ON THE THEORY OF TYPE II AND TYPE III SOLAR RADIO BURSTS

I. THE IMPOSSIBILITY OF NONTHERMAL EMISSION DUE TO COMBINATION SCATTERING OFF THERMAL FLUCTUATIONS

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Abstract

"Combination scattering" as proposed by Ginzburg and Zhelezniakov involves the coalescence of electron plasma waves from a nonthermal distribution with electron plasma waves from the distribution of thermal charge fluctuations. Reabsorption is neglected. Here it is shown that reabsorption is a major effect which limits the emitted radiation to that of a black body at twice the local electron temperature. Existing models for type II and type III bursts are discussed critically and found to be inadequate. An alternative model is indicated.

I. INTRODUCTION

Type II and type III solar radio bursts involve emission at two and only two harmonics (see Wild, Smerd, and Weiss 1963 and Kundu 1965 for reviews of the observations and basic theoretical interpretation). It is widely accepted that the emission processes involve the conversion of electron plasma waves into electromagnetic waves. Emission at the fundamental frequency is identified as due to the scattering of electron plasma waves into electromagnetic waves by (the shielding cloud of electrons around) plasma ions. Emission at the second harmonic is identified as due to the coalescence of two electron plasma waves into an electromagnetic wave.

By "combination scattering" we refer to the particular form of the coalescence process as first proposed by Ginzburg and Zhelezniakov (1958). In this case electron plasma waves generated through a nonthermal emission process coalesce with electron plasma waves from the thermal distribution of such waves. This particular process is invoked in many theories where the electron plasma waves are generated coherently, and so are confined to a small cone of angles of propagation. The relevant point is that for two electron plasma waves to coalesce they must be propagating in nearly antiparallel directions, and so the only electron plasma waves with which the coherently emitted waves can coalesce come from the (always present) thermal distribution of such waves.

Reabsorption is neglected in evaluating the power radiated due to combination scattering. In the present paper it is shown that reabsorption of the emitted waves in any coalescence process, reabsorption being the reverse process called the decay of waves, restricts the effective temperature of the emitted waves to about the

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minimum of the product of either of the effective temperatures of the two coalescing distributions of waves with the ratio of the frequency of the emitted waves to that of the initial wave. Effective temperatures are functions of the wave vector $\boldsymbol{k}$ in general; we refer to the effective temperatures at the values of each wave vector involved in the coalescence process. In combination scattering the effective temperature of the thermal electron plasma waves is the local electron temperature; reabsorption limits the emission at the second harmonic to an effective temperature which turns out to be twice the electron temperature. Thus significant nonthermal emission cannot result from combination scattering, and this process can play no significant role in the emission of type II or type III solar radio bursts. For there to be significant nonthermal emission at the second harmonic both coalescing electron plasma waves must come from nonthermal distributions of such waves.

In Section II use is made of the formalism reviewed by Tsytovich (1966*a*, 1966*b*, 1967) to substantiate the above statements. In Section III a number of theories for type II and type III bursts are discussed. None of these theories accounts for all the important features of the observations, and the modifications believed to be necessary to arrive at an acceptable theory are indicated.

Throughout this article the presence of any background magnetic field is ignored. The reader is referred to Tidman, Birmingham, and Stainer (1966) for a table of the, rather extreme, field strengths needed to have a significant effect.

II. Basic Equations

In this section we summarize the semiclassical description of plasma processes, following Tsytovich (1966*a*, 1966*b*, 1967). The proof of the assertions made in Section I above is an almost trivial consequence of the equations deduced in this formalism. We are concerned with the emission and absorption of waves by particles, the scattering of waves by particles, and the coalescence and decay of waves into other waves.

(*a*) *Wave Properties*

In an isotropic plasma, waves can exist in only three wave modes. These three modes are electron plasma waves, called l-waves for simplicity; electromagnetic waves, called t-waves; and ion–acoustic waves, called s-waves. We denote the wave mode by σ with $\sigma = l$, t, or s. Each wave in the mode σ with wave vector $\boldsymbol{k}$ has a specific frequency $\omega = \omega^\sigma(\boldsymbol{k})$, a particular (normalized) polarization vector $\boldsymbol{e} = \boldsymbol{e}^\sigma(\boldsymbol{k})$, and a particular ratio of the electrical energy $W_E^\sigma(\boldsymbol{k})$ to the total energy $W_T^\sigma(\boldsymbol{k})$ in the waves. For the three wave modes in question these properties are given by (leaving the $\boldsymbol{k}$ dependences understood)

$$\omega^l = (\omega_p^2+3k^2V_e^2)^{\frac{1}{2}}, \qquad e^l = \boldsymbol{\kappa}, \qquad \left[\frac{W_E}{W_T}\right]^l = \tfrac{1}{2}\left(\frac{\omega^l}{\omega_p}\right)^2; \tag{1}$$

$$\omega^t = (\omega_p^2+k^2c^2)^{\frac{1}{2}}, \qquad e^t . \boldsymbol{\kappa} = 0, \qquad \left[\frac{W_E}{W_T}\right]^t = \tfrac{1}{2}; \tag{2}$$

$$\omega^s = kv_s/(1+k^2\lambda_{De}^2)^{\frac{1}{2}}, \qquad e^s = \boldsymbol{\kappa}, \qquad \left[\frac{W_E}{W_T}\right]^s = \tfrac{1}{2}\left(\frac{\omega^s}{\pi_i}\right)^2. \tag{3}$$

