PLASMA EMISSION: A REVIEW*

(*Invited Review* presented by G. Dulk)

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Abstract. The theory of plasma emission is reviewed emphasizing general concepts rather than the details of the analysis. The generation of the Langmuir turbulence, its evolution due to nonlinear processes, and the plasma emission processes are described. Several outstanding difficulties in the applications to solar radio bursts are discussed, concentrating on those with implications on local density inhomogeneities in the sources.

1. Introduction

Plasma emission is the generation of escaping radiation near the local plasma frequency or its harmonics. Most meter-wave solar radio bursts are interpreted in terms of plasma emission at the fundamental (F) or second harmonic (H) (e.g., McLean and Labrum, 1985). Plasma emission is also observed ahead of the Earth's bow shock and of other interplanetary shocks. Plasma emission is necessarily a multistage process, unlike the direct emission processes such as cyclotron and synchrotron radiation. These stages include the generation of Langmuir turbulence, its nonlinear evolution and its partial conversion into escaping radiation. There are differences in detail in various treatments depending, for example, on whether or not phase-coherent or narrow-band effects are important, and whether or not strong turbulence effects are important. The simplest viewpoint is that these effects are not important, and this is the viewpoint adopted here. The generation of the Langmuir waves is then attributed to an instability involving negative absorption and saturation due to quasi-linear effects, and the evolution of the Langmuir turbulence is attributed to three-wave processes whose saturation may be discussed within the framework of weak-turbulence theory. The main emphasis in this review is in presenting the general ideas underlying this form of plasma emission. The phase-coherent and strong-turbulence effects are discussed only briefly.

Plasma emission remains at best a semi-quantitative theory. There have been some notable successes for the theory, e.g., in explaining the harmonic structure and the predominantly o-mode polarization, but there remain many difficulties. In discussing the applications here, it is argued that several difficulties can be overcome or alleviated by postulating that the sources are locally inhomogeneous. It seems that it is essential to take local inhomogeneities into account in order to reach a reasonable agreement between theory and observation.

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The review is set out as follows. In Section 2 the generation of Langmuir turbulence is discussed, Section 3 contains an overview of nonlinear processes involving Langmuir and ion sound turbulence, and in Section 4 specific plasma emission mechanisms are summarized. In Section 5 some of the outstanding problems concerning the applications of plasma emission are discussed: these are the observed harmonic ratio being less than two, the lack of agreement between theory and observation on the polarization, the apparent source sizes and directivities, and the clumpiness of Langmuir waves in the solar wind.

2. Generation of Langmuir Turbulence

The first step in any plasma emission mechanism is the generation of Langmuir turbulence.

2.1. Plasma turbulence

A nonthermal distribution of waves in any wave mode constitutes plasma turbulence in that mode. In plasma emission the magnetic field plays only a minor role, and considerable simplification results by neglecting it. (Magnetic effects need to be included when discussing the polarization and they may be important in causing certain fine structures.) Once the magnetic field is neglected, only three waves need to be considered. These are transverse (T) waves with dispersion relation $\omega = (\omega_p^2 + k^2c^2)^{1/2}$, or refractive index $n = kc/\omega = (1 - \omega_p^2/\omega^2)^{1/2}$, Langmuir ($L$) waves with dispersion relation $\omega \cong \omega_p + 3k^2 V_e^2/2\omega_p$ and ion sound (S) waves with dispersion relation $\omega = kv_s$, where the ion sound speed v_s is related to the thermal speed of electrons by $v_s \cong V_e/43$. Only transverse waves (with $\omega > \omega_p$) can escape, and an emission mechanism may be defined as a process which produces transverse waves.

A thermal distribution of waves in mode M corresponds to an energy density $W_M(\mathbf{k})\, \mathrm{d}^3\mathbf{k}/(2\pi)^3$ in the range $\mathrm{d}^3\mathbf{k}$ of wavevectors such that $W_M(\mathbf{k})$ is equal to the thermal energy. It is convenient to set Boltzmann's constant esual to unity so that temperatures are expressed in energy units. Then a thermal level of waves corresponds to $W_M(\mathbf{k}) = T_e$, where T_e is the electron temperature. It is convenient to write $W_M(\mathbf{k}) = T_M(\mathbf{k})$ and to interpret $T_M(\mathbf{k})$ as the effective temperature of the turbulence. The turbulence may be anisotropic, implying that $T_M(\mathbf{k})$ is a strong function of $\mathbf{k}$. The evolution of the turbulent spectrum involves transfer of energy from one region of $\mathbf{k}$-space to another, say from $\mathbf{k}$ to $\mathbf{k}'$ with $T_M(\mathbf{k}') < T_M(\mathbf{k})$.

