A PLASMA-EMISSION MECHANISM FOR TYPE I SOLAR RADIO EMISSION

D. B. MELROSE
Department of Theoretical Physics, University of Sydney, N.S.W. 2006, *Australia*

(Received 19 July; in revised form 3 November, 1979)

Abstract. A theory for type I emission is developed based on fundamental plasma emission due to coalescence of Langmuir waves with low-frequency waves. The Langmuir waves are attributed to energetic electrons trapped in a magnetic loop over an active region. It is argued that the low-frequency waves should be generated in connection with the heating of the region. The continuum can be explained in terms of Langmuir waves generated by a 'gap' distribution formed through collisional losses over a timescale of several tens of minutes. Bursts are attributed to local enhancements in the Langmuir turbulence associated with a loss-cone instability. No triggering mechanism for the bursts is identified. It is predicted that if the continuum is due to a 'large' source then its brightness temperature should rise over several tens of minutes to a value which is roughly independent of frequency and of position across the source and which should not exceed 3×10^9 K. For bursts, it is predicted that a fainter second harmonic component should accompany bright bursts.

1. Introduction

Despite the extensive and detailed observational data on type I solar radio emission (e.g. Elgarøy, 1977), no qualitative theory for the type I phenomenon has received widespread support. The essential ingredients in a qualitative theory are the identification of an emission mechanism and of an exciting agency. For type III and type II bursts, plasma emission is identified as the emission mechanism, and a stream of electrons and a shock wave, respectively, as the exciting agencies. The essential features of these qualitative theories were established in the early 1950's (Wild and McCready, 1950; Wild *et al.*, 1954). In contrast, there is no widespread agreement on either the emission mechanism or the exciting agency for type I emission. Existing theories for type I emission are discussed in Section 2. Briefly, the theories either involve plasma emission with type-III-like or type-II-like exciting agencies, or they involve amplified cyclotron emission generated by a stream of electrons. From a qualitative viewpoint, these ideas are unsatisfactory. Type I emission is quite unlike type II or type III emission in many ways. In particular there is no evidence that streaming motions play any essential role, and there is no direct evidence in support of the cyclotron hypothesis. More detailed criticisms of the existing theories are given in Section 2.

In this paper I propose to explore a theory for type I emission based on the following two assumptions:

(1) The emission mechanism is fundamental plasma emission and is due to coalescence of Langmuir waves with low-frequency waves (e.g. ion sound waves or lower-hybrid waves);

Solar Physics **67** (1980) 357–375. 0038/0938/80/0672–0357$02.85.

(2) The exciting agency for the Langmuir waves is a population of energetic particles trapped in a closed magnetic structure over an active region.

The first assumption has two parts. It is widely believed that type I emission is fundamental plasma emission (cf. Section 2). However, if one invokes plasma emission then one needs to identify a reason for the fundamental being observed exclusively in type I emission, in contrast with type II and type III emission which involve both fundamental and second harmonic emission, with a preference for the latter. The coalescence process invoked here leads to a strong preference for fundamental plasma emission. The second part is the coalescence of Langmuir waves and low-frequency waves to produce fundamental plasma radiation. This idea is not new. It was suggested by Sturrock (1965) and has been invoked in various connections by Melrose (1970), Smith (1970), Lacombe and Møller-Pedersen (1971), Melrose and Sy (1971), and Kuijpers (1974, 1975) amongst others. The suggested application to type I emission was proposed by Melrose (1977), and independently by R. van Hees (private communication, 1978; also reported by J. van Nieuwkoop, private communication, 1978). The particular feature emphasized in the present paper is that the coalescence process can proceed sufficiently rapidly for the process to saturate. Saturation occurs provided the level of the low-frequency turbulence is above a minimum level (which is identified) and is such that the effective temperature of the transverse waves becomes equal to that of the Langmuir waves. The quantitative discussion here is based on this equality. Another point emphasized in this paper is that the emission mechanism should be capable of accounting for both continuum and bursts; many of the existing theories cannot account for the continuum.

(Ideas similar to those discussed here have been presented very recently by Benz and Wentzel (1979).)

There are three obvious differences between type I and either type II or type III emission. First is the absence of harmonic structure in type I. This is readily explained in terms of the emission mechanism assumed here. Second is the high degree of polarization (in the sense of the *o*-mode) in type I. High polarization is implied for any fundamental plasma mechanism; the difference in degree of polarization between type I and fundamental type II or type III must be attributed to depolarization of the latter (Melrose, 1975a). Third is the absence of systematic streaming motions in type I. With the second assumption above no streaming motion is necessary.

The second assumption is based on the observational evidence that type I sources are associated with active regions (e.g. Elgarøy, 1977). In general type I emission consists of a continuum and bursts, either or both of which may be present. There is a lower cutoff frequency of type I emission from any given source, and this is usually interpreted as the emission frequency at the top of the trapping region. Many type I sources are known to be near the sites where acceleration of particles occurs, e.g. the type III electrons in type I–III storms must be accelerated in or near the type I source. Consequently, the assumption that the exciting agency involves trapped energetic particles seems a plausible one. A difficulty with it is that it does not account in any obvious way for bursts of emission, and type I bursts are a most important feature of type I emission. The problem of accounting for type I bursts is discussed in Section 6 below.

After the discussion of existing theories and of some relevant observational results in Section 2, the generation of Langmuir waves is discussed in Section 3 and the coalescence with low-frequency waves is discussed in Section 4. A theory for type I continuum is developed in Section 5.

2. Qualitative Discussion

2.1. Existing theories

In the first detailed theory for type I *bursts*, developed by Takakura (1963), it was postulated that the exciting agency is a slow stream of electrons with a speed of several thermal electron speeds. The postulated emission mechanism is scattering by thermal ions (the standard version of fundamental emission in type III and type II theories). The stream of electrons is attributed to local acceleration associated with the collision of two Alfvén waves. Trakhtengerts (1966) pointed out a difficulty in Takakura's acceleration mechanism and proposed a related alternative. Sy (1973) pointed out that induced scattering could be important in the conversion mechanism and that the resulting amplification favours the *o*-mode. Takakura's and Sy's treatments of the emission are, in effect, large-source and small-source models in the sense discussed in Section 2.3 below. Difficulties with the theory include: (i) the size of the acceleration region is incompatible with the size of the emission region; (ii) the efficiency of induced scattering is possibly too low (e.g., Melrose, 1977); (iii) the theory would predict observable second harmonic emission for an acceptable source size (Sections 4.2 and 6.1 below); and (iv) the theory does not account for the continuum at all. Point (i) requires further explanation. In Takakura's theory the volume required for emission is 2×10^{27} cm^3 which must be filled by Langmuir waves excited by electrons moving at a speed 5.5×10^8 cm s^{-1}, determined by the theory. This encounters the difficulty discussed in Section 3.3 below in an extreme form. In Sy's theory this difficulty is not so severe but with his estimated effective temperature $T^t\approx10^{16}$ K at the source one would predict clear harmonic structure (Section 6.1 below).

Both Zaitsev and Formichev (1973) and Vereshkov (1974) invoked exciting agencies which are MHD disturbances. Zaitsev and Formichev's theory is type-II-like and was applied specifically to type I chains. The theory does not account for the continuum, and there is a serious difficulty in accounting for the generation of the Langmuir waves (cf. Smith and Krall, 1974). Vereshkov's theory involves an MHD pulse overtaking a pulse of acoutic waves. This theory encounters the difficulties just mentioned. It is of interest that Vereshkov (1974) invoked coalescence of Langmuir waves and ion-sound waves as the emission mechanism.