The quantities introduced in (1), (2), and (3), namely

$$\boldsymbol{\kappa} = \boldsymbol{k}/k, \qquad \omega_{\mathrm{p}} = (4\pi n_{\mathrm{e}}\, e^2/m_{\mathrm{e}})^{\frac{1}{2}} = 5\cdot 63\times 10^4\{n_{\mathrm{e}}(\mathrm{cm}^{-3})\}^{\frac{1}{2}} \quad \mathrm{sec}^{-1}, \tag{4a, b}$$

$$\pi_{\mathrm{i}} = Z_{\mathrm{i}}(m_{\mathrm{e}}/m_{\mathrm{i}})^{\frac{1}{2}}\omega_{\mathrm{p}}, \qquad V_{\mathrm{e}} = (T_{\mathrm{e}}/m_{\mathrm{e}})^{\frac{1}{2}} = 3\cdot 9\times 10^5\{T_{\mathrm{e}}(^{\circ}\mathrm{K})\}^{\frac{1}{2}} \quad \mathrm{cm\,sec}^{-1}, \tag{4c, d}$$

$$V_{\mathrm{i}} = (T_{\mathrm{i}}/m_{\mathrm{i}})^{\frac{1}{2}}, \qquad \lambda_{\mathrm{De}} = V_{\mathrm{e}}/\omega_{\mathrm{p}} = 6\cdot 9\{T_{\mathrm{e}}(^{\circ}\mathrm{K})/n_{\mathrm{e}}(\mathrm{cm}^{-3})\}^{\frac{1}{2}} \quad \mathrm{cm}, \tag{4e, f}$$

$$v_{\mathrm{s}} = \pi_{\mathrm{i}}\,\lambda_{\mathrm{De}}, \tag{4g}$$

are a unit vector in the direction of $\boldsymbol{k}$, the electron plasma frequency, the ion plasma frequency ($Z_{\mathrm{i}}\,|\,e\,|$ being ionic charge and m_{i} ionic mass), the thermal velocity of electrons (temperatures being measured in ergs unless indicated otherwise), the ion thermal velocity, the Debye length for electrons, and the sound speed respectively.

As is well known the polarization of t-waves is arbitrary in the plane orthogonal to $\boldsymbol{k}$. In all formulae below we sum over the polarizations of t-waves, which sum is achieved by the rule of thumb replacement

$$e_i^t\, e_j^{t*} \to \delta_{ij} - \kappa_i\, \kappa_j\,.$$

For l-waves to be weakly (Landau) damped they must satisfy $k \ll \lambda_{\mathrm{De}}^{-1}$, that is, their phase velocities $v_\phi^l = \omega^l/k \approx \omega_{\mathrm{p}}/k$ must be much greater than V_{e}. The group velocity of l-waves,

$$v_{\mathrm{g}}^l = \partial\omega^l/\partial k \approx 3kV_{\mathrm{e}}^2/\omega_{\mathrm{p}} = 3V_{\mathrm{e}}^2/v_\phi^l,$$

is then less than V_{e} typically. For s-waves to be weakly damped the condition $T_{\mathrm{e}} \gg T_{\mathrm{i}}$ (or $Z_{\mathrm{i}} \gg 1$, which is not the case in the corona) must be satisfied.

(b) *Distribution Functions*

We describe particles of species α (α = e or i for electrons or ions respectively) by a distribution function $f_\alpha(\boldsymbol{p})$ normalized according to

$$\int \mathrm{d}^3\boldsymbol{p}\, f_\alpha(\boldsymbol{p}) = 1\,. \tag{5}$$

The number density of such particles is denoted by n_α.

We describe the waves semiclassically, i.e. for the purposes of discussion and of setting up a formalism we use quantum mechanical language while the actual calculations remain classical. In semiclassical language waves in the mode σ with wave vector $\boldsymbol{k}$ are regarded as a collection of wave quanta with energy $\hbar\omega^\sigma(\boldsymbol{k})$ and momentum $\hbar\boldsymbol{k}$, where $\hbar = h/2\pi$ is Planck's constant. For simplicity we refer to all wave quanta as photons.

Waves in the mode σ are described by a distribution function $N^\sigma(\boldsymbol{k})$ of photons normalized in the standard way in statistical mechanics, e.g. the total energy density in the waves is given by

$$\int \mathrm{d}^3\boldsymbol{k}\ (2\pi)^{-3}\hbar\omega^\sigma(\boldsymbol{k})\, N^\sigma(\boldsymbol{k})\,.$$

A thermal distribution of waves corresponds to a Planckian distribution, which in the classical limit $\hbar \to 0$ corresponds to

$$N^{\sigma}(\boldsymbol{k}) = [\exp\{-\hbar\omega^{\sigma}(\boldsymbol{k})/T_{\mathrm{e}}\}-1]^{-1} \to T_{\mathrm{e}}/\hbar\omega^{\sigma}(\boldsymbol{k}), \tag{6}$$

where the relevant temperature is T_{e} (rather than T_{i} for $T_{\mathrm{i}} \neq T_{\mathrm{e}}$) in all cases of interest to us here. More generally, when the waves are nonthermal it is convenient and physically relevant to define the effective temperature by analogy with (6), i.e. by writing

$$T^{\sigma}(\boldsymbol{k}) = \hbar\omega^{\sigma}(\boldsymbol{k})\, N^{\sigma}(\boldsymbol{k}). \tag{7}$$

Then for a thermal distribution of waves the effective temperature reduces to the electron temperature.

(c) *Particle–Wave Interactions*

The basic formalism we use applies to both particle–wave and wave–wave interactions. It is convenient to develop the formalism for the more familiar case of particle–wave interactions. The basic interaction between particles of species α and waves in the mode σ is described by the probability per unit time that a particle $(\alpha, \boldsymbol{p})$ spontaneously emit a photon $(\sigma, \boldsymbol{k})$. We denote this probability by $w_{\alpha}^{\sigma}(\boldsymbol{p}, \boldsymbol{k})$; it is given by (see e.g. Tsytovich 1966*a*; Melrose 1968)

$$w_{\alpha}^{\sigma}(\boldsymbol{p}, \boldsymbol{k}) = \big(2(2\pi)^2 q_{\alpha}^2/\hbar\omega^{\sigma}\big)[W_E/W_{\mathrm{T}}]^{\sigma}\, |\boldsymbol{e}^{\sigma}.\boldsymbol{v}|^2\, \delta(\omega^{\sigma}-\boldsymbol{k}.\boldsymbol{v}), \tag{8}$$

where the wave properties are given by (1), (2), or (3), q_{α} and m_{α} are the charge and mass of the particle respectively, and

$$\boldsymbol{v} = \boldsymbol{p}c^2/(m_{\alpha}^2 c^4 + p^2c^2)^{\frac{1}{2}}.$$

It should be noted that (8) vanishes identically for waves with phase velocities ω^{σ}/k greater than the particle velocity; in particular (8) is identically zero for t-waves in a plasma due to $\omega^{t}/k > c$.

