SESSION V

SOLAR BURSTS - METER-DECAMETER WAVELENGTHS

https://doi.org/10.1017/S0074180900036998 Published online by Cambridge University Press

CURRENT UNDERSTANDING OF THE PHYSICS OF TYPE III SOLAR RADIO BURSTS

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ABSTRACT

One of the most exciting plasma physics investigations of the recent years has been connected with the understanding of a new strong turbulent plasma state excited by propagating electron beams. This new state is initiated on the linear level by parametric instabilities (OTS, modulational, etc.) and results in a very dynamic state composed of collective clusters of modes called solitons, cavitons, spikons, etc. Introduction of these concepts to the classic beam plasma interaction problem has rendered quasilinear and weak turbulence theories inapplicable over most of the interesting parameter range, and helped explain many paradoxes connected with the propagation of beams in the laboratory and space. Following a brief review of these non-linear notions, we demonstrate how their application to type III solar radio-bursts has revolutionized our understanding of their propagation, radio-emission and scaling properties and has guided the in situ observations towards a more complete understanding. A particular burst (May 16, 1971) is analyzed in detail and compared with numerical predictions.

I. INTRODUCTION

Type III solar radiobursts are a form of sporadic solar radioemission originating throughout the solar corona and also in the interplanetary region, to heliocentric distances of 1 AU and beyond. They are associated with fluxes of moderately energetic electrons (10-100 keV) which are accelerated either in flares or in active storm regions and which escape along magnetic field lines that penetrate the high corona. The bursts appear to have a frequency corresponding to twice the local plasma frequency ($2\omega_e$) (sometimes, the plasma frequency); their frequency drifts at a rate suggesting a source velocity of $10^{10}\ \frac{cm}{sec}$. At a fixed frequency the duration ranges from < at the highest frequencies, to more than 1 hr at the lowest frequencies [Smith 1974].

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M.R. Kundu and T.E. Gergely (eds.), Radio Physics of the Sun, 287-297.

The observed frequency ranges from a few kHz to several hundred MHz. A vast literature of observational studies over the last thirty years has produced a fairly detailed morphological picture of the radio-emission, its scaling and decay properties. However, the interest in the phenomenon has not diminished and challenging theoretical problems have been encountered in constructing convincing models describing its behavior. It is a testimony to the increased sophistication of both the observational tools and of the theoretical computational modeling efforts that the last four years have shed light upon many of the paradoxes connected with type III radiobursts [Smith, et al., 1978].

Serious theoretical difficulties were 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 raised by Sturrock (1964) are the following three. First, why is the electron beam that excited the bursts not significantly decelerated. Second, why is the radiation predominantly emitted at the second harmonic of the local plasma frequency, ω_e. Finally, why does the beam have such a well defined velocity, typically between 0.2 and 0.3c.

In 1976 yet another curious observation was reported by Fitzenreiter, et al. (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 $(\text{cm}^2 \text{ sec ster})^{-1}$, the radio intensity, I, and the electron flux, J_E, were approximately linearly proportional. For larger electron fluxes $I \propto J_E^{2.4}$.

Two avenues of investigation were available at the time. We could assume that our understanding of the physics, based on quasilinear and weak turbulence theory was correct, and proceed to build complex computer models of the inhomogeneous propagation, or we could look for new physics. The first approach produced very limited success (Magelssen and Smith, 1977), while the second produced a major breakthrough not only for the type III bursts, but also for theoretical plasma physics (Zakharov (1972), Papadopoulos (1973, 1975)).

