Kodaikanal Obs. Bull. (1986) 6

A MODEL OF THE GENERATION MECHANISM OF TYPE II SOLAR RADIO BURSTS

G. Thejappa
INDIAN INSTITUTE OF ASTROPHYSICS
BANGALORE 560 034

Abstract

The proposed model for the generation mechanism of type II solar radio bursts is based on the fact that a small fraction of ions is reflected from the shock front and tend to evolve into a 'ring' in the downstream region behind the overshoot. This ring type ion distribution drive low frequency waves ($\Omega_i < \omega < \Omega_e$) unstable which effectively accelerate the magnetized electrons of the ambient plasma to very high energies along the field lines since their phase velocities parallel to magnetic fields are much higher than the electron thermal velocity in the solar corona. The distribution function of the accelerated electrons is calculated in the present paper assuming that Dory-Guest-Harris distribution correctly describes the 'ring' nature of the reflected ions. It is proposed that the frequency splitting in the observed radiation both at fundamental and at second harmonic is due to the non-linear scattering of electron-beam excited Langmuir and upper-hybrid waves parallel and perpendicular to the magnetic field on the ion-beam-excited whistler and lower-hybrid waves respectively into left-handed circularly polarized electromagnetic waves. The frequency splitting is approximately equal to electron cyclotron frequency.

1. Introduction

The type II solar radio bursts appear in the dynamic spectrum as narrow bands of emission that drift slowly to lower frequencies. Pickel'ner and Gitzburg (1963) were first to attribute type II radio bursts to outward moving shocks that excite radio emission at the fundamental and second harmonic of the local plasma frequency. Using the radial variation of the electron density in the corona, one can calculate the shock velocities from the observed frequency drift. They lie in the range from 200 km/s to 2000 km/s.

The observations of herringbone structures, and type III bursts associated with Type II bursts give evidence that a population of superthermal streaming electrons associated with the shock wave is necessary for generation of the type II radio bursts. Papadopoulos (1981) had proposed that the low frequency turbulence driven by reflected ions from the shock front can accelerate electrons to large energies producing an enhanced level of Langmuir waves.

Krasnoselskikh et al (1985) considered that the reflected ions behave like a beam in the shock front (foot and the ramp) and approximated them as a drifted Maxwellian. The growth rate of the low frequency waves was found to be $\gamma_{max} \sim \omega_{LH} (V_{Ai}/\Delta V_b)^2 n_b/n_o$ at $k \sim \omega_p/c$ and subsequently calculated the distribution function of the superthermal electrons and the energy density of the excited high frequency plasma waves.

Considering that Dory-Guest-Harris distribution (Dory et al. 1965) correctly represents the ring nature of the reflected ions in the downstream, we calculate the distribution function of the accelerated electrons. We show that these superthermal electron tails drive Langmuir (upper-hybrid) waves resonantly unstable. The excited Langmuir and upper hybrid turbulence explains the observed fine structure of the radiation when it is nonlinearly scattered on low frequency turbulence excited by the reflected ions in two limiting cases i.e., on Whistlers and lower hybrid waves respectively. It is also predicted that the polarization is in the ordinary sense.

2. Excitation of Lower-Hybrid Oscillations and Electron Acceleration

It is generally known that ion reflection occurs in a magnetosonic shock wave when its Mach number exceeds the so-called critical Mach number (Tidman and Krall, 1971). These ions form a beam of gyrating particles escapting along the magnetic field lines into the upstream region. The longitudinal phase velocity of the oblique Langmuir waves excited by the reflected ion beam can easily reach or even exceed the velocity of light, they can effectively accelerate electrons upto ultrarelativistic velocities (Vaisberg et a. 1983; Galeev, 1984; Krasnoselkikh et al. 1985).