The first suggestion that type I emission might be due to cyclotron emission was made by Twiss and Roberts (1958), who immediately dismissed the possibility on the ground that it would produce predominately x-mode radiation. Type I emission is predominantly o-mode, often reaching nearly 100% polarization (Payne-Scott and Little, 1951). Fundamental plasma emission is polarized in the sense of the o-mode (Takakura, 1963; Kai, 1970; Melrose and Sy, 1972; Dulk and Nelson, 1973; Sy, 1973). Fung and Yip (1966a, b) noted that emission at the fundamental gyrofrequency $\omega \approx \Omega_e$ is entirely in the o-mode for $\Omega_e > \omega_p$ and developed a theory for type I emission based on amplified fundamental cyclotron emission. Melrose (1973) pointed out that the second harmonic ($\omega \approx 2\Omega_e$) would grow faster than the fundamental ($\omega \approx \Omega_e$) and would strongly favour the x-mode. Mangeney and Veltri (1976a, b) argued that the o-mode could be produced under special circumstances and developed a detailed theory for type I bursts. The main theoretical objection to cyclotron theories is that they intrinsically favour the x-mode and it seems implausible that the x-mode would never be observed in the resulting emission. From an observational point of view there is strong circumstantial evidence that the emission frequency is related to the plasma frequency ω_p. The starting frequency for type III bursts in type I–III storms is approximately the same as the lowest frequency of type I emission (Malville, 1962; Hanasz, 1966; Stewart and Labrum, 1972). Granted that the type III emission occurs at ω_p or $2\omega_p$, this suggests that the frequency of the type I emission is related to ω_p.

2.2. Observed properties

The observed properties of type I emission have been reviewed by Elgarøy (1977). The properties considered particularly relevant here are: (i) the presence of two distinct components, continuum and bursts; (ii) the brightness temperature; (iii) the bandwidths (relative bandwidths of bursts are a few percent and continuum about 100% in general); (iv) the rise time of ≈ 0.1 s in bursts, which may last from ≈ 0.1 s to ≈ 10 s; and (v) the polarization, which implies fundamental plasma emission. Of these properties only the brightness temperature requires particular comment.

Estimates of the brightness temperature of both the continuum and of individual bursts are particularly important from a theoretical point of view, but very few reliable estimates are available. Stewart (1977, and private communication) estimated a brightness temperature of $\lesssim 10^{10}$ K, based on typical maximum flux values and on typical source sizes. Dulk and Nelson (1973) estimated $T_b \approx 7 \times 10^7$ to 2×10^8 K at 80 MHz. The brightness temperature of type I continuum at 43 MHz in one burst was measured as several times 10^9 K by Suzuki and Gary (1979). Suzuki and Gary (1979, and private communication) compared the brightness temperatures of (a) storm III's, (b) drift pairs, and (c) the type I continuum, all for the strong storm of February 18, 1979, and all at 43 MHz. They found the brightness temperatures of the III's and drift pairs between 2×10^9 K and 2×10^{10} K, with the continuum at 5×10^8 K. The size of the continuum source was considerably larger than the sizes of the III's and drift pairs, although the positions were identical. Recently G. A. Dulk (private communication) has analysed the same storm and found brightness temperatures of 4×10^9 K and 2×10^9 K for the continuum, of 4×10^{10} K and 5×10^{10} K for the brightest type I burst seen, and estimated the sizes to the $1/e$ point of 4.2 arc min and 7.2 arc min, where the numbers refer to 160 MHz and 80 MHz, respectively.

It may be concluded that type I bursts can have $T_b \gtrsim 10^{10}$ K, but as yet no type I continuum has been found to have $T_b > 4\times10^9$ K.

2.3. Source sizes

The apparent sizes of type I sources are two to three minutes of arc at 100 to 200 MHz being somewhat larger at lower frequencies and smaller at higher frequencies (e.g. Elgarøy, 1977, p. 135). A size of 2′ to 3′ corresponds to a linear dimension of 0.6 to .09 $\times 10^5$ km. A significantly smaller source (0$'$.7) was observed by Kerdraon (1979).

The observed source sizes could be the actual size or a scatter image of a much smaller source. It is important to decide whether the source is 'large' (actual size = apparent size), in which case the observed brightness temperature T_b may be equated to the effective temperature T^t of the escaping transverse waves at the source, or whether it is 'small' (actual size $\ll$ apparent size), in which case one has (neglecting free-free absorption)

$$T^t = \frac{T_b}{\eta}, \tag{1}$$

with η = (actual source area)/(apparent source area) $\ll 1$. This question is particularly important here because the coalescence process should result in $T^t = T^l$, where T^l is the effective temperature of the Langmuir waves. (The definition of T^l is such that T^l times Boltzmann's constant integrated over $d^3\mathbf{k}/(2\pi)^3$ is the energy density in the Langmuir waves; usually T^l peaks around some specific **k**-value, and the maximum value of T^l is implied). For a 'large' source one has $T_b \approx T^t = T^l \lesssim 10^{10}$ K: Langmuir waves with $T^l \lesssim 10^{10}$ K could be generated by a trapped distribution of electrons without any instability of the Langmuir waves (Section 3). For a 'small' source one requires $T^l \gg 10^{10}$ K.

There is no argument against a 'large' source for type I continuum. It is assumed in Section 5 below that the continuum is a 'large' source.

Type I bursts are confined to a relative bandwidth $\Delta\omega/\omega_p$ of a few percent. Let $L_N (= n_e\,|\text{grad}\, n_e|^{-1})$ be the characteristic distance over which the mean electron density in the corona changes. Then the emission in a given burst must come from a range of heights $(\Delta\omega/\omega_p)\, L_N/2 \approx 10^3$ km for $L_N \approx 10^5$ km. Thus a 'large' source for a burst is a flat disc $\approx(10^5\text{ km})^2 \times (10^3\text{ km})$. There is a strong argument against such a model. The rise time for a typical burst is ≈ 0.1 s (Elgarøy, 1977, p. 214) and to excite a region of dimensions 10^5 km in <0.1 s requires that the velocity of the exciting agency have a component across the source $>10^6$ km s^{-1}, i.e. $>3c$. Put another way, to excite a plane with cross-section $\approx 10^5$ km in <0.1 s due to a disturbance propagating at the Alfvén speed v_A would require a plane wave front (over $\gtrsim 10^5$ km) which deviates from the plane n_e = constant by no more than ($v_A/10^6$ km s^{-1}), which is $\ll 1°$ for any reasonable v_A. Thus for the bursts we must have $\eta \ll 1$ in (1). Assuming $v_A \lesssim 10^4$ km s^{-1} and a planar disturbance deviating from the plane by no more than ≈ 0.1 radians, one would estimate 10^4 km as the largest possible size of a type I burst. This is of the order of the size observed by Kerdraon (1979). A size $\lesssim 10^3$ km would be more easily compatible with the observed rise time. Then, with $T_b \lesssim 10^{10}$ K and $\eta \approx 10^{-4}$, (1) would imply $T^t \lesssim 10^{14}$ K.

The suggestion that type I sources are 'small' (Steinberg and Caroubalos, 1970; Steinberg *et al.*, 1971) requires that the apparent source be a scatter image, which would imply emission into a wide cone contrary to observations (Bougeret, 1973; Kerdraon, 1973; Steinberg *et al.*, 1974; Elgarøy, 1977, p. 110). Special assumptions seem to be required to account for the apparent size and the directivity simultaneously (Bougeret and Steinberg, 1977). Steinberg (1977) has discussed scattering-image sources and concluded that one can infer nothing useful concerning the actual size from the apparent size.

