A SELF-CONSISTENT MODEL FOR THE STORM RADIO EMISSION FROM THE SUN

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Abstract. A self-consistent theoretical model for storm continuum and bursts is presented. We propose that the Langmuir waves are emitted spontaneously by an anisotropic loss-cone distribution of electrons trapped in the magnetic field above active regions. These high-frequency electrostatic waves are assumed to coalesce with lower-hybrid waves excited either by the trapped protons or by weak shocks, making the observed brightness temperature equal to the effective temperature of the Langmuir waves.

It is shown that whenever the collisional damping (ν_c) is more than the negative damping ($-\gamma_A$) due to the anisotropic distribution, there is a steady emission of Langmuir waves responsible for the storm continuum. The type I bursts are generated randomly whenever the collisional damping (ν_c) is balanced by the negative damping ($-\gamma_A$) at the threshold density of the trapped particles, since it causes the effective temperature of Langmuir waves to rise steeply. The number density of the particles responsible for the storm radiation is estimated. The randomness of type I bursts, brightness temperature, bandwidth and transition from type I to type III storm are self-consistently explained.

1. Introduction

On dynamic-spectrum records a type I storm consists of many shortlived, narrow-band bursts scattered at random across the frequency-time plane. Usually these bursts are superposed on a background continuum of emission. Hey (1946) was first to recognize the association of noise storms with sunspots. Except near the limb, the storm radiation is generally strongly circularly polarized in the sense of the ordinary mode (Le Squeren, 1963). The observed polarization and heights are generally taken as evidence that the emission process is a plasma process at or near the local plasma frequency. Usually the type I storms are longlived, lasting from hours to days. The frequency of occurrence is in the range from 500 to 50 MHz, peaking around 100–150 MHz. The centroid of the burst emission is usually within the observed source region of the continuum (Daigne, Lantos-Jarry, and Pick, 1970). The most important characteristic of the storm radiation is the brightness temperature. Usually the brightness temperature of the continuum $< 10^{10}$ K, the maximum observed being $\sim 9.4 \times 10^9$ K at 169 MHz (Kerdraon and Mercier, 1983), whereas type I bursts are more intense and they can have $T_b \gtrsim 10^{10}$ K, even though there are no measurements at all the frequencies. At lower frequencies, there is a transition from type I storm to type III storm radiation. Kundu (1965) and Elgaroy (1977) have extensively reviewed the observations on solar noise storms:

Two emission mechanisms, namely cyclotron and plasma, were proposed to explain the storm radiation. The main objection to the cyclotron emission mechanism is that

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it favors the extraordinary mode whereas it was never observed (Melrose, 1980a). Therefore, it is now believed that the emission mechanism should be a plasma mechanism. The scenario of recent theories is as follows: (1) Langmuir waves are excited spontaneously or through a loss-cone type of instability by a trapped distribution of electrons (2) the conversion of these waves into radiation is by coupling with low-frequency waves, either lower-hybrid or ion-acoustic (Melrose, 1980a; Benz and Wentzel, 1981; Spicer, Benz, and Huba, 1981; Wentzel, 1981, 1986). All the above theories invoke spontaneous emission of Langmuir waves by isotropic distribution of electrons to explain the continuum and loss-cone type of instabilities to excite enhanced emission of high-frequency waves to explain the type I bursts. Even though the trapped distribution of electrons is identified as the responsible agent for type I storms, while calculating the effective temperature of the spontaneously emitted Langmuir waves, the damping due to trapped electrons is always taken as positive, since the distribution is considered to be isotropic. The Landau damping of Langmuir waves by the background plasma is estimated by considering that the electron density is independent of radial distance, and since it is extremely small, it has been neglected completely.

If one carefully observes the dynamic spectrum or intensity plots of type I emission, one sees that type I bursts occur as intensity fluctuations in frequency and time. Most probably the problem of type I emission is directly related to the problem of field fluctuations in a non-equilibrium plasma. The problem of critical fluctuations near the onset of a plasma instability is still an unsolved problem in plasma physics (see Ichimaru, 1973; Sitenko, 1982; Fung, Papagopoulos, and Wu, 1982). The study of spontaneous emission of Langmuir waves by the anisotropic distribution of electrons is similar to the study of critical fluctuations in a non-equilibrium plasma.

Here we propose that the spontaneous emission of Langmuir waves by the anisotropic distribution of electrons can explain both the continuum and bursts self-consistently. If the transformation of high-frequency electrostatic waves into radiation is assumed to be due to their coalescence with lower-hybrid waves excited by the trapped ions or by the weak shocks, the effective temperature is equal to the observed brightness temperature. We show that whenever collisional damping (ν_c) is dominant over the negative damping due to anisotropic electrons ($-\gamma_A$), there is a steady emission of Langmuir waves responsible for the continuum, and whenever $\nu_c = -\gamma_A$, i.e., at the threshold densities of the trapped particles, the effective temperature rises steeply, reaching values up to 10^{11} K, giving rise to bursts, which is similar to a critical fluctuation in the effective temperature of Langmuir waves near the onset of the instability. The threshold condition is satisfied randomly giving rise to random appearance of type I bursts superposed on the continuum, since the acceleration of electrons in the source regions is stochastic. The plan of the paper is as follows: in Section 2, we estimate the effective temperature of Langmuir waves spontaneously emitted by a loss-cone distribution in the coronal streamer. The marginally stable state of the electron distribution function is maintained because of the weak scattering and collisional losses of lower energy particles. In Section 3, we compare our estimates with those of quasi-linear estimates. We estimate the energy density of low-frequency waves in Section 4. In

Section 5 we discuss the scattering process, and in Section 6 we give a method to plot the brightness temperature versus density. In Section 7 we present the discussion, and in Section 8 we summarize our conclusions.

2. Emission of High-Frequency Langmuir Waves

Most probably the suprathermal electrons responsible for noise-storm radiation are trapped in the magnetic field structures above the active regions. In a trap, particles can be anisotropic due to the presence of the loss cones. For the formation and maintainance of the anisotropic distribution in the magnetic loop, the condition for weak diffusion $\nu_D \ll \frac{1}{2}\alpha_0^2 \nu_b$ should satisfy (Kennel and Petchek, 1966; Melrose and Brown, 1976). Here $\nu_D = 10^{-8} n_e E^{-3/2}\ \mathrm{s}^{-1}$ is the rate at which pitch-angle diffusion occurs where E is the electron energy in keV, α_0 is the pitch angle at the edge of the loss cone, ν_b is the bounce rate in the trap, and n_e is the ambient electron density. For the type I emission for the assumed value of $E = 28$ keV, the size of the trap $\approx 3 \times 10^{10}$ cm, $\alpha_0 \approx 0.1$ and $n_e \approx 10^8\ \mathrm{cm}^{-3}$, we obtain $\nu_D \approx 6.8 \times 10^{-3}\ \mathrm{s}^{-1}$ whereas $\frac{1}{2}\alpha_0^2 \nu_b \approx 1.7 \times 10^{-2}\ \mathrm{s}^{-1}$, indicating that the condition for weak diffusion is satisfied (see also, Melrose and White, 1979). Therefore, one can assume that the typical distribution responsible for type I emission is the loss-cone distribution

$$f_s(v) = \frac{1}{4\sqrt{2}\,\pi^{3/2}} \left(\frac{v_x}{v_{T_b}}\right)^5 \exp\left(-\frac{v_z^2 + v_x^2}{2v_{T_b}^2}\right), \tag{1}$$

where v_z and v_x are the velocity components parallel and perpendicular to the magnetic field, and v_{T_b} is the thermal spread in the velocities. The above distribution is normalized to unity.

