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A Manifestation of Quasilinear Diffusion in Whistlers in the Fine Structure of Type IV Solar Radio Bursts

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Abstract—A number of properties of the emission and absorption bands seen against the background of the type IV radio continuum (zebra structure) are explained by resonance switching in whistler instabilities for an anisotropic-beam-type distribution function (similar to the fan instability). There should be continuous diffusion in whistlers when there is a prolonged injection of particles. The instability can switch from a normal to an anomalous Doppler resonance in accordance with the shift of the distribution function maximum during a hollowing of the beam. This effect manifests itself as smooth or abrupt (depending on the whistler energy density) oscillations of the zebra-structure bands on the dynamical spectrum. Simultaneously, absorption appears at the low-frequency edge of the bands, since the instability of plasma waves within the whistler wave packet instantaneously decreases due to diffusion.

1. INTRODUCTION

The sources of type IV solar radio bursts are magnetic traps in which a conical distribution of energetic particles forms, exciting high-frequency (plasma) and low-frequency waves, in particular, whistlers. Emission and absorption bands against the background of the type IV continuum are conventionally divided into two types: zebra structure, observed at meter wavelengths (<300 MHz), with a predominantly positive frequency drift, and fiber bursts (intermediate drift bursts, IDB), with a negative drift [1].

The radioemission (t) of IDB fibers is unambiguously associated with the merging of plasma waves (l) with whistlers (w): $l + w \longrightarrow t$ [2]. The absorption accompanying every fiber at the low-frequency edge is explained by the efficient removal of plasma waves by the process $l + w \longrightarrow t$ or by screening of the radio emission by density inhomogeneities [2, 3]. The most common mechanism proposed for the zebra structure is emission in bands at twice the plasma frequency [4] and related schemes [5, 6]; in nearly all such mechanisms, the origin of the absorption between the bands of enhanced emission is not considered. The most substantial inconsistency of double-plasma-resonance mechanisms is that they require a rapid time variation of the magnetic field to explain the dynamics of the zebra structure bands, while very low magnetic field strengths are implied by the approximate equality of the frequency separation of the bands (Δf) and of the electron gyrofrequency (f_{ce}). As a rule, the implied magnetic field strengths (H) in the source vary irregularly with height in the corona [1], and are considerably lower than those required for the magnetic pressure to be greater than the kinetic pressure ($\beta < 1$), which is a condition for the existence of a magnetic trap.

In most of their basic parameters, however, zebra structure bands and IDB fibers display notable similarities: they have the same frequency separation between bands [1], location for the absorption at the low-frequency edge, and depth of continuum modulation. They also both have strong polarization, corresponding to the ordinary wave [7]. The collected observational data suggest that the zebra structure bands and IDB fibers are related, and that their origin can be explained by a common mechanism—the interaction of plasma waves with whistlers: $l + w \longrightarrow t$ [8]. The fibers are likely associated with channeled propagation of whistlers from the depths of the corona along a trap, while the zebra structure is associated with nonchanneled propagation oblique to the magnetic field, mainly at the top of the trap [9].

Long-term spectral fine structure observations reveal a number of features that cannot be easily described in the framework of proposed mechanisms for the zebra structure. For instance, in a number of events, we observed quasi-parallel bands with a wave-like frequency drift (Fig. 1); a continuous transition of zebra-structure bands with wave-like drifts to typical IDB fibers with constant negative drift; or strange fibers in type II bursts, similar to the zebra structure, but without absorption at the low-frequency edge. We do not yet have a satisfactory explanation for the absorption at the low-frequency edge of the enhanced-emission bands.

We show in this paper that these previously unexplainable features, as well as a number of basic properties of the zebra structure and fibers, can be understood in the framework of the $l + w \longrightarrow t$ mechanism, but only if we take quasilinear effects into consideration—the scattering of fast electrons on the whistlers excited by them and on plasma waves [10, 11].

