \sin^2\theta_t \left(\frac{1 + Y - \omega_p^2/\omega_t^2}{1 - \omega_p^2/\omega_t^2} \right)^2 \approx 4.5 \sin^2\theta_t . \tag{16}$$

The factor 4.5 is replaced by unity for an isotropic plasma (Smith and Spicer, 1979, following Tsytovich, 1970). The delta function in Eq. (14) yields

$$\delta(\omega_t - \omega_L) = \frac{\omega_p}{3 k_L v_e^2}\, \delta(k_L - k_{Lr}), \tag{17}$$

where k_{Lr} is the wavenumber fulfilling the resonance condition.

$$(k_{Lr}\lambda_D)^2=\frac{1}{3}\frac{k_t^2c^2}{\omega_p^2}-\frac{1}{3}Y\cos\theta_t<\frac{1}{3}Y\cos\theta_t. \tag{18}$$

The deletion of k_t in the delta function of Eq. (11) is justified by

$$\frac{k_t^3}{k_{Lr}^2}=3\frac{v_e^2}{c^2}\frac{r}{r-1}, \tag{19}$$

$$1<r\equiv\frac{k_t^2c^2}{\omega_p^2Y\cos\theta_t}<2. \tag{20}$$

We take as typical value $r=1.5$ yielding $k_t/k_{Lr}\simeq 0.03$. An interaction is also possible if $r\approx 1$ and $|k_s|\ll k_t\simeq k_L$. In fact this situation has been proposed by Vereshkov (1974). However, even if ion acoustic waves of small k_s exist, this condition strongly limits the range of possible resonance of k_t and thus of k_L. It can be shown (Vereshkov, 1974) that such an interaction makes use of a very small range of k_L values and less than 10^{-3} of the energy in Langmuir waves (assuming them to have a reasonable width) and therefore this interaction cannot lead to a high enough T_b unless a high level of Langmuir waves is assumed.

We obtain the absorption coefficient

$$\mu=\frac{\hbar e^2}{24\pi m}\frac{\omega_p^2}{k_tc^2v_e^2}\frac{c_s}{T_e}k_{Lr}^2[N_s(k_{Lr})+N_L(k_{Lr})]\Omega F(\theta_t,\Omega). \tag{21}$$

The physics of the waves is generally determined by the energy density, rather than the quantum density. The energy densities are defined through

$$U\equiv\hbar(2\pi)^{-3}\int\omega N(\boldsymbol{k})d^3k. \tag{22}$$

For the estimate of the thermal Langmuir wave level in Sect. 4 we used $\hbar\omega_pN=T_e$ for $k<k_L$,

$$\beta_L^{th}=\frac{k_L^{m3}}{6\pi^2n}. \tag{23}$$

As discussed earlier, the radio emission results from those ranges in wavenumber for which the two sets of waves overlap, $\boldsymbol{k}_L=-\boldsymbol{k}_s$. We write

$$U_L=\frac{\hbar}{(2\pi)^3}\omega_pN_Lk_L^2\Delta k_L\Omega_L, \tag{24}$$

$$U_s=\frac{\hbar}{(2\pi)^3}c_sN_sk_s^3\Delta k_s\Omega_s. \tag{25}$$

Upon collecting all terms, the two parts of the absorption coefficient become

$$\mu_L=\frac{\pi}{12}\left(3\frac{m}{M}\right)^{1/2}\left(\frac{T_i}{T_e}\right)^{1/2}\frac{\omega_p}{(k_t^2c^2-Y\omega_p^2\cos\theta_t)^{1/2}}\frac{\omega_p}{k_tc}\frac{k_L}{\Delta k_L}F\frac{\omega_p}{c}\beta_L \tag{26}$$

$$\approx 0.13\,F\frac{\omega_p}{c}\beta_L, \tag{27}$$

$$\mu_s=\frac{\pi}{12}\frac{\omega_p}{k_tc}\frac{k_s}{\Delta k_s}F\frac{1}{(k_s\lambda_D)^2}\frac{\omega_p}{c}\beta_s \tag{28}$$

$$\approx 5.0\,F\frac{\omega}{c}\beta_s. \tag{29}$$

These absorption coefficients exceed those of Smith and Spicer (1979) by a factor of about 10^2F and $10\,F$ respectively, in part because they set $k_t=\omega_p/c$ and assume isotropy, in part because they assume a wider wave spectrum and differently chosen numerical factors.