2.2. Bump-in-tail instability

The most important mechanism for the generation of Langmuir turbulence in solar radio bursts, particularly type III bursts, is a streaming instability. There are two versions of many instabilities and for a streaming instability these two versions are often called the bump-in-tail and the weak-beam instability, respectively. More generally these two versions may be called resistive and reactive instabilities, respectively, e.g., Melrose (1986a).

A resistive instability may be described in terms of negative absorption and attributed to a maser action. The saturation of a resistive instability is due to quasilinear relaxation when the particle distribution function is modified by the back-reaction of the waves on the particles as the free energy driving the instability drains away. For any resistive instability involving Langmuir waves and nonrelativistic electrons, the absorption coefficient $\gamma(\mathbf{k})$ is given by

$$\gamma(\mathbf{k}) = -\frac{\pi e^2 v_\phi^2}{\varepsilon_0 m \omega_p} \frac{\mathrm{d}F(v_\phi)}{\mathrm{d}v_\phi}, \quad (1)$$

with $v_\phi = \omega/k$, $\omega \cong \omega_p + 3k^2 V_e^2/2\omega_p$ and where $F(v_x)$ is the distribution function integrated over v_y and v_z, with the x-axis along $\mathbf{k}$ here. For semi-quantitative purposes (1) may be approximated by $\gamma \cong -\pi\omega_p(n_b/n_e)(v_b/\Delta v)^2$ for a stream with number density n_b, mean streaming velocity v_b and velocity spread Δv. Data on type III streams in the solar wind have been used to show that growth should occur according to (1) for the actual electron streams observed (Lin *et al.*, 1981).

Quasi-linear relaxation of the bump-in-tail instability involves formation of a plateau and its gradual spread to lower velocities, as illustrated in Figure 1. For type III streams advection is important; advection involves faster particles overtaking slower particles

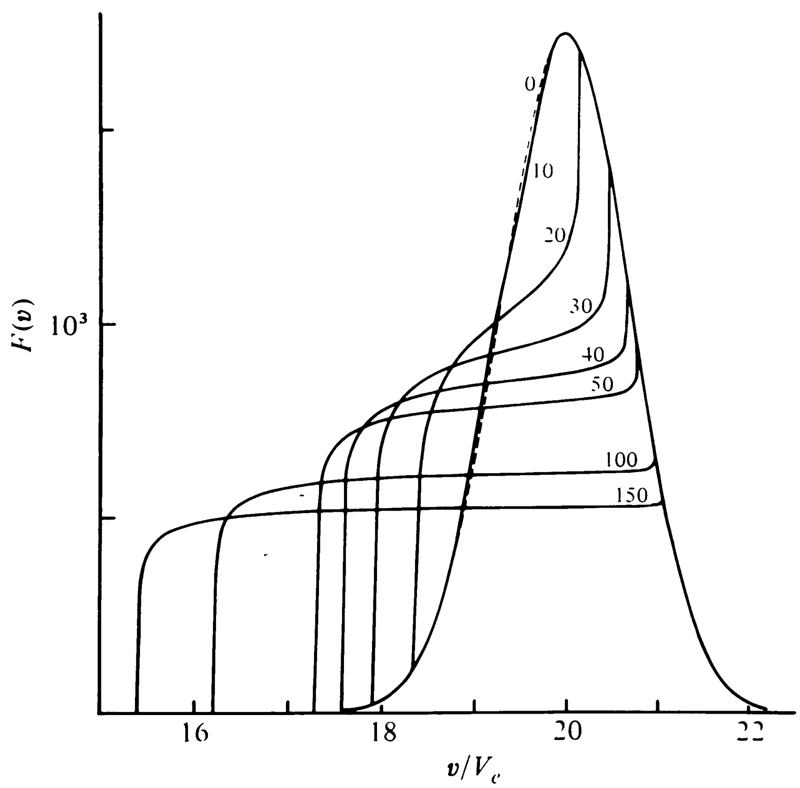

Fig. 1. The evolution of an initial gaussian beam centered at $v = 20V_e$ is illustrated. The labels on the curves refer to the number of (initial) growth times which has elapsed. Note the formation of a plateau which moves to lower velocities due to quasilinear relaxation (Grognard, 1975).