2.1 Wave Excitation

Leroy et al. (1981, 1982) showed that the reflected ions tend to form a gyrating stream in the down stream region behind the overshoot and evolve into a 'ring' with a significantly large velocity spread. A realistic representation of the ring type distribution of reflected ions which includes all thermal effects is given by the Dory-Guest-Harris distribution (Dory et al. 1965).

$$f_b(v) = \frac{1}{\pi^{3/2}(N+1)!}\,\frac{1}{V_{Tz}V_{T\perp}^2}\left(\frac{V_\perp}{V_{T\perp}}\right)^{2N} e^{-\left(\frac{V_\perp}{V_{T\perp}}\right)} e^{-\left(\frac{V_\perp}{V_{T\perp}}\right)^2} \tag{1}$$

$V_{T\perp}(V_{Tz})$ is the thermal velocity $\perp$ ($\parallel$) to $\vec{B}$, N is a parameter which measures the anisotropy of the distribution. n_b is density of the reflected ions in the background plasma with ion density n_0. The distribution function resembles a ring or torus in V-space, when N=0 it reduces to a Maxwellian; and when N (the ring anisotropy) is large, the perpendicular energy is concentrated near the maximum.

$$V_{\perp max} = N^{1/2}\, V_{T\perp} \equiv V \;. \tag{2}$$

We take V to be the perpendicular speed of reflected ions, approximately 2000 km/s $\approx 2MV_A$, where M is Mach Number, V_A is the Alfven velocity. The distribution function has many of the features expected for monoenergetic ions reflected from the shock front.

The general dispersion equation for the waves excited by the reflected ion beam and absorbed by the background suprathermal electron beam is

$$\epsilon(\omega,\vec{k}) = 1 + \frac{\omega_{pe}^2}{\Omega_e^2}\left(1+\frac{\omega_{pe}^2}{k^2c^2}\right)\frac{k_\perp^2}{k^2} - \frac{\omega_{pi}^2}{\omega^2} - \frac{\omega_{pe}^2}{\omega^2}\,\frac{k_\parallel^2}{k^2}\,\frac{1}{\left(1+\frac{\omega_{pe}^2}{k^2c^2}\right)} + \frac{\omega_{pi}^2}{k^2 n_0}\int \frac{\vec{k}\cdot\frac{\partial f_b}{d\vec{v}}\, d^3\vec{v}}{\omega - \vec{k}\cdot\vec{v} + i0} + \frac{\omega_{pe}^3}{k^2 n_0}\int \frac{k_\parallel \frac{\partial f_e}{\partial v_\parallel}\, dV_\parallel}{\omega - k_\parallel v_\parallel + i0}\;, \tag{3}$$

where $\omega_{pj}^2 = \frac{4\pi n_0 e^2}{m_j}$ is the plasma frequency of j-sort particles; ω is the frequency, $\vec{k}$ is the wave vector; $K_\parallel$ and $K_\perp$ are the components of the above vector, along and across the undistributed magnetic field, $f_b(v)$ is the ion-beam distribution function; $f_e(v_\parallel)$ is the distribution function of suprathermal electrons over longitudinal velocities. The above equation is derived under the assumption that the oscillation frequencies and wave vectors are in the range:

$$\Omega_e \gg \omega \gg \Omega_i \ , \qquad I_m\omega \gg \Omega_i \ ,$$

$$K_\perp r_{Li} \gg 1 \gg K_\perp r_{Le} \ , \qquad \omega \gg K V_{Ti} \ , \qquad \omega \gg K_\parallel V_{Te} \ , \tag{4}$$

where $V_{Tj} = \sqrt{2T_j/m_j}$ is the thermal velocity of J-sort particles, $r_{Lj} = V_{Tj}/\Omega_j$ is the Larmor radius of J-sort particles. The above inequalities make it possible to neglect the effect of the magnetic field on the motion of ions and employ a drift approximation to describe electrons. The dielectric constant computed for the background plasma is

$$\varepsilon_0 = 1 + \frac{\omega_{pe}^2}{\Omega_e^2}\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right)\frac{K_\perp^2}{K^2} - \frac{\omega_{pi}^2}{\omega^2} - \frac{\omega_{pe}^2}{\omega^2}\,\frac{K_\parallel^2}{K^2\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right)} \tag{5}$$

The solution of $\varepsilon_0 = 0$ is

$$\omega^2 = \frac{\omega_{LH}^2}{\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right)}\left[1 + \frac{m_i}{m_e}\,\frac{K_\parallel^2}{K^2\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right)}\right] \tag{6}$$

In solar coronal conditions where $\omega_{pe}^2 \gg \Omega_e^2$ is satisfied, the short-wave oscillations $kc \gg \omega_{pe}$ are electrostatic and their frequency is equal (at $k_\parallel =$) to higher (at $k_\parallel \neq 0$) than, the lower-hybrid resonance frequency ω_{LH}. With wave-length increasing the non-potentiality of oscillations becomes essential, and they turn into well-known Whistlers.