3. Generation of the Langmuir Waves

As already mentioned, the proposed emission mechanism produces $T^t \approx T^l$. In a large-source ($\eta \approx 1$ in (1)) one requires Langmuir waves with $T^l \lesssim 10^{10}$ K. It is conceivable that such a distribution of Langmuir waves could be maintained in a steady state through emission and absorption. This possibility is examined first. For bursts ($\eta \ll 1$ in (1)) one requires $T^l \gg 10^{10}$ K. The generation of such a distribution of waves is discussed in subsections (c) and (d).

3.1. Steady-state level

A steady-state level of Langmuir turbulence can be maintained when emission and absorption are in balance. Assuming that the emission and absorption is due to energetic electrons (e) and energetic ions (i) in a thermal plasma, the steady-state level is (e.g. Melrose, 1975b)

$$T^l(\mathbf{k}) = \frac{\alpha_e(\mathbf{k}) + \alpha_i(\mathbf{k}) + T_e\{\gamma_c(\mathbf{k}) + \gamma_L(\mathbf{k})\}}{\gamma_e(\mathbf{k}) + \gamma_i(\mathbf{k}) + \gamma_c(\mathbf{k}) + \gamma_L(\mathbf{k})}. \tag{2}$$

In (2) $\alpha_{e,i}(\mathbf{k})$ and $\gamma_{e,i}(\mathbf{k})$ are emission and absorption coefficients, respectively, for the energetic electrons and ions, and $\gamma_c(\mathbf{k})$ and $\gamma_L(\mathbf{k})$ are the collisional absorption coefficient and Landau damping coefficient, respectively, for the thermal electron gas at temperature T_e. Landau damping is negligible in comparison with collisional damping for $k\lambda_{\text{De}} \lesssim \frac{1}{5}$ (Tidman and Dupree, 1965), where $\lambda_{\text{De}} := V_e/\omega_p$ is the Debye length and $V_e := (T_e/m_e)^{1/2}$ is the thermal speed of the electrons. For $\gamma_e(\mathbf{k})$ to exceed $\gamma_c(\mathbf{k})$ in magnitude requires that there be more than about one energetic electron with $v \gtrsim \omega_p/k$ per Debye sphere (Melrose, 1975b). Both $\gamma_e(\mathbf{k})$ and $\gamma_i(\mathbf{k})$ can be negative when the distribution of energetic particle s is anisotropic and is also a 'gap' or 'plateau' distribution. Also (2) applies only if the denominator is positive.

The spectrum was evaluated explicitly by Melrose (1980b, p. 140) for the case of an isotropic power-law distribution

$$f_e(p) = K_e p^{-n}\,. \tag{3}$$

Let us define

$$v_\phi = \omega_p/k\,, \qquad p_\phi = m_e v_\phi (1 - v_\phi^2/c^2)^{-1/2} \tag{4}$$

and v_0, and the corresponding p_0, such that there are less than one electron per Debye sphere for $v > v_0$ or $p > p_0$. The result is

$$T^l = \frac{\alpha_e(\mathbf{k})}{\gamma_e(\mathbf{k})} \approx \begin{cases} \dfrac{m_e v_\phi^2}{n-2}\left[\dfrac{1-(v_\phi/v_0)^{n-2}}{1-(v_\phi/v_0)^n}\right] & \text{for} \quad v_0 \ll c\,, \\[2ex] \dfrac{(n-2)p_\phi c}{n(n-3)}\left[\dfrac{1-(p_\phi/p_0)^{n-3}}{1-(p_\phi/p_0)^{n-2}}\right] & \text{for} \quad p_0 \gg m_e c\,. \end{cases} \tag{5}$$

The maximum value of $T^l(\mathbf{k})$ is for $p_\phi \approx p_0$. For example, for $n = 5$ and $v_0 = c/3$ (5) gives $(T^l)_{\max} \approx 10^8$ K, and for $n = 5$ and $p_0 = 2m_e c$ it gives $(T^l)_{\max} \approx 2 \times 10^9$ K.

It follows that an isotropic power-law distribution of electrons can produce $T^l \approx 10^{10}$ K only if the electrons have energies $\gg$1 MeV, when one would expect them to radiate observable gyro-synchrotron emission. It may be concluded that one cannot account for an observed $T_b \approx 10^{10}$ K in terms of fundamental plasma emission from isotropic energetic electrons.

The value of $T^l(\mathbf{k})$ is enhanced in the presence of energetic ions. If the ions have the same power-law distribution function (3) as the electrons but with normalization coefficient K_i, then the emission and absorption coefficients are related by

$$\alpha_i(\mathbf{k}) = \alpha_e(\mathbf{k})\frac{K_i}{K_e}\,, \qquad \gamma_i(\mathbf{k}) = \gamma_e(\mathbf{k})\frac{m_e}{m_i}\frac{K_i}{K_e}\,, \tag{6}$$

where $v_0 \gg m_e c$ is assumed. Then one has

$$T^l(\mathbf{k}) = \frac{\alpha_e(\mathbf{k}) + \alpha_i(\mathbf{k})}{\gamma_e(\mathbf{k}) + \gamma_i(\mathbf{k})} = \frac{1 + K_i/K_e}{1 + m_e K_i/m_i K_e}\,\frac{\alpha_e(\mathbf{k})}{\alpha_i(\mathbf{k})}\,,$$

i.e. enhancement above the level (5) by the factor $(1 + K_i/K_e)/(1 + m_e K_i/m_i K_e)$ (Melrose, 1975b). To have a significant enhancement requires $K_i \gtrsim K_e$. Let N_e and N_i be the numbers (or number densities) of electrons and ions *above a fixed energy*. Then one has $K_i/K_e = (N_i/N_e)\ (m_e/m_i)^{(n-1)/2}$. In other words K_i/K_e is the rato of the number of ions to the number of electrons *above a fixed speed*. It is implausible that the factor K_i/K_e would be greater than unity. Hence one may conclude that a steady-state $T^l(\mathbf{k}) \gtrsim 10^{10}$ K cannot be achieved under plausible conditions due to isotropic distributions of energetic particles.

3.2. Collisional losses: gap distribution

Coulomb interactions ('collisions') in a trap cause an initial distribution of electrons to flatten at low energies (Takakura and Kai, 1966; Benz and Kuijpers, 1976). The actual form of the change in a magnetic trap (where electrons are lost through scattering into the loss cone in addition to slowing down) is that $f_0(v)$ at $t = 0$ evolves into (Melrose and Brown, 1976)

$$f(v, t) = \left(\frac{v}{v_0(t)}\right)^{2t_E/t_D} f_0(v_0(t)) \tag{7}$$

with

$$v_0(t) = \left(v^3 + \frac{t}{2t_E} V_e^3\right)^{1/3}, \tag{8}$$

and where t_E and t_D are the energy-loss and deflection times, respectively, for an electron with $v = V_e$. Taking $t_D = t_E/2$ (Melrose and Brown, 1976), (5) with (6) implies that $f(v, t)$ is an increasing function of v for $v^3 \lesssim (t/2t_E)\ V_e^3$. Such a distribution has a positive slope in velocity space. If there is less than one particle per Debye sphere in the range $V_e \ll v \ll (t/2t_E)^{1/3} V_e$, then it is a 'gap' distribution (Tidman and Dupree, 1965; Melrose, 1975b). The maximum effective temperature for Langmuir waves with phase speeds in the gap (i.e. for $v_\phi \ll v_0(t)$) is (Robinson, 1977; Melrose, 1980b, p. 142)

$$\kappa T^l(\mathbf{k}) \approx \tfrac{1}{2} m_e c^2 \approx 3 \times 10^9\ \mathrm{K}, \tag{9}$$

where the electrons are assumed isotropic and non-relativistic.