II. NUMERICAL RESULTS BASED ON STRONG TURBULENCE THEORY

It was shown (Zakharov, 1972) that our concept of electrostatic turbulence uniformly distributed in space is invalidated even for rather low wave energy levels, (i.e., $W^{\ell}/nT < (k\lambda_D)^2$ where k is the typical wave-number of the e.p.o. spectrum). The physical reason for this can be seen by noting that the presence of high frequency (ω_e) waves exerts a low ponderomotive force (i.e., radiation pressure) on the plasma, which results in a modification of the local density n, in which the change in pressure $p = nT + \frac{1}{2} E^2/8\pi = nT + \frac{1}{2} W$ is zero ($\delta p = 0$). Therefore $\delta n/n \approx -\frac{1}{2} W/nT$. The dispersion relation for e.p.o.'s thus becomes (Abdulloev, et al., 1975)

$$\omega_{ek} = \omega_e\left(1 + \frac{3}{2}(k\lambda_D)^2 + \tfrac{1}{2}\frac{\delta n}{n}\right) \approx \omega_e\left(1 + \frac{3}{2}(k\lambda_D)^2 - \tfrac{1}{4}\frac{W}{nT}\right). \qquad (1)$$

Eq. (1) has a simple physical interpretation if the e.p.o.'s are viewed as quasiparticles subject to attractive and repulsive forces and capable of emitting sound waves. For nonrelativistic velocities, eq. (1) can be viewed as the definition of energy of the quasiparticle (with $h = 1$), with an effective mass of $m_{eff} = \omega_e/3V_e^2$, a momentum of k, and the quantity corresponding to the velocity of light is $c^2 = 3V_e^2$. The last term corresponds to the potential energy of the quasiparticle in the field of others. Since its sign is negative it implies attraction. As long as the kinetic energy (i.e., $\frac{3}{2} k^2V_e^2/\omega_e$) is larger than the attractive potential (i.e., $\frac{1}{4}\omega_e / W^{\ell}/nT$), the plasma waves behave in the usual sense described by the weak turbulence theory. However, when $W^{\ell}/nT > 6(k\lambda_D)^2$ they start collapsing to smaller and smaller sizes, and form localized clumps of wave energy which have been given the name solitons, spikons, or cavitons. As shown by Manheimer and Papadopoulos (1975), this process is equivalent to the oscillating two stream instability (O.T.S.) known from parametric interactions. The inequality

$$\frac{W^{\ell}}{nT} > (k\lambda_D)^2 \qquad (2)$$

is usually considered as the condition for invalidation of weak turbulent theory, and has two consequences. The first is that the dispersive term in eq. (1) becomes negligible, thereby radically modifying the real part of $\varepsilon_L(\underline{k},\omega)$. The second is that instead of uniformly distributed turbulence, we end up with a series of highly intense and localized wave packet-like structures.

It was first shown (Papadopoulos, et al., 1974) that effects of strong plasma turbulence can readily account for the observed fact that the electron streams associated with the 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, et al., 1976, 1978; Goldstein, et al., 1978; Nicholson, et al., 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 $2\omega_e$), the correlation between the radio and electron fluxes, and for the observed decay times of the radiation. The full impact of the theory was, however, due to the results of the numerical modeling. Rate equations including strong turbulence mode coupling effects, reabsorption and collisionless damping, were utilized to model actual observations. The complete set of equations can be found in Smith, et al. (1976, 1978) and will not be repeated here. The input to the code was a beam distribution based on *in situ* particle observations at 1 AU. The numerical computations can be performed at any point in space at which the density and temperature of the ambient solar wind can be estimated. Typically, distances between 0.1 and 1.0 AU were chosen, and it was 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 ($\simeq 2\omega_e/2\pi$). The radio noise increased in intensity until 1335 UT, and little further evolution was observed in the electron spectrum after that time. Their results show 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, η, estimated to be 5×10^{-6}. The results of the model are shown in Fig. 1 as a function of time, where the logarithm of the electron distribution function $f_T(\underline{v})$, the electron plasma wave energy level W^{ℓ} (normalized to nT) and the amplitude of the density fluctuations ($\delta n/n$) are plotted as a function of $V_p = \omega_e/k$.