The derivative of ε_0 with respect to ω is given by

$$\frac{\partial \varepsilon_0}{\partial \omega} = \frac{2}{\omega}\left[\frac{\omega_{pi}^2}{\omega^2} + \frac{\omega_{pe}^2}{\omega^2}\,\frac{k_\parallel^2}{k^2}\,\frac{1}{\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right)}\right] \tag{7}$$

and we get using the equation (6)

$$\omega \frac{\partial \varepsilon_0}{\partial \omega} = 2\,\frac{\omega_{pe}^2}{\Omega_e^2}\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right) \tag{8}$$

When the DGH distribution of reflected ions are added to the background plasma the above oscillations are driven unstable when $N > 0$.

The imaginary part of the dielectric constant due to the presence of DGH distribution of ions in the ambient plasma is given by

$$\mathrm{Im}\,\varepsilon_i = \frac{4\omega_{Pb}^2 R^3 e^{-R^2}}{(N+1)!\,\omega^2} \int_0^\infty dz e^{-z^2}(R^2+z^2)^N \left(1 - \frac{N}{R^2+z^2}\right) \quad , \tag{9}$$

where $R = \frac{\omega}{k_\perp V_{T\perp}}$. [Barbosa et al. 1985]

Therefore

$$\mathrm{I_m}\varepsilon_i(N=1) = 2\sqrt{\pi}\ \frac{\omega_{Pb}^2}{\omega^2} R^3 e^{-R^2}(R^2 - \tfrac{1}{2}) \quad , \tag{10}$$

$$\mathrm{I_m}\varepsilon_i(N=2) = \sqrt{\pi}\ \frac{\omega_{Pb}^2}{\omega^2} R^3 e^{-R^2}(R^4 - R^2 - \tfrac{1}{2}) \quad . \tag{11}$$

For small growth or damping such that $\gamma/\omega \ll 1$, we may use the expression

$$\gamma = -\frac{\mathrm{Im}\varepsilon}{\partial\varepsilon_0/\partial\omega} \tag{12}$$

Here, a necessary condition for growth is

$$\frac{\omega}{k_\perp} < N^{1/2}\, V_\perp = V \quad , \tag{13}$$

for instability on the positive slope of the beam, which is

$$\frac{\gamma_b}{\omega} = -\frac{2}{(N+1)!}\ \frac{n_b}{n_o}\ \frac{m_e}{m_i}\ \frac{\Omega_e^2}{\omega^2}\ \frac{1}{\left(1 + \frac{\omega_{pe}^2}{k^2c^2}\right)} R^3 e^{-R^2} \int dz e^{-z^2} \tag{14}$$

2.2 Electron Acceleration

The imaginary part of the dielectric constant due to damping of the waves by the ambient electrons is given by

$$\mathrm{Im}\varepsilon_e = -\frac{\pi\omega_{pe}^2}{n_o k^2}\ \frac{\partial f_e}{\partial V_\parallel} \tag{15}$$

The damping of waves due to the non Maxwellian tail of the electron distribution function is given by

$$\frac{\gamma_e}{\omega} = \frac{\pi}{2}\ \frac{\Omega_e^2}{\left(1+\frac{\omega_{pe}^2}{k^2c^2}\right)}\ \frac{1}{n_o k^2}\ \frac{\partial f_e}{\partial V_\parallel}\Bigg|_{V_\parallel = \omega/k_\parallel} \tag{16}$$

In order to find the main parameters of accelerated electrons it is sufficient to onsider the wave balance equation i.e.

$$(\vec{V}_{gz}-\vec{V}_o)\cdot\frac{\partial}{\partial z}E_k^2 = 2E_k^2(\gamma_b+\gamma_e+\gamma_i) \quad , \tag{17}$$

here $\vec{V}_{gz}$ is the group velocity of the excited waves along the magnetic field, and V_o is the velocity of the shock. $|E_k^2/8\pi|$ is the energy density of the excited low freuency waves.