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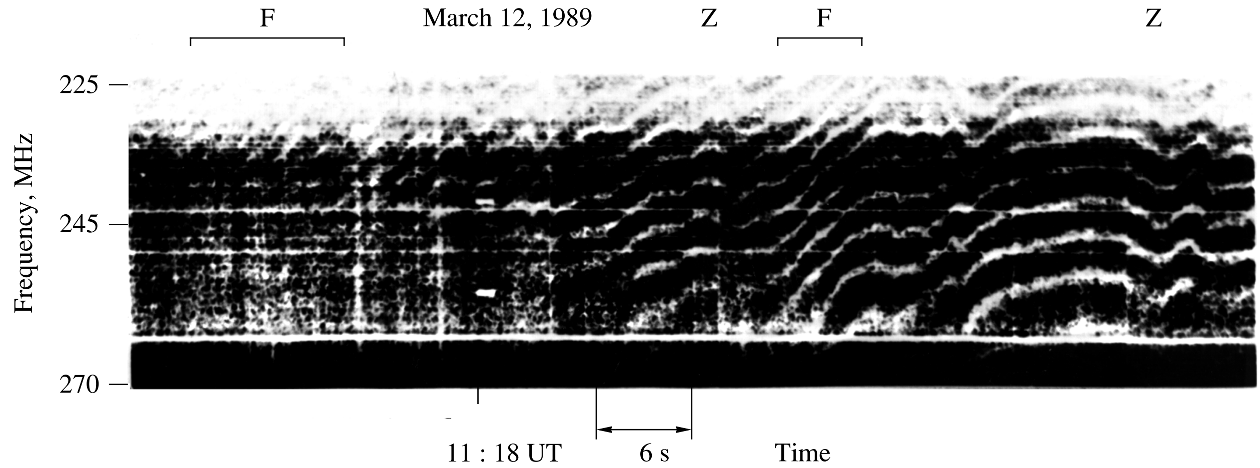

Fig. 1. Dynamical spectrum of zebra structure bands with wave-like frequency drift in the type IV burst of March 12, 1989. The "F" labels above the spectrum refer to the times when zebra structure bands (Z) with a constant drift toward low frequencies ($d_f/d_t < 0$) become similar to fiber bursts.

In models for the merging of plasma waves with whistlers for fibers with intermediate frequency drift (IDB fibers), quasilinear effects have been used to explain the periodic excitation of whistlers [2]. In some models, the passage of electrons into the loss cone during interaction with whistlers and the formation of an electron distribution with a "gap" have been employed to explain pulsations of the radio emission [12]. For the moderate and strong diffusion that is most likely to occur in solar magnetic traps, however, finer features of the dynamics of the loss-cone distribution function should also appear in the radio emission. The particles could be accelerated both by electric fields in the explosive phase of a flare and by shock waves at large heights in the corona. At the top of the trap (especially for small bottle ratios), a distribution of electrons can form in which the longitudinal velocities exceed the transverse velocities ($v_\parallel > v_\perp$) and the loss-cone angles are moderate. Such distributions excite whistlers at arbitrary angles to the magnetic field [13, 14]. The possibility of nonlongitudinal wave propagation considerably broadens the energy range of the electrons interacting with whistlers in a given frequency band; this is so because, in this case, diffusion takes place not only in the normal Doppler effect (conventionally used for longitudinal propagation), but also in the anomalous Doppler effect and Cherenkov resonance [15, 16]. Real hollow-beam distributions have been directly measured at the fronts of interplanetary shock waves [17].

2. FORMULATION OF THE PROBLEM

Numerous data have been accumulated from theoretical and numerical studies of wave–particle interaction, mainly aimed at explaining satellite and ground-based observations of magnetospheric whistlers [10, 11, 13, 15, 16, 18–21]. In fact, these results have not been applied to the interpretation of solar phenomena, although whistlers are widely used to explain the fine structure of solar radio bursts.

In this work, we explore the possibility of applying known theoretical and numerical results on wave-particle interactions to the conditions in the solar corona. The most complete numerical simulations, which included both whistlers and electrostatic (plasma) waves, were carried out by Omura and Matsumoto [11] using a distribution function in the form of a monoenergetic and warm beam with a loss cone. Such beams may also be typical for energetic particles captured in solar magnetic traps—sources of type IV bursts. Under the conditions of moderate or strong diffusion occurring in whistlers (typical for the solar corona), the lifetime of fast electrons is ~10 min, which is comparable to the duration of series of the radio emission fine structure [15]. Omura and Matsumoto [11] showed that diffusion in whistlers and plasma waves is separated in time. Initially, there is energy diffusion in plasma waves (hollowing of the beam); however, it is known that nonlinear effects can damp the plasma wave instability [22]. At the second stage, there is diffusion in whistlers, primarily in angle. It is important to note here that only a small fraction of the particles falls into the loss cone [11, 20], and the instability may be supported by a series of particle bounces between the ends of the magnetic traps.