thereby tending to reform a region of positive slope $dF(v_\phi)/dv_\phi > 0$ in opposition to quasilinear relaxation tending to decrease this slope. A steady state with very specific properties is possible in principle due to a balance between these two effects. It is therefore highly significant that measured distributions of type III electrons in the solar wind exhibit these specific properties, as shown by Grognard (1982) by inserting the observed distributions in a numerical code that follows the evolution, and showing that essentially the same distribution emerges as the self-consistent solution. Evidently other effects are unimportant in influencing the evolution.

2.3. Weak-beam instability

The weak-beam instability is a reactive version of the bump-in-tail instability. This instability was first recognized by Bohm and Gross (1949), and it was invoked in Ginzburg and Zheleznyakov's (1958) original treatment of the plasma emission mechanisms. The weak-beam instability applies when the velocity spread is small. It is attributed to phase bunching of the streaming electrons. The growth rate is given approximately by $|\gamma| \cong \omega_p(\pi n_b/n_e)^{1/3}$; this growth rate applies for $\Delta v/v_b \lesssim (\pi n_b/n_e)^{1/3}$, with the bump-in-tail instability applying when this inequality is reversed. By way of illustration, parameters reported by Lin *et al.* (1986) for a type III event in the solar wind imply $\Delta v/v_b \cong 10^{-1}$ and $(n_b/n_e)^{1/3} > 10^{-2}$.

The weak-beam instability, like other reactive instabilities saturates due to wave trapping. Wave trapping is important only for phase-coherent waves; reactive instabilities involve phase-coherent processes, and resistive instabilities involve phase-random processes. The oscillatory or bouncing motion of the trapped particles can lead to the development of side bands, and this effect has been invoked in at least one theory for fine structures in solar radio bursts (Smith and de la Noë, 1976).

2.4. Alternative sources of Langmuir turbulence

Other proposed mechanisms for generation of Langmuir turbulence include the following.

(1) An isotropic *gap distribution* has no particles in a range $V_e \ll v \ll v_0$, and then the particles at $v \gtrsim v_0$ emit Langmuir waves with phase speeds in the gap but do not absorb them (Tidman and Dupree, 1965; Melrose, 1980b). A relativistic effect in the absorption (Robinson, 1978) limits the emission to $T_L(\mathbf{k}) \lesssim 3 \times 10^9$ K for phase speeds in the gap.

(2) A loss-cone distribution with a gap can form naturally (eg., Melrose and Brown, 1976) and then Langmuir waves can grow due to a loss cone instability, cf. Hewitt and Melrose (1985) for a review of the literature. Such a mechanism is favoured for type I emission (e.g., Melrose, 1980c; Benz and Wentzel, 1981).

(3) Recently, Winglee and Dulk (1986) have proposed an electron cyclotron mechanism (driven by $\partial f/\partial v_\perp > 0$) for generating Langmuir (or upper hybrid) waves for type V emission.

(4) Also, recently Cairns (1986) has shown that in a region (the foreshock) ahead of a curved shock an electron distribution with $\partial f/\partial v_\parallel > 0$ develops naturally, and may account for the continuum (non-herringbone) type II emission.

(5) A further alternative but controversial mechanism is turbulent bremsstrahlung (Tsytovich *et al.*, 1975), which I claim does not exist (Kuijpers and Melrose, 1985).

3. Nonlinear Evolution of Langmuir Turbulence

Nonlinear effects modify a spectrum of Langmuir turbulence by transferring energy from one range of wavevectors **k** to another. Relevant processes include three-wave interactions involving ion sound waves, scattering off thermal particles, scattering off density fluctuations, parametric instabilities and modulational instabilities.

3.1. Kinematic restrictions on three-wave interactions

Three waves, described by ω_i, $\mathbf{k}_i$ with $i = 1$, 2, and 3, can beat together when the conditions

$$\omega_1 + \omega_2 = \omega_3\,, \qquad \mathbf{k}_1 + \mathbf{k}_2 = \mathbf{k}_3\,, \tag{2a, b}$$

are satisfied. If we denote Langmuir waves, transverse waves and ion sound waves by L, T, and S, respectively, then the following processes are allowed: $L + S \to L$, $L + S \to T$, $T + S \to L$, $T + S \to T$, and $L + L \to T$. The first process is of interest here, the next two are of interest for fundamental plasma emission, the fourth is of interest for scattering of transverse waves and the last is of interest for second harmonic plasma emission.