In this formalism, induced emission by particles $\boldsymbol{p} \to \boldsymbol{p}-\hbar\boldsymbol{k}$ and absorption by particles $\boldsymbol{p}-\hbar\boldsymbol{k} \to \boldsymbol{p}$ are described by the probability

$$w_{\alpha}^{\sigma}(\boldsymbol{p}, \boldsymbol{k})\, N^{\sigma}(\boldsymbol{k}).$$

By adding the effects of spontaneous emission (sp) and the induced processes (ind) and expanding in $|\hbar\boldsymbol{k}| \ll |\boldsymbol{p}|$, the rates of change of $N^{\sigma}(\boldsymbol{k})$ and $f_{\alpha}(\boldsymbol{p})$ are found to be given by

$$\left[\frac{\partial N^{\sigma}}{\partial t}\right]^{\mathrm{sp}} = n_{\alpha}\int \mathrm{d}^3\boldsymbol{p}\, w_{\alpha}^{\sigma} f_{\alpha}, \tag{9}$$

$$\left[\frac{\partial N^{\sigma}}{\partial t}\right]^{\mathrm{ind}} = -\gamma^{\sigma} N^{\sigma}, \qquad \gamma^{\sigma} = -n_{\alpha}\int \mathrm{d}^3\boldsymbol{p}\, w_{\alpha}^{\sigma}\hbar\boldsymbol{k}.\frac{\partial f_{\alpha}}{\partial \boldsymbol{p}}, \tag{10}$$

$$\frac{\partial f_{\alpha}}{\partial t} = \int \frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3}\, \hbar\boldsymbol{k}.\frac{\partial}{\partial \boldsymbol{p}}\left\{w_{\alpha}^{\sigma}\left(1+N^{\sigma}\hbar\boldsymbol{k}.\frac{\partial}{\partial \boldsymbol{p}}\right)f_{\alpha}\right\}, \tag{11}$$

where we leave all arguments understood, suppose that only one species of particle and waves in only one mode are involved, and combine spontaneous and induced processes in (11). The quantity γ^σ is called the absorption coefficient.

It follows from (9) and (10) that when the waves come into equilibrium due to emission and reabsorption by particles we have

$$N^\sigma(\boldsymbol{k}) = -\int \mathrm{d}^3\boldsymbol{p}\, w_\alpha^\sigma(\boldsymbol{p},\boldsymbol{k})\, f_\alpha(\boldsymbol{p}) \Big/ \int \mathrm{d}^3\boldsymbol{p}\, w_\alpha^\sigma(\boldsymbol{p},\boldsymbol{k}) \hbar\boldsymbol{k} \cdot \frac{\partial f_\alpha}{\partial \boldsymbol{p}}(\boldsymbol{p}), \tag{12}$$

which is a statement of Kirchhoff's law. Inserting a thermal distribution of electrons in (12), one rederives (6) for l-waves and s-waves; the effect of thermal ions can be shown to be negligible by summing over all species of particles in (9) and (10) and rewriting (12) appropriately.

Provided (12) leads to a positive value of $N^\sigma(\boldsymbol{k})$ (as negative values imply coherent emission and so no equilibrium), it can be shown that this equation implies that particles with velocity $v \leqslant v_0$ generate l-waves with an effective temperature, for $v_\phi^l = \omega_\mathrm{p}/k \leqslant v_0$, which never exceeds a value of the order of the kinetic energy of particles with velocity v_0.

(*d*) *Scattering Processes*

In the same manner we treat scattering processes. Let $w_\alpha^{\sigma\sigma'}(\boldsymbol{p},\boldsymbol{k},\boldsymbol{k}')$ be the probability per unit time that a particle $(\alpha,\boldsymbol{p})$ scatter a photon $(\sigma',\boldsymbol{k}')$ into another photon $(\sigma,\boldsymbol{k})$. The reverse process is described by the same probability. The rates of change of N^σ and $N^{\sigma'}$ are then described by (see e.g. Tsytovich 1966*b*, 1967)

$$\left[\frac{\partial N^\sigma}{\partial t}\right]^{\mathrm{sp}} = n_\alpha \int \mathrm{d}^3\boldsymbol{p} \int \frac{\mathrm{d}^3\boldsymbol{k}'}{(2\pi)^3} w_\alpha^{\sigma\sigma'} N^{\sigma'} f_\alpha, \tag{13}$$

$$\left[\frac{\partial N^{\sigma'}}{\partial t}\right]^{\mathrm{sp}} = -n_\alpha \int \mathrm{d}^3\boldsymbol{p} \int \frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3} w_\alpha^{\sigma\sigma'} f_\alpha, \tag{14}$$

$$\left[\frac{\partial N^\sigma}{\partial t}\right]^{\mathrm{ind}} = N^\sigma n_\alpha \int \mathrm{d}^3\boldsymbol{p} \int \frac{\mathrm{d}^3\boldsymbol{k}'}{(2\pi)^3} w_\alpha^{\sigma\sigma'} N^{\sigma'} \hbar(\boldsymbol{k}-\boldsymbol{k}') \cdot \frac{\partial f_\alpha}{\partial \boldsymbol{p}}, \tag{15}$$

where all arguments are left understood and the remaining equations describing the change in f_α are omitted. The rate of change of $N^{\sigma'}$ due to the induced processes follows by interchanging primed and unprimed quantities in (15).

The scattering probabilities are cumbersome expressions in general due to the large number of effects which contribute to the scattering. The scattering ascribed to either an individual electron or an individual ion involves three effects. Firstly the particles scatter the waves through Compton scattering (negligible for ions). Secondly the scattering can be due to charge density inhomogeneities due to the shielding cloud of plasma particles around the scattering particle. Thirdly the scattering can be due to transverse (orthogonal to $\boldsymbol{k}$) currents associated with the shielding cloud. Furthermore, the effect of the shielding cloud of particles can be separated into the effects due to the shielding cloud of electrons and that of ions.

In most applications the effect of transverse currents can be neglected and the effect of the shielding cloud of ions neglected compared to that of electrons. The remaining effects, Compton scattering and the scattering off the shielding cloud of electrons, tend to cancel for electrons. Thus the dominant effect is usually the scattering off the shielding cloud of electrons around individual ions.