Initially, the linearly unstable beam produces resonant plasma waves (indicated by cross hatching in Fig. 1) that grow until the modulational threshold is reached (Fig. 1a). Periodic ion waves are then excited (gray shading) as are shorter wavelength "daughter" Langmuir waves (Fig. 1b-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 (Fig. 1d-f). In the 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 $f_T(\underline{v})$.

It is important to note that the total elapsed time between the onset of OTSI and its final stabilization was little more than 0.1 sec, during which the electron distribution is 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 periodic modulational instability (i.e., oscillating two stream).

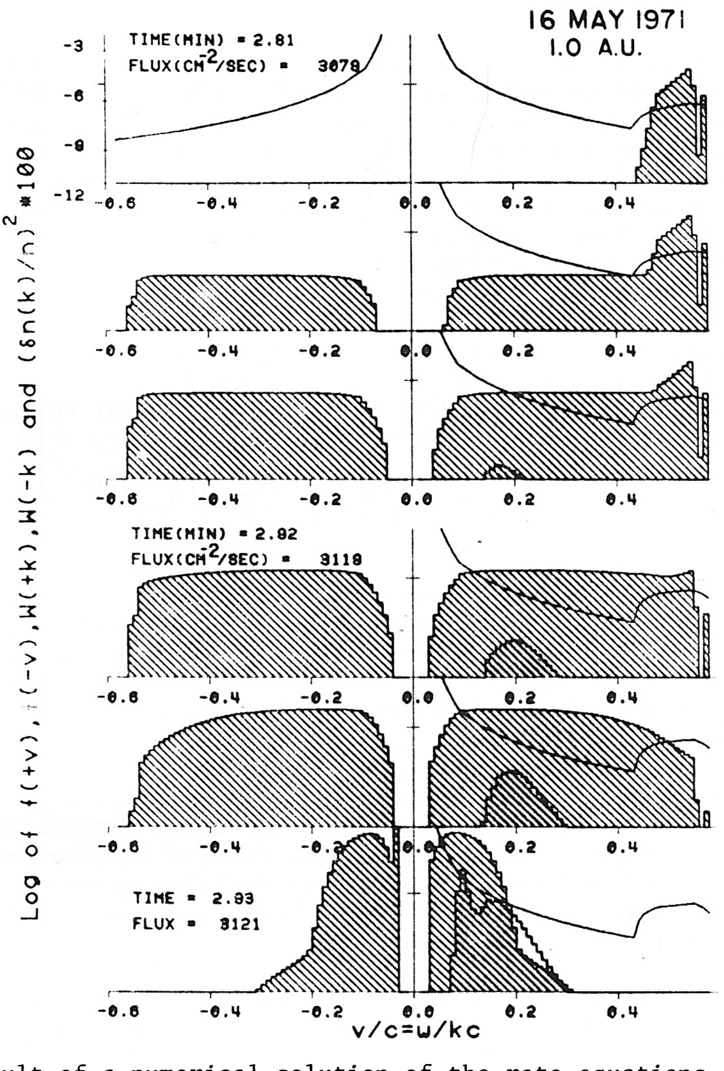

Fig. 1 - Result of a numerical solution of the rate equations. Parameters were chosen to model the May 16, 1971 event at 1 AU. The top panel (a) shows the distribution function, f_r, of the solar wind plus the linearly unstable beam. Langmuir waves (diagonally striped histograms) are shown near $W_T(a)$, and during subsequent stages of excitation and stabilization (b-f). 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 17a, d, and f.

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 Fig. 1a 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 $50\lambda_D$, with an energy density, W/nT, of nearly 10^{-2}. Such structures are very difficult to observe with present spacecraft instrumentation. In a 400 km/s solar wind, a 350m ($50\lambda_D$) soliton is convected past a 30m 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

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

where Δz 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 $2\omega_e$, E_o is the electric field in the soliton, and L is the dimension of the soliton transverse to the magnetic field. Eq. (3) 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 ΔR 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 (Fig. 1f), $I(2\omega_e) \cong 1 \times 10^{-17}\,Wm^{-2}sec^{-1}$, close to the peak intensity observed at 55 kHz. Finally, we should note that using the results of computations such as shown in Fig. 1 and eq. (3), excellent detailed agreement was found concerning the exponent α of the $I \sim J_E{}^{\alpha}$ dependence between the radio intensity and electron flux (Fig. 2).