After discussing general aspects of the quasilinear theory in Section 3, we chose a distribution function in the form of an anisotropic beam (8) and investigated its possible dynamics during diffusion in whistlers; this investigation is based on [19], the results of which are fully applicable to conditions in the solar corona. Below, this analysis, together with the numerical simulations of [11], the analytical results of [10], and the theory of fan instability in whistlers [16], allows us to qualitatively show that it is possible for a whistler instability to switch from a normal Doppler resonance to an anomalous resonance during diffusion. Thus, without solving a new quasilinear problem, we will attempt to estimate possible manifestations of these effects in the fine structure of solar radio bursts.

In Section 4, we apply the necessary numerical estimates to the conditions in the solar corona. It is usually assumed that the energy release in sources of type IV bursts is extended, which eases the problem of replenishment of the distribution function during diffusion by both bounces in the trap and periodic injection of new particles.

In Section 5, we show that only by taking into account the diffusion of fast particles in whistlers and plasma waves is it possible to explain naturally the most controversial observational fact—the formation of absorption in the continuum between bands of enhanced emission.

3. SOME ASPECTS OF THE QUASILINEAR THEORY

It is known that the spectra of naturally-excited whistler waves satisfy the validity condition of the quasilinear approximation, in which particles wander in a multiple-wave field. Essentially, the spectrum ($\Delta\omega^w$) must be rather broad and the relative amplitude of the wave must be modest [15, 18, 23]:

$$\Delta\omega^w \gg [(H^w/H_0)k\omega_{ce}v_\perp]^{1/2}, \tag{1}$$

where H^w/H_0 is the relative amplitude of the magnetic field of the wave, k is the wave number, ω_{ce} is the electron gyrofrequency, and $v_\perp$ is the electron velocity component transverse to the magnetic field. Condition (1) is better satisfied by solar whistlers than by magnetospheric whistlers. For example, for typical values of $H^w/H_0 \sim 10^{-4}$, $\omega^w/\omega_{ce} \sim 0.1$, and an approximate equality of the wave phase velocity ω^w/k and $v_\perp$, we obtain from (1) $\Delta\omega^w/\omega^w > 3\times10^{-2}$. For the inverse inequality to (1), the particle is completely trapped in the wave field [18]. Note that conical instability yields whistler spectra bounded above, $\Delta\omega^{\omega}/\omega^w < 0.1$–$0.2$, excluding significant group spreading of the whistlers over several seconds of propagation in the corona.

If a fast particle interacts with whistlers at the cyclotron resonance

$$\omega^w - k_\parallel v_\parallel - s\omega_{ce} = 0 \tag{2}$$

($k_\parallel$ and $v_\parallel$ are components of the wave vector and velocity parallel to the magnetic field), it moves along the diffusion curves

$$v_\perp^2 + (v_\parallel - \omega^w/k_\parallel)^2 = \text{const} \tag{3}$$

in the direction of decrease of the distribution function $F(v_\perp, v_\parallel)$ [19]. If the resulting flux of particles in velocity space is directed toward increasing particle energy, energy in the given range of velocities will be transferred from resonant waves to particles, and the wave will weaken. This physical process finds its theoretical reflection in the identity operator $\hat{\Lambda}$ [which has the meaning of a derivative along the curve (3)] in formulas for the whistler instability increment γ^w and the distribution function diffusion equation [19, 24]:

$$\hat{\Lambda} = (s\omega_{ce}/\omega v_\perp)(\partial/\partial v_\perp) + (k_\parallel/\omega)(\partial/\partial v_\parallel)|v_\parallel = (\omega - s\omega_{ce})/k_\parallel. \tag{4}$$