The probability for scattering of l-waves into t-waves due to the last effect reduces to the relatively simple expression

$$w_{\mathrm{i}}^{tl}(\boldsymbol{p},\boldsymbol{k}_t,\boldsymbol{k}_l) = \frac{(2\pi)^3 Z_{\mathrm{i}}^2 e^4}{m_{\mathrm{e}}^2\,\omega^t\,\omega^l}\,|\,\boldsymbol{\kappa}_t\times\boldsymbol{\kappa}_l\,|^2\,\delta\{\omega^t-\omega^l-(\boldsymbol{k}_t-\boldsymbol{k}_l)\,.\,\boldsymbol{v}\}\,. \tag{16}$$

The δ function implies that $\omega^t \approx \omega^l \approx \omega_{\mathrm{p}}$. The corresponding probabilities for scattering of l-waves into l-waves and t-waves into t-waves differ from (16) only by an obvious relabelling plus the replacement of $|\,\boldsymbol{\kappa}_1\times\boldsymbol{\kappa}_2\,|^2$ by $|\,\boldsymbol{\kappa}_1\,.\,\boldsymbol{\kappa}_2\,|^2$ and $\frac{1}{2}(1+|\,\boldsymbol{\kappa}_1\,.\,\boldsymbol{\kappa}_2\,|^2)$ respectively, where unscattered and scattered waves are labelled by 1 and 2.

In any scattering process involving nonthermal unscattered waves σ' and thermal particles, the scattering process proceeds in the direction $\sigma' \rightarrow \sigma$ if and only if the effective temperature of the scattered waves σ remains less than that of the unscattered waves σ'. This follows directly from (13), (14), and (15) and the corresponding equations for $\sigma' \rightarrow \sigma$.

(e) *Coalescence and Decay Processes*

Similarly let $w_{\sigma}^{\sigma'\sigma''}(\boldsymbol{k},\boldsymbol{k}',\boldsymbol{k}'')$ be the probability per unit time that two photons $(\sigma',\boldsymbol{k}';\,\sigma'',\boldsymbol{k}'')$ spontaneously coalesce into a single photon $(\sigma,\boldsymbol{k})$, which probability is also that of the reverse process. In the presence of distributions N^{σ}, $N^{\sigma'}$, and $N^{\sigma''}$ of waves the total probabilities of the processes $\sigma'+\sigma'' \rightarrow \sigma$ and $\sigma \rightarrow \sigma'+\sigma''$ are given respectively by

$$w_{\sigma}^{\sigma'\sigma''}\,N^{\sigma'}\,N^{\sigma''}(1+N^{\sigma}), \qquad w_{\sigma}^{\sigma'\sigma''}(N^{\sigma'}+1)(N^{\sigma''}+1)N^{\sigma}\,.$$

The unit terms correspond to the spontaneous processes which are of no significance in themselves. It is then trivial to show that the rates of change of N^{σ}, $N^{\sigma'}$, and $N^{\sigma''}$ are given respectively by

$$\frac{\partial N^{\sigma}}{\partial t} = \int\frac{\mathrm{d}^3\boldsymbol{k}'}{(2\pi)^3}\int\frac{\mathrm{d}^3\boldsymbol{k}''}{(2\pi)^3}\,w_{\sigma}^{\sigma'\sigma''}(N^{\sigma'}\,N^{\sigma''}-N^{\sigma'}\,N^{\sigma}-N^{\sigma''}\,N^{\sigma})\,, \tag{17}$$

$$\frac{\partial N^{\sigma'}}{\partial t} = -\int\frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3}\int\frac{\mathrm{d}^3\boldsymbol{k}''}{(2\pi)^3}\,w_{\sigma}^{\sigma'\sigma''}(N^{\sigma}\,N^{\sigma''}-N^{\sigma'}\,N^{\sigma}-N^{\sigma''}\,N^{\sigma})\,, \tag{18}$$

$$\frac{\partial N^{\sigma''}}{\partial t} = -\int\frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3}\int\frac{\mathrm{d}^3\boldsymbol{k}'}{(2\pi)^3}\,w_{\sigma}^{\sigma'\sigma''}(N^{\sigma'}\,N^{\sigma''}-N^{\sigma'}\,N^{\sigma}-N^{\sigma''}\,N^{\sigma})\,, \tag{19}$$

where again all arguments are left understood, and the terms corresponding to the spontaneous decay $\sigma \rightarrow \sigma'+\sigma''$ are omitted.

The probability for the process $l+l \to t$ is given by (see Tsytovich 1967)

$$w_t^{ll}(\boldsymbol{k}_t, \boldsymbol{k}'_l, \boldsymbol{k}''_l) = \frac{(2\pi)^5}{8}\frac{\hbar e^2}{m_e^2}\frac{(k_l'^2-k_l''^2)^2}{\omega^t k_t^2}\,|\,\boldsymbol{\kappa}'_l\times\boldsymbol{\kappa}''_l\,|^2\,\delta^3(\boldsymbol{k}_t-\boldsymbol{k}'_l-\boldsymbol{k}''_l)\,\delta(\omega^t-\omega^{l'}-\omega^{l''})\,, \tag{20}$$

while that for $t+l \to t$ is given by

$$w_t^{tl}(\boldsymbol{k}_t, \boldsymbol{k}'_l, \boldsymbol{k}'_t) = \frac{(2\pi)^5}{8}\frac{\hbar e^2}{m_e^2}\frac{\omega_p k_l'^2}{\omega^t\omega^{t'}}(1+|\,\boldsymbol{\kappa}_t\cdot\boldsymbol{\kappa}'_t\,|^2)\,\delta^3(\boldsymbol{k}_t-\boldsymbol{k}'_l-\boldsymbol{k}'_t)\,\delta(\omega^t-\omega^{l'}-\omega^{t'})\,. \tag{21}$$

We also note that the probability for the process $l+s \to t$ is

$$w_t^{ls}(\boldsymbol{k}_t, \boldsymbol{k}'_l, \boldsymbol{k}'_s) = \frac{(2\pi)^5}{4}\frac{\hbar e^2}{m_e^2}\frac{\omega_p^2}{\pi_i^2}\frac{(\omega^s)^3}{k_s'^2 V_e^4}\,|\,\boldsymbol{\kappa}_t\times\boldsymbol{\kappa}'_l\,|^2\,\delta^3(\boldsymbol{k}_t-\boldsymbol{k}'_l-\boldsymbol{k}'_s)\,\delta(\omega^t-\omega^{l'}-\omega^{s'})\,, \tag{22}$$

while the probability for $l+s \to l$ follows by a relabelling of t by l and that for $t+s \to t$ follows by an obvious relabelling plus the replacement of $|\,\boldsymbol{\kappa}_t\times\boldsymbol{\kappa}_l'\,|^2$ by $(\omega_p^2/\omega^t\omega^{t'})\frac{1}{2}(1+|\,\boldsymbol{\kappa}_t\cdot\boldsymbol{\kappa}_t'\,|^2)$. These probabilities are again approximate but the approximations in (20) and (21) involve only expansions in the ratios V_e/v_ϕ and m_e/m_i.