Absorption is dominated by ion-acoustic waves for $U_L<40\,U_s$. For radio bursts at 100 MHz, $\mu_s=0.1\,F\beta_s$. With $\beta_s\approx 10^{-2}$, optical thickness unity may well be achieved over a distance of a few meters! We have neglected the stimulated decay process $L\to t+s$ which is only important behind the optically thick region.

Brightness Temperature

The source function of the optically thick source is given by

$$S=2\pi\frac{J}{\mu}=\left(\frac{k_t}{2\pi}\right)\hbar\omega_t\frac{\int wN_LN_sd\boldsymbol{k}_Ld\boldsymbol{k}_s}{\int w[N_L+N_s]d\boldsymbol{k}_Ld\boldsymbol{k}_s} \tag{30}$$

$$=\frac{T_b^s}{\lambda_t^2}, \tag{31}$$

where T_b^s is the brightness temperature measured at the surface of the source in energy units. Upon cancellation of common factors

$$T_b^s=\hbar\omega_p\frac{\int F(\theta_t,\theta_L)N_LN_s\delta(k_L-k_{Lr})d\boldsymbol{k}_L}{\int F(\theta_t,\theta_L)[N_L+N_s]\delta(k_L-k_{Lr})d\boldsymbol{k}_L} \tag{32}$$

$$=\hbar\omega_p\frac{N_LN_s}{N_L+N_s}. \tag{33}$$

The brightness temperature is determined by the set of waves with the lower quantum density, although the absorption occurs via the waves with the higher density. More specifically

$$T_b^s=\hbar\omega_pN_L=\frac{(2\pi)^3}{\Omega}\frac{U_L}{k_L^2\Delta k_L} \tag{34}$$

when $N_L\ll N_s$. This corresponds roughly to $U_L\ll 400\,U_s$, and is certainly satisfied for our theory.

Equation (34) says that $T_b^s\gtrsim 2\,10^{11}$ from a source at $T_e=2\,10^7$ must have a Langmuir wave energy density $>10^4$ times the thermal value *within* the range of k overlapping with the spectrum of ion-acoustic waves. If this range is very small, the total Langmuir wave energy need not be huge. However, this presumes optical thickness exceeding unity, and if the source is merely 0.1 km thick it requires a substantial overlap. Moreover, if the Langmuir waves are nearly isotropic in the presence of the ion-acoustic waves, then Eq. (34) translates into T_b^s/T_e=ratio of actual to thermal Langmuir wave density within the range of k_L common to both. With a thermal $\beta_L=0.5\,10^{-12}$, and $\beta_L<10^{-6}$ as set by the lack of harmonic emission, it is possible to obtain $T_b^s\leqq 4\,10^{13}$ K.

The main accomplishment of this section, then, is to demonstrate that optical thickness unity can be accomplished over merely a few meters when β_s is of order 10^{-2}.

The interaction $s+L\to t$ takes place where there are enhanced levels of both sets of waves. An important condition for achieving high T_b is that absorption of the radio waves outside the overlapping regions, by interaction $t+L\to s$, is small. For a level of $\beta_L=10^{-6}$, an *e*-fold reduction takes of the order of 4000 km [Eq. (27)] and more for parallel direction [cf. Eq. (16)]. It can thus be neglected. Once the local plasma frequency is sufficiently different (say 2%) from the radio frequency, no such interaction is possible and there remains only free-free absorption.