According to [15, 24], the general expressions for γ^w and $\partial F/\partial t$ have the form

$$\gamma^w = 8\pi^3(\omega/k_\parallel)(n^h/n^c) \times \sum_{s=0,\mp1}\int dv^3 G_{ks}\delta(\omega - k_\parallel v_\parallel - s\omega)\hat{\Lambda}F \tag{5}$$

$$\partial F/\partial t = v_\perp^{-1}\hat{\Lambda}\{v_\perp G_{ks}\hat{\Lambda}F\}, \tag{6}$$

where n^h/n^c is the relative number density of hot and cold particles and G_{ks} is a weighting function defining the emission contribution per particle per whistler wave. The sign of the increment depends on the sign of the operator $\hat{\Lambda}$. For example, in the normal Doppler resonance, when $s = +1$ in (2), the waves and particles are oppositely directed ($k_\parallel v_\parallel < 0$ or $\omega/k_\parallel < 0$), and $\hat{\Lambda}$ should be negative for positive values of the operator γ^w. Therefore, the dominant contribution to the emission comes from the part of the distribution where $\partial F/\partial v_\parallel < 0$ and $\partial F/\partial v_\perp > 0$. The contribution to the anomalous Doppler instability [$s = -1$ in (2)] comes from particles in the part of the distribution where the signs of these derivatives are opposite. Therefore, whether the predominant instabilities are in normal or anomalous resonances will depend on whether there is an excess of particles with $v_\perp$ or $v_\parallel$, respectively.

The operator $\hat{\Lambda}$ can be simply expressed in terms of the derivative of the distribution function with respect to the particle energy $\partial F/\partial E$ [19]:

$$\hat{\Lambda} \sim 2v_\perp v_p \partial F/\partial E \mid E \sim v_\perp^2 + v_\parallel^2;\ v_p = \omega/k_\parallel. \tag{7}$$

This relationship implies that, for positive increments, the derivative $\partial F/\partial E$ should be positive for motion in the direction of decreasing F. This in turn means that the diffusion curve (D) should lie between the constant-density curve (F) and the constant-energy circle (E) [19]. This relation and the physical meaning of the diffusion interaction are visually depicted by the graphical determination of the increment signs for different resonances presented in Fig. 2 for an anisotropic beam distribution function

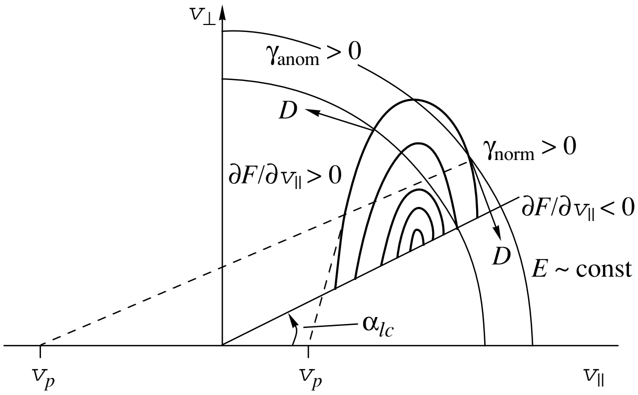

Fig. 2. An anisotropic-beam type (8) distribution function (F) in the ($v_\parallel$, $v_\perp$) plane with loss cone α_{lc}. Thin circular lines denote levels of constant particle energy E. Arrows labelled D indicate the direction of particle diffusion toward decreasing F. Thick lines are levels of equal density of F. v_p is the phase velocity, $\gamma_{anom} > 0$ is the domain of anomalous Doppler whistler instability, and $\gamma_{norm} > 0$ is the region of normal Doppler resonance instability.

$$F(v_\perp, v_\parallel) = C_n v_\perp^j \exp[-(m_e v_\perp^2/2kT_\perp) - (m_e(v_\parallel - v_\perp)^2/2kT_\perp)], \quad (8)$$

where C_n is a normalization constant and j is the loss cone index.

Quasilinear diffusion establishes a plateau along the diffusion line, i.e., $\hat{\Lambda} = 0$ (by analogy with the plateau state in the quasilinear theory of plasma oscillations [25]). Interaction in the normal resonance leads to the passage of particles into the loss cone [20, 26]. In the anomalous Doppler effect, the longitudinal velocities are hollowed out of a beam-type distribution (8), but the particles do not fall into the loss cone; instead, they diffuse in pitch angle toward high $v_\perp$ [16, 27].