Suppose that in (2a) one sets $\omega_1 = \omega_p + 3k^2 V_e^2/2\omega_p$ for an initial Langmuir wave, $\omega_3 = \omega_p + 3k'^2 V_e^2/2\omega_p$ for a final Langmuir wave and $\omega_2 = \pm k_s v_s$ for an ion sound wave, and that (2b) is then replaced by $\mathbf{k} + \mathbf{k}_s = \mathbf{k}'$. The following kinematic restrictions are implied:

$$k^2 - k'^2 = \mp k_0 k_s\,, \qquad \cos\theta = \frac{k_0 \mp k_s}{2k}\,, \tag{3a, b}$$

with

$$k_0 = \frac{2\omega_p v_s}{3V_e^2} = \frac{\omega_p}{65V_e}\,, \tag{4}$$

and where θ is the angle between **k** and $\mathbf{k}_s$.

In applications to solar radio events the inequality $k > k_0$ usually applies. Then the three-wave interactions are possible for k_s ranging from arbitrarily small k_s to $k_s \cong 2k$. Interactions with $k_s \ll k$ cause a small change in the direction and magnitude of the wave vector **k**, and this effect may be described as a diffusion in **k** space (cf. Section 3.3). Interactions with $k_s \cong 2k$ lead to a backscatter with $\mathbf{k}'$ approximately antiparallel to **k** and with $k' - k \cong \pm k_0$. Interactions in the intermediate range $k_s \cong k$ lead to $\mathbf{k}'$ at an oblique angle to **k**.

3.2. Saturation of growth in three-wave interactions

Examination of the kinetic equations for a three-wave process leads to two simple conclusions (e.g., Melrose, 1980b, d, 1986a). Let $N_M(\mathbf{k}) = T_M(\mathbf{k})/\omega_M(\mathbf{k})$ be the wave action for any wave mode M. One conclusion is that in a coalescence process $A + B \to C$ in which waves in mode C are generated, the process saturates at

$$N_C(\mathbf{k}_C) = \frac{N_A(\mathbf{k}_A)N_B(\mathbf{k}_B)}{N_A(\mathbf{k}_A) + N_B(\mathbf{k}_B)}\,, \tag{5}$$

due to the absorption-like effect of the inverse process $C \to A + B$. This applies to the process $L + S \to L'$, which saturates at

$$T'_L(\mathbf{k}') = \frac{T_L(\mathbf{k})T_s(k_s)}{T_s(\mathbf{k}_s) + T_L(\mathbf{k})\omega_s(\mathbf{k}_s)/\omega_p}\,. \tag{6}$$

The other conclusion is that if modes A or B are generated by the decay process $C \to A + B$ then when the limit (5) implied by (2) on N_A or N_B is exceeded the absorption-like process is negative, implying growth of the waves A and B. This can occur for the secondary Langmuir waves L' in the decay process $L \to L' + S$. This kind of nonlinear wave growth is an example of a parametric instability.

3.3. Scattering processes

The conversion of Langmuir waves at one $\mathbf{k}$ into another $\mathbf{k}'$ is called scattering when it is attributed either to thermal particles or to density fluctuation. Both these can be treated as three-wave processes, with the ion-sound wave replaced by the relevant fluctuation.

For scattering off thermal ions the fluctuations have $k'' \gtrsim 1/\lambda_D = \omega_p/V_e$ and $\omega'' \lesssim k''V_i$. For $k'' = 1/\lambda_D$ and $\omega'' = k''V_i$ the relations (2a, b) apply with $k_0 = 2\omega_p V_i/3V_e^2$. The decay $L \to L' + F$, where F denotes the fluctuation ω'', $\mathbf{k}''$, leads to an instability called induced scattering (e.g., Kaplan and Tsytovich, 1968). This growth occurs only when a threshold analogous to (6) is exceeded, and the specific limit $T'_L \gtrsim T_i v_\phi/V_i$ corresponds to $T'_L \gtrsim$ few $\times\ 10^9$ K for type III bursts. The growth rate for this induced scattering instability is almost identical to that for the parametric decay $L \to L' + S$ mentioned above. For semiquantitative one may regard induced scattering off ions as being included in the parametric decay process $L \to L' + S$, and ignore it.

