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# A Theory of Solar Type III Radio Bursts

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## ABSTRACT

A theory of type III bursts is reviewed. Energetic electrons propagating through the interplanetary medium are shown to excite the one dimensional oscillating two stream instability (OTSI). The OTSI is in turn stabilized by anomalous resistivity which completes the transfer of long wavelength Langmuir waves to short wavelengths, out of resonance with the electrons. The theory explains the small energy losses suffered by the electrons in propagating to 1 AU, the predominance of second harmonic radiation, and the observed correlation between radio and electron fluxes.

## INTRODUCTION

Solar type III radio bursts have been studied for more than 30 years. The persistent interest in this phenomenon has been due in no small part to the theoretical difficulties encountered in constructing a convincing interpretation of many of the most striking properties of the bursts. Several basic questions were posed by Sturrock fifteen years ago, and are only now beginning to be answered. Among the issues he raised (Sturrock, 1964) are three that will be discussed in some detail in this brief review. First, why is the electron beam that excites the bursts not significantly decelerated; second, why is the radiation predominately emitted at the second harmonic of the local plasma frequency, $f_{pe} = \omega_e/2\pi$; and third, why does the beam have such a well defined velocity, typically between 0.3 and 0.2c.

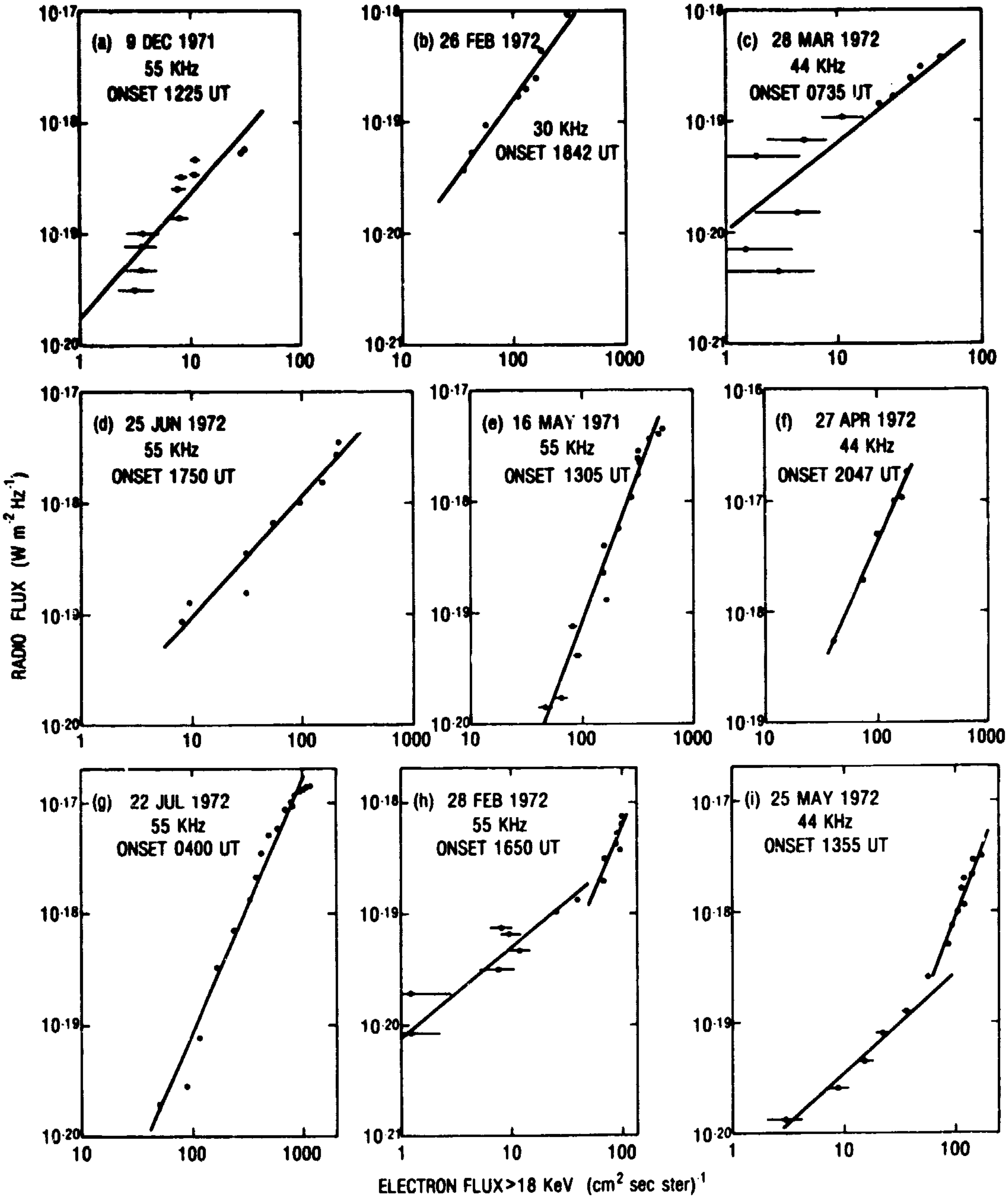


Figure 1. After Fitzenreiter et al. (1976), showing the correlation between radio and electron fluxes for nine type III bursts observed at 1 AU. The dots represent data taken approximately two minutes apart.

In 1976 yet another curious observation was reported by Fitzenreiter, Evans, and Lin (1976). In looking at simultaneous observations of both the electron and radio fluxes of type III bursts that had traveled out to 1 AU, they found that for electron fluxes less than about 100 $(cm^2 \cdot s \cdot ster)^{-1}$, the radio intensity I and the electron flux $J_E$ were approximately linearly proportional. For larger electron fluxes I $\propto J_E^{2.4}$ Figure 1 shows this power law dependence for nine events analyzed by Fitzenreiter, et al. (1976).