Thus far it was 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

$$E = \int_{R_o}^{R} dr A(r) \int_{t_1(r)}^{t_2(r)} dt \frac{d\tilde{W}(r,t)}{dt} \tag{4}$$

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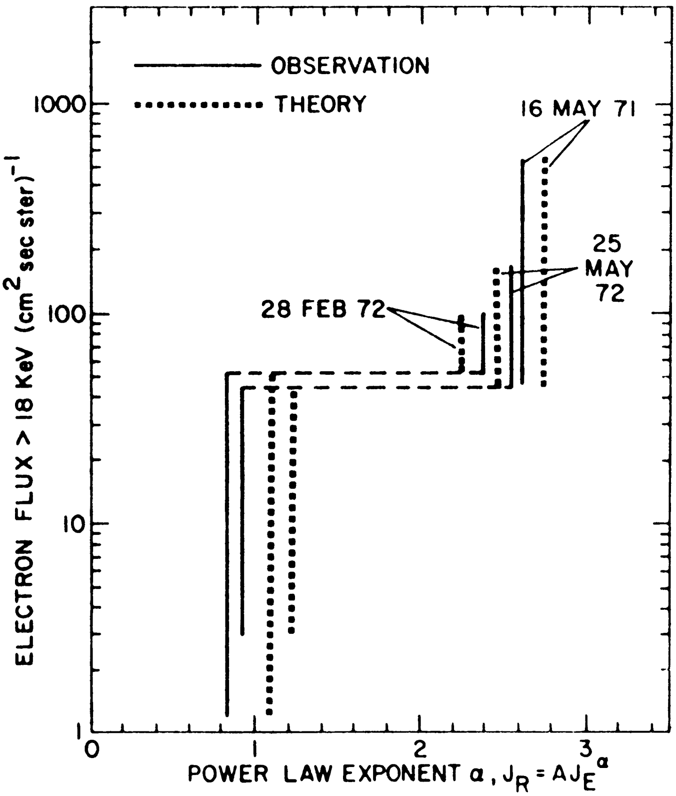

Fig. 2 - After Fitzenreiter et al. (1976). The electron flux and power law exponent, α, from the relationship $I \propto J_E^\alpha$ are shown for the three events for which numerical calculations could be performed. Observed and computed values of α are plotted.

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 collapse, one can assume that it takes place at the steady rate dW/dt = W_T/τ_o, where W_T is taken to be $W_o \exp(\gamma_L \tau_o)$.

When eq. (4) was evaluated, Smith, et al. (1978) found that ~90% of the energy loss occurred in the inner corona, and that $\Delta E = 10^{30}$ W (ergs). With $W \cong 10^{-4}$, the exciter loses some 10^{26} ergs in leaving the corona. The total energy in the type III exciter will typically lose only a few percent of its energy.

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One additional consequence of this energy-loss calculation was that it proved an explanation for why the electron streams appear to have such well-defined velocities, or order c/3 at high frequencies, decreasing to c/2 or less at low frequencies.

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 do not necessarily imply that the exciter is decelerated between 0.005 AU - 1 AU, but rather reflects the decrease in the temperature of the solar wind with increasing heliocentric distance.

In concluding, we should note that the most important lesson from the above is the fact that plasma theory supported by computation has reached the level of sophistication where detailed predictions can be derived even in complex systems.