Scattering off density fluctuations may be treated as a limit of the three-wave interaction in which the low-frequency waves have negligible frequency and are described by a spectrum in k''. Suppose the fluctuations are described by an r.m.s. level δn_e and a mean wave number $\bar{k}''$. One may relate δn_e to $T_s(\mathbf{k}'') = k''v_s N_s(\mathbf{k}'')$ by writing

$$\int \frac{d^3\mathbf{k}''}{(2\pi)^3}\, T_s(\mathbf{k}'') = \left(\frac{\delta n_e}{n_e}\right)^2 n_e T_e\,, \tag{7}$$

$$D_\theta \cong \frac{\pi}{12}\, \frac{\omega_p}{(k\lambda_D)^2}\, \frac{k''}{k} \left(\frac{\delta n_e}{n_e}\right)^2. \tag{8}$$

This diffusion can suppress the bump-in-tail instability by removing the Langmuir waves from the range $\Delta\theta$ of forward angle where they grow to angles θ where they damp: suppression occurs when the scattering rate $\cong D_\theta(\Delta\theta)^2$ exceeds the growth rate of the bump-in-tail instability. Using observational data on type III events Muschietti *et al.* (1985) estimated that this suppression should be effective, thereby posing a serious problem for the conventional theory of type III bursts.

3.5. Parametric instabilities, modulational instabilities, and strong turbulence

Besides the processes mentioned so far, there is a rich and confusing variety of other nonlinear processes. The following is a brief summary of some relevant points.

(1) Parametric instabilities have both resistive and reactive versions (e.g., Melrose, 1986b).

(2) The resistive version of the parametric three-wave decay is the process most likely to be relevant in solar radio bursts. It has a nonlinear growth rate

$$\gamma_{NL} \cong \frac{\omega_p \omega_s}{4\Delta\omega} \frac{W_L}{n_e T_e}, \tag{9}$$

with $k_s \cong 2k_L$, and where the Langmuir waves have an energy density W_L over a frequency range $\Delta\omega \cong 3\omega_p \Delta v_\phi V_e^2/v_\phi^3$. A condition on the growth is that γ_{NL} exceed the damping rate $\cong \omega_s/50$ of the ion sound waves.

(3) Induced scattering off ions need not be distinguished from this decay process for semi-quantitative purposes.

(4) The parametric decay instability usually discussed in the literature (e.g., Liu and Kaw, 1976) is a reactive counterpart to this decay instability (Melrose, 1986b). The reactive version applies for $\Delta\omega < \gamma_{NL}$ and has a growth rate $\gamma_{PD} \cong [\gamma_{NL}\Delta\omega]^{1/2}$.

(5) Modulational instabilities are due to the ponderomotive force (e.g., Goldman, 1983, 1984). They cause a Langmuir wave packet to shrink and hence transfer some of the energy to larger k. The most familiar such instability is OTSI (e.g., Papadopoulos *et al.*, 1974, Bardwell and Goldman, 1976).

(6) Langmuir collapse occurs when this shrinking becomes dynamically unstable and the wave packet shrinks to much smaller dimensions. Collapse has yet to be explored fully other than in one dimension (e.g., Goldman, 1984). A relevant point is that even if collapse does occur, the transfer of the energy to larger k is not necessarily more important for the generation of plasma emission than the transfer to smaller k, which necessarily continues even during collapse.

4. Plasma Emission Mechanisms

The most likely forms for fundamental (F) and second harmonic (H) plasma emission involve three-wave processes. These are the only processes discussed here. Alternative forms of plasma emission were discussed by Melrose (1980b, d).

4.1. Fundamental emission

Fundamental plasma emission can occur due to the three-wave processes $L + S \rightarrow T$ and $L \rightarrow T + S$. In the presence of suitable ion sound turbulence both processes saturate at $T_T(\mathbf{k}_T) = T_L(\mathbf{k}_L)$. The requirements on the ion sound turbulence are twofold. First, the wave vectors $\mathbf{k}_s$ must satisfy $\mathbf{k}_L \pm \mathbf{k}_s = \mathbf{k}_T$, and $k_T \ll k_L$ then implies $\mathbf{k}_s \cong \mp \mathbf{k}_L$. Second, the level of the ion sound turbulence must be such that saturation occurs. If there is just one localized region where the ion sound and Langmuir turbulence coexist, then T_T must increase to reach T_L in less than the time it takes a transverse wave to propagate across the source. For a source with linear dimensions L, this condition reduces to (Melrose *et al.*, 1986)