The effect of these particles on the Langmuir wave distribution may be described by the transfer equation

$$\frac{\mathrm{d}T^{\mathrm{L}}(k)}{\mathrm{d}t} = \alpha^{\mathrm{L}}(k) - \gamma^{\mathrm{L}}(k)T^{\mathrm{L}}(k)\,. \tag{2}$$

Here T^{L} is the effective temperature of the Langmuir waves, α^{L} is the emission coefficient, and γ^{L} is the effective absorption coefficient, which is the sum of negative absorption due to anisotropic electrons (γ_A), collisional damping (ν_c) and Landau damping by ambient electrons (γ_L). Therefore, it can be written as

$$\gamma^{\mathrm{L}} = \gamma_A + \nu_c + \gamma_{\mathrm{L}}\,. \tag{3}$$

To estimate the effective temperature of Langmuir waves, one should calculate the emission coefficient α^{L} and the effective damping γ^{L}.

2.1. Emission coefficient

The emission coefficient α^{L} for an electrostatic wave in the case of magnetized electrons; i.e., energy per unit volume and wave number interval per second is given by (Melrose, 1980c; Wentzel, 1985)

$$\alpha^{\mathrm{L}}(k) = 16\pi^3 e^2 R \frac{\omega^2}{k^2} n_s \sum \int v_x \,\mathrm{d}v_x \,\mathrm{d}v_z \, J_s^2\left(\frac{k_x v_x}{\Omega}\right) f_s(v)\delta(\omega - s\Omega - k_z v_z). \tag{4}$$

Here R is the ratio of the energy of the electrostatic wave to the total energy and it is $\approx \frac{1}{2}$ for Langmuir waves, e is the electronic charge, J_s is the Bessel function of order s; Ω is the cyclotron frequency; k_z and k_x are the components of the wave vector along and across the magnetic field respectively, and n_s is the number density of the suprathermal particles. The Cherenkov resonance condition for the magnetized electrons is

$$\omega = s\Omega + k_z v_z \,. \tag{5}$$

By substituting the function f_s given by Equation (1) in Equation (4) and by integrating it over v_z, using the resonance condition (5), we obtain

$$\alpha^{\mathrm{L}}(k) = 2\sqrt{2}\,\pi^{3/2} \frac{e^2}{v_{T_b}^5} R \frac{\omega^2}{k^2} \frac{n_s}{k_z} \sum_s \exp\left(-\frac{(\omega - s\Omega)^2}{2k_z^2 v_{T_b}^2}\right) \times$$

$$\times \int_0^\infty v_x^3 \,\mathrm{d}v_x J_s^2\left(\frac{k_x v_x}{\Omega}\right) \exp(-v_x^2/2v_{T_b}^2)\,. \tag{6}$$

The integration over v_x can be carried out by using the relation (Erdelyi, 1953)

$$I = \int_0^\infty e^{-a^2x^2} J_m^2(bx) x \,\mathrm{d}x = \frac{1}{2a^2} e^{-b^2/2a^2} I_m\left(\frac{b^2}{2a^2}\right), \tag{7}$$

where I_m is the modified Bessel function of the first kind for an imaginary argument. But

$$-\frac{\mathrm{d}I}{\mathrm{d}a^2} = \int_0^\infty x^3 \,\mathrm{d}x \, J_m^2(bx) \exp(-a^2 x)\,. \tag{8}$$

In the limit of high frequencies, $\omega \gg \Omega$, the effect of the magnetic field is negligible. Then, absorption must also be possible for $k_z = 0$, since this limit corresponds to the isotropic plasma without an external magnetic field. This is physically obvious since the Larmor radius of the particles greatly exceeds the wavelength when $\Omega \to 0$. The formal transition to the limit $\Omega \to 0$ is nontrivial and connected with the problem of the asymptotic representation of the Bessel functions of high order in large arguments. For $\Omega \to 0$, the arguments of the Bessel functions become large. Then all terms with $|m| < s$ contribute to the same order, whereas the terms with $|m| > s$ are exponentially small where $s = b^2/2a^2 = k_x^2 v_{T_b}^2/\Omega^2$.

Therefore, by differentiating the right-hand side of (7) with respect to a^2 and by using the asymptotic relation for $I_m(z)$ for large z which is

$$I_m(z) = \frac{\exp(z)}{\sqrt{2\pi z}}\left(1 - \frac{m^2 - 0.25}{2z}\right), \tag{9}$$

we obtain

$$\alpha^{\mathrm{L}}(k) = \frac{1}{4}\,\frac{n_s}{n_e}\,\frac{m_e}{v_{T_b}^2}\,\frac{\omega_p^2}{k^2}\,\frac{\omega^2}{k_x k_z}\,\Omega \sum_s \times$$

$$\times \exp\left[-\left(\frac{\omega - s\Omega}{\sqrt{2}\,v_{T_b} k_z}\right)^2\right]\left(1 + \frac{(s^2 - 0.25)}{2k_x^2 v_{T_b}^2}\,\Omega^2\right). \tag{10}$$

Here we substituted $R = \frac{1}{2}$. The main contribution is by the term $\omega = s\Omega$. The s can take values from 5 to 20 for a typical noise storm. And for $\omega_p \approx \omega$, $\omega/k_x \approx v_{T_b}$, $k_z \approx 0.7k_x$, the emission coefficient is reduced to

$$\alpha^{\mathrm{L}}(k) \approx 0.1 m_e \frac{n_s}{n_e}\,\omega v_{T_b}^2\,. \tag{11}$$

2.2. Effective damping

The negative damping of Langmuir waves ($-\gamma_A$) excited by the anisotropic distribution of electrons as given by Equation (1) can be written as (Zaitsev and Stepanov, 1975)

$$-\gamma_A = \frac{1}{4}\sqrt{\frac{\pi}{2}}\,\frac{n_s \omega_p^4}{n_e k^3 v_{T_b}^3}\,\frac{\left(1 - \dfrac{\omega^2}{k^2 v_{T_b}^2} - 2\cot\theta\right)}{(1 + \cot\theta)}\,e^{-\omega^2/2k^2 v_{T_b}^2}, \tag{12}$$

where θ is the angle between the magnetic field and the wave vector. The γ_A is negative only for Langmuir waves with phase velocities

$$\left(\frac{\omega}{k}\right)^2 < v_{T_b}^2 (1 - 2\cot^2\theta)\,. \tag{13}$$

In other words the system is unstable only for $54^\circ\!.74 \leqq \theta < \pi/2$.

The waves are actually z-modes since $\mathbf{k}$ is almost perpendicular to $\mathbf{B}$. When $\cot^2\theta \ll 1$, $-\gamma_A$ reaches its maximum:

$$-\gamma_A^{\max} \approx 4.4 \times 10^{-2}\,\omega_p\,\frac{n_s}{n_e}\,. \tag{14}$$

Therefore, the negative damping is directly proportional to the frequency as well as the number density of the nonthermal particles.