(*f*) *Consequences of Reabsorption*

The coalescence process proceeds in the direction $\sigma'+\sigma'' \to \sigma$, rather than in the direction of the decay process $\sigma \to \sigma'+\sigma''$, if and only if the integrands in (17), (18), and (19) are positive, i.e. if and only if we have

$$N^{\sigma'}(\boldsymbol{k}')\,N^{\sigma''}(\boldsymbol{k}'') > N^{\sigma}(\boldsymbol{k})\{N^{\sigma'}(\boldsymbol{k}')+N^{\sigma''}(\boldsymbol{k}'')\}\,, \tag{23}$$

$$\boldsymbol{k} = \boldsymbol{k}'+\boldsymbol{k}''\,, \qquad \omega^{\sigma}(\boldsymbol{k}) = \omega^{\sigma'}(\boldsymbol{k}')+\omega^{\sigma''}(\boldsymbol{k}'')\,.$$

Using (7), condition (23) obviously implies an inequality for the effective temperatures. In particular for the coalescence process $l+l \to t$ with

$$N^l(\boldsymbol{k}) = N_1^l(\boldsymbol{k})+N_2^l(\boldsymbol{k})\,,$$

where $N_1^l(\boldsymbol{k})$ is a nonthermal distribution which gives zero emission at $2\omega_p$ due to angular limitations (e.g. for $N_1^l(\boldsymbol{k})$ a unidirectional distribution) and $N_2^l(\boldsymbol{k})$ a thermal distribution, (23) implies

$$N_1^l(\boldsymbol{k}')\,N_2^l(\boldsymbol{k}'') > N^t(\boldsymbol{k})\{N_1^l(\boldsymbol{k}')+N_2^l(\boldsymbol{k}'')\}\,,$$

where we use the symmetry in the $\boldsymbol{k}$ integrals in (17) and neglect the terms $N_1^l(\boldsymbol{k}')\,N_1^l(\boldsymbol{k}'')$ (assumed to be zero) and $N_2^l(\boldsymbol{k}')\,N_2^l(\boldsymbol{k}'')$ (which can lead only to thermal emission). Then using (7) with $\omega^l \approx \omega_p$ and $\omega^t \approx 2\omega_p$ we find

$$T_1^l(\boldsymbol{k}')\,T_2^l(\boldsymbol{k}'') > \tfrac{1}{2}T^t(\boldsymbol{k})\{T_1^l(\boldsymbol{k}')+T_2^l(\boldsymbol{k}'')\}\,. \tag{24}$$

For $T_2^l = T_e$ and $T_1^l \gg T_e$, as is the case in combination scattering off thermal fluctuations, (24) leads to

$$T^t < 2T_e\,.$$

Thus in combination scattering as formulated by Ginzburg and Zhelezniakov (1958) reabsorption limits the effective temperature of the radiation emitted at the second harmonic to that of a black body at twice the local electron temperature.

Zaitsev and Kaplan (1966, Section 5) consider an analogous process in which nonthermal s-waves coalesce with thermal electron plasma waves. In this case reabsorption limits the effective temperature of the emitted radiation, now at the fundamental, i.e. at $\omega^t = \omega_p$, for this process to that of a black body at the local electron temperature.

More generally reabsorption in any coalescence process $\sigma'+\sigma'' \rightarrow \sigma$ limits the effective temperature T^σ of the waves so produced to less than about the *minimum* of $(\omega^\sigma/\omega^{\sigma'})T^{\sigma'}$, $(\omega^\sigma/\omega^{\sigma''})T^{\sigma''}$.

Returning to the process $l+l \rightarrow t$, the wavenumber of the emitted t-wave is given by (see equations (2))

$$(2\omega_p)^2 = \omega_p^2+k_t^2 c^2\,, \qquad k_t = \sqrt{3}\,\omega_p/c\,. \tag{25}$$

On the other hand, all l-waves generated by particles with velocity v have $v_\phi^l \leqslant v$ and so

$$k_l \geqslant \omega_p/v > \omega_p/c\,.$$

Even for $v = c$ the δ function for wave vectors in (20) requires that the two coalescing l-waves be at an angle of 60°, that is, $\boldsymbol{k}_l'\cdot\boldsymbol{k}_l'' = \frac{1}{2}|\boldsymbol{k}_l'||\boldsymbol{k}_l''|$. For $k_l \gg \omega_p/c$ the two coalescing l-waves must be approximately antiparallel.

The implication is that a unidirectional distribution of l-waves generated by particles, as results from a two-stream instability, leads to only approximately thermal emission at the second harmonic. Nonthermal emission at the second harmonic occurs only if the two coalescing l-waves come from nonthermal distributions. Thus there must be either scattered l-waves or scattered particles generating l-waves over a wide cone of angles for nonthermal emission at the second harmonic to be possible in theories based on the two-stream instability. The scattering of l-waves occurs (see above) and the scattering of fast electrons due to collisions is effective enough to be considered to be of relevance in placing a lower limit on the velocity of the electrons causing type III bursts (see Kundu 1965 and references therein). Thus the objection raised here applies to particular models but is not difficult to overcome by modifying the models to include the above effects.

III. Existing Models

In this section we consider several of the models found in the literature for type II and type III bursts with a view to determining an acceptable class of model. A number of theories are open to theoretical objections, including the objection raised above. We find that the other theories discussed below lead to predictions in conflict with observation. It is suggested that the observations point to a particular class of model but only at the expense of invoking physical processes of whose validity we are uncertain.