In 1974 we proposed a theory (Papadopoulos, Goldstein, and Smith 1974) which demonstrated that effects of strong plasma turbulence can readily account for the observed fact that the electron streams associated bursts are able to travel large distances without significant deceleration. In contrast, conventional weak turbulence plasma theory predicts that all the streaming energy should be dissipated within a few kilometers of the injection site.

The strong turbulence theory also suggested an explanation for the dominance of second harmonic radiation. During the last several years, that theory has been expanded in a series of papers (Smith, Goldstein, and Papadopoulos 1976, 1978; and Goldstein, Smith, and Papadopoulos 1978). In its present version the theory not only accounts for the minimal energy losses suffered by the electrons, but also is able to account for the observed intensities of electromagnetic radiation (at $2\omega_e$), and the correlation between the radio and electron fluxes.

The essential features of the theory will be reviewed in this paper. The reader is referred to the original papers mentioned above for additional details, and to the review by Nicholson and Smith (1979) for a discussion of other aspects of type III burst observations and theory.

ELEMENTS OF THE THEORY

Following the injection into the solar atmosphere of a power law distribution of electrons, the faster particles will begin to out pace the slower ones as the beam propagates to higher altitudes. Sufficiently far from the injection point only the fastest particles will at first be detected. The distribution function $f_T$ will have the form of a "bump" on the tail of the thermal component, and will be unstable to the excitation of Langmuir waves with phase velocities equal to the velocities of the fast electrons. In the conventional weak turbulence theory, as the slower particles arrive, they continue to amplify the Langmuir waves ("oscillation pileup") until the energy flux of the waves equals that of the beam $[v_g(E^2/8\pi) \simeq (1/2)n_b m v_b^3(\Delta v_b/v_b)]$, where $v_g$ and E are the group velocity and electric field of the Langmuir waves; and $n_b$, $v_b$,

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and $\Delta v_b$ are the density, velocity, and velocity spread of the energetic electrons, respectively. In the solar corona a beam could propagate only some 60 km before losing all of its energy to Langmuir waves. Clearly this does not happen because energetic electrons are often observed at 1 AU in association with type III bursts (Figure 2).

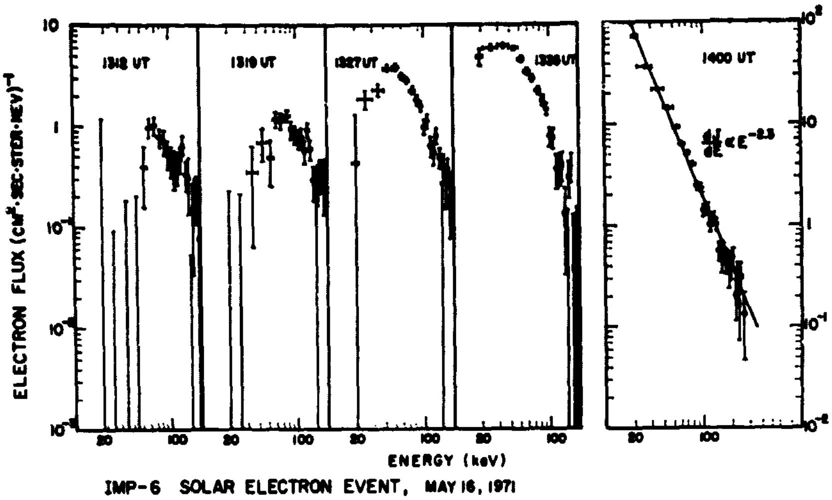


Figure 2. After Lin et al. (1973). Electron spectra are shown at various times during the type III burst of May 16, 1971.

The strong turbulence theory can be formulated in terms of a two-fluid hydrodynamic model of the plasma in which the motions are separated into fast plasma oscillations and quasineutral slow ion oscillations. If one writes the ion and electron densities as $n_i = n_o + \delta n_i$ ($\delta n_i/\delta n \ll 1$), and $n_e = n_i + \delta n_e$ ($\delta n_e/n_i \ll 1$), then the fundamental equations describing the coupling of high frequency plasma waves to low frequency density fluctuations are

$$\frac{\partial \delta n_e}{\partial t} - n_o \nabla\cdot\delta \underline{v}_e = -\nabla\cdot(\delta n_i \delta \underline{v}_e)$$

$$n_o \frac{\partial}{\partial t}\delta \underline{v}_e + \frac{\gamma_e T_e}{m}\nabla\delta n_e = \frac{e}{m} n_o \underline{\varepsilon} - \nu_e n_o \delta \underline{v}_e \qquad (1)$$

where $\underline{\varepsilon}$ satisfies $\nabla\cdot\underline{\varepsilon} = -4\pi e\delta n_e$, $\gamma_e$ is the polytropic index, and $\nu_e$ is a phenomenological damping decrement. In writing equations (1) we have assumed that $(k\lambda_e)^2 \ll 1$ and $\nu_e/\omega_e \ll 1$. $C_s$ is the speed of sound, defined by $C_s^2 = (\gamma_e T_e + \gamma_i T_i)/M$. The brackets, < >, denote averages over the fast time scale. Introducing the slowly varying quantity $\underline{E}(x,t)$

$$\underline{\varepsilon}(\underline{x},t) = \frac{1}{2}\ \{\underline{E}(\underline{x},t)e^{-i\omega_e t} + \underline{E}^*(\underline{x},t)e^{i\omega_e t}\}$$

results in the following simplification

$$\{i\frac{\partial}{\partial t} + \frac{\gamma_e T_e}{m\omega_e}\nabla\nabla + i\nu_e\}\underline{E} = \frac{\omega_e}{2n_o}\,\delta n_i \underline{E}$$

$$\{\frac{\partial^2}{\partial t^2} + \nu_i\frac{\partial}{\partial t} - C_s^2\nabla^2\}\delta n_i = \frac{1}{16\pi M}\nabla^2|\underline{E}|^2 \qquad (2)$$