The damping due to electron-ion collisions in the coronal plasma is given by

$$\nu = 5.5(n_e/T_e^{3/2}) \ln(10^4 T_e^{2/3}/n_e^{1/3}) , \quad (15)$$

where T_e is electron temperature in degrees of Kelvin. The collisional damping rate is given by

$$\nu_c = \nu/2 . \quad (16)$$

The absorption coefficient due to Landau damping in the background plasma is

$$\mu_{\mathrm{L}} = \frac{1}{3} \sqrt{\frac{\pi}{2}} \frac{\omega_p^2 \omega^3}{k^4 v_T^5} \exp\left(- \frac{\omega^2}{2 v_T^2 k^2}\right), \quad (17)$$

where v_T is the thermal speed of electrons in the ambient plasma. The above formula for μ_{L} is valid only in the case of weak absorption, i.e., when $\mu_{\mathrm{L}} \ll k$ in a transparent medium. The condition of weak absorption is fullfilled, if $\omega^2/2k^2 v_T^2 = v_{ph}^2/2v_T^2 \gg 1$, where v_{ph} is the phase velocity of the waves. We should note that in the region of weak absorption of plasma waves, their frequency $\omega \approx \omega_p$. If $v_{ph} \sim v_T$, the expression (17) is no longer valid, because in that region $\gamma_{\mathrm{L}} \sim \omega$, i.e., the waves are damped in $t \sim \gamma^{-1} \sim \omega^{-1}$ s, or in a distance $L \sim \lambda/2\pi$ which means that the plasma waves with $v_{ph} \sim v_T$ cannot propagate in the plasma due to a strong Landau damping. Therefore, as a result of Landau damping, plasma waves cannot go out of the region $\omega \approx \omega_p$ into the rarefied plasma (e.g., from solar corona into interplanetary medium). Actually during the propagation of waves with frequency ω in a smoothly varying plasma, their phase velocity $v_{ph} = \sqrt{3}\, v_T/\sqrt{\varepsilon}$ is decreased with the plasma frequency ω_p (since $\varepsilon = 1 - \omega_p^2/\omega^2$). Therefore, when v_{ph} is decreased to v_T, strong Landau damping takes place leading to a complete absorption of the plasma waves. (For details, see Zheleznyakov, 1977; Thejappa, Gopalswamy, and Kundu, 1990). It is, therefore, useful to estimate the length $\Delta\rho$ travelled by a plasma wave before it is completely Landau damped. Since the absorption coefficient is related to absorption length $\Delta\rho$ as

$$\mu_{\mathrm{L}} = \frac{1}{\Delta\rho} , \quad (18)$$

we can use Equation (17) as an equation for the absorption length $\Delta\rho$. From the dispersion relation for Langmuir waves, the dielectric constant in the cold approximation $\varepsilon = 1 - \omega_p^2/\omega^2$ can be written as

$$\varepsilon = 3 \frac{v_T^2}{v_{ph}^2} . \quad (19)$$

By Taylor expanding ε around $\omega \approx \omega_p$ we obtain

$$\varepsilon = \frac{\mathrm{d}\varepsilon}{\mathrm{d}\rho} \Delta\rho = - \frac{1}{\omega^2} \frac{\partial \omega_p^2}{\partial \rho} \Delta\rho . \quad (20)$$

Since the noise storm radiation occurs mainly above the active regions, i.e., in the streamers, the most appropriate density model for such a streamer is the 5 × Newkirk's density model

$$N = 2.1 \times 10^5 \times 10^{4.32/\rho}, \tag{21}$$

where ρ is the radial distance in units of solar radii. The plasma frequency ω_p is related to the density N as

$$\omega_p^2 = 3.18 \times 10^9 N. \tag{22}$$

By using Equation (21), differentiating Equation (22) with respect to ρ, and equating $\omega \approx \omega_p$ we obtain

$$\varepsilon = 9.95 \frac{\Delta\rho}{\rho^2}. \tag{23}$$

By substituting Equation (23) into (17) in place of $3v_T^2/v_{ph}^2$, according to the relation (19), we obtain an equation for $\Delta\rho$:

$$\Delta\rho = 3.8 \times 10^{-2} \frac{R_\odot}{v_T} \omega_p \rho^4 \exp(-0.15\rho^2/\Delta\rho), \tag{24}$$

where $R_\odot$ is a solar radius. The radial distance ρ is related to the plasma frequency f_p in MHz in the density model given by Equation (21) as

$$\rho = \frac{4.32}{\log f_p^2 - 1.23}. \tag{25}$$

For $T = 10^6$ K, $v_T \approx 3.89 \times 10^8$ cm s^{-1} and $R_\odot = 6.96 \times 10^{10}$ cm, one finds for the absorption length $\Delta\rho$:

$$\ln \Delta\rho - \left[23.42 - \ln\left(\frac{f_p}{(\log f_p^2 - 1.23)^4}\right)\right] + \frac{2.82}{\Delta\rho(\log f_p^2 - 1.23)^2} = 0. \tag{26}$$

One can solve the above equation for the absorption length $\Delta\rho$ for different frequencies. The Landau damping can be estimated as

$$\gamma_{\mathrm{L}} = v_g/\Delta\rho. \tag{27}$$

Here v_g is the group velocity, estimated as $v_g \sim 3v_T^2/v_{ph}$. For $v_T \approx 3.89 \times 10^8$ cm s^{-1} and $v_{ph} \sim v_{T_b} \sim 10^{10}$ cm s^{-1} one obtains $v_g \sim 4.54 \times 10^7$ cm s^{-1}. A tabulation of the absorption length and Landau damping for different frequencies is given in Table I.

TABLE I

The effective absorption length and Landau damping of the Langmuir wave estimated for different plasma frequency levels in the corona

Frequency in MHz	Absorption length ($\Delta\rho$) in $R_\odot$	Landau damping γ_L in s^{-1}
300	1.05×10^{-2}	6.2×10^{-2}
250	1.19×10^{-2}	5.48×10^{-2}
167	1.39×10^{-2}	4.69×10^{-2}
80	2.11×10^{-2}	3.09×10^{-2}
50	2.9×10^{-2}	2.25×10^{-2}
30	4.92×10^{-2}	1.33×10^{-2}

2.3. Effective temperature of Langmuir waves

From Equation (2) it is clear that the effective temperature of the spontaneously emitted Langmuir waves in the stationary case is given by:

$$T^{\mathrm{L}}(k) = \frac{\alpha^{\mathrm{L}}(k)}{\gamma^{\mathrm{L}}(k)} . \tag{28}$$

In the theory of fluctuations, the effective fluctuation temperature of a non-equilibrium plasma is determined by a similar expression. When the negative damping ($-\gamma_A$) due to an anisotropic distribution of electrons is balanced by the collisional damping (ν_c), the effective damping $\gamma^{\mathrm{L}}(k)$ is equal to the Landau damping γ_{L}. This situation corresponds to the critical fluctuation level where the effective temperature rises steeply before the onset of the instability. We call the density of the suprathermal electrons when $-\gamma_A = \nu_c$ the threshold density $(n_s/n_e)_{th}$, which is different for different frequencies. In Table II, we give the collisional damping (ν_c), the negative damping ($-\gamma_A$) and $(n_s/n_e)_{th}$ at different frequencies, assuming $T_e = 10^6$ K.