The theoretical objection raised above applies particularly to the model proposed by Ginzburg and Zhelezniakov (1958) and to subsequent modifications of this model (e.g. Smerd, Wild, and Sheridan 1962; Wild, Smerd, and Weiss 1963; Zheleznjakov 1964; Kundu 1965). Other objections to these and further models based on a two-stream instability are discussed below.

The first class of model to be considered applies only to type II bursts and avoids all difficulties encountered with the two-stream instability. This class of model is based on the work of Tidman and Dupree (1965) and modifications thereof. The basic idea of the model (see Tidman 1965) is that suprathermal electrons are either produced by or are comoving with the shock front associated with type II bursts and that these electrons cause l-waves to be excited well above the thermal equilibrium value in a large volume behind the shock front.

(a) *Incoherently Emitted l-waves*

In the model proposed by Tidman (1965) the l-waves are generated due to incoherent emission and reabsorption by suprathermal electrons. For isotropically distributed electrons, (12) with (7), (8), and (1) for $k \ll \lambda_{\mathrm{De}}^{-1}$ gives

$$T^l(k) = -\int_{v>v_\phi} \frac{p^2\,\mathrm{d}p}{v} f_\mathrm{e}(p) \Big/ \int_{v>v_\phi} \frac{p^2\,\mathrm{d}p}{v^2} \frac{\partial f_\mathrm{e}}{\partial p}(p)\,, \quad v_\phi = \omega_\mathrm{p}/k\,. \tag{26}$$

On separating $f_\mathrm{e}(p)$ into a thermal (T) and suprathermal (F) part by writing

$$f_\mathrm{e}(p) = \beta f_\mathrm{e}^\mathrm{T}(p) + (1-\beta) f_\mathrm{e}^\mathrm{F}(p)\,, \tag{27}$$

with

$$f_\mathrm{e}^\mathrm{T}(p) = \{(2\pi)^{3/2} m_\mathrm{e}^3 V_\mathrm{e}^3\}^{-1} \exp(-v^2/2V_\mathrm{e}^2)\,,$$

where $(1-\beta)/\beta \ll 1$ is the ratio of the number density of suprathermal to thermal electrons, (26) leads to $T^l(k)$ determined solely by $f_\mathrm{e}^\mathrm{F}(p)$ for v_ϕ sufficiently in excess of V_e. If we choose $f_\mathrm{e}^\mathrm{F}(p)$ to be a thermal distribution with temperature $T^\mathrm{F} \gg m_\mathrm{e} V_\mathrm{e}^2$ $(= T_\mathrm{e})$ then (26) gives

$$T^l(k) = T^\mathrm{F}\,, \quad \alpha V_\mathrm{e} \lesssim v_\phi < c\,; \qquad \alpha^2 = 2\ln\left(\frac{V_E}{V_\mathrm{e}} \frac{\beta}{1-\beta}\right), \tag{28}$$

where V_E is given by $T^\mathrm{F} = m_\mathrm{e} V_E^2$ for $V_E \ll c$ and by $V_E \approx c$ for $T^\mathrm{F} \gtrsim m_\mathrm{e} c^2$. Any other choice of $f_\mathrm{e}^\mathrm{F}(p)$ leads to $T^l(k)$ a decreasing function of k, e.g. for

$$\begin{aligned} f_\mathrm{e}^\mathrm{F}(p) &= (2\pi/m_\mathrm{e}^3 v_0^3)(v/v_0)^{-1} \qquad &&\text{for} \qquad v \leqslant v_0 \leqslant c\,, \\ &= 0\,, && \qquad\qquad v > v_0\,, \end{aligned}$$

we have

$$T^l(k) = m_\mathrm{e} v_0 \omega/k\,, \qquad \alpha V_\mathrm{e} \lesssim \omega_\mathrm{p}/k \leqslant v_0\,,$$

where α is a logarithmic factor similar to that in (28).

Under no circumstances can this class of model lead to emission of t-waves with effective temperature in excess of the maximum value of $T^l(k)$ as a function of k. The observed effective temperatures of type II bursts range up to $T^t \approx (0\cdot1\text{–}1) m_\mathrm{e} c^2$ $(T = m_\mathrm{e} c^2$ corresponds to $T = 0\cdot6 \times 10^{10}\,{}^\circ\mathrm{K})$ so that at least mildly relativistic electrons must be involved. Furthermore, we argue that we must have $T^t \ll \max T^l(k)$

because as T^t approaches T^l the process $l+t \to t$ becomes comparable in rate to the process $l+l \to t$ (cf. equations (20) and (21)) thereby predicting significant emission at $3\omega_p$ contrary to observation. (It is noted that Tidman (1965) obtains spurious agreement with observation for $T^F = m_e V_E^2 \approx 0 \cdot 1\, m_e c^2$ due to a numerical slip in evaluating the radiated power.)

Papadopoulos (1969) extended the analysis of Tidman and Dupree (1965) to include relativistic particles. The results of Papadopoulos show that, apart from numerical factors, the radiated power is given by intuitively obvious generalizations of formulae given by Tidman and Dupree (1965). Let us rederive these formulae from the formalism presented in Section II above.

The emission at the fundamental frequency ($\omega^t \approx \omega_p$) follows from (13) with (16). For $|(\boldsymbol{k}_t - \boldsymbol{k}_l)\cdot\boldsymbol{v}| \ll |\omega^t - \omega^l|$ in (16), the δ function gives $\omega^t \approx \omega_p$ with

$$\delta(\omega^t - \omega^l) \approx (\omega_p / k_t c^2)\, \delta\{k_t - \sqrt{3}\, k_l (V_e / c)\}, \tag{29}$$

where we use (1) and (2). The above indicated inequality with $v \sim V_i$ then requires

$$|(\boldsymbol{k}_t - \boldsymbol{k}_l)\cdot\boldsymbol{v}| \approx k_l V_i \ll 3 k_l^2 V_e^2 / 2\omega_p^2,$$

$$v_\phi^l / V_e \ll V_e / V_i \approx 43.$$

This may be violated in the corona ($V_e \approx 10^{-2} c$) for $v_\phi^l \approx c$ but it turns out that the emission at ω_p is dominated by l-waves with $v_\phi^l \ll c$ (see below) so that (29) remains valid.