$$\frac{v_g}{L} \lesssim \frac{\pi}{6}\,\omega_p\,\frac{v_\phi^2}{V_e^2}\,\frac{W_s}{n_e T_e}\,, \tag{10}$$

where $v_g = nc \cong \sqrt{3}\,V_e c/v_\phi$ is the group speed of the transverse waves, v_ϕ is the phase speed of the Langmuir waves, and W_s is the energy density in the (assumed isotropic) ion sound waves in a range $\Delta k_s \cong k_s$ about $k_s = \omega_p/v_\phi$.

The crucial question concerning the efficacy of this process is the presence of ion sound waves with $\mathbf{k}_s \cong \pm \mathbf{k}_L$. It may be that the broad spectrum of ion sound waves in the solar wind (e.g., Gurnett and Frank, 1978) includes the required waves, or it may be that they are generated through a decay instability, as suggested by Lin *et al.* (1986). It may also be that the required waves are not present, and that F emission is due to some other process.

An attractive feature of this mechanism for F emission is that it allows one to determine the value of the brightness temperature T_L at which saturation occurs. Saturation may be treated using (5) with $A = L$, $B = S$, and $C = T$, giving $T_T \cong T_L$ at saturation for $T_s \gg T_L v_s/v$ with $v_\phi = v_b$ for type III emission. For Langmuir waves confined to a range of angle $\Delta\theta$, the bump-in-tail instability starts to saturate when an energy $W_L \cong \frac{1}{2} n_b m v_b \Delta v$ is reached, and then the implied value of T_L and $T_T \lesssim T_L$ gives a saturation value

$$T_T \lesssim \frac{\frac{1}{2} n_b m v_b^2}{\pi(\Delta\theta)^2}\left(\frac{v_b}{f_p}\right)^3. \tag{11}$$

4.2. Second harmonic emission

H emission is attributed to the coalescence $L + L \rightarrow T$, which also saturates at $T_T(\mathbf{k}_T) = T_L(\mathbf{k}_L)$. The kinematic conditions require $\mathbf{k}_L + \mathbf{k}_L' = k_T$ and $k_T = k_H \cong \sqrt{3}\,\omega_p/c$ for transverse waves at $\omega \cong 2\omega_p$. With $v_\phi \ll c/\sqrt{3}$ one has $k_L \gg k_T$, one requires $\mathbf{k}_L' \cong \mathbf{k}_L$ and, hence, the two Languir waves must coalesce nearly head on. This head-on condition would seem to be easily satisfied when a backscatter of the Langmuir waves has occurred. However, this is not necessarily the case: a backscatter changes k_L by k_0, cf. Equation (4), and for $k_0 > k_H$ the coalescence of k_L and $k_L' = K_L \pm k_0$ into $k_T = k_H$ is not possible (Melrose, 1983). Three possibilities arise:

(i) The condition $k_0 > k_H$ may not be satisfied. According to Equation (4), k_0 depends only on the electron temperature and for $T_e > 5 \times 10^5$ K the coalescence is possible. The condition $T_e > 5 \times 10^5$ K is satisfied in the solar corona but not in the solar wind.

(ii) The initial spectrum of Langmuir waves covers a relatively broad range Δk about k_L. The spectrum formed in one backscatter then covers a range Δk about $k_L - k_0$ in the backward direction. In this case, provided one has $\Delta k > k_H - k_0$, H emission is possible for $k_0 > k_H$.

(iii) Scattering of the Langmuir waves (Section 3.3) may produce a roughly isotropic distribution which contains nearly anti-parallel waves. The coalescence $L + L \rightarrow T$ is then possible between two such waves.

An important qualitative point is that the coalescence processes considered here for F and H emission saturate at the same value of T_T, which may be estimated using (11).

5. Some Outstanding Difficulties in Applications of Plasma Emission

Despite approximately three decades of theoretical work on plasma emission, there is no case where there is convincing quantitative agreement between theory and observation. Here several difficulties with the applications are discussed, concentrating on those with implications on local density structures in sources.