Thus a gap distribution could produce a Langmuir spectrum with an effective temperature of the order required. If the gap distribution forms as implied by (9), then there must be sufficient time for the collisional losses to remove the low-energy electrons.

3.3. Loss-cone anisotropy

To produce $T^l \gg 10^{10}$ K requires an anisotropic distribution of electrons. In a trap particles can be anisotropic due to the presence of the loss cones, and a loss-cone anisotropy is known to lead to growth of Langmuir waves under certain conditions (e.g. Stepanov, 1973; Kuijpers, 1974). However, an important point which has not been raised previously in this connection is whether the scattering is 'weak' or 'strong' (Kennel and Petschek, 1966), i.e. whether the scattering time is less or greater, respectively, than the bounce time in the trap. If the scattering is 'strong' then the loss cone is nearly full and there is essentially no loss-cone anisotropy. Melrose and White (1979) discussed this point in connection with a trap model for solar hard X-ray sources. They found that the scattering is 'weak' only for energies $\gtrsim$1 keV and $\gtrsim$5 keV for two different models. The parameters in type I sources are not radically different from those in these hard X-ray models. Consequently, a loss-cone anisotropy should be present only for electrons with energies of several keV or greater.

Estimates of the growth rate for a loss-cone instability have been made for conditions relevant to the solar corona by Zaitsev and Stepanov (1975), Benz and Kuijpers (1976), Melrose and Stenhouse (1977), Melrose (1977), and Robinson (1978). Effective growth occurs at θ around $\pi/2$ for phase speeds in the gap ($V_e \ll v_\phi \ll v_0$) with a growth rate of the form (Melrose, 1977)

$$\gamma_e(k, \theta) = -\pi \frac{n_1}{n_e} \omega_p G(\theta) , \tag{10}$$

where n_1 is the number density of energetic electrons and $G(\theta)$ is a function which depends on the details of the distribution function. The magnitude of the growth rate (10) is large, in comparison with say the collisional damping rate, and if growth occurs one would expect it to saturate. The energy density in Langmuir waves should then be comparable with the energy density in the trapped electrons.

Several difficulties arise with a theory based on the hypothesis that the Langmuir waves are generated by an instability. The most serious difficulty is that one would expect to see clear harmonic structure. The threshold for saturation of the second harmonic would be exceeded (cf. Section 6.1 below). Also the timescale for saturation of the instability to occur is several tens of growth times, i.e. $\approx 10^{-8}\, n_e/n_1$ s from (10), and this is likely to be much shorter than the rise time ≈ 0.1 s for type I bursts. Finally, there is no ready explanation for the observed brightness temperatures. The saturation value of T^l is uncertain by several orders of magnitude around 10^{16} K. Then $T_b \lesssim 10^{10}$ K in (1) implies $\eta \lesssim 10^{-6}$ and hence the size of the region where the instability satures would need to $\lesssim 100$ km, with an uncertainty of an order of magnitude or so. There is no apparent reason for localization of the instability in a region of such size.

3.4. Marginal stability

A more favourable idea for the generation of the Langmuir waves is that the distribution of electrons is maintained in a marginally stable state. The idea is that $\gamma_e(\mathbf{k})$ is negative and of order $\gamma_c(\mathbf{k})$ so that the denominator in (2) is small, e.g. of order $\gamma_c(\mathbf{k})$.

A marginally stable state can be maintained when two opposite tendencies balance. In the present case collisions tend to increase the growth rate, i.e. to make $\gamma_e(\mathbf{k})$ more negative. This occurs due to the combination of two effects. The first effect is the increase in the size of the loss-cone with increasing speed due to the scattering becoming 'weaker' with increasing speed, cf. Section 3.3 above. Growth requires that the negative contributions to $\gamma_e(\mathbf{k})$ from the anisotropic higher energy particles not be offset by a positive contribution from the less anisotropic lower energy particles. Collisional loss of lower energy particles then favours growth. The tendency to suppress the growth comes from quasilinear relaxation, i.e. the back-reaction of the Langmuir waves on the distribution of particles.

Consider the denominator in (2) at phase speeds well above the thermal speed, such that $\gamma_L(\mathbf{k})$ is negligible. The contribution of any energetic ions is negligible except under extreme circumstances, as discussed above. Thence (2) reduces to

$$T^l(\mathbf{k}) \approx \frac{\alpha_e(\mathbf{k})}{\gamma_e(\mathbf{k}) + \gamma_c(\mathbf{k})}. \tag{2'}$$

Marginal instability then corresponds to $\gamma_e(\mathbf{k})$ of order the collisional damping rate $\gamma_c(\mathbf{k})$. The background level of the Langmuir waves would then be maintained at the value (cf. Melrose, 1975b)

$$\kappa T^l(\mathbf{k}) \approx 2\pi n_1 \lambda_{\mathrm{De}}^3 \frac{m_e v_\phi^3}{v_0} \tag{11}$$

for $V_e \ll v_\phi \lesssim v_0$. For example, for $v_\phi \approx v_0$ and $\frac{1}{2} m_e v_0^2 = 5$ keV at $\omega_p/2\pi = 100$ MHz, (11) gives

$$T^l(\mathbf{k}) \approx \left(\frac{n_1}{n_e}\right) 10^{17}\,\mathrm{K}. \tag{11'}$$

Using (11) one can estimate the timescale over which quasi-linear relaxation occurs. One finds that the rate of quasi-linear relaxation is of the order of $n_1 \lambda_{\mathrm{De}}^3$ times the collisional relaxation rate of the particles. This rate is of the order of γ_c for $n_1 \lambda_{\mathrm{De}}^3 \approx (v/V_e)^3$, which approximate equality should be roughly satisfied in the marginally stable state.

In the marginally stable state one would expect T^l to increase over a characteristic time of order $\gamma_c^{-1} \approx 0.1$ to 1 s over a region of order $v_0 \gamma_c^{-1} \approx 10^3$ to 10^4 km, i.e. over the distance the electrons would propagate in the time γ_c^{-1}. In the discussion in Section 6 it is assumed that type I bursts are due to localized enhancement of T^l to between 10^{10} K and 10^{13} K over such timescales and distances.

4. The Emission Mechanism

4.1. Basic Equations

The conversion of Langmuir waves (l) into transverse waves (t) due to coalescence with low-frequency waves (σ) occurs when the parametric equations

$$\mathbf{k}' \pm \mathbf{k}'' = \mathbf{k}, \qquad \omega^l(\mathbf{k}') \pm \omega^\sigma(\mathbf{k}'') = \omega^t(\mathbf{k}) \tag{12a, b}$$

are satisfied with either sign. The +sign implies 'up-conversion' and the −sign implies 'down-conversion'. The kinetic equations in semi-classical form (e.g. Tsytovich, 1970, p. 89; Melrose, 1980a, p. 173) are

$$\frac{\mathrm{d}N_\pm^t(\mathbf{k})}{\mathrm{d}t} = \int \frac{\mathrm{d}^3\mathbf{k}'}{(2\pi)^3} \int \frac{\mathrm{d}^3\mathbf{k}''}{(2\pi)^3} u_\pm^{tl\sigma}(\mathbf{k}, \mathbf{k}', \mathbf{k}'') [N^l(\mathbf{k}') N^\sigma(\mathbf{k}'') \mp$$
$$\mp N_\pm^t(\mathbf{k}) \{N^l(\mathbf{k}') \pm N^\sigma(\mathbf{k}'')\}] \tag{13}$$

with (cf. Melrose, 1980b, p. 172)

$$u_{\pm}^{tl\sigma}(\mathbf{k},\mathbf{k}',\mathbf{k}'')=\frac{(2\pi)^5}{2}\frac{\hbar e^2}{m_e^2}\frac{\omega^{\sigma}(\mathbf{k}'')}{k''^2V_e^4}R_E^{\sigma}(\mathbf{k}'')\,\delta^3(\mathbf{k}-\mathbf{k}'\mp\mathbf{k}'')\times$$
$$\times\,\delta(\omega^t(\mathbf{k})-\omega^l(\mathbf{k}')\mp\omega^{\sigma}(\mathbf{k}'')),\tag{14}$$

where $R_E^{\sigma}(\mathbf{k}'')$ is the ratio of electric to total energy in the waves (Melrose, 1980a, p. 47).