Details of this derivation can be found in Papadopoulos, *et al.* (1974), Papadopoulos (1975), Manheimer and Papadopoulos (1975), and Smith *et al.* (1978). Equations (2) describe modulational instabilities including the OTSI, or soliton formation, depending on ones point of view.

Upon taking the Fourier transform of equations (2) the general dispersion relation for the OTSI and other modulation instabilities can be written as

$$[-\omega(\omega+i\nu_i) + k^2C_s^2] = -\frac{k^2\omega_e}{32\pi Mn_o}\int d\underline{k}'|\underline{E}(k')|^2 \times$$

$$\frac{[\underline{k}'\cdot(\underline{k}-\underline{k}')]^2}{|\underline{k}'|^2|\underline{k}-\underline{k}'|^2}[-(\omega+i\nu_{ek}) - \frac{3}{2}\omega(k'\lambda_e)^2 + \frac{3}{2}(\underline{k}-\underline{k}')^2\lambda_e^2\omega_e]^{-1} +$$

$$\frac{[\underline{k}'\cdot(\underline{k}-\underline{k}')]^2}{|\underline{k}'|^2|\underline{k}+\underline{k}'|^2}[-(\omega+i\nu_{ek}) - \frac{3}{2}\omega(k'\lambda_e)^2 + \frac{3}{2}(\underline{k}+\underline{k}')^2\lambda_e^2\omega_e]^{-1} \qquad (3)$$

where we have assumed that a zero-order spectrum of "pump" Langmuir waves exists between $\underline{k}_1 < \underline{k} < \underline{k}_2$. Outside that region the waves are taken to be small perturbations, $\delta\underline{E}(\underline{k},t)$ and $\delta n(\underline{k},t)$. The analysis can be further simplified by noting that a full three dimensional treatment is not necessary (contrast Bardwell and Goldman 1976 and Nicholson *et al.* 1978). In reducing these equations to one dimension, one must separately consider the question of whether the pump waves are predominately one dimensional, whether the parametric instabilities preferentially give rise to daughter waves also aligned predominately along the magnetic field direction, and whether the decay instability has a lower threshold than the OTSI.

The pump waves are excited by the beam plasma instability which in the presence of a magnetic field, for arbitrary propagation direction, has a growth rate given by:

$$\gamma_b = \frac{n_b}{n_o}\sqrt{\pi}\ \exp(-\lambda) I_o(\lambda)\left(\frac{\omega - \underline{k}\cdot \underline{V}_b}{k\Delta V_b}\right)\times \exp\left(\frac{\omega - \underline{k}\cdot \underline{V}_b}{k\Delta V_b}\right)^2$$

where $\lambda \simeq k_\perp \Delta v_b/\Omega_e$ and we assume that the beam velocity spread, $\Delta v_b$, is isotropic. Because $\gamma_b(k_\perp) \neq 0$ only if $\lambda \ll 1$, the angular spread of the spectrum $\theta$ will be given by $\sin(\theta) < (\Omega_e/\omega_e)\cdot(V_b/\Delta V_b) \simeq (1/3)\cdot(\Omega_e/\omega_e)$. With $\Omega_e/\omega_e \simeq 10^{-2}$ the angular spread of the instability spectrum is less than $1^o$. Thus the pump waves are one dimensional.

One can now approximate the pump spectrum, $W(k')$ in equation (3) as a one dimensional Lorentzian in which the pump wave spectrum is centered at $k_o$ with width $\Delta k_o$. The dispersion relation then becomes

$$[\omega(\omega + i\nu_i) - \mu k^2 T] + \frac{\frac{3}{4}\mu k^2 W_o F(k,k_o,\psi)}{\frac{9}{4}k^4 - [(\omega + i\nu_e) + 3i|k|\Delta k_o\cos\psi - 3|k|k_o\cos\psi]^2} = 0 \qquad (4)$$

where, for convenience, we have used a dimensionless notation in which $k \to (k\lambda_e)$, $\omega \to (\omega/\omega_e)$, $\nu \to \nu/\omega_e$, $\mu \equiv m/M$, $W_o \equiv \int dk'^2 |E(\underline{k}')|^2/8\pi n T_e$ and $T \equiv 1 + (3/2)\cdot(T_i/T_e)$. In equation (4) $\psi$ is the angle between $\underline{k}$ and the magnetic field direction, $F(\psi) \simeq \cos^2\psi$, $k \gg k_o$ and $F(\psi) \simeq 1 - k^2\sin^2\psi/k_o^2 + 4k^2\cos^2\psi/k_o^2$ for $k \ll k_o$. In the limit $(\Delta k_o/k_o) \to 0$ the dispersion relation reduces to the one found by Papadopoulos (1975) and Bardwell and Goldman (1976). In the interplanetary medium $W_o < \mu$ and $\omega(\omega + i\nu_i) < \mu k^2 T$ with the result

$$\frac{9}{4} k^4 - [\omega+3i|k|\Delta k_o\cos\psi-3|k|k_o\cos\psi]^2 - \frac{3}{4} k^2 \frac{W_oF(\psi)}{T} = 0$$