6. Storm-type Continuum and Type III Bursts

The type I bursts, the storm continuum and storm-type III bursts are correlated in time (Boischot et al., 1970). In this section we propose that the current-driven instability of Sect. 4 ejects fast electrons onto the magnetic field lines passing through the site, and that these electrons if trapped yield the type I burst and continuum, if untrapped yield the storm-type III bursts. The requirement of trapping is important because the same electrons added to a preexisting distribution of suprathermal electrons yield a different set of plasma instabilities than if they are released into a region with purely thermal electrons.

6.1. Langmuir Waves in the Continuum Source

The brightness temperature of the continuum is usually below 10^8 K, but up to $5\,10^9$ K has been observed (Kerdraon, 1978). The time structure is more than one minute (Mätzler et al., 1978). Therefore, there is no reason to believe that the continuum source is much brighter than the observed T_b. Moreover, the time scale suggests that emission is caused by trapped electrons. Fast electrons of energy E lose their energy by collision with ions in a time of order $10^3\ E_{\rm keV}^{3/2}\ (10^8/n)$ s. Thus electrons exceeding 1 keV are useful candidates.

The similar position, height and polarization of bursts and continuum (Sect. 2) suggest that both processes occur at the same frequency. Since the polarization is high, the emission is expected to be near ω_p.

The important observation of high polarization also places strong limits on the ratio of harmonic to fundamental plasma-frequency emission of a continuum, since harmonic emission is generally far less polarized. Partially polarized sources are preferentially observed near the limb, which strongly suggests a propagation effect. In the following we discuss radiation for which the ratio of harmonic to fundamental brightness temperatures does not exceed 5% (Benz and Jaeggi, 1980, in preparation).

Limits on β_L can be estimated in two ways. First, we may again apply Eq. (1) for harmonic emission from isotropic Langmuir waves with a uniform spectrum for $k<k_L$. If we again use $n=10^8$ cm^{-3} and $k_L^m\lambda_D=0.3$, but now $T_e=2\,10^6$ K appropriate to the ambient corona, then $P=12\,\beta_L^2$ erg cm^{-3} s^{-1}. We choose an emitting volume for the continuum (here and in all subsequent estimates) with a depth of $10^{4.5}$ km and an area 10^5 km $\hat{=}\,2'$ on a side, that is, $10^{29.5}$ cm^3. The volume may follow along a magnetic loop, and may encompass a large range in plasma density. If we divide the emission into a band of 100 MHz, the resulting flux at Earth is $4\,10^{-5}\,\beta_L^2$ erg cm^{-2} s^{-1} Hz^{-1}. If the maximum harmonic flux based on the high polarization is taken to be 10^{-19} erg cm^{-2} s^{-1} Hz^{-1}, we obtain $\beta_L<0.5\,10^{-7}$. The thermal level for $k_L^m\lambda_D=0.3$ is $\beta_L=10^{-11}$. Thus, if Langmuir waves determine T_b, the corresponding upper limit is $T_b=0.5\,10^4\,T_e=10^{10}$ K, which is well in excess of observed T_b.

A second and direct measurement of Langmuir wave energy density has recently been attempted by Benz and Fitze (1979). In a radar experiment at 2380 MHz they probed type I source regions for Langmuir waves at 220 ± 50 MHz. The interaction of a transverse photon with a plasmon produces an echo photon at the sum and the difference of the two frequencies. Contrary to the first method, this measurement is not based on antiparallel Langmuir waves, but refers to waves in the line of sight towards the observer. However, no echo has been found. The upper limit to the Langmuir wave energy density was determined to be $\beta_L<6\,10^{-4}$, resp. $5\,10^{-3}$ in two possible type I source regions. Unpublished results of the latest measurements confirm that such values apply also for very strong sources. The upper limits correspond to a constant level, i.e. to the continuum source.