5.1. The Harmonic Ratio

Theory implies that the harmonic ratio for plasma emission should be very close to two, i.e., the ratio of the frequencies of H and F emission should differ insignificantly from 2 : 1. However, observations show that it is systematically less than 2 : 1. The only plausible explanation for this is that the F and H components have moderately broad bandwidths and that the lower part of the F band does not escape (Wild *et al.*, 1954). However, this bandwidth cannot be intrinsic to plasma emission itself, and must be due to a range of plasma frequencies in the source (Roberts, 1959). To prevent the lower half of the F band from escaping it may be necessary to assume that the source is in an underdense duct (Stewart, 1974). Thus to explain harmonic ratios different from two, it seems necessary to appeal to density structures in the source, with the F emission being able to escape only from the density peaks. This argument is independent of other arguments for ducting (Duncan, 1979), and combining the two arguments implies a model in which the source is an underdense duct, with local variations in the plasma density across the duct. The observed value of the harmonic ratio does not provide evidence on the sizes of the density irregularities, but it does imply a minimum value for the density variations. For example, to account for an harmonic ratio of 1.8 : 1 (Stewart, 1974) with only half the F bnd escaping requires a density variation of at least 20%. Observed fluctuations in the solar wind (Celnikier *et al.*, 1983) are considerably smaller than this.

Alternative mechanisms which could produce an harmonic ratio less than two seem implausible. The production of F emission could involve a coalescence with a low frequency wave with a higher frequency than the ion sound waves. The required frequency to produce a ratio 1.8 : 1, say, is $\omega = 0.1\omega_p$. However, all the candidates, such as electron cyclotron waves ($\omega \cong \Omega_e$ = cyclotron frequency) and electron acoustic waves (e.g., Gary and Tokar, 1985) seem quite implausible. Another possibility is that once the F emission is generated it is scattered e.g., off ion sound waves, many times before it escapes. Each such scattering changes the wave frequency by an amount equal to the frequency of the ion sound wave. The transverse waves then diffuse in frequency causing a spread in the bandwidth of the escaping F emission. However, the frequency of the ion sound waves involved is $\cong 10^{-4}\omega_p$ and, hence, $\cong 10^6$ scatterings would be required for the bandwidth to increase by 10% of the plasma frequency. These alternative mechanisms are much less plausible than that involving density variations.

5.2. Polarization of plasma emission

Theory implies that F emission should be essentially completely polarized in the sense of the o-mode for plausible values of the coronal magnetic field, and that H emission should be weakly polarized in the sense of the o-mode for streaming electrons (e.g., Melrose *et al.*, 1980). The argument for F emission is that the frequency of the Langmuir waves lies between the cutoff frequencies, ω_p for the o-mode and $\cong \omega_p + \frac{1}{2}\Omega_e$ for the x-mode, so that only o-mode emission can be generated. Although observed plasma emission is polarized in the sense of the o-mode, there are at least three difficulties relating to the magnitude of the polarization.

(i) The polarization of F emission in type III and type II bursts is always less than 100%. The degree of polarization does not exceed about 70% for type III bursts (Suzuki and Sheridan, 1977; Dulk and Suzuki, 1980).

(ii) The H emission in type III events is only modestly polarized (Suzuki and Sheridan, 1977), but the degree of polarization is higher than can be explained by simple theory (Melrose *et al.*, 1980).

(iii) Type I emission, which is accepted as being F emission, is typically 100% polarized. However, the degree of polarization can have a characteristic value less than 100%, especially for sources away from the central meridian (Zlobec, 1975).

There are two possibilities for explaining the $<100\%$ polarization of F emission in type III bursts: the lower polarization is intrinsic or the radiation is depolarized as a propagation effect. Lower intrinsic polarization can occur only if the frequency of the F component on generation exceeds the cutoff frequency for the x-mode, and this seems highly implausible. Depolarization due to propagation effects is more favorable. The evidence for type I bursts on variation of the degree of polarization with viewing angle (Zlobec, 1975) is strongly suggestive of a propagation effect. However, the details of the depolarizing mechanism remain unclear. There are three suggested mechanisms for depolarization, and in one sense all three are equivalent: these are mode coupling (Melrose, 1975), scattering by low-frequency waves (Wentzel, 1984) and reflection at density inhomogeneities (Hayes, 1985). All these mechanisms involve small scale lengths and, hence, large spatial Fourier coefficients in plasma inhomogeneities, allowing large angle changes for the scattered (or reflected) waves. After a large angle deflection an initial o-mode component is composed of a mixture of o-mode and x-mode components. Depolarization results whenever a large angle deflection occurs for waves above the cutoff frequency for the x-mode. The conditions under which the required degree of depolarization occurs are very specific (Hayes, 1985).