4.2. Properties of the conversion process

The properties of the conversion processes have been discussed by Melrose (1980c) and may be summarized as follows:

(i) The energy density W^t in the transverse waves builds up with distance l along the ray path according approximately to

$$\frac{\mathrm{d}W^t}{\mathrm{d}l}\approx\sigma_{\mathrm{T}}n_eW^l\left(\frac{T^{\sigma}}{T_e}\right),\tag{15}$$

where σ_{T} is the Thomson cross-section, W^l is the energy density in the Langmuir waves. Equation (15) follows approximately from (13) when only the term $N^l(\mathbf{k}')$ $N^{\sigma}(\mathbf{k}'')$ is retained.

(ii) For $N^{\sigma}(\mathbf{k}'')\ll N^l(\mathbf{k}')$ the down-conversion process amplifies for $N^t(\mathbf{k})>$ $N^{\sigma}(\mathbf{k}'')$. The up-conversion process never amplifies.

(iii) The processes saturate at $N^t(\mathbf{k})\approx N^l(\mathbf{k}')$, i.e. at

$$T^t\approx T^l.\tag{16}$$

(iv) Provided the source is of depth $\gtrsim L_N\omega^{\sigma}/\omega_p$, saturation occurs for

$$\frac{W^{\sigma}}{n_e\kappa T_e}\gtrsim\frac{6\sqrt{3}}{\pi}\frac{V_ec}{\omega_pL_Nv_{\phi}};\tag{17}$$

where W^{σ} is the energy density in the low-frequency waves. The right hand side of (17) is of order 10^{-9} in the solar corona. (These results apply specifically to ion sound waves and to lower-hybrid waves and are slightly modified for other waves.)

(v) The up- and down-conversions lead to transverse waves with frequency

$$\omega_{\pm}\approx\omega_p\pm\omega^{\sigma}.\tag{18}$$

Thus split-line emission results provided $2\omega^{\sigma}$ is less than the frequency spread $\Delta\omega^l(\approx 3(V_e/v_{\phi})^2\,(\Delta v_{\phi}/v_{\phi})\omega_p)$ in the Langmuir waves.

(vi) The emission is polarized in the sense of the o-mode with the degree of polarization the same as for scattering $l\to t$ by thermal ions (cf. Kai, 1970; Melrose and Sy, 1972).

(vii) Second harmonic emission may be described approximately by an equation of the form (15),

$$\frac{\mathrm{d}W^t}{\mathrm{d}l} \approx \sigma_T n_e W^l \left(\frac{T^l}{m_e c^2}\right), \tag{19}$$

where the Langmuir waves are assumed isotropic. The second harmonic saturates at $T^l \approx T^t$. The constitution analogous to (17a) for saturation to occur is $T^l > T_0^l$ with

$$\frac{T_0^l}{m_e c^2} \approx \frac{5\sqrt{3}}{2} \frac{v_\phi^3}{c\omega_p^2 L_N r_0} \left(\frac{v_\phi}{\Delta v_\phi}\right) \left(\frac{4\pi}{\Delta\Omega}\right), \tag{20}$$

where the Langmuir waves are assumed to be confined to a range Δv_ϕ of phase speed and $\Delta\Omega$ of solid angle. Under coronal conditions for $\Delta v_\phi \approx v_\phi$ and with $\Delta\Omega$ in steradians, (20) implies

$$T_0^l \approx \left(\frac{v_\phi}{0.1\mathrm{c}}\right)^3 \frac{10^{15}\,\mathrm{K}}{\Delta\Omega}. \tag{20$'$}$$

4.3. The low-frequency waves

The required low-frequency waves must have wavenumbers nearly equal to those of the Langmuir waves. Suitable waves can be ion-sound waves, lower-hybrid waves, ion-cyclotron waves, ion-Bernstein waves and certain drift waves with large wavenumbers. Hydromagnetic waves and whistlers are not suitable (Melrose, 1975c), except near resonant frequencies. In particular, whistlers near the electron cyclotron resonance have been invoked by Kuijpers (1974, 1975) and by Benz and Kuijpers (1976). (The resonance in the whistler mode is at the electron gyrofrequency for parallel propagation, the lower-hybrid frequency for perpendicular propagation and at intermediate frequencies at intermediate angles.) We cannot probe the solar corona directly and so have no direct evidence on the existence of these waves in solar radio sources. We have: (i) direct evidence for their existence from probes in the solar wind and the magnetosphere; (ii) some indication of their presence from radar echoes from the corona; (iii) indirect evidence of their existence from some solar radio bursts; and (iv) theoretical arguments that they are likely to be present in regions of coronal heating.

(i) Electrostatic turbulence, interpreted as ion-sound turbulence, has been observed in the interplanetary medium from Helios-1 and 2 (Gurnett and Anderson, 1977; Gurnett and Frank, 1978; Gurnett *et al.*, 1979; Gurnett, 1979). The turbulence is bursty and its mean intensity decreases with heliocentric distance. The peak intensity corresponds to $W^s/n_e\kappa T_e \gtrsim 10^{-9}$. The origin of the ion-sound turbulence is not known definitely; the instability associated with heat conduction (Forslund, 1970) is one plausible mechanism. Low-frequency turbulence has been observed in the far magnetotail (Gurnett *et al.*, 1976). The peak relative energy density is higher than in the solar wind ($W^\sigma/n_e\kappa T_e \approx 10^{-5}$). The turbulence has been attributed to a lower-hybrid-drift instability by Huba *et al.* (1978). A third example of observed low frequency turbulence is above the auroral zone (Gurnett and Frank, 1977). It is associated with field-aligned currents and 'inverted V' electron precipitation events. The peak relative energy density is high ($W^\sigma/n_e\kappa T_e \approx 10^{-5}$). The mode has not been identified, but it may be near the resonance in the whistler mode mentioned above (Gurnett and Frank, 1977).

It may be concluded that low-frequency turbulence is common in naturally occurring plasmas and can be excited under a variety of conditions. This adds plausibility to the suggestion that low-frequency turbulence may play a significant role in some solar radio bursts.

(ii) Evidence for the presence of low-frequency turbulence in the solar corona is provided by experiments on radar echoes from the Sun (James, 1966). The interpretation (Gordon, 1973) provides general support for the suggestion that plasma turbulence is present. More recent attempts to detect turbulence above active regions have not yet been successful (Benz and Fitze, 1979).

(iii) Indirect evidence for the presence of low-frequency turbulence in the solar corona is provided by the interpretation of certain radio bursts. An event discussed by Laconbe and Møller-Pedersen (1971) suggested the presence of ion-sound turbulence over a large distance ($\approx 10^5$ km) in the wake of a shock front. The apparent existence of type III bursts in absorption suggests the presence of low-frequency turbulence (Melrose, 1974).