The maximum growth rate as a function of angle is given by $\partial(\mathrm{Im}\omega)/\partial\psi = 0$, or

$$-3|k|\Delta k_o\sin\psi + \frac{\sqrt{3}}{4}\left[ \frac{W_oF(\psi)}{T} - 3k^4 \right]^{-(1/2)} \partial F(\psi)/\partial\psi = 0$$

which can be satisfied _only_ if $\psi = 0$ for both $k \gtrless k_o$. Note that we have neglected ion inertia so that one must have $k_o^2 < \mu$ as well as $W_o < \mu$; but both conditions are easily satisfied in the interplanetary medium. Therefore we conclude that the daughter waves are also one dimensional and the one dimensional treatment of parametric instabilities developed in Papadopoulos _et al_. (1974) is completely justified for type III bursts.

Finally, the question arises as to whether the decay instability is important when $W_o \approx 10^{-5}$, the threshold for the OTSI. The decay instability is a three wave interaction whose threshold is found by setting $F(\psi) = 1$ and $\cos\psi = 1$ in equation (4). If the decay instability were excited, the daughter waves would have _smaller_ wave numbers than the pump wave. Stabilization of the beam plasma instability could not be achieved in that case. The threshold condition is $W^{thr}_{decay} \approx (8\nu_i/\mu kT) \cdot (3|k|\Delta k_o)$. For $T_e/T_i = 1$, $\nu_i/\mu kT \approx 2$ and $W^{thr}_{decay} \approx 50\,|k|k_o >> 50k_o\Delta k_o \approx 10\,k_o^2 >> 10^{-5}$, and thus the OTSI has the lower threshold.

The foregoing discussion has shown that both the pump and daughter waves are produced with wavenumbers aligned along the direction of the electron beam. In addition, none of the competing modulational instabilities will have lower thresholds than the OTSI, at least for the parameter range appropriate for type III bursts. The derivation of the OTSI growth rate, $\gamma_{OTS}$, has been given in Papadopoulos _et al_. (1974) and Smith _et al_. (1976). If $\delta \equiv \omega_{ek} - \omega_{ek'}$ is the frequency shift between the pump wave with wavenumber $k_o$ and the daughter wave with wavenumber $k'$, then for the OTSI $\delta < 0$ and $\gamma_{OTS}$ is given by

$$\gamma^2_{OTS}(k_o,k') = -(1/2)\cdot(\omega_A^2 + \delta^2) + \{(\omega_A^2 + \delta^2)^2 - 4\delta^2\omega_A^2[1 + W(k_o)\omega_e^2/(2\delta\omega_{ek'})]\}^{(1/2)} \quad (5)$$

where $(\omega_A/\omega_e)^2 = (m/M)\cdot(k'\lambda_e)^2/[1+(k'\lambda_e)^2]$. In the absence of collisions the threshold for the OTSI IS $W_T = -2\delta_{ek'}/\omega_e^2$.

Once the OTSI is excited, the pump waves in resonance with the electron beam will couple to daughter waves having larger wavenumbers. These daughter waves will no longer resonate or exchange energy with the electrons. It is important to note that other modulational instabilities, e.g. direct collapse, do not have this propert_ of scattering the pump waves out of resonance with the electrons. Thus the beam plasma instability would continue to amplify the pump waves until the threshold for the OTSI is reached. Once the OTSI begins to dominate the wave-wave interactions, stabilization of the electrons against catastrophic energy loss will proceed as outlined.

The OTSI is not by itself adequate to explain the lack of energy loss of the electron streams. Excitation of the OTSI itself would result in a marginally stable situation in which wave energy densities of order $W_T \approx 10^{-5}$ remain resonant with the beam, and over large distances could decelerate it. However, in the presence of such energy densities, an approximately isothermal plasma will produce density depressions $\delta n/n \approx \tilde{W}/nT$ due to pressure balance $(nT + \tilde{W} = \text{const.})$. (The ~ denotes dimensioned variables.) These will modify the local plasma frequency, and hence the local Bohm-Gross dispersion relation of both the pump and daughter Languir waves. The nonlinear dispersion relation will be approximately given by $\omega_k = \omega_e[1 + (3/2)k^2\lambda_e^2 - (1/2)\delta n/n] \approx \omega_e[1 + (3/2)k^2\lambda_e^2 - (1/2)\tilde{W}/nT]$. When $W \approx (k\lambda_e)^2$ this correction term will become important and will cause $\delta$ and $W_T$ to decrease. Then virtually all the energy initially resonant with the energetic electrons will be rapidly transferred to larger wavenumbers, out of resonance. A more rigorous derivation of these effects can be found in Kaw, Lin, and Dawson (1967) and Smith _et al_. (1978).

An additional effect which must be included is anomalous resistivity (Dawson and Oberman 1963, and Dawson 1968). Once nonthermal levels of ion fluctuations are excited, a high frequency anomalous resistivity is produced which causes long wavelength Langmuir waves to cascade to shorter wavelengths. In the presence of a three-dimensional spectrum of correlated, nonthermal ion fluctuations, the scalar impedance at frequencies near $\omega_e$ is

$$Z(\omega \approx \omega_e) \approx \frac{(2\pi)^3(\delta n_i/n_o)^2}{36\pi k_i^2 \Delta k}\,\frac{\omega(k_i\lambda_e)}{\omega_e^2\lambda_e^3}$$

The effective electron-ion collision time, $\tau_c$, is related to Z by

$$Z(\omega) = \frac{4\pi i\omega}{\omega_e^2}(1 - i/\omega\tau_c)$$

so that

$$\gamma_{NL}/\omega_e \equiv 1/(\omega_e \tau_c) \approx \frac{\pi}{2} \frac{(\delta n_i/n_o)^2}{(k_i \lambda_e)^2}$$

$$= \Sigma \frac{[S(k_i)-S_o(k_i)]}{(k_i \lambda_e)^4}$$

where $S_o(k)$ is the thermal noise level of the ion spectrum and $S(k)$ is the total electric field energy density in ion waves at wavenumber k.

