INTERPRETATION OF A SPECIAL FINE STRUCTURE IN TYPE-IV SOLAR RADIO BURSTS

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(Received in revised form 8 November, 1988)

Abstract. A special fine structure (slowly drifting chains of narrowband fiber bursts), firstly observed during the solar type-IV radio burst on April 24, 1985, is interpreted as the radio signature of whistler waves periodically excited by a switch-on/switch-off process of a loss-cone instability in a localized wave packet of the fast magnetoacoustic mode.

1. Introduction

For the first time, Aurass *et al.* (1987) reported on a new fine structure observed during the solar type-IV radio burst on April 24, 1985. Figure 1 shows the spectrographic features of three samples of these fine structures. They are characterized by a slowly positively or negatively drifting chain of narrowband, always negatively drifting fiber bursts. The fibers are enclosed by an 'envelope' accompanied with an absorption ridge at the low-frequency side. The drift rate of the envelope is approximately 1 MHz s^{-1} similar to that of the type-II bursts (Krüger, 1979). The instantaneous bandwidth of the envelope is roughly 2 MHz. The fibers have typical drift rates of -5 MHz s^{-1} and instantaneous bandwidths of 0.5 MHz. These data represent typical features of fiber bursts in the frequency range around 200 MHz. Their small frequency extent of 2 MHz, however, is unusual in this spectral range (Kuijpers, 1975; Elgarøy, 1982; Bernold, 1983). The individual fibers in the envelope have approximately the same characteristics (drift rate, instantaneous bandwidth, frequency extent) and are connected with an absorption ridge at the low-frequency side, typical for fiber bursts (Kuijpers, 1975; Elgarøy, 1982; Bernold, 1983). The time distance between the fibers in the chain is roughly 0.4 s. The duration of these chains can be up to 7 s.

In this paper, we propose a model which is capable of explaining the appearance of this special fine structure. The principal scheme of this model is the following: during a flare process, high-energy electrons and protons are injected into a coronal loop, in which they are partly trapped. We assume the presence of a localized wave packet of the fast magnetosonic mode excited by energetic protons (Meerson, Sasorov, and Stepanov, 1978), which propagates at a finite angle to the loop magnetic field. This wave packet is accompanied with an enhanced magnetic field, and it is this local magnetic

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Solar Physics **120** (1989) 383–391.