During the lifetimes of particles in a trap in the moderate and strong diffusion whistler modes (several tens of seconds for solar magnetic traps [15, 28]), the loss cone is emptied twice during each bounce period. Therefore, after a short period of particle injection, the instability for both normal and anomalous resonance will be periodic, with a period (t_p) of the order of a second: for a trap length $l \sim 10^{10}$ cm and velocity $v \sim 10^{10}$ cm/s, we obtain $t_p \sim (1/2)(l/v) \sim 0.5$ s.

The abnormal-resonance instability favors the confinement of particles in the trap. This confinement probably explains the longer sequences of zebra structure caused by anomalous-resonance whistler instability compared to those of IDB fibers, due to normal-resonance whistlers.

The above effects allow us to estimate the importance of alternating switch-on of whistler instabilities due to the normal and anomalous Doppler effects, if the distribution function can support this dynamic transition. Two such distribution functions are an anisotropic high-energy beam (8), which can form in shock fronts [29], or a ring distribution [30]. The conical anisotropy of these distributions is determined by the acceleration mechanism—reflection from magnetic mirrors (shock fronts). Whistler instability has primarily been considered for longitudinal propagation in the normal Doppler effect [2, 31, 32]. It is known, however, that the contributions from all three resonances ($s = 0, \pm 1$) to the instability are comparable for oblique propagation [13]. Shapiro and Shevchenko [16] considered a turn of the beam in transverse velocity for the anomalous resonance, termed the fan instability. An initial distribution function in the form of a beam in longitudinal velocity turns in transverse velocity as a result of the excitation of oscillations and quasilinear diffusion; a new distribution function in the form of a beam in transverse velocity is established, if $\omega/k_\parallel < v_\parallel$. In our case, this last inequality is fulfilled even more strictly. In addition, this is one of the conditions for excitation of whistler waves in the anomalous resonance.

An initial distribution function in the form of a beam in longitudinal velocity is most probable under the conditions of multiple, repeated flare processes at large coronal heights during large-scale events. The unusual properties of the zebra structure (Fig. 1) can be understood in the framework of a mechanism of periodic whistler instability. The analog of fan instability considered above provides smooth (or rapid) switching of the instability from the anomalous to the normal Doppler resonance, which implies a smooth (or abrupt) change of the frequency drift of the zebra structure bands. The frequency drift is determined by the direction of the whistler group velocity. In the anomalous resonance, the directions of the waves and particles coincide, while in the normal resonance, they move in opposite directions. The character of the fan instability will depend on the specific parameters of the plasma and beam.

4. NUMERICAL ESTIMATES

The complete solution of the quasilinear problem for the three resonances with particle sources and losses is very difficult, and has not yet been obtained either analytically or numerically. Usually, a simplified problem is posed, and quasilinear equations are solved numerically for one resonance and one oscillation mode with a conical distribution function [21, 26]. In addition to the known effects of the passage of slower particles into the loss cone and the formation of a "gap" in the distribution, the calculations in [20, 21] also showed a simple shift of the distribution function maximum toward the initial loss cone for large velocities.

Periodic fan instability in a self-oscillatory mode was first considered in [33] in the framework of a theory of microwave spikes with millisecond periods; this analysis was based on instability switching from the Cherenkov to the normal resonance for tokamak plasma parameters, which are far from cosmic parameters. The condition for switch-on of the self-oscillatory mode obtained in [33] indicates that the increments for both resonances are comparable if there is sufficient beam temperature anisotropy and if the longitudinal velocity source power is moderate, so that particles cannot instantaneously leave the emission region. In our case, these conditions are fulfilled for a hollow-beam distribution function.

It is known that whistler diffusion occurs mainly in the pitch angle.

Transforming the expression for $\partial F/\partial t$ to a dependence of the distribution function F on pitch angle α, we obtain a diffusion equation in which the pitch angle diffusion coefficient D_α significantly exceeds the velocity diffusion coefficients when $v \gg \omega/k_\parallel$, which is satisfied in our case. According to [34], the coefficient D_α is proportional to the wave energy:

$$D_\alpha \sim (H^w/H_0)(\omega_{ce}/\cos\alpha). \quad (9)$$

For a reasonable wave amplitude relative to the background field $H^w/H_0 \sim 10^{-4}$ and $\omega_{ce} \sim 10^7$ s^{-1}, we obtain a simple estimate for the timescale for pitch angle diffusion by $\Delta\alpha \sim 1$: $t_D = D_\alpha^{-1} \sim 1$–2 s. These values are consistent with the timescales for sign alternation of the zebra structure band frequency drift in Fig. 1. Thus, we indeed observe the diffusion of fast electrons in whistlers in the radio emission fine structure.