Landau damping of Langmuir waves by the thermal solar wind electrons has also been included. The total electron distribution function, $f_T$, is the sum of the solar wind and beam components. Thus Landau damping, the linear beam plasma instability, and reabsorption caused by evolution of the electron beam to lower velocities can all be described in terms of a single damping decrement, $\gamma_L$, where

$$\frac{\gamma_L}{\omega_e} = \pi \frac{\beta^3}{|\beta|} \frac{\partial f_T}{\partial \beta} \bigg|_{\beta = \omega/kc}$$

The evolution of the growth, spectral transfer, and damping of the Langmuir and ion waves in the presence of the time evolving electron beam can now be described in terms of the various transfer rates $\gamma_L$, $\gamma_{OTS}$, and $\gamma_{NL}$. The complete set of rate equations has been given before (Smith et al. 1976, 1978) and we will not repeat them here.

Before discussing the results of our calculations, it is necessary to briefly review some properties of the model we constructed for the electron beam (v. Smith et al. 1976, 1978 for more details). The beam distribution model is a semi-empirical one based on in situ particle observations at 1 AU (Lin 1974, Lin, Evans, and Fainberg 1973). In general it is very misleading to attempt to construct a beam model based solely on its interaction with the self-consistently produced Langmuir turbulence because of the importance of scattering by magnetic irregularities even for events called "scatter-free". One such event was observed on May 16, 1971 (Figure 2). Of particular importance is the fact that it took more than 15 minutes for the peak of the spectrum to evolve from 80 eV to 33 keV, implying that reabsorption of Langmuir waves by the electron stream is unimportant. Theories which do not accurately model this behavior will, of necessity, conclude that reabsorption is important. As an example of the latter, Magelssen and Smith (1977) use a model for the electron beam evolution for the May 16 event which takes less than 5 minutes to evolve from 80 to below 20 keV. We will return to this subject below when discussing the results of our numerical calculations.

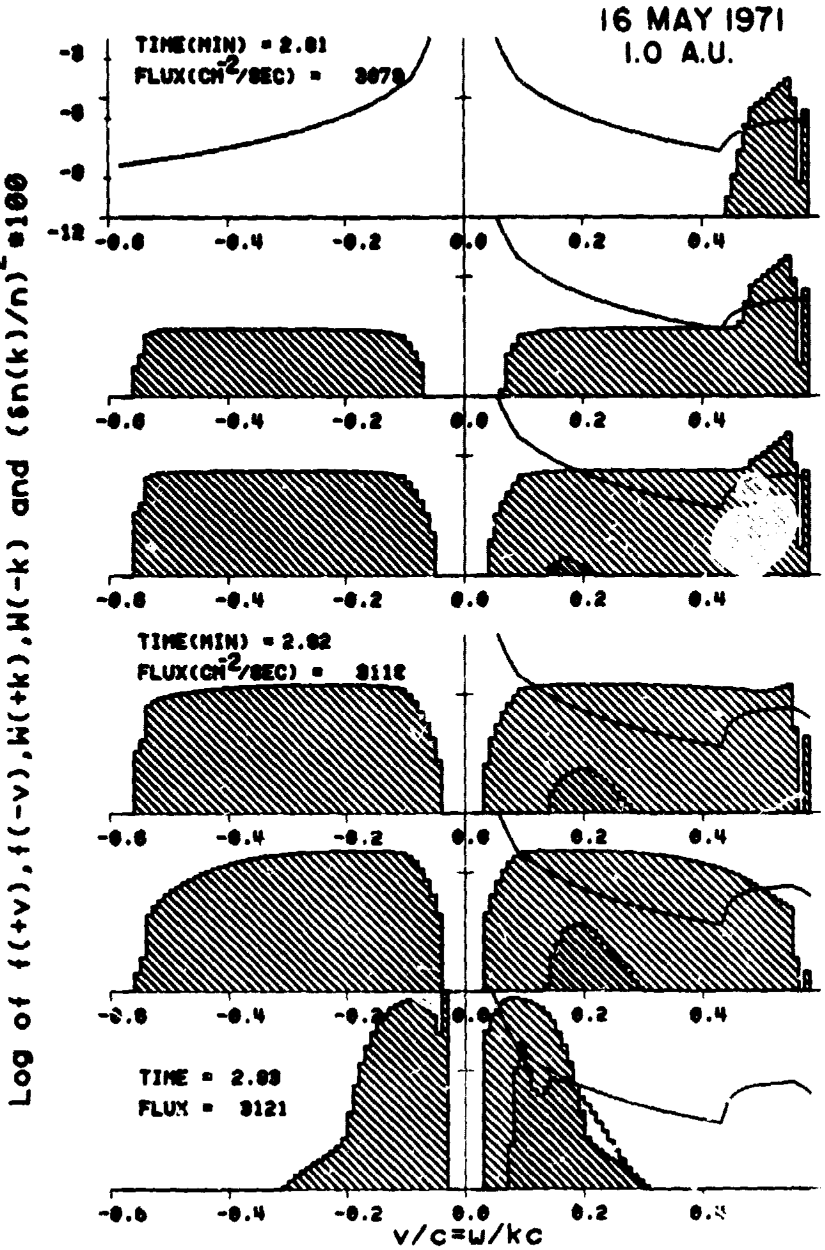


Figure 3. Results of a numerical solution of the rate equations that describe the OTSI. Parameters were chosen to model the May 16, 1971 event at 1 AU. The top panel (3a) shows the distribution function $f_T$ of the solar wind plus the linerarly unstable beam. Langmuir waves (diagonally striped histograms) are shown near $W_T$ (3a), and during subsequent stages of excitation and stabilization of the OTSI (3b-3f). Ion oscillations are depicted by the gray shading. Times computed from the start of the numerical calculations and the calculated values of the electron flux are given in 3a, 3d, and 3f.