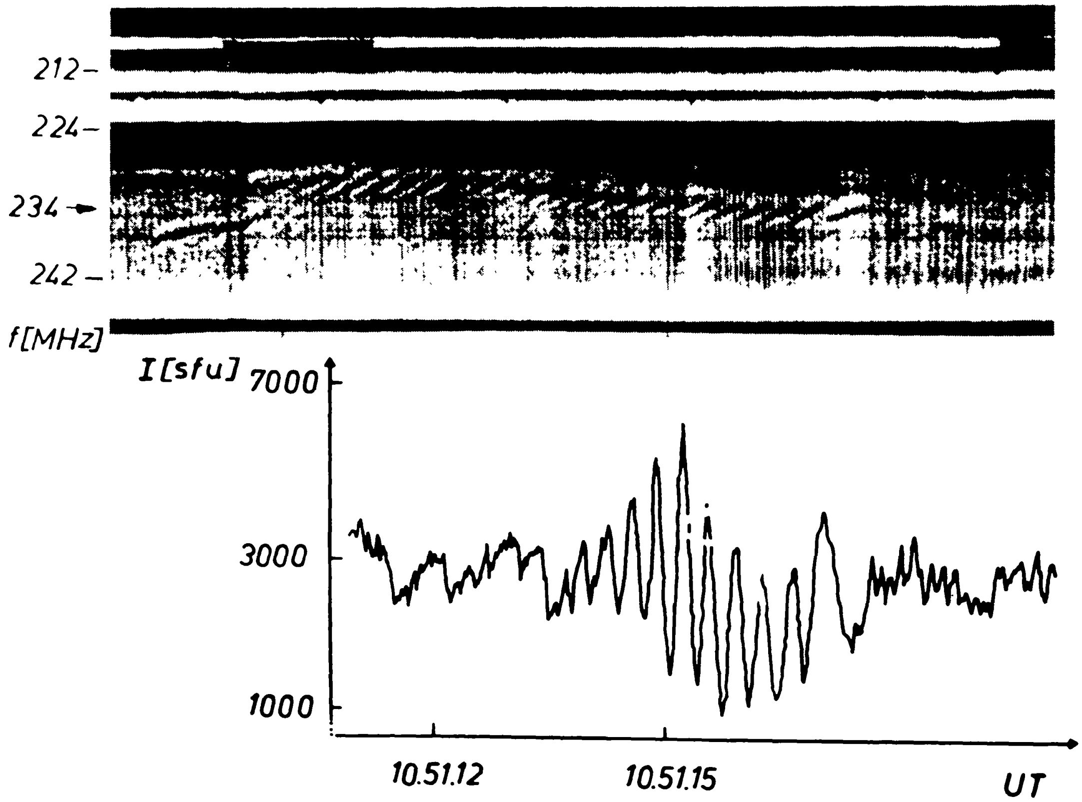

Fig. 1a.

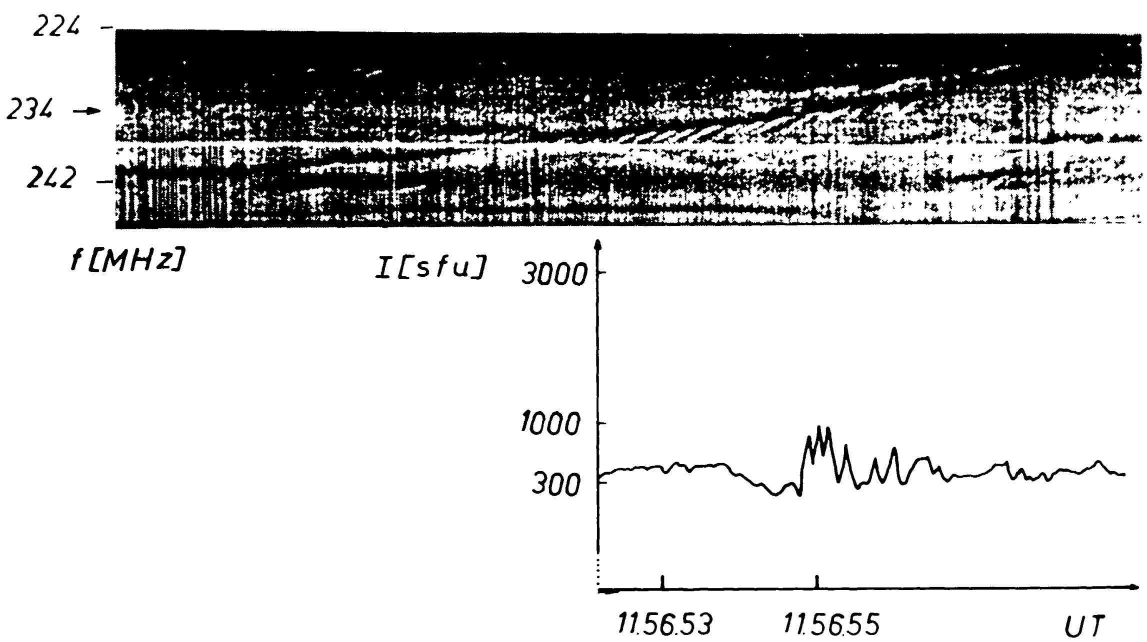

Fig. 1b.

Fig. 1. IZMIRAN spectrograms in the 200–247 MHz range and the corresponding radio flux registration at 234 MHz of the Tremsdorf Observatory. The spectrograms and the single-frequency records have identical time-scales. The slowly drifting chains of narrow band fiber bursts appear near (a) 10 : 51 : 10 UT, (b) 11 : 56 : 53 UT, and (c) 11 : 58 : 55 UT (from Aurass *et al.*, 1987).