With (29) the $\boldsymbol{p}$ integral in (13) is trivial. We then find

$$\frac{\partial N^t}{\partial t}(\boldsymbol{k}_t) = \sum_i \frac{Z_i^2 n_i e^4}{m_e^2 \omega_p} \int \mathrm{d}^3 \boldsymbol{k}_l\, N^l(\boldsymbol{k}_l)\, |\boldsymbol{\kappa}_t \times \boldsymbol{\kappa}_l|^2 \frac{\omega_p}{k_t c^2}\, \delta\!\left(k_t - \sqrt{3}\, \omega_p \frac{V_e}{c}\right), \tag{30}$$

where we sum over all ionic species. For isotropically distributed l-waves (30) predicts a power radiated per unit volume at ω_p given by

$$P_{\omega_p} = \int \frac{\mathrm{d}^3 \boldsymbol{k}_t}{(2\pi)^3} \hbar\omega_p \frac{\partial N^t}{\partial t}(\boldsymbol{k}_t) = \frac{e^2 \omega_p V_e}{\sqrt{3}\, \pi^2 c} \int \mathrm{d}k\, k^3 \frac{T^l(k)}{m_e c^2}, \tag{31}$$

where we suppose all ions to be protons. A formula analogous to (31) was found by Tidman and Dupree (1965).

For the emission at the second harmonic, $\omega^t = 2\omega_p$, we use (17), ignoring reabsorption, with (20). The δ function for the frequency in (20) reduces to (c.f. equation (25))

$$\delta(\omega^t - \omega^{l'} - \omega^{l''}) \approx (2\omega_p / k_t c^2)\, \delta(k_t - \sqrt{3}\, \omega_p / c),$$

while for $v_\phi^l \ll c$ we have

$$\{(k_l'^2 - k_l''^2)^2 / k_t^2\}\, |\boldsymbol{\kappa}_l' \times \boldsymbol{\kappa}_l''|^2 = \{(k_l'^2 - k_l''^2)^2 / k_l''^2\}\, |\boldsymbol{\kappa}_t \times \boldsymbol{\kappa}_l'|^2$$

$$\approx 4 k_t^2\, |\boldsymbol{\kappa}_t \cdot \boldsymbol{\kappa}_l'|^2\, |\boldsymbol{\kappa}_t \times \boldsymbol{\kappa}_l'|^2.$$

This approximation, $v_\phi^l \ll c$, leads to numerical underestimation of the power radiated at $2\omega_p$ for waves with $v_\phi^l \approx c$ but gives the correct functional form (see Papadopoulos and Lerche 1969*a*, 1969*b*). The power radiated per unit volume at $2\omega_p$ then reduces to

$$P_{2\omega p} = \frac{2\sqrt{3}}{5\pi^2}\frac{e^2\omega_p^2}{c}\int k^2\,dk\left(\frac{T^l(k)}{m_e c^2}\right)^2, \tag{33}$$

where we again assume the l-waves to be isotropically distributed. Apart from a factor of two, Tidman and Dupree (1965) found a formula analogous to (33). Their underestimation by a factor of two is due to terms omitted in the formulation of Dupree (1964) (see Papadopoulos and Lerche 1969*a*).

On comparing (31) with (33) we draw two important qualitative conclusions. Firstly any overall increase in T^l causes an increase in the ratio $P_{2\omega p}/P_{\omega p}$ because T^l occurs squared in (33). Secondly if we compare $P_{\omega p}/P_{2\omega p}$ for T^l constant and for T^l decreasing with k, more slowly than $k^{-3/2}$, from the same constant maximum value, then we conclude that $P_{\omega p}/P_{2\omega p}$ is a maximum for the situation where T^l remains constant (for T^l decreasing faster than $k^{-3/2}$ the overall power is greatly diminished).

Observationally the intensities at the two harmonics are about equal. In view of the greater difficulty which the fundamental has in escaping from the corona (see Roberts 1959), any model which predicts $P_{\omega p} \ll P_{2\omega p}$ is in direct conflict with observation. The comments above imply that the most favourable choice of $f_e^F(p)$ in (27) is one when it is a thermal (but high temperature) distribution. Papadopoulos (1969) considered such a situation with $T^F = m_e c^2$ in (28) and included all the numerical enhancements which result when l-waves with $v_\phi^l \approx c$ are involved. He purported to find $P_{\omega p} \approx P_{2\omega p} \approx 10^{-12}$ erg cm^{-3} sec^{-1} and so, with an emitting volume of 10^{30} cm^{-3}, a total power of 10^{18} erg sec^{-1} (actually up to 10^{19} erg sec^{-1} can be radiated in type II bursts). However, the choice of parameters made by Papadopoulos (1969) actually leads to $P_{\omega p} \approx 0\cdot 1\,P_{2\omega p}$, $P_{2\omega p} \approx 10^{-12}$ erg cm^{-3} sec^{-1}. Thus even for the most favourable case the model predicts $P_{\omega p} \ll P_{2\omega p}$.

This discrepancy seems to be inescapable in the model as formulated, i.e. the model either leads to too little power radiated or leads to $P_{\omega p} \ll P_{2\omega p}$. It might be thought that inclusion of the process $l+s \rightarrow t$ could enhance the emission at ω_p. For thermal s-waves reabsorption limits the effective temperature of the emission at the fundamental to $\omega_p T_e/\omega^s \lesssim 10^2 T_e \sim 10^8\,$°K, where we set $\omega^s = kv_s$ and $k \sim \omega_p/aV_e$. Because the effective temperature observed is greater than $10^8\,$°K, the process $l+s \rightarrow t$ requires that nonthermal s-waves be present throughout the emitting volume. The mere presence of suprathermal electrons leads to no significant enhancement in the level of s-waves. Coherently generated s-waves are rapidly Landau damped and so could not possibly be excited over a volume of the order necessary in this model.

A further discrepancy with observation is in the frequency dependence predicted. Papadopoulos and Lerche (1969*a*) point out that the functional form gives $P_{\omega p} \propto \omega_p^5$ for any reasonable choice of $f_e^F(p)$. The intensity of the emitted radiation is proportional to (emitting volume) times $(P_{\omega p}/\omega_p)$. The most extreme assumption that can be made about the emitting volume is that it expands spherically, i.e. as (time elapsed)^3. Observationally the radiated frequency decreases roughly linearly with time so that this extreme assumption corresponds to the intensity being proportional to $\omega_p^{-3}(P_{\omega p}/\omega_p) \propto \omega_p$. Observationally the intensity increases with *decreasing frequency*.