6.2. Number of Trapped Electrons

In the following we wish to estimate the number of trapped electrons due to type I sources, the manner of trapping, and the relevant radiation processes. While the results are reasonable, it is clear that they are based on theoretical papers that have not actually been written with the present application in mind and involve somewhat disparate assumptions, for instance for the velocity distribution function of the trapped electrons.

We first discuss the escape of electrons from the hot region of the current instability. This is somewhat similar to the situation recently considered for flares (e.g. Vlahos and Papadopoulos, 1979), except that our temperature is lower and, most importantly, the source is much thinner. If the smallest dimension of the source is 0.1 km and $j_\perp=0.1\,j$, the width of the source measured along a field line is 1 km. Now thermal electrons at $T_e=2\,10^7$ K with a collision frequency $10^{-2}\,\omega_p=6\,10^6$ s^{-1} have a mean free path of 3 m. They diffuse out of the source by a random walk, involving 10^5 steps, in roughly 10^{-2} s. They cannot simply leave this rapidly, since an unstable return current is created. (This would be an additional source of ion-acoustic waves). Clearly, electrons at $v>3\,v_e$ escape practically freely. We assume that most of the electrons in the source are replaced by a stable return current once during the burst and that most of the escaping electrons have $v\gtrsim3\,v_e$, that is, energies exceeding 10 keV and collisional loss times in the ambient corona exceeding 10^4 s. Compared to the ambient corona, the electrons are "hotter" by a factor $T_h/T_e\approx10^2$, but there is no claim that the energy distribution is in any sense Maxwellian.

The electrons are injected onto a magnetic flux tube with a cross section that depends on the angle between the magnetic fields and the current sheet. With $j_\perp=0.1\,j$, the fluxtube has a crossection of 0.1 (6000 km)$^2=10^{16.5}$ cm^2, or 10^{-3} of the crossection of the emission volume we adopted above. Therefore, if the trapping time of suprathermal electrons exceeds the time needed for some 10^3 type I bursts to occur in the region, i.e. roughly 10^3 s, then suprathermal electrons from any one burst site are injected onto a filled flux tube. For later reference, we also estimate the fraction of trapped relative to thermal electrons, Δ. If all electrons in a type I burst source are replaced once, each burst creates $10^{29.5}$ electrons, and if the trapping time includes 10^3 bursts, then $10^{32.5}$ fast electrons in a volume $10^{29.5}$ cm^3 results in $n_T=10^3$ cm^{-3}, $\Delta=n_T/n=10^{-5}$. Since some electrons are lost, $\Delta=10^{-6}$ will be assumed below for numerical estimates.

6.3. The Trapping of Particles

How are the fast electrons trapped? Since most of the escaping fast electrons are likely to have $v_\parallel>v_\perp$, guiding-center motion would lead to the loss of many of the electrons by collisions in the chromosphere. However, an anomalous-cyclotron instability may develop which scatters the electrons and lets them become trapped (see Papadopoulos and Palmadesso, 1976, and references therein). We follow the work of Vlahos (1979). When the suprathermal tail has a distribution $f(v)\simeq f(v_\perp)$ nearly independent of $v_\parallel$, then the dispersion for electrostatic waves (Krall and Trivelpiece, 1973) can be simplified for lower-hybrid waves, and, requiring that they are not Landau-damped in the ambient corona, one has

$$\omega=\omega_c^e\frac{k_\parallel}{k},\qquad \frac{\omega}{k_\parallel}\gtrsim3\,v_e,\qquad \omega_c^e\equiv\frac{eB}{mc},\tag{35}$$

satisfying a single cyclotron resonance for electrons somewhat slower than their maximum velocity v_M,

$$\omega + \omega_c < v_M k_{\|} \tag{36}$$

to yield a growth rate

$$\gamma \approx \frac{\pi}{4}\left(\frac{k_{\|}}{k}\right)^3 \frac{\omega_c}{k_{\|} v_M} \frac{n_T}{n} \omega_c \approx 10^{-2} \Delta\omega_c. \tag{37}$$