(iv) Theoretical ideas on the heating of the solar corona have moved recently in favour of dissipation of currents in the form of non-potential magnetic fields (the review by Withbroe and Noyes, 1977). Dissipation of currents due to anomalous conductivity in magnetic loops in the corona (Rosner *et al.*, 1978) is now the favoured heating mechanism (Wentzel, 1978).

If the mechanical energy F (erg cm^{-2} s^{-1}) supplied from below is dissipated in a height h due to transfer of energy to the plasma by damping of low-frequency waves at a rate γ, then the energy density in the low-frequency waves is maintained at the value

$$\frac{W^\sigma}{n_e\kappa T_e} = \frac{F}{\gamma h n_e \kappa T_e}. \tag{21}$$

Estimates of F range between about 10^6 and 10^8 erg cm^{-2} s^{-1} (Osterbrock, 1961; Kuperus, 1965; Athay, 1966; Boland *et al.*, 1973; Ulmschneider, 1974). One has $\gamma \approx 10^5$ s^{-1} for ion-sound waves in the corona, and dissipation should occur over a height $h \approx 10^5$ to 10^6 km. These values in (19) lead to $W^\sigma/n_e\kappa T_e \gg 10^{-9}$. Thus heating of corona should result in a level of turbulence in excess of the value implied by (17b).

Further support for the presence of low-frequency turbulence in type I storms comes from recent work on coronal reconnections following coronal transients (Pneuman, 1979). This reconnection is related to post-flare loops and possibly to the storm continuum (the start of a type I storm). A steady reconnection of the form envisaged would provide a steady supply of energy in dissipating currents.

5. Type I Continuum

The simplest model for type I continuum based on the foregoing ideas is as follows: The type I continuum source has an actual size equal to the apparent size. The Langmuir waves are generated by a trapped distribution of energetic electrons; no instability is involved. Conversion into escaping *o*-mode waves is due to coalescence with low-frequency (e.g. ion-sound) waves generated and maintained at a steady level in association with some process in the corona, possibly the heating of the region.

This model is a 'large-source' model. One could consider a 'small-source' model as an alternative. However, in the absence of any evidence to the contrary, the 'large-source'model seems the simpler. Implications of the model include the following.

(i) A trapped distribution of energetic electrons will generate a steady uniform level of Langmuir turbulence with T^l up to about 3×10^9 K. One would not expect continua with brightness temperature greater than 3×10^9 K. (Free-free absorption is neglected here, it will reduce the limiting brightness temperature to $\ll 3\times10^9$ K except for $T_e \gtrsim 10^7$ K.)

(ii) The value of T^l can increase with time following an injection of electrons due to collisional effects (Section 3.2 above). Conversely, a new injection of energetic particles can fill the 'gap' and cause a sudden reduction in T^l. Hence one expects the brightness temperature of the continuum to rise slowly with occassional abrupt decreases associated with new injections. The timescale of the rise is that for producton of a gap at energies $\lesssim$ few keV due to collisions, which timescale is tens of minutes, i.e. timescale for the start of storm continua after flares.

(iii) The trapped distribution of electrons should have roughly the same energy spectrum throughout the trapping region, and hence T^l should be roughly constant across the source. This implies that the brightness temperature should be roughly independent of frequency. More specifically, the brightness temperatures in bipolar regions should be roughly the same in the two sub-sources. (This prediction is not dependent on any assumed symmetry of the flux loop, but like the other predictions it does presuppose that the level of low-frequency turbulence is everywhere above the threshold (17a, b).)

(iv) Alternatively, if the continuum is due to a 'small' source, then one would expect it to fluctuate on a timescale comparable to that of the bursts. Further, the brightness temperature should vary with frequency roughly as does the brightness temperature of bursts.

6. Type I Bursts

6.1. Lower limit to the size of the source

An important question concerning the theory of how type I bursts are generated is the intrinsic size of the emitting region. In Section 2.3 above it was pointed out that the rise time of the bursts implies an upper limit on the size of the emitting region. The absence of a second harmonic implies a lower limit to the size. To see this, consider a burst with apparent brightness temperature T_b whose actual size is a fraction η of its apparent size. Then we have, cf. (1),

$$T^l \approx T_1^t = \frac{T_{1b}}{\eta_1}, \tag{22}$$

where subscripts 1 and 2 refer to the fundamental and second harmonic respectively. In (22) free-free absorption has been neglected. Free-free absorption has a large effect on the escape of fundamental plasma radiation from a smoothly varying corona at a temperature $\approx 10^6$ K. The neglect of free-free absorption here is justified only if the source region is hot (e.g., $\gtrsim 10^7$ K), which is plausible if the region is being heated, or inhomogeneous, e.g. along the lines suggested by Bougeret and Steinberg (1977). To include free-free absorption one includes a factor $e^{-\tau}$ in the definition of η_1, where τ is the optical depth.

Suppose T^l exceeds the values (22). Then the second harmonic saturates at $T_2^t \approx T^l$. However, the fundamental also saturates at $T_1^t \approx T^l$. Consequently, provided a burst is due to a local enhancement in T^l, one has $T_2^t \approx T_1^t$ and hence $T_{2b}/\eta_2 \approx T_{1b}/\eta_1$. On the other hand, for $T^l < T_0^l$, the value of T_2^t is smaller than the saturation value by $(T^l/T_0^l)^2$, cf. (19). Hence one has

$$\frac{T_{2b}/\eta_2}{T_{1b}/\eta_1} = \frac{T_2^t}{T_1^t} = \frac{T^l}{T_0^l}. \tag{23}$$

Then (22) and (23) imply

$$\frac{\eta_1^2}{\eta_2} \approx \frac{T_{1b}}{T_0^l}\left(\frac{T_{1b}}{T_{2b}}\right). \tag{24}$$

Absence of a second harmonic implies T_{2b}/T_{1b} less than some small value, and (24) then implies a lower limit to η_1^2/η_2.

The value of T_0^l implied by (20) or (20′) is uncertain to the extent that it depends on the details of the Langmuir spectrum, i.e. on v_ϕ, Δv_ϕ, and $\Delta\Omega$. For Langmuir waves generated by a loss-cone gap distribution we should have $\Delta v_\phi \approx v_\phi$ in the 'gap', i.e. for $v_\phi \lesssim 0.1c$ for ≈ 5 keV electrons, and the waves should fill a relatively large range of angles about $\theta = \pi/2$, say $\Delta\Omega \approx 1$. Then (20′) implies $T^l \approx 10^{15}$ K. For unresolved sources, which is the case here, η is determined by the resolution of the instrument used and we may set $\eta_1 = \eta_2$, provided again that free-free absorption can be neglected. For a moderately bright burst, $T_{1b} \approx 10^9$ K say, with a much fainter second harmonic, $T_{2b} \lesssim 10^7$ K say, $\eta_1 = \eta_2$ in (24) implies $\eta_1 \gtrsim 10^{-4}$. This implies a linear dimension greater than about 10^3 km. Arguments based on the rise time (Section 2.3) suggest a size less than about 10^3 km.

It is worth emphasizing that even with a size of 10^3 km the second harmonic of bright type I bursts (at 10^9 K to 10^{10} K) should be observable (at 10^7 K to 10^8 K).