NUMERICAL CALCULATIONS

The numerical evaluation of the rate equations can be carried out at any point in the interplanetary medium at which the density and temperature of the ambient solar wind can be estimated. Typically, we chose distances between 0.1 and 1.0 AU, and assumed that the ambient density varied as $r^{-2}$. At a given location the calculation began ($t = 0$) with the arrival of energetic electrons with velocities of about 0.7c. The exact velocity distribution being given by the beam evolution model. As an example, consider the burst on May 16, 1971. The local plasma frequency at 1 AU on that date was about 30 kHz and electrons with energies above 100 keV were first observed at 1305 UT when the radiometer on IMP-6 first detected radio noise at 55 kHz ($\approx 2\omega_e/2\pi$). The radio noise increased in intensity until 1335 UT, and little further evolution was observed in the electron spectrum after that time. From Figure 2 we can see that the distribution function had a positive slope below the peak energy. The other parameters needed for the numerical model were the path length traversed by the electron beam, taken to be 1.5 AU; the ratio of the beam to solar wind density $\eta$ estimated to be $5 \times 10^{-6}$; and the spectral index $\zeta \approx 4.6$ of the power-law portion of $f_T$.

The solution is shown in Figure 3, where the logarithms of $f_T(\beta)$, $W(\pm k)$, and $(\delta n/n)^2 = S(k)/(k\lambda_e)^2$ are plotted against $v_{ph} = \omega_{ek}/kc$ at various times.

Initially, the linearly unstable beam produces resonant plasma waves (indicated by cross hatching in Figure 3a) that grow until the OTSI threshold is reached (Figure 3a). Aperiodic ion waves are then excited (gray shading) at the rate $\gamma_{OTS}$, as are shorter wavelength "daughter" Langmuir waves (Figure 3b-d). The combined effects of nonlinear changes in the Bohm-Gross dispersion relation and anomalous resistivity then complete the decoupling of the electron beam from the Langmuir turbulence (Figure 3d-f). In our calculations the collapse to short wavelengths ceases when Landau damping by the thermal solar wind electrons balances the spectral transfer. No further energy exchange will then take place. Gradually the ion fluctuations and Langmuir waves will simultaneously decay back to thermal levels whereupon the linear instability will again be excited, and the process will cyclically repeat until the electron beam has merged with the ambient solar wind distribution and no positive slope exists to $f_T(\beta)$.

It is important to note that the total elapsed time between the onset of the OTSI and its final stabilization was little more than 0.1s, during which the electron distribution was essentially constant. Therefore, neither reabsorption nor quasilinear relaxation can be important.

Similar calculations were performed at 0.5 and 0.1 AU and for the type III bursts observed on May 25, 1972 and February 28, 1972; the results are similar to those described here and are reported in

Goldstein et al. (1978). In all cases stabilization and decoupling of the electron beam from the Langmuir turbulence is due to excitation of the OTSI.

## ADDITIONAL THEORETICAL RESULTS

We now turn to the question of why type III bursts are preferentially observed at the second harmonic of the local plasma frequency. Much of this discussion is based on a recent paper by Papadopoulos and Freund (1978).

From a comparison of Figure 5a and f, one sees that the long wavelength pump waves have collapsed into shorter wavelength daughter waves. In configuration space these short wavelength structures are solitons (Manheimer and Papadopoulos 1975), whose spatial extent in the direction parallel to the magnetic field can be estimated to be about $50\lambda_e$, with an energy density, W, of nearly $10^{-2}$. Such structures are very difficult to observe with present spacecraft instrumentation. In a 400 km/s solar wind, a 350 m ($50\lambda_e$) soliton is convected past a 30 m dipole antenna in little more than a millisecond. This must be compared to the electronic response times of plasma wave experiments typically no faster than 20 ms (Gurnett, private communication).

Papadopoulos and Freund (1978) found that the total volume emissivity of a soliton, integrated over solid angle is

$$J(2\omega_e) = \frac{3\sqrt{3}}{8}\left(\frac{v_e}{c}\right)^4\left(\frac{cE_o^2}{8\pi\Delta Z}\right)\left(\frac{1}{k_oL}\right)^2 \qquad (6)$$

where $\Delta z$ is the parallel dimension of the linearly unstable wave-packet, $k_o = \sqrt{3}\omega_e/c$ is the wavelength of the electromagnetic wave at $2\omega_e$, $v_e$ is the thermal electron velocity, $E_o$ is the electric field in the soliton, and L is the dimension of the soliton transverse to the magnetic field. (Papadopoulos and Freund argue that collapse is likely only in the parallel dimension, and that L should be greater than the electron Larmor radius.) Equation (6) is valid for $k_o^2L^2 \gg 4$, a good approximation throughout the interplanetary medium. The intensity of emission outside a spherical shell of radius R and thickness $\Delta R$ centered on the sun is (Gurnett and Frank 1975) $I = JR(2\omega_e/2\pi)$. For the May 16 burst at the time of soliton formation (Figure 3f), $I(2\omega_e) \approx 1 \times 10^{-17}$ W m$^{-2}$ s$^{-1}$, close to the peak intensity observed at 55 kHz.