ACKNOWLEDGMENTS

It is a pleasure to acknowledge significant contributions to the above work by my colleagues Drs. R. Smith, M. Goldstein, H. Freund, H. Rowland and P. Palmadesso. The work was performed at the Naval Research Laboratory, and was supported by ONR and NASA.

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DISCUSSION

Nakagawa: I would like to point out the similarity of your soliton formation and shock formation in compressible fluid turbulence. In fluid turbulence the statistics of the magnitude and period of shocks become the subject of theoretical study. What controls the statistics of solitons?

Papadopoulos: A statistical theory of strong or spiky turbulence, is an ongoing subject of investigation and probably one of the most challenging current theoretical plasma problems. We still do not have any definite models. Some limited success has been achieved in Manheimer and Papadopoulos, Phys. Fl. 18, 1397 (1975) and Tsytovich's book on plasma turbulence.

Gergely: I would like to know how do you define inner corona, when you say that beams with energy less than 10^{26} ergs do not get out. In other words, I would like to know at what frequency should we see these bursts to cut off?

Papadopoulos: It would correspond to a level of 15 MHz.

Dulk: Would you expect this strong turbulence to develop in the inner corona, say below 1 $R_\odot$, or only in the solar wind? i.e. What is the criterion for strong turbulence to develop?

Papadopoulos: The criterion for strong turbulence is $\frac{W}{nT} > (k_o \lambda_D)^2 = (\frac{V_b}{V_e})^2$.

Kuijpers: The plateau in velocity space you talked of in the beginning is only an equilibrium situation in the one-dimensional case and you understand also that your strong turbulence calculations are one-dimensional. How does three-dimensionality influence your results?

Papadopoulos: The fact that the maximum growth of the beam plasma instability is in the direction of the stream, forces both the linear and non-linear stage of the collapse to behave in a one dimensional fashion. Namely $L_\parallel \ll L_\perp$ ($L_\parallel$, $L_\perp$ are parallel and transverse scale lengths).

Melrose: What pitch-angle distribution of the beam did you choose in estimating the angular distribution of the growth rate of Langmuir waves?

Papadopoulos: Our analysis was one dimensional, so that no assumption on the pitch angle was made.

D. Smith: Papadopoulos and coworkers made an important contribution in pointing out the possible effects of strong turbulence. However, it is not clear that the problem can be treated in one dimension. Nicholson and Goldman find that for $\omega_{ce}/\omega_{pe} \gtrsim 100$ at 0.45 A.U. two dimensional

https://doi.org/10.1017/S0074180900036998 Published online by Cambridge University Press

effects are important. The dominant energy transfer mechanism is direct two-dimensional collapse which is a highly nonlinear state of the modulational instability. The amount of time spent in the linear stage is completely negligible.

Papadopoulos: The claim that the direct two dimensional collapse can stabilize the beam plasma instability, has caused a confusion in the literature. It is very simple to see by elementary and quite general physical arguments, that even if the transverse collapse rate is much faster than the beam growth rate, it cannot stabilize the beam-plasma instability till $\frac{W}{nT} > \frac{1}{10}$, a rather large value. This is because all the wave numbers $k_\perp \lambda_D < \frac{1}{3}$ are unstable. Therefore transverse collapse results in redistribution of the unstable wave-spectrum from peaking at the maximum growth, to other less faster growing modes. This results in a somewhat lower overall growth but not stabilization till the collapse has reached stable $k_\perp$ wave-numbers (i.e. $k_\perp \lambda_D > \frac{1}{3}$) when $\frac{W}{nT} > \frac{1}{10}$. The stabilization on the basis of field aligned modulation process occurs at levels $\frac{W}{nT} \lesssim 10^{-3}$-$10^{-4}$, which is much earlier. Note that this is independent of whether transverse modulational instabilities or direct collapse are faster than parallel. The fact, however, that two dimensional collapse cannot stabilize the beam instability, does not imply that it is not an interesting and potentially important aspect of the phenomenon on a long time scale and on the potential decay of the turbulence.