Typical values might be $v_M = 10^{10}$ cm s^{-1}, $k_{\|} = 1.5\ \omega_c/v_M = 10^{-2}$ cm^{-1}, $k/k_{\|} = 4$, and $\omega_p/\omega = 40$. With $B = 3$ g and $\Delta > 10^{-4}$ among the fast electrons that have just escaped, one obtains $\gamma < 10^2$ s^{-1}. It is likely that there is an ambient spectrum of such waves (see below). In one growth time, the waves travel normal to the field by less than the width of the fluxtube containing new electrons (600 km). Since the new electrons take longer than some 10^2 s to pass any one point, they can indeed excite the instability, which will scatter the electrons and trap them.

The trapped electrons involve a loss-cone (and perhaps an anti-loss-cone) instability. They are too slow to generate whistlers, which need energies above roughly 100 keV. They generate the upper-hybrid instability. The saturation of this instability has been evaluated by Benz (1980) in the hydrodynamic approximation. For typical parameters (cf. Berney and Benz, 1978) such as $Y = 0.06$, $k_u v/\omega_c = 25$, $T_h/T_e = 100$ (thus $k_u \lambda_D = 0.25$), Benz arrives at a wave energy density $\beta_u = 10^{-2}\ \Delta = 10^{-8}$ which is permitted by our estimate concerning emission of harmonic radiation, and a trapping time of roughly $5\ 10^{-6}\ \Delta^{-3/2}$ s, or about one hour, which is consistent with the duration of the continuum.

6.4. Radiation Processes

How is the continuum created? It is clear that the waves and particles which so far have entered the picture outside the burst source cannot be responsible for suitable wave-wave interactions. The reason is fairly simple: waves at frequencies near ω_p have k much larger than the k of the available low frequency waves. Thus the combination cannot result in the rather small $\mathbf{k}$ associated with the radio emission at $k_t \lambda_D = Y^{1/2} v_e/c \approx 3\ 10^{-3}$ [cf. Eq. (9)]. Loss-cone driven upperhybrid waves have $k_u \lambda_D \approx 0.25$, and Langmuir waves resonant with fast electrons have $k_L \lambda_D = v_e/v_M > 0.04$. The beam-driven lower-hybrid waves satisfy $k\lambda_D = 1.5\ Y v_e/v_M = 4\ 10^{-3}$. Whistlers also have a low k (Melrose, 1975).

The explanation of the continuum requires two ingredients: fast electrons causing Langmuir waves and low-frequency fluctuations that combine with the Langmuir waves to yield radio emission. Isotropic Maxwellian electrons can yield a brightness temperature no higher than the electron temperature. If the electrons are at about 15 keV, they can cause only $T_b < 2\ 10^{-8}$ K. However, electrons with a "gap" distribution can yield $T_b < 10^{10}$ K ($\simeq mc^2$) and such a distribution may well exist in a trapping region (Melrose, 1980). If the electrons have a loss-cone, still higher T_b might be reached. However, the maximum observed, $T_b = 5\ 10^9$ K, cannot be reached if the Langmuir waves are merely scattered off ion thermal fluctuations because these fluctuations cannot yield an optical depth of order unity in the corona (Melrose, 1980). Thus one is faced either with invoking very intense Langmuir waves or some non-thermal level of low-frequency waves. The former, however, is not possible without exceeding the observational limit on the harmonic. To reach $T_b > 10^9$ K by intense Langmuir waves without high optical depth requires the Langmuir wave temperature to be substantially above 10^{10} K. But even isotropic Maxwellian electrons produce a ratio of harmonic to fundamental of $2\alpha(E_h/mc^2)$ where $\alpha \simeq 6$ (Tidman and Dupree, 1965) and the ratio is larger for anisotropic electrons since these produce anisotropic waves. However, the second and most likely explanation for the continuum is that low-frequency waves with low k decay into other low-frequency waves with a k sufficiently high to combine with the Langmuir waves. Whistlers and lower-hybrid waves may represent pump waves. They stay in the continuum source long enough to have time to decay, and the threshold for decay tends to be quite low (e.g. Berger and Perkins, 1976). The coalescence of lower and upper hybrid modes (the two most likely waves) yields a transverse wave below the x-mode cut-off for $\omega_p/\omega_c > 1$. The radiation is thus fully polarized in the same sense as the burst radiation.