In the limit $\omega_{pe}^2/k^2c^2 \ll 1$,

$$(V_{gz}-V_A)\cdot\frac{\partial}{\partial z}E_k^2 \approx \frac{\alpha V_o\Omega_i}{V_o}E_k^2 = \alpha\Omega_i E_k^2 \quad , \tag{18}$$

where $\alpha \leq 1$.

Therefore the equation (16) can be written as

$$(\gamma_b+\gamma_e+\gamma_i-\alpha\Omega_i)E_k^2 = 0 \tag{19}$$

Since $E_k^2 \neq 0$, and $(\gamma_i-\alpha\Omega_i) \ll \gamma_b+\gamma_e$, we have

$$(\gamma_b+\gamma_e) = 0 \quad . \tag{20}$$

Actually $\gamma_e \geq \gamma_b$ so that the waves are absorbed in the plasma. The equality is satisfied only for one K where the energy spectrum of the oscillations is streamer-type and is concentrated in a narrow band near the line $k=k_{max}$ on the phase plane $(K_\perp, K_z)$. Since the growth rate and damping of waves due to ion beam and accelerated electrons are given respectively by

$$\gamma_b = -\frac{\mathrm{Im}\varepsilon_i}{\omega(\partial\varepsilon_0/\partial\omega)} \;;\quad \gamma_e = -\frac{\mathrm{Im}\varepsilon_e}{\omega(\partial\varepsilon_0/\partial\omega)} \quad , \tag{21}$$

By equating them we get,

$$\mathrm{Im}\varepsilon_0 = \mathrm{Im}\varepsilon_i \quad . \tag{23}$$

Using the equations (9) and (16), we get

$$\frac{\partial f_e}{\partial V_\parallel} = -\frac{n_0k^2}{\pi\omega_{pe}^2}\,\mathrm{Im}\varepsilon_i \quad . \tag{24}$$

By integrating the above equation we obtain

$$f_e = -\frac{n_0k^2}{\pi\omega_{pe}^2}\,\mathrm{Im}\varepsilon_i[V_{max}-|V_\parallel|] \quad . \tag{25}$$

By substituting the expressions (10) and (11) in the above equation, we obtain

$$f_e(N=1) = \frac{n_{a1}}{V_{T_\perp}^2}\,(V_{max} - |V_\parallel|); \qquad f_e(N=2) = \frac{n_{a2}}{V_{T_\perp}^2}\,(V_{max} - |V_\parallel|) \;, \tag{26}$$

where

$$n_{a_1} = \frac{2}{\sqrt{\pi}}\left(\frac{m_e}{m_i}\right) n_b R e^{-R^2}(R - \tfrac{1}{2}) \;,$$

$$n_{a_2} = \frac{1}{\sqrt{\pi}}\left(\frac{m_e}{m_i}\right) n_b R e^{-R^2}(R^4 - R^2 - \tfrac{1}{4}) \;, \tag{28}$$

The value $\alpha = n_{a_2}/n_{a_1}$ is given by

$$\frac{n_{a_2}}{n_{a_1}} = \frac{1}{2}\,\frac{R^4 - R^2 - 1/4}{R^2 - 1/2} \;, \tag{29}$$

The maximum velocity acquired by the accelerated electrons V_{max} is calculated by condition that the energy flux of ion beam should be more than the energy flux of the accelerated electrons; i.e.,

$$n_b m_i V_b^2 \,\Delta V_b \le n_a \varepsilon_e (\varepsilon_e/m_e)^{1/2} \;; \qquad \frac{n_b}{n_o}\left(\frac{V_{Ai}}{\Delta V_b}\right)^2 \simeq \frac{n_a}{n_o}\,\frac{m_e V_{Ae}^2}{\varepsilon_e} \;, \tag{30}$$

where ε_e is kinetic energy of electrons, which is

$$\varepsilon_0 = \frac{m_e V_{max}^2}{2} = \left(\frac{m_e}{m_i}\right)^{1/5} (m_i V_b^2)^{2/5} (m_i \,\Delta V_b^2)^{3/5} \;. \tag{32}$$

For $V_b = 2000$ km/s; $\Delta V_b = 0.2\,V_b$; $V_{max} = 1.55 \times 10^9$ cms/sec.