The correlation between the radio and electron fluxes also has a straightforward explanation. In Figure 4 we plot $\gamma_{OTS}$ against $W(k_o)$--

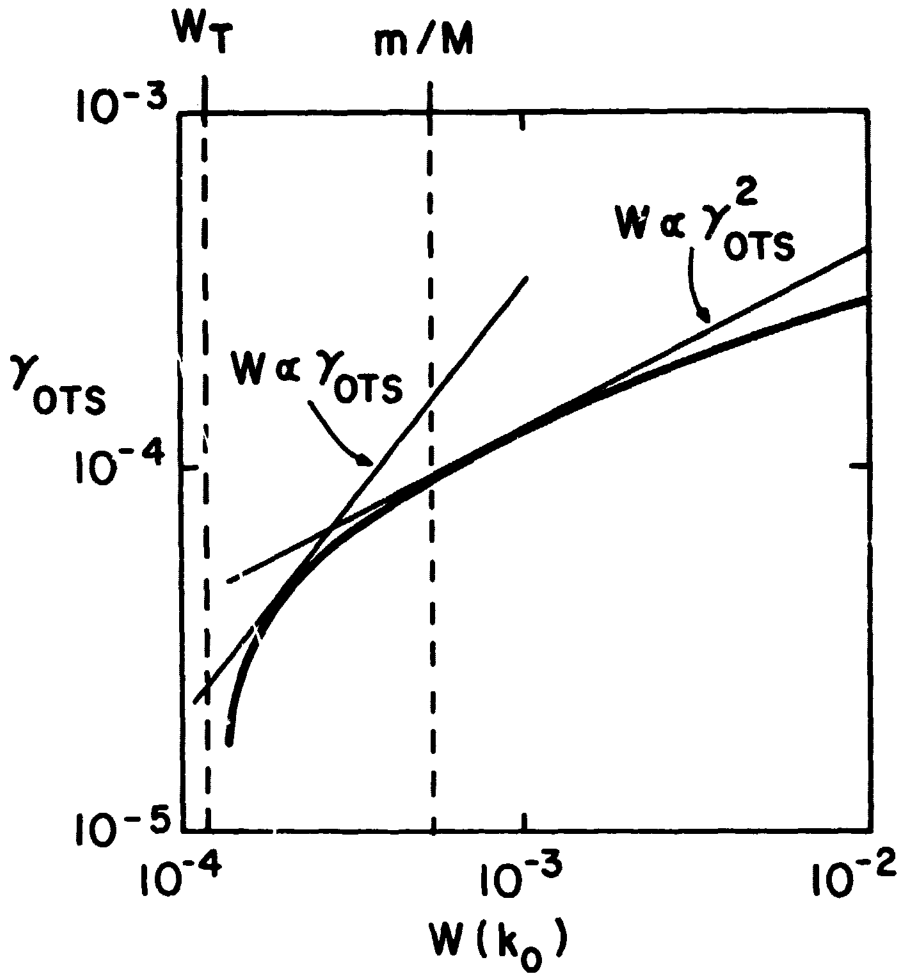


Figure 4. $\gamma_{OTS}$ versus $W(k_o)$ from equation (5). Note the approximate scaling relations for $W_T < W(k_o) <$ m/M and $W(k_o) >$ m/M.

equation (5). Papadopoulos (1975) and Rowland (1977) have shown that one can find an approximate relationship between $W(k_o)$ and $\gamma_{OTS}$ near and above the threshold:

$$W(k_o) \sim \gamma_{OTS} \qquad m/M > W(k_o) > W_T$$

$$W(k_o) \sim \gamma_{OTS}^2 \qquad W(k_o) > m/M$$

These regimes are noted in Figure 4.

When the OTSI stabilizes the linear instability, $\gamma_{OTS} \propto \gamma_L$, and $\gamma_L \propto n_b$. In addition, from equation (6), $I(2\omega_e) \propto W(k_o)$ the electron flux, $J_E \propto n_b \langle v \rangle$, where $\langle v \rangle$ is the <u>peak</u> of the electron distribution, so that with $n \propto \langle v \rangle^{-\zeta+1}$ one has

$$I(2\omega_e) \propto J_E^{(1-\zeta)/(2-\zeta)} \qquad m/M > W(k_o) > W_T$$

$$I(2\omega_e) \propto J_E^{2(1-\zeta)/(2-\zeta)} \qquad W(k_o) > m/M \qquad (7)$$

In Figure 5 we compare the observations of Fitzenreitr r *et al.* (1976) with the predictions of the theory from equations (7). The three bursts