The most complete numerical simulation of whistler diffusion of a beam with conical anisotropy in the normal Doppler resonance has been carried out by Omura and Matsumoto [11]. Figure 5 in [11] shows the initial values of the volume and planar distribution functions $F(v_\perp, v_\parallel)$, the linear increment γ, and the distribution height at the beam center. This figure also shows the hollowing of F in pitch angle. It is important to note that, when F is hollowed in $v_\perp$, a beam in the longitudinal velocities is formed in a time $t \sim 1300\omega_{ce}^{-1}$: $\partial F/\partial v_\parallel > 0$. Thus, the term $\partial F/\partial v_\parallel$ in the operator $\hat{\Lambda}$ in (4) remains positive, the operator $\hat{\Lambda}$ in (5) and (6) will be nonzero, and diffusion toward increasing $v_\perp$ will continue, but in the anomalous Doppler resonance. As can be seen from Fig. 5 in [11], the wave gradually weakens ($\hat{\Lambda}$ decreases with spreading of the beam in pitch angle). Since only a small fraction of the slowest particles falls into the loss cone, however, the resonance switching can repeat many times, as in the tokamak spike mode [33]. Particle replenishment will occur through bounce motions in the trap; if the particles arrive in phase with the growth of the longitudinal velocity bump (i.e., at the moment of alternation of the sign of the frequency drift), the instability will be sharply enhanced. We probably observe this effect when we detect short zebra-structure spikes in the dynamical spectra in some events, seen as herringbone structures over a broad frequency range.

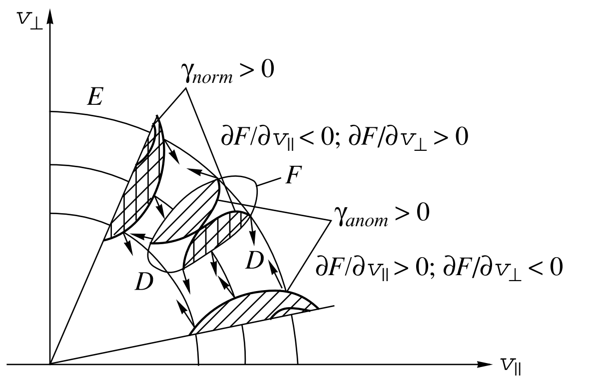

Fig. 3. Schematic presentation of the fan instability—switching of whistler instability from normal Doppler cyclotron resonance (cross-hatched F regions) to anomalous resonance (single-hatched regions) due to the shift of the maximum (bump) of the distribution function F during diffusion along the diffusion curves D (arrows) from large values of $v_\perp$ (where the operator $\hat{\Lambda} < 0$) to large $v_\parallel$ (where $\hat{\Lambda} > 0$).

The beam relaxation timescale $t_{\rm rel} \sim \gamma^{-1}$ in [11] is of the order of $3000\omega_{ce}^{-1}$. Correcting for coronal plasma and beam parameters (plasma frequency ratio $\omega_{pe}/\omega_{ce} \sim 20$, $n^h/n^c \sim 10^{-6}$), this timescale should be about 3×10^7 gyroperiods (in [11], $n^h/n^c \sim 10^{-2}$, $\omega_{pe}/\omega_{ce} \sim 2$). For a gyrofrequency of $\sim 5 \times 10^7$ Hz, we obtain $t_{\rm rel} \sim 0.8$ s, which is close to our estimate of the diffusion timescale $t_D \sim 1$ s.

An initial distribution in the form of a beam with conical anisotropy (8) will be unstable to both normal and anomalous resonances. The regions of positive-increment instability determined from Fig. 2 are shaded in Fig. 3. The cross-hatched areas refer to particles that are unstable to the normal resonance ($\gamma_{\rm norm} > 0$, where $\partial F/\partial v_\parallel < 0$, and $\partial F/\partial v_\perp > 0$), and the single-hatched areas refer to those unstable to the anomalous Doppler resonance ($\gamma_{\rm anom} > 0$ where $\partial F/\partial v_\parallel > 0$, and $\partial F/\partial v_\perp < 0$). According to the calculations of [11], the beam will smear in pitch angle. The gradual diffusive accumulation of particles with maximum F along $v_\parallel$ (later along $v_\perp$) will mean a gradual switching of the instability from anomalous to normal resonance. Figure 3 schematically shows the shapes of the distribution at small and large pitch angles for two such times. Pitch angles close to 90° cannot be reached when there is particle motion with conservation of the first adiabatic invariant. At a qualitative level, we can say that when there is a large particle number density in the beam, diffusion will be rapid and resonance switching (alternation of the sign of the frequency drift) will be abrupt.