By using relation (11) for the emission coefficient and relation (28) when $\gamma^{\mathrm{L}}(k) = \gamma_{\mathrm{L}}(k)$, we can compute the effective temperature of the Langmuir waves at threshold densities. In Table III, we give the emission coefficient and the effective

TABLE II

The ambient electron density, collisional damping, negative damping due to trapped electrons, and the density ratio of the trapped particles to the ambient electrons at threshold for different plasma frequency layers in the corona

Frequency in MHz	n_e cm^{-3}	ν_c	$-\gamma_A^{\max}$	$(n_s/n_e)_{th}$
300	1.11×10^9	34.73	$8.29 \times 10^7 n_s/n_e$	4.19×10^{-7}
250	7.74×10^8	24.69	$6.9 \times 10^7 n_s/n_e$	3.57×10^{-7}
167	3.45×10^8	11.26	$4.62 \times 10^7 n_s/n_e$	2.44×10^{-7}
80	7.92×10^7	2.69	$2.21 \times 10^7 n_s/n_e$	1.23×10^{-7}
50	3.1×10^7	1.08	$1.38 \times 10^7 n_s/n_e$	7.83×10^{-8}
30	1.1×10^7	0.40	$8.29 \times 10^6 n_s/n_e$	4.82×10^{-8}

TABLE III

The emission coefficient and the effective temperature of the Langmuir waves emitted by the trapped electrons at threshold, i.e., when negative damping is balanced by the collisional damping at various frequencies

Frequency in MHz	α_{th}^{L}	T_{th}^{L} in K
300	6.53×10^{-6}	8.39×10^{11}
250	7.11×10^{-6}	6.75×10^{11}
167	2.33×10^{-6}	3.6×10^{11}
80	5.63×10^{-6}	1.32×10^{11}
50	2.24×10^{-7}	7.2×10^{10}
30	8.27×10^{-8}	4.05×10^{10}

brightness temperature of Langmuir waves at threshold beam densities for various frequencies. Here we assumed $v_{T_b} \sim 10^{10}$ cm s^{-1}.

3. Comparison with Quasi-Linear Analysis

In Section 2, we have estimated the effective temperature of the spontaneously emitted Langmuir waves, by an anisotropic electron distribution. Above the threshold values, i.e., when $-\gamma_A > v_c$, the linear analysis does not hold since the condition for quasi-equilibrium between the emission and absorption does not hold. Therefore one should estimate the T^{L} by solving the two-dimensional quasi-linear equations. Breizmann (1986) has solved such a problem for a typical trapped distribution of electrons. The maximum energy density of the Langmuir waves excited by such a distribution is estimated as

$$\frac{W_{\mathrm{L}}}{n_e T_e} \approx \frac{n_s m v_b^2}{2 \Lambda n_e T_e} q , \tag{29}$$

where Λ is the Coulomb logarithm, v_b – velocity of the suprathermal particles and q is a dimensionless function whose maximum value is ~ 0.1. The form of the distribution function used by Breizmann (1986) is different from what we have used, however, for estimating the energy density, the form of the distribution function is not important. The effective temperature of the plasma waves is given by (Melrose, 1974)

$$\langle T^{\mathrm{L}} \rangle_{\Delta\Omega} = W^{\mathrm{L}} \left(\frac{v_{ph}}{f_p} \right)^3 \frac{v_{ph}}{\Delta v_{ph}} \frac{1}{\Delta\Omega} , \tag{30}$$

where Δv_{ph} is the spread in phase velocities and $\Delta\Omega$ is the solid angle of the emission. For $v_{ph} \sim \Delta v_{ph} \sim 10^{10}$ cm s^{-1}; and $\Delta\Omega \sim 4\pi$, we can estimate the effective temperature of Langmuir waves at different frequencies. In Table IV we give the Coulomb logarithm, the energy density of Langmuir waves as given by Equation (29) and the effective temperature of Langmuir waves at different frequencies.

The effective temperature of Langmuir waves computed by solving the quasi-linear equation T_{QL}^{eff} as given in Table IV, is always more than the threshold effective tempera-

TABLE IV

The Coulomb logarithm, the ratio of the total energy density of Langmuir waves to the thermal energy and the effective temperature of Langmuir waves estimated by quasi-linear analysis for different frequencies

Frequency in MHz	$\Lambda = \ln(10^4 T_e^{3/2}/n_e^{1/3})$	$W_L/n_e T_e$	T_{QL}^{eff}
300	11.48	1.2×10^{-6}	3.89×10^{12}
250	11.6	1.02×10^{-6}	4.02×10^{12}
167	11.87	6.78×10^{-7}	3.99×10^{12}
80	12.36	3.28×10^{-7}	4.03×10^{12}
50	12.67	2.04×10^{-7}	4.03×10^{12}
30	13.01	1.22×10^{-7}	3.99×10^{12}

ture of spontaneously emitted Langmuir waves (see Table III). One should note that the effective temperature T_{QL}^{eff} remains constant, $\sim 4 \times 10^{12}$ K, independent of frequency at the threshold values of n_s/n_e.

4. Low-Frequency Waves

In addition to the electrons, if one assumes that the fast ions are injected across the field into the magnetic confinement system with a velocity v_{bi}, the distribution function for such fast ions can be assumed to be close to a δ-function: $f(\mathbf{v}) \sim n_b \delta(v - v_{bi})$, which we call a ring. For a sufficiently large plasma density, i.e., $\omega_{pi} \gg \Omega_i$ and when the ratio of the beam to plasma density is not too small, i.e., $n_b/n_e \gtrsim m_e/m_i$, the main instabilities are those in the frequency range $\Omega_i \ll \omega \ll \Omega_e$ with the growth rate $\gamma > \Omega_i$. Here $m_{e,i}$ and $\Omega_{e,i}$ are the masses and cyclotron frequencies of electrons and ions, respectively, and n_b is the density of ions in the ring.

When studying instabilities with $\gamma > \Omega_i$, the ions can be considered to be moving along straight lines. In this case the magnetic field does not influence the quasi-linear relaxation of the velocity distribution of the fast ions so that the quasi-linear equations for their distribution function are the same as that in the absence of magnetic field. The quasi-linear relaxation of such a beam due to the excitation of the lower-hybrid oscillations proceeds in exactly the same way as the quasi-linear relaxation of the electron beam due to high-frequency Langmuir waves (Shapiro and Shevchenko, 1988). In the general case this situation corresponds to the time variation of the particle distribution function in three-dimensional velocity space. For the case of $k_\parallel \to 0$, general quasi-linear equations can be integrated over $v_\parallel$ and $k_\parallel$ and the problem becomes that of investigating the time behaviour of a two-dimensional distribution $f(v_\perp, t)$ and a two-dimensional spectrum of oscillations $W_{\text{LH}}(k_\perp, t)$. Therefore, $f = f(v_\perp, t)$, i.e., the ion distribution function remains isotropic during the whole process of relaxation (see, Kulygin, Mikhailovskii, and Tsapelkin, 1971). The energy density of lower-hybrid waves excited by the isotropic distribution of trapped ions can be estimated (Breizmann, 1986) as

$$\frac{W_{\text{LH}}}{n_e T_e} \approx \frac{n_b m_i v_b^2}{n_e 2\Lambda T_e} q . \tag{31}$$

For $n_b/n_e \approx m_e/m_i$; $v_b = 5v_{T_i}$; $\Lambda \approx 12$, $q = 0.1$, and $T_e = T_i = 10^6$ K, we obtain $W_{\text{LH}}/n_e T_e \sim 10^{-4}$.

5. Conversion of Plasma Waves into Transverse Waves

5.1. Rayleigh scattering

As we noted earlier, the polarization and heights are generally taken as evidence that the emission process is a plasma process at or near the local plasma frequency. The emission at the second harmonic should be negligibly small. If the effective temperature of Langmuir waves $T_{\text{eff}}^{\text{L}} < T^* = 4 \times 10^{13}$ K, the emission at the second harmonic is predominant over that of the fundamental in the case of Rayleigh scattering of plasma waves on ion density fluctuations (Zaitsev and Stepanov, 1983).

Since the effective temperature of Langmuir waves emitted spontaneously by an anisotropic electron distribution is less than 4×10^{13} K (see Table III) and the radiation should be at the fundamental, most probably the conversion of Langmuir waves into radiation by scattering on ion density fluctuations does not play any role in the case of storm radiation, and the low-frequency waves in the source region may play the role of scatterers of these longitudinal waves into transverse waves (see Melrose, 1980a).