6.2. Generation of the bursts

Let us now assume $\eta_1 = 10^{-4}$ and explore the conditions required to account for the brightest bursts implied by current data, i.e. $T_{1b} \approx 10^{10}$ K and hence $T_1^t \approx 10^{14}$ K. Such a burst requires $T^l \approx 10^{14}$ K. Using (iii), such waves could be generated by energetic electrons with $n_1/n_e \approx 10^{-3}$. If these electrons have an energy $\approx$10 keV then their energy density is $\approx$10% of the thermal energy density.

It may be concluded that one can account for $T_{1b} \approx 10^{10}$ K for a source of size $\approx 10^3$ km with Langmuir waves generated by a trapped distribution of energetic electrons without invoking an instability. If one were to invoke an instability, one would expect $T^l \gg 10^{14}$ K and then a definite harmonic structure ($T_{2b} \approx T_{1b}$) would be predicted.

The generation of the bursts must involve an increase in the level of the Langmuir waves. That is, it is not acceptable to attribute a burst to an increase in the level of low-frequency turbulence. This is implicit in the discussion is Section 3.4. The reason involves the second harmonic. To explain a burst with $T_{1b} \approx 10^{10}$ K by an increase in the low-frequency turbulence (from below to above threshold implied by (17a)) would require that T^l be at a steady level of $\approx 10^{14}$ K. However, then (23) implies a steady level of the second harmonic at $T_{2b} \approx 10^9$ K. This second harmonic continuum would be polarized in the sense of the x-mode (e.g. Melrose *et al.*, 1978). Such a continuum has never been observed. It is reasonable to conclude that a burst is due to an increase in the Langmuir turbulence and not in the low-frequency turbulence. However, in the absence of firm observational limits on a second harmonic x-mode continuum, this conclusion is not compelling.

As pointed out in Section 3.4, while there is no obvious trigger for the generation of a burst of Langmuir waves with T^l between 10^{10} K and 10^{14} k, if such a burst were to occur, a timescale of $\approx$0.1 s and a size of $\approx 10^3$ km would be reasonably consistent with theory. In a more detailed theory one would need to account for the triggering of the bursts of Langmuir waves. Chains of type I bursts (implied speeds $\approx v_A$) presumably offer a clue to the triggering mechanism. However, this point will not be pursued further here.

6.3. Frequency splitting

Frequency splitting is observed in some type I bursts. Elgarøy (1961; 1977, p.170) observed splitting of 5 MHz to 12 MHz at 200 MHz with a mean separation of 8 MHz. Sometimes a third component midway between the other two was observed.

These observations find a ready qualitative explanation in terms of the present mechanism. The splitting between the up- and down-conversion processes leads to a line separation of twice the frequency of the low-frequency waves. A third line could be due to the presence of another kind of low-frequency turbulence with an even lower frequency. The two, or three, lines should all be at the same brightness temperature ($T^t = T^l$) and should all be similarly polarized.

If the low-frequency waves are ion-sound waves then the splitting would correspond to

$$\frac{\Delta\omega}{\omega_p} \approx \frac{2}{43} \frac{V_e}{v_\Phi}, \tag{25}$$

where v_Φ is the phase speed of the Langmuir waves. For a reasonable value of v_Φ, say $v_\Phi \approx 10\, V_e$, this splitting is much less than that observed and is likely to be smaller than the intrinsic width determined by the frequency spread in the Langmuir waves, i.e. $3\omega_p(V_e/v_\Phi)^2\,(\Delta v_\Phi/v_\Phi)$. If the low-frequency waves are lower-hybrid waves, then an observed splitting of $\Delta\omega$ implies

$$\frac{\Omega_e}{\omega_p} \approx \frac{43}{2} \frac{\Delta\omega}{\omega_p}. \tag{26}$$

For $\Delta\omega/\omega_p \approx 1/25$, (26) implies $\Omega_e/\omega_p \approx 0.9$. This is a high but not unacceptable value at the 200 MHz level in a loop in the corona.

It may be concluded that: (i) the present mechanism provides a ready qualitative explanation for frequency splitting in type I bursts; (ii) the observed splitting suggests that the low-frequency turbulence is near the lower-hybrid frequency; and (iii) the third component sometimes observed could be explained in terms of the simultaneous presence of lower frequency (e.g. ion-sound) turbulence.

7. Conclusions

The basic idea proposed in this paper is that type I emission is fundamental plasma emission due to coalescence of Langmuir waves and low-frequency waves. The required level of low-frequency turbulence is likely to be present in regions where the corona is being heated, and, in particular, is likely to be present over active regions. Granted the presence of the low-frequency waves, the effective temperature of the transverse waves is maintained equal to that of the Langmuir waves ($T^t = T^l$).

Type I continuum can be explained in terms of emission from a source equal in size to the apparent source size with the Langmuir waves generated by trapped electrons. However, the electrons need to have a 'gap' distribution to achieve $T^l \gtrsim 10^9$ K, and such distribution can be formed through collisional losses on a timescale of several tens of minutes. Specific predictions of this model are made in Section 5.

Type I bursts are not so easily explained. One difficulty concerns the size. The short rise time suggests a size $\lesssim 10^3$ km and the absence of a readily observable second harmonic suggests $\gtrsim 10^3$ km. Assuming a size $\approx 10^3$ km a second harmonic component should be detectable for bright bursts. (The same applies to the continuum if it is attributed to a 'small' source, i.e. if the continuum is assumed to consist of unresolved bursts.) A second difficulty is the absence of a plausible triggering mechanism for bursts. It is plausible that Langmuir waves can build up to the level required in the observed rise time and over $\approx 10^3$ km without invoking an instability, but it is not obvious why such a burst should occur. Generation of

Langmuir waves through an instability is unfavourable because clear harmonic structure should result.

The attractive features of the proposed mechanism are that it accounts automatically for the main features of type I emission, notably the polarization and the absence of clear harmonic structure. It also accounts naturally for frequency splitting in type I bursts. However, the magnitude of the splitting is too large to be explained in terms of ion sound waves and requires lower-hybrid waves. (Ion sound waves or lower-hybrid waves are equally as effective otherwise.) The main implications of the theory are limits on the brightness temperature of the continuum and on the size and harmonic structure in the bursts.

Acknowledgement

I would like to thank Dr G. A. Dulk for helpful comments on the manuscript.