5. FORMATION OF ABSORPTION BETWEEN EMITTING FILAMENTS

The origin of the absorption usually accompanying the fibers and zebra structure at the low-frequency edge is a very nontrivial question, often bypassed by many authors. For example, the double-plasma-resonance mechanism produces only emission bands [2, 4, 6]. In the whistler fiber model, it was assumed without proper detailed analysis that plasma waves within the whistler wave packet volume are efficiently evacuated in the process of merging with whistlers, $l + w \longrightarrow t$, and only residual l-waves scattered on thermal ions result in a decreased continuum level [2, 8].

A more concrete basis for this effect can only be associated with the evacuation of resonant l-waves with small wave numbers k^l in the process $1 + w \longrightarrow t$, which are achieved in the course of the differential scattering of l-waves. The maximum transformation to electromagnetic radiation takes place precisely at low k^l, since, according to Tsytovich [35], the rate of growth of the radio emission is

$$\partial W^t/\partial t = \alpha^{li} W^t[(W^l/k^l) + (\partial W^l/\partial k^l)], \qquad (10)$$

where W^t and W^l are the energy densities of electromagnetic and plasma waves and α^{li} is a transformation coefficient. For small k^l, the dominant term in the brackets is W^l/k^l. Small k^l are also resonant in the process $1 + w \longrightarrow t$, however. Turning to quantitative estimates, only $\sim 10^{-2}$–$10^{-3} W^l$ is transformed to electromagnetic waves in a source that is optically thick to the process $1 + w \longrightarrow t$. Detection of such weak absorption is at the resolution limit of spectrographs.

Another explanation of absorption is based on the assumption that the radio emission is screened by plasma-density enhancements created by the whistler wave packet [3, 37]. However, simple estimates show that the creation of a one-percent plasma density enhancement requires the improbably high relative density of whistler energy $W^w/n_eT_e \sim 10^{-3}$, while the values achieved by weak turbulence are five orders of magnitude lower [2, 38].

In relation to this, we will consider the role of quasilinear effects. First of all, it is natural to adopt a Maxwellian distribution function with conical anisotropy for the electrons responsible for type IV continuum emission [2, 6]. Since this distribution is unstable to whistlers, we must take into account the influence exerted by the diffusion of fast particles in whistlers to plasma wave excitation.

For large loss cone angles ($v_\perp > v_\parallel$), the main contribution to whistler instability is given by the normal Doppler resonance for longitudinal propagation ($k^w \parallel H_0$). In this case, plasma waves with the maximum increment are excited at double the plasma resonance frequency for quasitransverse propagation ($k^w \parallel H_0$) [22, 39]. This situation is characteristic for the bases of magnetic traps (for heights at decimeter emission wavelengths). Because of the diffusion of fast particles in whistlers, the distribution function maximum instantaneously shifts toward greater longitudinal velocities [11] and smaller pitch angles. As a result, the plasma-wave instability sharply weakens, because the l-wave excitation increment at the upper hybrid frequency is several times more sensitive to changes of the loss-cone angle α_{lc} than are whistlers. A comparison of Figs. 3 and 8 in [32] shows that the increments γ^w and γ^l are nearly equal for $\alpha_{lc} \approx 10°$, while for $\approx 2 \times 10^{-5}\omega_{ce}$, we have $\gamma^l \approx 5\gamma^w$. This decrease of γ^l with decreasing α_{lc} is the main cause of continuum suppression within the whistler wave packet volume. As in the first hypothesis, the zebra emission band will lie at a higher frequency $\omega^t = \omega^l + \omega^w$.