5.2. Coalescence of Langmuir waves with lower-hybrid waves

The transformation of Langmuir waves (L) into electromagnetic radiation (t) due to scattering on lower-hybrid waves (LH) takes place when the resonance conditions

$$\mathbf{k}_{\text{L}} + \mathbf{k}_{\text{LH}} = \mathbf{k}\,, \qquad \omega_{\text{L}}(\mathbf{k}_{\text{L}}) + \omega_{\text{LH}}(\mathbf{k}_{\text{LH}}) = \omega(\mathbf{k})\,, \tag{32}$$

are satisfied. The kinetic equations for the waves can be written in the form of transfer equations. The equation for the electromagnetic waves can be written as (Zheleznyakov, 1977, p. 413)

$$\frac{\partial W_{\mathbf{k}}}{\partial t} + v_g\,\frac{\partial W_{\mathbf{k}}}{\partial L} = \int \omega\, W\left(\frac{W_{\mathbf{k}_{\text{L}}} W_{\mathbf{k}_{\text{LH}}}}{\omega_{\text{L}}\omega_{\text{LH}}} - \frac{W_{\mathbf{k}_{\text{L}}}}{\omega_{\text{L}}\omega} - \frac{W_{\mathbf{k}_2} W_{\mathbf{k}}}{\omega_{\text{LH}}\omega}\right)\frac{\mathrm{d}^3\mathbf{k}_{\text{L}}\,\mathrm{d}^3\mathbf{k}_{\text{LH}}}{\hbar(2\pi)^3}\,. \tag{33}$$

Here $W_{\mathbf{k}}$, $W_{\mathbf{k}_{\text{L}}}$ and $W_{\mathbf{k}_{\text{LH}}}$ are the energy densities of t, L, and LH waves respectively, $\hbar$ is Planck's constant, and W is the probability of L + LH $\rightarrow t$ process and it is given as (Melrose, 1980b)

$$W = \frac{(2\pi)^5}{4}\,\frac{\hbar e^2}{m_e^2}\,\frac{\omega_{\text{LH}}^3(\mathbf{k}_{\text{LH}})}{k_{\text{LH}}^2 V_T^4}\,\frac{m_i}{m_e}\,\sin^2\theta\,\delta(\mathbf{k} - \mathbf{k}_{\text{L}} - \mathbf{k}_{\text{LH}})\,\delta(\omega - \omega_{\text{L}} - \omega_{\text{LH}})\,. \tag{34}$$

Here θ is the angle between $\mathbf{k}$ and $\mathbf{k}_{\text{L}}$. The first term on the right-hand side of Equation (33) corresponds to the emission coefficient of t-waves with wave vector $\mathbf{k}$ and can be written as

$$\alpha_{\mathbf{k}} = \int W W_{\mathbf{k}_{\text{L}}} \mathrm{W}_{k_{\text{LH}}}\,\frac{\omega}{\omega_{\text{L}}\omega_{\text{LH}}}\,\frac{\mathrm{d}^3 k_{\text{L}}\,\mathrm{d}^3 k_{\text{LH}}}{\hbar(2\pi)^3}\,. \tag{35}$$

The damping of t-waves due to decay into L and LH waves is given by the second and third terms as

$$2\gamma_{\mathbf{k}} = \int W\left(\frac{W_{\mathbf{k}_{\mathrm{L}}}}{\omega_{\mathrm{L}}} + \frac{W_{\mathbf{k}_{\mathrm{LH}}}}{\omega_{\mathrm{LH}}}\right)\frac{\mathrm{d}^3k_{\mathrm{L}}\,\mathrm{d}^3k_{\mathrm{LH}}}{\hbar(2\pi)^3} . \tag{36}$$

In a quasi-stationary large source, the energy density $W_{\mathbf{k}}$ can be estimated as follows. If one assumes that $W_{\mathbf{k}}$ and $W_{\mathbf{k}_{\mathrm{LH}}}$ change very little during $\Delta t \sim 1/\gamma_{\mathbf{k}}$, in a distance $\Delta L \sim 1/\mu_{\mathbf{k}}$ (where $\mu_{\mathbf{k}} = 2\gamma_{\mathbf{k}} v_g^{-1}$ is the absorption coefficient along the ray due to decay $t \to \mathrm{L} + \mathrm{LH}$), the energy density W_k is maintained at a constant level. Therefore,

$$\frac{\partial W_{\mathbf{k}}}{\partial t} + v_g \frac{\partial W_{\mathbf{k}}}{\partial L} \approx 0 . \tag{37}$$

Here v_g is the group velocity of the electromagnetic waves. If we express the energy densities in terms of temperatures as

$$W_{\mathbf{k}} = \frac{\chi T_{\mathrm{eff}}}{(2\pi)^3} , \qquad W_{\mathbf{k}_{\mathrm{L}}} = \frac{\chi T^{\mathrm{L}}}{(2\pi)^3} , \qquad W_{\mathbf{k}_{\mathrm{LH}}} = \frac{\chi T^{\mathrm{LH}}}{(2\pi)^3} , \tag{38}$$

where χ is Boltzmann's constant, and by equating the right-hand side of Equation (33) to zero, we obtain a relation

$$T_{\mathrm{eff}} = \frac{\omega T^{\mathrm{L}} T^{\mathrm{LH}}}{\omega_{\mathrm{L}} T^{\mathrm{LH}} + \omega_{\mathrm{LH}} T^{\mathrm{L}}} . \tag{39}$$

In the case of coalescence of Langmuir waves with lower-hybrid waves with $\omega_{\mathrm{L}} \gg \omega_{\mathrm{LH}}$, the effective temperature of electromagnetic waves T_{eff} with $\omega = \omega_{\mathrm{L}} + \omega_{\mathrm{LH}} \approx \omega_{\mathrm{L}}$ can have a maximum value equal to the effective temperature of the Langmuir waves (T^{L}) in a wide range of values $T^{\mathrm{LH}}/T^{\mathrm{L}} \geq \omega_{\mathrm{LH}}/\omega_{\mathrm{L}}$. It does not exceed T^{L} for any value of $T^{\mathrm{LH}}/T^{\mathrm{L}}$.

If one assumes that the spectrum of plasma waves $W_{\mathbf{k}}$ is isotropic and if the source is stationary and homogeneous, the effective temperature of the radiation T_{eff} is equal to the effective temperature of Langmuir waves T^{L} only if the condition $\mu L \gg 1$ is satisfied. Here L is the linear dimension of the source. In other words, T_{eff} raises to the maximum value equal to T^{L} if the source is optically thick with respect to the decay of electromagnetic wave (t) into Langmuir (L) and lower-hybrid (LH) waves. In the case of the optically thin case T_{eff} is determined by $\alpha_{\mathbf{k}}$. Therefore, to find T_{eff}, one has to estimate $\alpha_{\mathbf{k}}$ and $\gamma_{\mathbf{k}}$.