References

Athay, R. G.: 1966, *Astrophys. J.* **146**, 223.
Benz, A. O. and Fitze, H. R.: 1979, *Astron. Astrophys.* **76**, 354.
Benz, A. O. and Kuijpers, J.: 1976, *Solar Phys.* **46**, 275.
Benz, A. O. and Wentzel, D. G.: 1979, 'Solar Type I Radio Bursts: An Ion Acoustic Model', preprint.
Boland, B. C., Engstrom, S. F. T., Jones, B. B., and Wilson, R.: 1973, *Astron. Astrophys.* **22**, 161.
Bougeret, J. L.: 1973, *Astron. Astrophys.* **24**, 53.
Bougeret, J. L. and Steinberg, J. L.: 1977, *Astron. Astrophys.* **61**, 777.
Dulk, G. A. and Nelson, G. J.: 1973, *Proc. Astron. Soc. Australia* **2**, 211.
Elgarøy, Ø.: 1961, *Astrophys. Norv.* **7**, 123.
Elgarøy, Ø.: *Solar Noise Stroms*, Pergamon Press, Oxford.
Forslund, D. W.: 1970, *J. Geophys. Res.* **75**, 17.
Fung, P. C. W. and Yip, W. K.: 1966a, *Australian J. Phys.* **19**, 759.
Fung, P. C. W. and Yip, W. K.: 1966b, *Planetary Space Sci.* **14**, 1139.
Gordon, I. M.: 1973, *Space Sci. Rev.* **15**, 157.
Gurnett, D. A.: 1979, in H. Rosenbauer (ed.), *Solar Wind* 4, Springer Verlag, New York-Heidelberg-Berlin, (in press).
Gurnett, D. A. and Anderson, R. R.: 1977, *J. Geophys. Res.* **82**, 632.
Gurnett, D. A. and Frank, L. A.: 1977, *J. Geophys. Res.* **82**, 1031.
Gurnett, D. A. and Frank, L. A.: 1978, *J. Geophys.* **83**, 58.
Gurnett, D. A., Frank, L. A., and Lepping, R. P.: 1976, *J. Geophys. Res.* **81**, 6059.
Gurnett, D. A., Marsch, E., Pilipp, W., Schwenn, R., and Rosenbauer, H.: 1979, *J. Geophys. Res.* **84**, 2029.
Hanasz, J.: 1966, *Australian J. Phys.* **19**, 653.
Huba, J. D., Gladd, N. T., and Papadopoulos, K.: 1978, *J. Geophys. Res.* **83**, 5217.
James, J. C.: 1966, *Astrophys. J.* **146**, 356.
Kai, K.: 1970, *Solar Phys.* **11**, 456.
Kennel, C. F. and Petschek, H. E.: 1966, *J. Geophys. Res.* **71**, 1.
Kerdraon, A.: 1973, *Astron. Astrophys.* **27**, 361.
Kerdraon, A.: 1979, *Astron. Astrophys.* **71**, 266.
Kuijpers, J.: 1974, *Solar Phys.* **36**, 157.
Kuijpers, J.: 1975, *Solar Phys.* **44**, 173.
Kuperus, M.: 1965, *The Transfer of Mechanical Energy in the Sun and the Heating of the Corona*, D. Reidel Publ. Co., Dordrecht, Holland.

Lacombe, C. and Møller-Pedersen, B.: 1971, *Astron. Astrophys.* **5**, 406.
Mangeney, A. and Veltri, P.: 1976a, *Astron. Astrophys.* **47**, 165.
Mangeney, A. and Veltri, P.: 1976b, *Astron. Astrophys.* **47**, 181.
Malville, J. M.: 1962, *Astrophys. J.* **136**, 266.
Melrose, D. B.: 1970, *Australian J. Phys.* **23**, 885.
Melrose, D. B.: 1973, *Australian J. Phys.* **26**, 229.
Melrose, D. B.: 1974, *Australian J. Phys.* **27**, 271.
Melrose, D. B.: 1975a, *Solar Phys.* **43**, 79.
Melrose, D. B.: 1975b, *Solar Phys.* **43**, 211.
Melrose, D. B.: 1975c, *Australian J. Phys.* **28**, 101.
Melrose, D. B.: 1977, *Izv. Vyssh. Uchebn. Zaved. Radiofiz.* **20** 1369; *Radiophysics and Quantum Electronics* **20**, 945.
Melrose, D. B.: 1980a, *Plasma Astrophysics*, Volume 1, Gordon and Breach, New York.
Melrose, D. B.: 1980b, *Plasma Astrophysics*, Volume 2, Gordon and Breach, New York.
Melrose, D. B.: 1980c, *Space Sci. Rev.* **26**, 3.
Melrose, D. B., and Brown, J. C.: 1976, *Monthly Notices Roy. Astron. Soc.* **176**, 15.
Melrose, D. B. and Stenhouse, J. E.: 1977, *Australian J. Phys.* **30**, 481.
Melrose, D. B. and Stenhouse, J. E.: 1979, *Astron. Astrophys.* **73**, 151.
Melrose, D. B. and Sy, W. N.: 1962, *Australian J. Phys.* **25**, 387.
Melrose, D. B., Dulk, G. A., and Smerd, S. F.: 1978, *Astron. Astrophys.* **66**, 315.
Osterbrock, D. E.: 1961, *Astrophys. J.* **134**, 347.
Payne-Scott, R. and Little, A. G.: 1951, *Australian J. Sci. Res.* **A4**, 508.
Pneuman, G. W.: 1979, in E. Jensen *et al.* (eds.), 'Physics of Solar Prominences', *IAU Symp.* **44**, 281.
Robinson, R. D.: 1977, Dissertation, University of Colorado.
Robinson, R. D.: 1978, *Astrophys. J.* **222**, 696.
Rosner, R., Golub, L., Coppi, B., and Vaiana, G. S.: 1978, *Astrophys. J.* **222**, 317.
Smith, D. F.: 1970, *Adv. Astron. Astrophys.* **7**, 147.
Smith, D. F. and Krall, N. A.: 1974, *Astrophys. J.* **194**, L163.
Steinberg, J. L.: 1977, in M. A. Shea *et al.* (eds.), '*Study of Travelling Interplanetary Phenomena*/1977', D. Reidel Publ. Co., Dordrecht, Holland, p. 65.
Steinberg, J. L. and Caroublas, C.: 1970, *Astron. Astrophys.* **9**, 329.
Steinberg, J. L., Aubier-Giraud, M., Leblanc, Y., and Boischot, A.: 1971, *Astron. Astrophys.* **10**, 362.
Steinberg, J. L., Caroubalos, C., and Bougeret, J. L.: 1974, *Astron. Astrophys.* **37**, 109.
Stepanov, A. V.: 1973, *Astron. Zh.* **50**, 1243; *Soviet Astron.-AJ* **17**, 781.
Stewart, R. T.: 1977, *Izv. Vyssh. Uchebn. Zaved. Radiofiz.* **20**, 1338; *Radiophysics and Quantum Electronics* **20**, 922.
Stewart, R. T. and Labrum, N. R.: 1972, *Solar Phys.* **27**, 192.
Sturrock, P. A.: 1965, *Phys. Fluids* **8**, 281.
Suzuki, S. and Gary, D.: 1979, *Proc. Astron. Soc. Australia* **3**, 379.
Sy, W. N.: 1973, *Proc. Astron. Soc. Australia* **2**, 215.
Takakura, T.: 1963, *Publ. Astron. Soc. Japan* **15**, 462.
Takakura, T. and Kai, K.: 1966 *Publ. Astron. Soc. Japan* **18**, 57.
Tidman, D. A. and Dupree, T. H.: 1965, *Phys. Fluids* **8**, 1860.
Trakhtengerts, V. Yu.: 1966, *Astron. Zh.* **43**, 436; *Soviet Astron.-AJ* **10**, 281.
Tsytovich, V. N.: 1970, 'Nonlinear Effects in Plasma,' Plenum, New York.
Twiss, R. Q. and Roberts, J. A.: 1958, *Australian J. Phys.* **11**, 424.
Ulmschneider, P.: 1974, *Solar Phys.* **39**, 327.
Vereshkov, G. M.: 1974, *Astron. Zh.* **51**, 261; *Soviet Astron.-AJ* **18**, 152.
Wentzel, D. G.: 1978, *Rev. Geophys. Space Phys.* **16**, 757.
Wild, J. P. and McCready, L. L.: 1950, *Australian J. Sci. Res.* **A3**, 387.
Wild, J. P., Roberts, J. A., and Murray, J. D.; 1954, *Nature* **173**, 532.
Withbroe, G. L. and Noyes, R. W.: 1977, *Ann. Rev. Astron. Astrophys.* **15**, 363.
Zaitsev, V. V. and Fomichev, V. V.: 1973, *Astron. Zh.* **49**, 817; *Soviet Astron.-AJ* **16**, 666.
Zaitsev, V. V. and Stepanov, A. V.: 1975, *Astron. Astrophys.* **45**, 135.