(b) *Coherently Emitted l-waves*

It would appear that the class of model discussed above proves inadequate for type II bursts. This class of model is not directly applicable to type III bursts. If models based on incoherent emission and reabsorption of l-waves are inadequate then we must appeal to the generation of l-waves directly by the propagating disturbance itself. For type III bursts the disturbance appears to be a bunch of fast electrons (see e.g. Wild, Smerd, and Weiss 1963). Incoherent emission by a monoenergetic bunch of fast electrons is examined in Part II (present issue pp. 885–903) and there found to be incapable of accounting for the observed power radiated in type III bursts (here we note that about 10^{38} electrons per bunch are required to account for the power radiated in type III bursts if the electrons radiate l-waves only incoherently). Thus we are led to return to theories based on coherent emission. Firstly we consider the possibility that type II bursts are generated in an analogous way to type III bursts, i.e. by bunches of fast moving electrons.

Smerd, Wild, and Sheridan (1962) noted that type II bursts include series of rapidly drifting bursts so reminiscent of type III bursts that they conjectured that the slow moving shock front was the seat of a continuous ejection of fast electrons. Some theoretical support for this view is provided by the acceleration mechanism invoked by Lacombe and Mangeney (1969). These authors were concerned with the acceleration of electrons to mildly relativistic energies in order to account for type IV emission (synchrotron radiation). However, the acceleration mechanism invoked, which involves acceleration by s-waves coherently emitted due to a current instability in the shock front, produces fast non-relativisitic electrons copiously if it produces any relativistic electrons at all. It seems reasonable to identify these electrons with those whose presence was conjectured by Smerd, Wild, and Sheridan (1962), and then to regard type II bursts as a superposition of many localized type III bursts.

Besides the objection raised in the present article (this objection is relatively easily overcome) two other strong objections to the view that type III bursts are due to a stream of fast electrons coherently emitting l-waves are presented in the literature. Firstly, Sturrock (1964) points out that the two-stream instability is such an efficient instability that the streaming motion is stopped after propagating absurdly short distances compared to those over which type III bursts are observed to be generated. Thus any acceptable model must invoke processes which are highly efficient in suppressing the two-stream instability.

Secondly, Tidman, Birmingham, and Stainer (1966) point out that it is expected, on theoretical grounds, that the damping time for l-waves generated through a two-stream instability be the Landau damping time rather than the collision time. These authors refer in particular to models proposed by Pikel'ner and Gintsburg (1964) and Zheleznakov (1965) where the damping time is assumed to be the collision time. Observationally, the damping time, if identified with the duration of type III bursts at fixed frequency, corresponds so well with the collisional time that it is even used to estimate the temperature in the emitting region (see Kundu 1965 and references therein). If the l-waves are collisionally damped after the passage of the stream of electrons, then this implies that any two-stream instability is so effectively suppressed that the velocity distribution of the streaming electrons is not significantly affected. (In the two-stream instability the electron distribution tends to form a "plateau" (see e.g. Shapiro 1963) such that the velocity gap between a few times V_e and the streaming velocity is filled. The slower electrons filling this gap are outpaced and so remain behind to Landau damp the l-waves. Evidently this plateau cannot form if the l-waves are to be only collisionally damped.)

One model which overcomes some of the objections raised above is that proposed by Kaplan and Tsytovich (1967). These authors argue that the two-stream instability is suppressed due to nonlinear processes, in particular due to the coherent version of the scattering of l-waves into l-waves by plasma particles (see equation (15) and the comments after (16) in Section II(d) above). This process causes scattering of l-waves as fast as they are generated by the instability thereby stopping the exponential growth of the l-waves. These authors also assume that the scattering of l-waves into t-waves is coherent (see (15) and (16)). Not surprisingly the emission region is then optically thick to the t-waves so generated due to the high efficiency of the conversion of l-waves into t-waves.

An embarrassingly intense emission results in this model. This is because the effective temperature of the t-waves approaches that of the l-waves, which effective temperature is very large whenever coherent emission is involved. Although Kaplan and Tsytovich (1967) do not discuss the emission at $2\omega_p$ in detail it is clear from their model that the effective temperature of the t-waves at $2\omega_p$ also approaches that of the l-waves. As already pointed out, one then expects emission at $3\omega_p$, $4\omega_p$, . . . to arise due to the process $l+t \rightarrow t$. Indeed Colgate (1967) carries this process to its logical conclusion in his theory for the infrared emission from quasi-stellar objects; by analogy we might expect the maximum radiated power to be at many times ω_p in the model of Kaplan and Tsytovich.

It would appear that any two-stream instability needs to be more strongly suppressed than by the process considered by Kaplan and Tsytovich (1967). For coherent scattering processes to be unimportant, the energy density in l-waves must remain much less than the thermal energy density (of particles) (see Tsytovich 1966b). That coherent scattering does not play an important role is indicated both by the absence of higher harmonics in the emission and by the relaxation time being the collision time.

All these arguments suggest that in order to retain the widely held and well-founded view that type III bursts are due to the passage of a bunch of fast electrons we must find an effective way of suppressing the two-stream instability. One suggestion which does not seem to have been considered is the following. Consider a stream of electrons with velocity $u \gg V_e$ and number density $n_s \ll n_e$. In order for charge neutrality to be maintained thermal electrons must stream in the opposite direction with velocity U where

$$U = (n_s/n_e)u. \tag{34}$$

If U exceeds the sound speed v_s, that is, if the thermal electrons stream relative to the ions at greater than the sound speed, then a current instability develops. The

s-waves generated in the current instability are favourably directed with respect to the l-waves generated in the two-stream instability for the two instabilities to suppress each other. It seems that if the s-waves grow faster due to the current instability than do the l-waves due to the two-stream instability then the latter instability can be very effectively suppressed.

If the two-stream instability is sufficiently strongly suppressed then one expects the l-waves generated by a bunch of electrons to be collisionally damped. The possibility of accounting for type III and type II bursts in this way is explored in Part II.

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