In summary, the observed polarization of plasma emission appears to have been strongly influenced by propagation effects. Models to account for this require relatively sharp gradients in plasma density involving scalelengths comparable with the wavelength of the radiation. Further investigation of detailed models is required.

5.3. Scattering and directivity

It is thought that the apparent sizes of sources of plasma emission are scatter images of smaller actual sources (Steinberg *et al.*, 1971). Such scattering is attributed to density fluctuations in the source. The sources remain highly directive, however, and it has been assumed tha this is difficult to reconcile with the amount of scattering required to explain the apparent sizes. The directivity has been explained in terms of specific structures in the corona. For type I bursts Bougeret and Steinberg (1977) suggested a fibrous structure, involving parallel columns of overdense plasma. For type III bursts Duncan (1979) suggested that the sources are in underdense ducts. There is no other supporting evidence for either of these suggestions.

It is worth pointing out that the directivity may not be an actual problem, and the perceived difficulty in reconciling the scattering and the directivity may be an artefact of assumptions made in scattering models. The point is the following. Snell's law implies that radiation with refractive index $n \ll 1$ initially becomes confined to a cone with half angle $\theta \cong n/n'$ at another point where the refractive index is $n' \gg n$. It is assumed that scattering destroys this directivity. However, it is not clear that this is the case, and in fact it requires specific assumptions on the scattering for this directivity to be destroyed.

One way of seeing this point is to note the existence of a conserved quantity, called the 'generalized étendue' in geometric optics (e.g., Welford and Winston, 1978). Consider a collection of rays confined to direction cosines dL and dM about the z-axis and to an area dx dy at some point P where the refractive index is n. At any other point P' where the refractive index is n', suppose the same rays are confined to ranges dL' and dM' of direction cosines about the new ray direction along the z'-axis, and to an area dx' dy'. Then one has

$$n^2 \, \mathrm{d}x \, \mathrm{d}y \, \mathrm{d}L \, \mathrm{d}M = n'^2 \, \mathrm{d}x \, \mathrm{d}y' \, \mathrm{d}L' \, \mathrm{d}M' \; . \tag{12}$$

Scattering, in the sense used here, involves only propagation in accord with the laws of geometric optics. The path from the actual to the apparent source can then be regarded as equivalent to an optical system to which (12) applies. The generalized étendue them implies that the range of angles to which the radiation is confined must shrink not only as the refractive index increases (as is well known) but also as the area increases. From this viewpoint the difficulty is in understanding why the radiation is not much more highly directive than it appears to be. To be compatible with the generalized étendue it appears that radiation filling a range of angles much greater than $\theta \cong n$ (for

$n' = 1$) must be such that rays at a particular angle arise only from a small fraction of the area of the apparent source. This has implications on the relation between the apparent brightness temperature and the value of T_T in the source.

These points require more detailed discussion. What the generalized étendue does is provide a different viewpoint on the problem of scattering by density inhomogeneities in the corona. Qualitatively, an implication is that it is much more difficult to destroy directivity by scattering than has been assumed hitherto.

5.4. Clumpy Langmuir Waves

Direct evidence on inhomogeneities in sources is available for type III sources in the solar wind. It is found that the Langmuir waves are distributed in a highly inhomogeneous way (Gurnett and Anderson, 1977; Lin *et al.*, 1986), referred to as 'spikes' or 'clumps' of Langmuir waves. A plausible explanation for the clumpy distribution is the effect of local density inhomogeneities (Smith and Sime, 1979). As mentioned in Section 3.3, scattering of Langmuir waves off density inhomogeneities appears to e strong enough to suppress growth entirely (Muschietti *et al.*, 1985) and specific assumptions are required to explain any Langmuir waves at all (Melrose *et al.*, 1986). Investigation of this problem is continuing.

The evidence from the solar wind provides qualitative support for the roles of inhomogeneities considered above for emission from the corona. However, a potentially serious problem is that the observed density fluctuations (Celnikier *et al.*, 1983) seem to be isotropic (Muschietti, 1986, private communication). How can one explain ducting without any anisotropic density structures? Either the anisotropic density structures are present and have escaped detection for some reason, or ducting effectively occurs without ducts. The remarks on the generalized étendue in Section 5.3 suggest that the latter possibility cannot be excluded.

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