6.5. Relation to Storm Type III Bursts

The number of fast electrons invoked above, $10^{29.5}$, is comparable to the number needed for a (small) storm-type III burst. Probably these electrons do cause a type III burst if the ion-acoustic instability occurs on an open field line, but if the same phenomenon occurs on a closed field line, either a type I burst occurs (sufficient trapped electrons) or nothing is observed. The boundary in frequency between storm-type III and type I bursts corresponds to a height where predominantly closed magnetic field lines tend to change over to open field lines at greater heights (Newkirk, 1974). We expect that newly injected hot electrons in a region with trapped electrons cannot yield a type III burst because the bump-in-tail distribution of the new electrons occurs only after they have traveled over a large distance and, in addition, it would be exceeded by the steep velocity distribution of the trapped electrons.

7. Conclusions

The primary ingredients in our theory are the following: (i) Type I bursts involve a nonthermal phenomenon: association with the evolution of a large sunspot suggests coronal disequilibrium and complex magnetic situations, leading to a localized plasma phenomenon – one dimension of such plasma phenomenon is inevitably very small, of the order of a km or less. (ii) We assume radio emission by wave-wave interaction involving a Langmuir wave, consistent with polarization and frequency width of the bursts. The lack of harmonic emission combined with the high intensity and the required small physical thickness of the source requires that the Langmuir wave be the less intense of the two kinds of waves. (iii) We have proposed ion-acoustic waves as the low-frequency waves because they have a sufficiently high saturation level, they have a short enough decay time compared to observed time structures, and their excitation by electrical currents due to coronal disequilibrium appears plausible. (iv) Nonthermal electrons trapped in the coronal magnetic fields have been evoked to provide the Langmuir waves. Their presence in noise-storm sources is well demonstrated by the emergence of storm type III bursts. In our model they are generated by the sites of type I bursts. The plasma phenomenon at sites of type I bursts and the associated coronal evolution may occur even when there is no radiation signature (insufficient trapped electrons). (v) The same electrons may give rise to the continuum emission. We have found the conversion of the associated Langmuir waves on thermal ion-acoustic fluctuations to be unlikely, but interactions of the Langmuir waves with lower-hybrid or other low-frequencies decay-product waves are an attractive possibility. The presence of upper-hybrid waves at a saturated or marginally stable level is most likely for a loss-cone velocity distribution expected for trapped electrons.

The field is wide open for future theoretical work, particularly for MHD studies of other possible localized current situations, for analysis of fields and particles in regions of a high level of ion-acoustic waves, and for the generation of low-frequency waves for the continuum radiation.

This leads us, finally, to note that improved observations would be useful in distinguishing among some of the theoretical questions. For the problem with the polarization: is there any evidence of less-than-complete polarization for continuum of high flux and/or low source area at frequencies above say 200 MHz. For the trapping: Is a continuum with short (minutes) time variation associated with a high T_b, less than complete polarization, and/or a high rate of type I (or even storm type III) bursts? Is there a minimum level of continuum before type I bursts appear (due to trapped electrons)?

Much of our analysis has depended on the lack of harmonic emission from type I bursts. Would such emission become observable with higher resolution? For intense bursts at $T_b = 10^{10}$ K, the predicted flux ratio (harmonic/fundamental) is 10^{-6} $(1/\delta)$ $(0.2/n_t)^2$, where δ is the reduction due to free-free absorption. For $\delta = 0.1$ and $n_t = 0.2$ this ratio is two orders of magnitude below the currently observed limit. The ratio of harmonic and fundamental brightness temperature as measured by an ideal interferometer (better than 40″ n_t resolution) is $5\ 10^{-4}$ $(1/\delta)$ $(0.2/n_t)^2$, or with the adopted values T_b(harmonic) $= 5\ 10^7$ K. For an angular resolution of 1′ (typical for modern radioheliographs) this value is reduced to an apparent temperature of 10^6 K, which may in certain cases be observable against the continuum.