By substituting Equation (34) in Equations (35) and (36) and by assuming that $\omega_{\mathrm{LH}}/k_{\mathrm{LH}} \sim v_{bi}$, since the lower-hybrid waves are excited due to the resonance condition, $\omega_{\mathrm{LH}} \approx \mathbf{k}_{\mathrm{LH}} \cdot \mathbf{v}_{bi}$ is satisfied, and by writing $\mathrm{d}^3\mathbf{k}_{\mathrm{L}} = k_{\mathrm{L}}^2\,\mathrm{d}k_{\mathrm{L}} \sin\theta\,\mathrm{d}\theta\,\mathrm{d}\psi$, we can easily perform the integration in Equations (35) and (36), to obtain

$$\alpha_{\mathbf{k}} = \frac{(2\pi)^3}{9}\,\frac{e^2}{m_e^2}\,\frac{m_i}{m_e}\,\frac{v_{bi}^2\omega_p}{v_T^6}\,W_{\mathbf{k}_{\mathrm{L}}} W_{\mathbf{k}_{\mathrm{LH}}} k_{\mathrm{L}} , \tag{40}$$

and

$$\gamma_{\mathbf{k}} = \frac{(2\pi)^3}{18} \frac{e^2}{m_e^2} \frac{m_i}{m_e} \frac{v_{bi}^2}{v_T^4} \frac{\omega_p}{v_T^2} k_{\mathrm{L}} \left[\frac{\omega_{\mathrm{LH}}}{\omega_{\mathrm{L}}} W_{\mathbf{k}_{\mathrm{L}}} + W_{\mathbf{k}_{\mathrm{LH}}} \right]. \tag{41}$$

The emissivity α_ω and absorption coefficient μ are given by

$$\alpha_\omega = k^2 \alpha_{\mathbf{k}} (|\cos\theta|\, v_g)^{-1}, \tag{42}$$

and

$$\mu = 2\gamma_{\mathbf{k}}/v_g. \tag{43}$$

By taking for the plasma waves $k_{\mathrm{L}} \approx \omega_p/v_{ph}$ and for the electromagnetic waves $\omega \approx \omega_p$, $k \approx \omega_p/c$ and $v_g \approx c\sqrt{\Omega_e/\omega_p}$ and by noting that in an isotropic plasma $\cos\theta \approx 1$, we obtain

$$\alpha_\omega = \frac{(2\pi)^3}{9} \frac{e^2}{m_e^2} \frac{m_i}{m_e} \sqrt{(\Omega_e/\omega_p)} \frac{\omega_p^4}{c^3 v_{ph}} \frac{v_{bi}^2}{v_T^6} W_{\mathbf{k}_{\mathrm{L}}} W_{\mathbf{k}_{\mathrm{LH}}}, \tag{44}$$

and

$$\mu = \frac{(2\pi)^3}{9} \frac{e^2}{m_e^2} \frac{m_i}{m_e} \sqrt{(\omega_p/\Omega_e)} \frac{v_{bi}^2}{v_T^6} \frac{\omega_p^2}{c v_{ph}} \left[\frac{\omega_{\mathrm{LH}}}{\omega_{\mathrm{L}}} W_{\mathbf{k}_{\mathrm{L}}} + W_{\mathbf{k}_{\mathrm{LH}}} \right]. \tag{45}$$

By using the relations (38) we can write

$$\mu L = \frac{1}{9} \frac{e^2}{m_e^2} \frac{m_i}{m_e} \chi \sqrt{(\omega_p/\Omega_e)} \frac{v_{bi}^2}{v_T^6} \frac{\omega_p^2}{v_{ph}} \frac{L}{c} \left[\frac{\omega_{\mathrm{LH}}}{\omega_{\mathrm{L}}} T_{\mathrm{L}} + T_{\mathrm{LH}} \right]. \tag{46}$$

For $(\omega_{\mathrm{LH}}/\omega_{\mathrm{L}})T_L \ll T_{\mathrm{LH}}$, and for $\omega_p/\Omega_e \sim 5$, $v_{bi}^2 \sim 5v_{T_i}^2$, $v_T \sim 3.86 \times 10^8 \text{ cm s}^{-1}$, $v_{ph} \sim 10^{10} \text{ cm s}^{-1}$, $\omega_p = 2\pi \times 10^8 \text{ s}^{-1}$, and by taking into account that the absorption of electromagnetic waves due to decay into two longitudinal waves is significant only in a layer of thickness $L \sim 10^8$ cm (see Zheleznyakov and Zlotnik, 1974), we can estimate the lower limit of lower hybrid waves when $T_{\mathrm{eff}} \approx T^{\mathrm{L}}$ as 3.53×10^9 K. It agrees with the estimates given by Melrose (1980b) and Wentzel (1986). If $T^{\mathrm{L}} < 3.53 \times 10^9$ K, the effective temperature T_{eff} is given by

$$T_{\mathrm{eff}} = (2\pi)^3 \frac{c^2}{\chi \omega_p^2} a_\omega L. \tag{47}$$

By substituting Equation (44) in (47) and using the relations (38), we obtain

$$T_{\mathrm{eff}} = \frac{1}{9} \frac{e^2}{m_e^2} \frac{m_i}{m_e} \sqrt{\frac{\omega_p}{\Omega_e}} \frac{\omega_p^2}{c v_{ph}} \frac{v_{bi}^2}{v_T^6} L \times T^L T^{\mathrm{LH}}. \tag{48}$$

For the same parameters as above and for $T_{\mathrm{LH}} \sim 10^8$ K we obtain $T_{\mathrm{eff}} \sim T_{\mathrm{L}}/10$. Since the effective temperature estimated from Equation (31) exceeds 3.5×10^9 K by several orders of magnitude the effective temperature of radiation, in other words the brightness temperature is equal to the effective temperature of Langmuir waves (T^{L}) in the present case.

6. The Brightness Temperature of the Radiation

As we stated earlier, since the optical depth in the scattering process is very large, the brightness temperature (T_B) is equal to the effective temperature of the Langmuir waves. Therefore, we can write

$$T^{\mathrm{L}} = T_B = \frac{\alpha^{\mathrm{L}}}{\gamma^{\mathrm{L}}} \, . \tag{49}$$

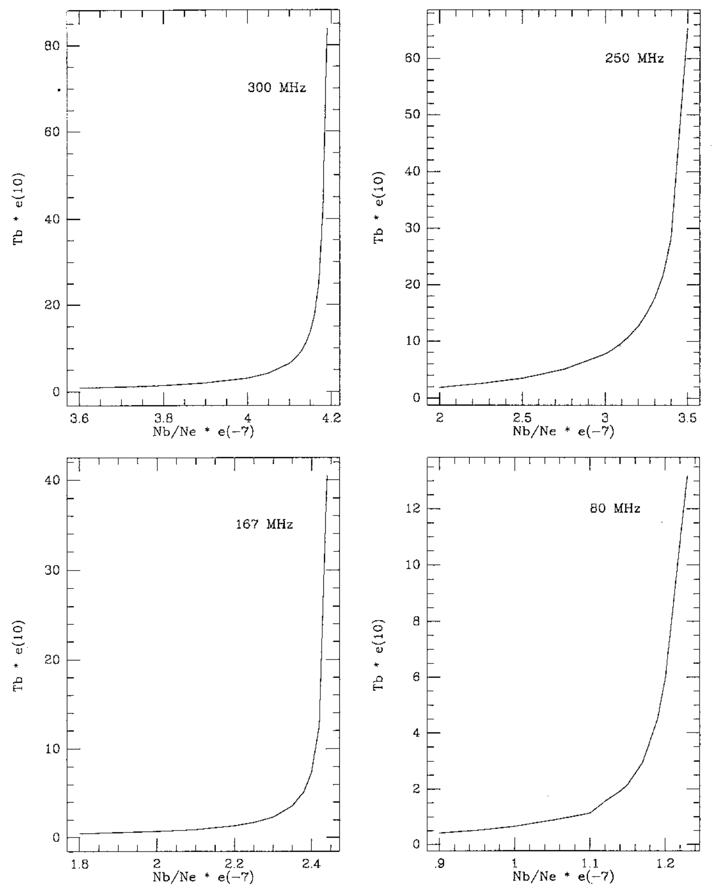

Fig. 1.