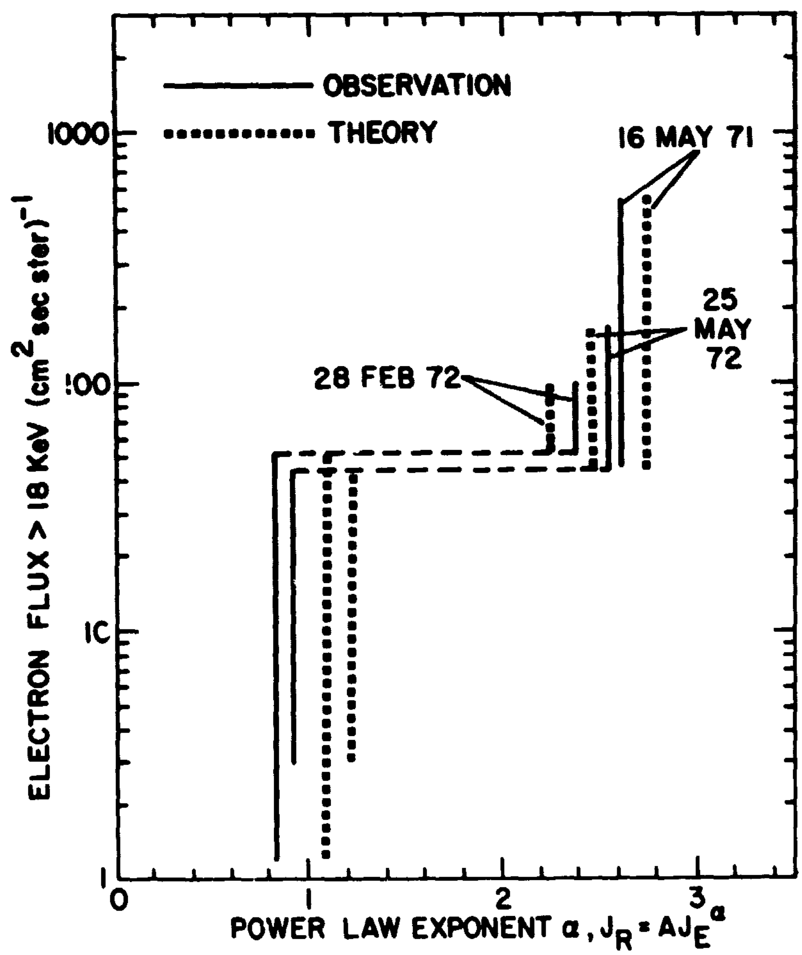


Figure 5. After Fitzenreiter *et al.* (1976). The electron flux and power law exponent $\alpha$ from the relationship $I \propto J_E^{\alpha}$ are shown for the three events for which numerical calculations could be performed. Observed and computed values of $\alpha$ are plotted.

shown in Figure 5 are the only ones for which the electron observations were sufficiently detailed to permit estimation of $\zeta$, L, and $\eta$--the parameters used in our beam evolution model. It should be emphasized that the correlation between $J_E$ and I cannot be explained using weak turbulence theories. The excellent agreement shown in Figure 5 is confirmation that type III bursts are stabilized by the OTSI and that solitons radiate electromagnetic radiation proportional to W and not $W^2$.

Although the scaling $W(k_o) \sim \gamma_{OTS}$ would seem to imply that $\alpha$ should not be less than 1, Figure 4 indicates that very close to threshold, $W(k_o) \sim \gamma_{OTS}^{\nu}$ with $\nu < 1$. Thus, for bursts such as the ones on February 28 and May 25, 1972 which initially only weakly excite the OTSI, values of $\alpha < 1$ are quite reasonable.

Thus far we have tacitly assumed that because the electron beam becomes decoupled from the radiation field, no significant energy loss will occur. Smith et al. (1978) have investigated this in some detail; we only summarize that discussion here.

If the beam is injected near the solar surface, the total energy lost by the beam in propagating to the point R is given by

$$\Delta E = \int_{R_o}^{R} dr A(r) \int_{t_1(r)}^{t_2(r)} dt \frac{d\tilde{W}(r,t)}{dt} \tag{8}$$

where A(r) is the source area at r, and $t_1(r)$, and $t_2(r)$ are the times at which the instabilities at r begin and end. Because all the beam energy loss occurs in the resonant region until the onset of the OTSI, one can assume that it takes place at the steady rate $dW/dt = W_T \tau_o$, where $W_T$ is taken to be $W_o \exp(\gamma_L \tau_o)$.

When equation (8) was evaluated, Smith et al. (1978) found that $\sim$90% of the energy loss occurred in the inner corona, and that $\Delta E = 10^{30}$ W (ergs). With $W \approx 10^{-4}$, the exciter loses some $10^{26}$ ergs in leaving the corona. The total energy in the type III exciter has been estimated to be $\approx 10^{28}$ ergs (Lin 1971). Thus, the exciter will typically lose only a few percent of its energy.

One additional consequence of this energy-loss calculation was that it provides an explanation for why the electron streams appear to have such well-defined velocities, of order c/3 at high frequencies (Wild and Smerd 1972), decreasing to c/2 or less at low frequencies (Evans, Fainberg and Stone 1973).

The peak intensity at any frequency is reached just before the linear beam-plasma instability stops at that frequency, for at that time the density in the energetic electron beam is maximum. It is this peak velocity which is directly deduced from the observed frequency drift rates as being the nominal velocity of the beam.

Smith et al. (1978) found that in the inner corona the peak velocity when the linear instability stopped was $v_p = 0.3c$, while near 1 AU, because the ambient solar wind is cooler, $v_p$ was about 0.2c. This suggests that the nominal velocity (c/3) is not characteristic of electron acceleration, but rather reflects the evolution of the particle spectrum. In addition, the observations of Evans et al. (1973) do not necessarily imply that the exciter is decelerated between 0.05 AU - 1 AU, but rather reflects the decrease in the temperature of the solar wind with increasing heliocentric distance.

We have reviewed a theory of type III bursts which is able to account for many seemingly diverse aspects of the phenomenon, in particular we have offered an explanation of the small energy losses of the exciter, the predominance of radiation at $2f_{pe}$, the characteristic exciter velocities of 0.2-0.3c, and the correlation between electron and radio fluxes.

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