The calculations of [11, 20, 26] obviously show that a gap distribution with positive derivative $\partial F/\partial v_\parallel$ is formed in the course of diffusion in the normal Doppler resonance. Melrose [40] was the first to stress the importance of this effect in astrophysical plasmas. The minimum velocity at which the beam forms is determined by the minimum frequency in the whistler spectrum, in accordance with the resonance condition (2). For typical values $k_\parallel \sim \omega_{pe}/c$ (c is the velocity of light) and $\omega/\omega_{ce} \sim 0.1$, we obtain $v_\parallel \approx 43 v_A$ (v_A is the Alfvenic velocity). A beam with such velocities efficiently excites plasma waves along the field ($k^l \parallel H_0$) in the Cherenkov resonance. Therefore, the emission of plasma waves (and, consequently, of radio continuum) will be enhanced in regions adjacent to the whistler wave packet, amplifying the contrast of the continuum depression along the low-frequency edge of the fiber. This scheme offers a natural explanation for the continuum enhancement usually observed in a series of fibers and zebra structure [41].

If the initial whistler diffusion proceeds in the anomalous Doppler resonance in "oblique" whistlers (when $v_\parallel > v_\perp$) and the plasma waves have $k^l \parallel H_0$, the plasma-wave level will also be reduced because of the beam turn in transverse velocity. The plasma-wave instability at the upper hybrid frequency will gradually grow, also increasing the absorption contrast.

Thus, in any case, the reduction of the plasma-wave level due to electron diffusion in whistlers can explain the absorption at the low-frequency edge of a fiber. Another important observational fact is simultaneously explained: the small delay of the beginning of the emission band relative to that of the absorption in fibers (see the Fig. 2 in [37]).

According to the arguments above, the weakening of plasma waves should begin instantaneously, and the merging mechanism at the combined frequency $\omega^t = \omega^l + \omega^w$ requires, according to conservation laws, either oppositely directed or nearly aligned wave vectors $\bar{k}^l$ and $\bar{k}^w$ in order for decay at the difference frequency $\omega^t = \omega^l - \omega^w$ to take place [42]. Therefore, in any case, some time is needed for the isotropization of the wave vectors. While the plasma waves become isotropic nearly immediately, whistlers undergo comparatively slow scattering on thermal electrons, and especially on thermal ions [36]. According to estimates given in [43], this time should be ~0.2–0.3 s. Radio flux fluctuations in both emission and absorption on profiles along fibers occur on nearly the same timescales; these fluctuations could also be associated with the quasilinear diffusion of fast particles in whistlers, which proceeds at different rates at different levels in the trap during the lifetime of a fiber.

In a number of events, we have also observed fiber structure in type II bursts (see, e.g., Fig. 3 in [44]), associated with the propagation of shock waves through the solar corona. The main peculiarities of these fibers are the absence of any evident absorption at the low-frequency edge and their long duration (up to 40–50 s). Observations of interplanetary shock waves [45, 46] testify to the presence of plasma waves and whistlers in front of the shock. It is therefore natural to interpret the fiber structure of type II bursts as a manifestation of whistler propagation through clusters of preshock plasma waves. Plasma waves and whistlers can be generated only by electrons reflected from the shock and therefore possessing conical anisotropy. However, there cannot be trapped fast particles in front of an oblique, collisionless (in a general case) shock wave, and therefore the interaction of whistlers with fast particles cannot be extended. After each pulse injection of particles, the excited whistlers and fast electrons disperse in space. Judging from the negative frequency drift of fibers, whistlers propagate forward from the front (or remain nearly standing). Therefore, whistlers should be excited in the anomalous Doppler resonance, when the particles and waves propagate in the same direction, but at large angles. Thus, a beam of particles with conical anisotropy escapes along the magnetic lines of force, while whistlers propagate at a large angle to the field and, after a certain isotropization, interact with plasma waves. The lack of strict periodicity of fibers is also a consequence of the absence of quasilinear effects. Without the diffusion of fast particles in whistlers, there should not be any absorption at the low-frequency edge of fibers, which is confirmed by observations.

Thus, we have shown that it is natural and necessary to take into account the quasilinear effects of the diffusion of fast electrons with conical anisotropy in whistlers; this allows us to explain certain very important properties of fibers with intermediate frequency drifts (IDB or fiber bursts) and the zebra structure in type II and IV solar radio bursts. This analysis demonstrates the suitability of mechanisms for the merging of plasma waves with whistlers to explain the structures considered above.

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