Appendix

We consider the threshold for the current-driven ion-acoustic instability. The tearing mode has a threshold that is lower than the ion-acoustic current-driven instabilities. However, its main signature is the creation of suprathermal particles, which have already been shown to be insufficient to explain bursts. We expect that the motions creating intense currents (Sect. 2) make the currents sufficiently narrow for current-driven instabilities. For $T_e = T_i$, the electrostatic ion-cyclotron wave has the lower threshold. It saturates at a low level, insufficient by far to account for bursts without requiring Langmuir wave intensities that would yield strong harmonic emission or a huge source volume.

We imply that the driving phenomenon not only causes very narrow current sheets or filaments but also, at some time, heats the electrons so that $T_e \gtrsim 5\,T_i$. One such process which may heat the electrons is the lower-hybrid drift instability associated with the current running normal to the magnetic field (Huba et al., 1978).

A second possibility (due to Spicer, private communication) are tearing modes. The third and most probable process (also due to Spicer) invokes classical resistivity for a gradually increasing current density. When the current-carrying electrons drift at 4 km s^{-1}, the electrons heat more rapidly than the ions (neglecting electron heat conduction). If the coronal disequilibrium then raises the current to the ion-acoustic threshold fast enough, such that the electrons have been heated while the ions are still at their original temperature, then the threshold for the electron drift will be at 300 km s^{-1}. Once the threshold is exceeded, the effect of the anomalous resistance spreads rapidly to trigger the ion-acoustic instability elsewhere. We believe that sufficient arguments have been given for the occurence of the ion-acoustic instability that the incomplete explanation of the triggering at this stage is not vital.

The linear phase of the ion-acoustic instability is well known (Fried and Gould, 1961). Theoretically, one expects saturation at an ion-acoustic wave energy density relative to nT_e (Boltzmann constant = 1) of $\beta_s \simeq 10^{-2}$ (Papadopoulos, 1977). Also according to the laboratory experiments, the energy density due to a collisionless current-driven instability has typically about $\beta_s = 10^{-2}$ (Gekelmann and Stenzel, 1978), but may reach $\beta_s = 0.2$ (Hollenstein, 1979). The angular distribution seems to depend significantly on experimental conditions, as indicated by Illic (1977), Mase and Tsukishima (1975), and Gekelmann and Stenzel (1978), all dealing with parallel currents.

The complex behavior of Langmuir waves in the presence of intense ion-acoustic waves can be seen from the following two extrapolations from better-known situations. If the ion-acoustic spectrum has a small width $\Delta k_s \ll k_L^m$, then there occurs a diffusion in the k_L. For $\Delta k_s/k_L^m = 0.1$, $k_L^m \lambda_D = 0.05$, and $\beta_s = 10^{-2}$, the diffusion time is about 60 Langmuir wave periods (Davidson, 1972). If the k_s and k_L are comparable, the diffusion time is an order of magnitude shorter. Of course, diffusion then is no longer the appropriate approximation, but this indicates that the Langmuir waves are likely to be isotropic and k_L assumes most of the allowed values ($k_L \lambda_D \lesssim 0.3$). Perhaps more disturbing is that the momentum of the electron oscillations is transferred to the ions at a rate of order $\omega_p \beta_s/(k_L^m \lambda_D)^2$ (Dawson, 1968).

Acknowledgements. We wish to thank K. Appert and J. Vaclavik from Centre de Recherche de la Physique des Plasmas, Lausanne, for discussions and reading the manuscript.

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