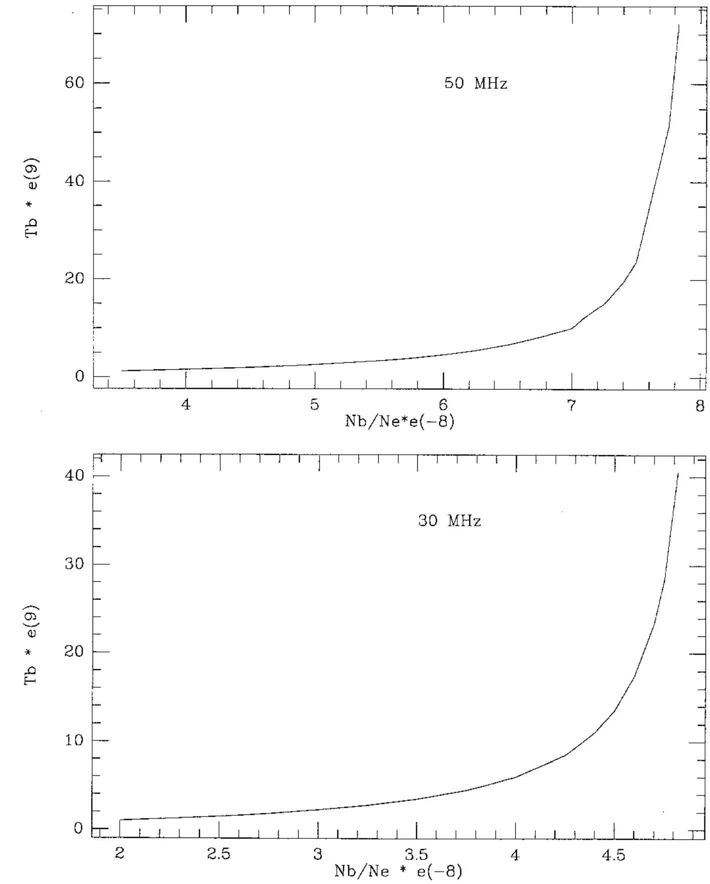

Fig. 1. The brightness temperature as a function of number density of the suprathermal electrons in the loss-cone distribution at 300, 250, 167, 80, 50, and 30 MHz frequencies.

By using Equations (3), (11), and (14) we can rewrite the above equation as

$$T_B \approx \frac{0.1 \dfrac{n_s}{n_e} \omega v_{T_b}^2 m_e}{-4.4 \times 10^{-2} \dfrac{n_s}{n_e} \omega_p + \nu_c + \gamma_L} . \tag{50}$$

For $v_{T_b} = 10^{10}$ cm s^{-1}, and using the Tables I and II for ν_c and γ_L respectively, we can plot in Figure 1, T_B as a function of n_s/n_e. From Figure 1, one can notice that T_B rises sharply when n_s/n_e approaches the threshold values. For an increase of 0.5×10^{-7} in n_s/n_e near the threshold, the increase in T_B is by two orders of magnitude. Equation (50) is valid only for $n_s/n_e \leqq (n_s/n_e)_{th}$. Beyond threshold values, one should solve the quasi-linear equations to obtain T_B, since the linear analysis is not valid for $-\gamma_A > \nu_c$.

One should note that the solution for the transfer equation, in the limit of the optically thick case ($\mu L \gg 1$) Equation (28) is valid even at threshold. For a source of linear dimensions $L \sim 5 \times 10^{-2} R_{\odot}$, which is typical for type I bursts, the condition $\mu L \gg 1$ is always satisfied for all the frequencies. One should also note that as n_s/n_e increases, the effective damping γ approaches zero, and solution (28) is not valid. There is a sort of stimulated emission of Langmuir waves at this limit even though the net damping is greater than zero, similar to the critical fluctuation temperature in a non-equilibrium plasma. We have shown that the brightness temperatures of type I bursts up to 10^{11} K can be explained by mere spontaneous emission. As the density of the beam is increased, making $-\gamma_A > \nu_c$, the emission is no longer incoherent but coherent, and it can give high brightness temperatures explaining the very bright type I bursts.

7. Discussion

As seen in Figure 1, there is a steady emission for number densities n_s/n_e less than 4×10^{-7}, 3×10^{-7}, 2.2×10^{-7}, 1×10^{-7}, 6×10^{-8}, and 4×10^{-8} at 300, 250, 167, 80, 50, and 30 MHz, respectively, giving rise to steady emission of storm continuum. As the density n_s/n_e approaches the threshold values at these frequencies, the brightness temperature T_B increases sharply by two orders of magnitude giving rise to intense type I bursts. The threshold condition may be satisfied randomly in space and time which depends on the acceleration or injection mechanism of electrons into the source region.

Weak shocks driven by emerging magnetic flux are supposed to be responsible for the high energy electrons in the source region as well as the low-frequency lower-hybrid waves (Spicer, Benz, and Huba, 1981; Wentzel, 1981). The acceleration of electrons by collisionless shocks through the lower-hybrid waves was discussed by Vaisberg *et al.* (1983), Krasnoselkikh *et al.* (1985), Thejappa (1987), Lampe and Papadopoulos (1977), and Benz and Thejappa (1988). The acceleration is a stochastic process, i.e., the injection of particles into the source region takes place stochastically making both n_s and v_b stochastic.

The low-frequency lower-hybrid waves can be excited easily either by shocks due to the modified two-stream instability whose energy density for typical type I parameters is $W_{\mathrm{LH}}/n_e T_e \simeq 10^{-6}$ (Spicer, Benz, and Huba, 1981), or as stated earlier, by the trapped distribution of energetic ions.

In Figure 2, we plot the brightness temperature (T_B) of the type I burst (i.e., at threshold) as a function of the frequency as predicted in this model for a strong storm. One can see from Figure 2 that the brightness temperature increases with the frequency which is one of the predictions of our model. It can be tested by measuring T_b of type I bursts at various frequencies. The brightness temperature of the continuum can lie anywhere from 10^7 K to 10^{10} K at all the frequencies. By knowing the brightness temperature of the continuum at a given frequency, one can estimate the number density of the suprathermal electrons trapped in the field lines by using Equation (50). In Figure 3, we plot the threshold beam density n_s/n_e as a function of frequency. One can easily see the linear relationship. By knowing the frequency of observations of type I

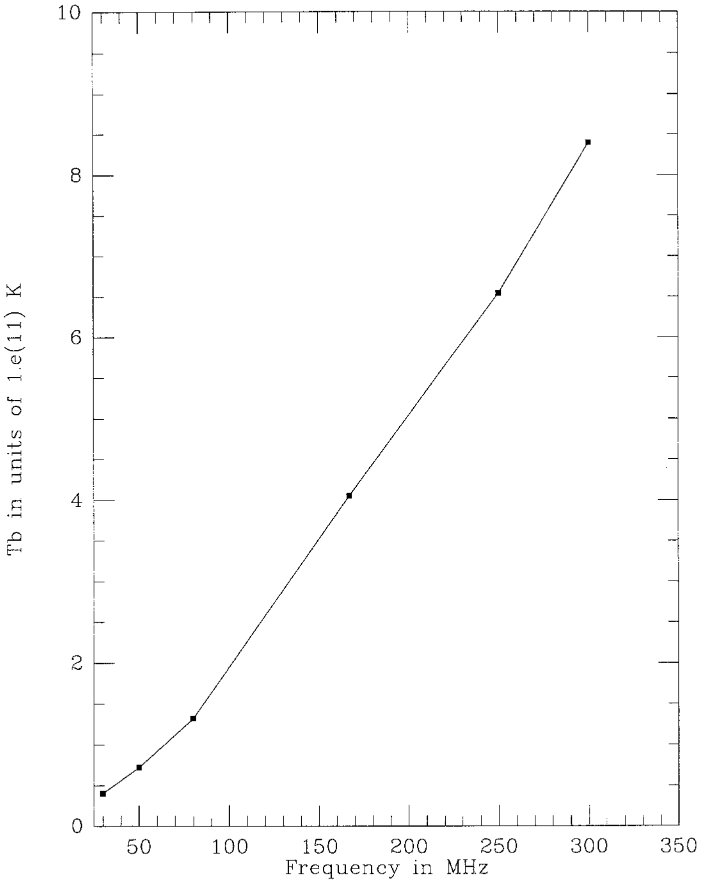

Fig. 2. The brightness temperature spectrum for type I bursts.

bursts, we can estimate the threshold density of the beam when $\nu_c = -\gamma_A$, and can estimate their brightness temperatures since it is difficult to measure them because of the difficulty in estimating the source size.