3. Radiation Process

In the process of electron acceleration it is only the energy of their motion along magnetic field lines that increases, the resulting electron distribution turns to be very anisotropic. Such non Maxwellian electron tails drive Langmuir waves and upper-hybrid waves resonantly unstable in the parallel and perpendicular directions of the ambient magnetic field.

We propose that the nonlinear scattering of the Langmuir waves and UH waves into Left handed polarized electromagnetic waves on Whistlers and Lower hybrid waves explain the observed type II radiation at the fundamental plasma frequency. The frequency difference inthe above two processes will be of the order of electron cyclotron frequency (Ω_e), which explains the observed frequency splitting of the type II radiation at the fundamental frequency.

The radiation at second harmonic of the plasma frequency and the fine structure are explained due to the combination scattering of Langmuir waves and upper hybrid waves on themselves.

There is an equilibrium between the electromagnetic and Langmuir waves, and the brightness temperature of the radio-emission is equal to the effective temperature of Langmuir waves:

$$T_b \sim T_{eff} \sim T_e(n_o\lambda_D^3)\frac{W_L}{n_oT_e} \quad . \tag{33}$$

For $n_o=10^8$, $T_e=10^6$ K, one obtains $T_b \sim 10^{11}$ K.

The optical depth in the case of second harmonic emission also is very large. Therefore the brightness temperature of the harmonic emission is appeared to be approximately equal to that of fundamental emission. The wave-wave interactions favour the ordinary electromagnetic waves with left-handed polarization for the fundamental. The polarization of the second harmonic in the region of generation is connected with the coupling of two Langmuir waves.

4. Conclusions

The proposed model of the type II radio emission consists of the following steps

1. The reflected ions behind the overshoot are described by the ring distribution.
2. The oblique Langmuir waves excited by these reflected ions accelerate portion of background electrons along the field lines due to the anisotropy in their phase velocities.
3. The accelerated electrons drive Langmuir waves and upper hybrid waves resonantly whose energy levels may be

$$\left(\frac{W_{L(UH)}}{n_oT_e} \sim 10^{-4} \div 10^{-5}\right)$$

4. The Langmuir waves are scattered by ion beam excited Whistler waves (lower hybrid waves) and converted into the ordinary left-handed polarized waves at the same frequency $\omega \approx \omega_{pe}$. The radiation at second harmonic is due to the combination scattering of Langmuir waves (UH-waves). The observed frequency splitting is naturally explained by the above process. For the typical parameters of the coronal plasma, the brightness temperature at the fundamental and second harmonic is of the order of 10^{11} K.

References

Barbosa, D.D., Coroniti, F.V., Kurth, W.S. and Scarf, F.L. 1985. Astrophys. J. 289, 392.
Dory, R.A., Guest, G.I. and Harris, E.G. 1965. Phys. Rev. Letters, **14**, 131.
Galeev, A.A. 1984. Exper. Teor. Fiz., **86**, 1655.
Krasnosels'kikh, V.V., Kruchina, E.N., Thejappa, G. and Volokitin, A.S. 1985. Astron. Astrophys. **149**, 323.
Leroy, M.M., Goodrich, C.C., Winske, D., Wu, C.S. and Papadopoulos, K. 1981. Geophys. Res. Lett. **8**, 1268.
Leroy, M.M., Goodrich, C.C., Winske, D., Wu, C.S. and Papadopolos, K. 1982. JGR, **87**, 5081.
Papadopolous, K. in Plasma Astrophysics, ESA Sp. 161:1981, p.313.
Pikelner, S.B. and Ginzburg, M.A. 1963. Astron. Zh. **40**, 842.
Tidman, D.A. and Krall, N.A. 1971. Shock waves in collisionless plasmas, Wiley-Interscience, Newyork.
Vaisberg, O.L., Galeev, A.A., Zastenker, G.N., Klimov, S.I., Norzdrachev, M.M., Sagdeev, R.Z., Sokolov, A. and Shapiro, A. 1983. Zhur. Exper. Teor. Fiz., **85**, 1232.