As we noted earlier, since the electrons are accelerated stochastically by the shocks, they are injected into the source region above the active regions stochastically. The electrons are also lost due to collisions and also escape into the loss-cone. Therefore, the threshold condition for the number density of particles n_s/n_e in the source region is satisfied randomly, explaining in a natural way the random appearance of type I bursts superposed on continuum. At all other time, the suprathermal electrons emit steady background continuum.

As we have noted from Figure 3, the threshold density of the suprathermal particles decreases steadily with frequency. The threshold condition is purely dependent on the number density (n_s/n_e) and a slight decrease in n_s/n_e violates the threshold conditions, making T_b equal to the temperature of the continuum. Most probably the fluctuations in n_s or n_e, i.e., δn_s or δn_e, determine the bandwidth and lifetime of type I bursts. In the

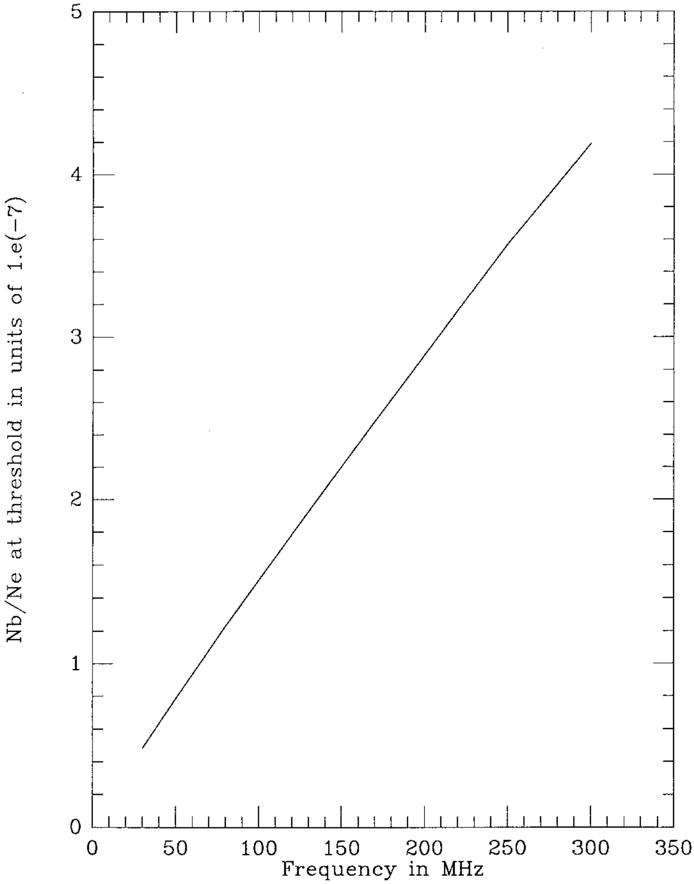

Fig. 3. Threshold density of the energetic electrons in the loss-cone distribution as a function of frequency.

absence of an electric field, we can write the kinetic equation for the nonthermal particles as

$$\frac{\partial f_s}{\partial t} = -\nu_p f_s, \tag{51}$$

which can be written as

$$\frac{(-\partial n_s)}{n_s} \sim \frac{(-\partial f_s)}{f_s} \sim \nu_p \, \delta t, \tag{52}$$

where the effective collisions

$$\nu_p = 3\sqrt{\frac{\pi}{2}\frac{v_T^3}{v_z^2}}\;\nu_c. \tag{53}$$

At 167 MHz frequency, the fluctuations in the density of nonthermal particles are $\delta n_s/n_s \sim 0.02$ when the brightness temperature of the burst is 10 times that of the continuum. For $v_z \approx 10^9$ cm s^{-1}, $v_T \approx 3.89 \times 10^8$ cm s^{-1}, and $\nu_c \sim 11.26$ s^{-1}, we obtain $\nu_p = 2.47$ s^{-1}. Therefore the lifetime of the fluctuation at 167 MHz is $\sim 1.0 \times 10^{-2}$ s, and the distance travelled by the density enhancement δn_s is $S \sim 100$ km, which corresponds to the bandwidth of $\partial f_p/f_p \sim 0.2\%$. Here we have used the relation $\partial f_p/f_p L_N/2 \sim S$, where L_N is the characteristic distance over which the mean electron density in the corona changes and it is $\sim 10^5$ km. Therefore, by knowing the continuum level and the peak T_B of the burst, we can estimate the bandwidth and the lifetime of the type I bursts.

For type I bursts, n_s/n_e is not always necessarily at threshold values; it can be anywhere near the threshold values. At decameter wavelengths the suprathermal particles escape and excite the type III storms. Since the observed brightness temperature of type III storm bursts is lower than the flare related type III bursts, they may be microbursts (see Kundu *et al.*, 1986; White, Kundu, and Szabo, 1986; Thejappa, Gopalswamy, and Kundu, 1990). As reported by Thejappa, Gopalswamy, and Kundu (1990) the microbursts are due to spontaneous emission of Langmuir waves by the electron beams with the threshold density. The conversion is most probably due to ion-density fluctuations in the case of microbursts. As discussed by Thejappa, Gopalswamy, and Kundu (1990), there is a clear distinction between the storm type III's and flare-related type III bursts. In the case of normal type III bursts, Langmuir waves are emitted coherently through the two-stream instability whereas microbursts are exclusively due to incoherent emission of Langmuir waves by the electron beams.

8. Conclusions

(1) The spontaneous emission of Langmuir waves by the anisotropic distribution of electrons in the closed magnetic fields above active regions is most probably the emission mechanism for the storm radiation.

(2) The conversion of Langmuir waves into radiation is most probably by their coalescence with lower-hybrid waves excited either by weak shocks or by fast ions injected across the field lines into the magnetic loops. The background ion density fluctuations do not play the role of scatterers in the case of storm radiation, since in this case the second harmonic emission will be predominant over the fundamental, contrary to observations.

(3) Whenever the source region is filled with suprathermal particles, there is a steady enhancement of Langmuir waves due to spontaneous emission, explaining the wideband steady continuum. For typical parameters, the brightness temperature of the continuum is in the range 10^7 K to 10^{10} K and it can be explained with the number density of the suprathermal particles $n_s/n_e \sim 10^{-9}$–10^{-7} at all the observed frequencies. Usually the brightness temperature decreases with frequency for a particular storm at a given time.

(4) Whenever the collisional damping is equal to the negative damping due to anisotropic distribution of electrons, there is a steep rise in the brightness temperature giving

rise to type I bursts. The random nature of the threshold condition naturally explains the randomness of the type I bursts. The brightness temperature spectrum of type I bursts shows that T_B increases with frequency.

(5) The finite bandwidth and short lifetime of type I bursts are the spatial and temporal scales of the random density fluctuations of the suprathermal particles in the source region.

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