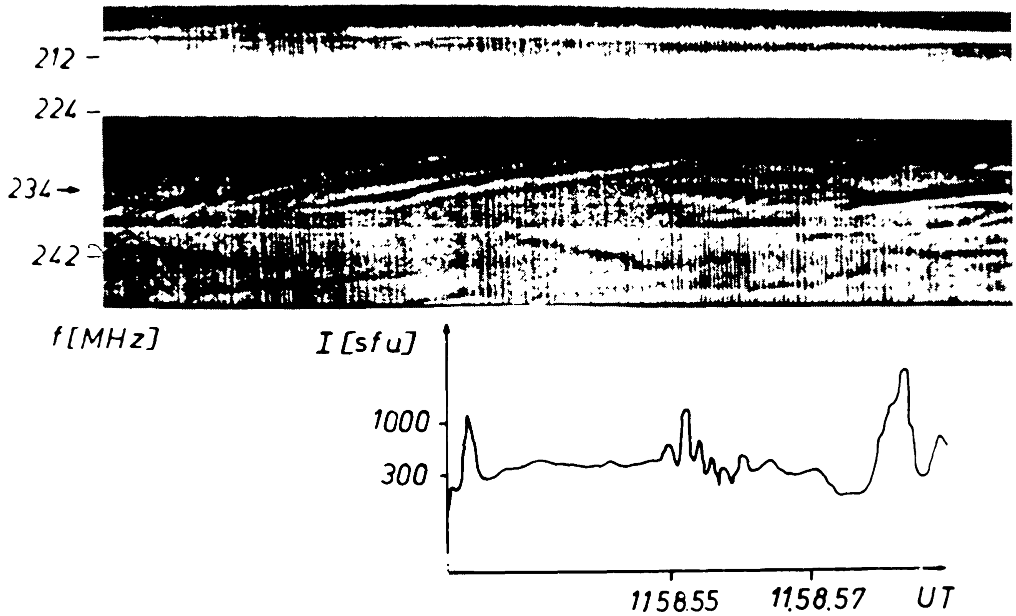

Fig. 1c.

field perturbation which is suggested to trigger loss-cone instability driven whistler waves via a local exceeding of the threshold of this instability.

Because the generated whistler waves exhibit a group velocity greater than the Alfvén speed for coronal conditions, they can leave the unstable region (i.e., the localized wave packet) and are damped outside. This explains the small frequency extent of the fibers generated by coalescence of the whistlers with Langmuir waves into escaping electromagnetic waves (Kuijpers, 1975). Furthermore, the wave packet is connected with a local density enhancement, i.e., the whistlers propagate towards decreasing densities. That always leads to a negative drift rate of the fibers, as observed. The slowly drifting envelope is interpreted as the radio signature of this localized wave packet. Because shock waves which induce type-II bursts propagate approximately with the Alfvén velocity in the solar corona, the drift rate of the envelope agrees with that of type-II bursts. If the wave packet travels in direction of an increasing/decreasing density, a positive/negative drift rate of the envelope is produced.

According to Kuijpers (1975) the recurrence of the fiber bursts in an individual envelope is explained by a periodic renewal of the loss-cone distribution owing to the bounce movement of the superthermal particles after the quenching of the instability by quasi-linear pitch angle diffusion of the fast particles into the loss-cone.

In Section 2, the stability of the fast magnetoacoustic wave packet is investigated. The properties of the instability induced by a highly energetic loss-cone distribution of electrons superimposed on a Maxwellian distribution is described in Section 3. In Section 4, these results are applied to explain the special fine structure discussed above. In addition, some estimates of plasma parameters in the source region are given.

2. Stability of the Fast Magnetoacoustic Wave Packet

The fast magnetoacoustic wave packet may be affected by dispersion, nonlinear wave breaking, and scattering at the boundary of the coronal loop. These processes are briefly discussed in this section.

The nonlinear behaviour of a fast magnetoacoustic wave propagating along the z-axis at an angle φ to the external magnetic field B_0 is described by a Korteweg–de Vries-type equation

$$\frac{\partial v_z}{\partial t} + \frac{3}{2} v_z \frac{\partial v_z}{\partial \xi} - \frac{1}{2} v_A d^2 \cot^2 \varphi \frac{\partial^3 v_z}{\partial \xi^3} = 0 , \tag{1}$$

with $\xi = z - v_A t$. Here, v_z, $v_A = B_0/(4\pi\rho_0)^{1/2}$ and $d = c/\omega_{pi}$ denote, respectively, the z-component of the velocity disturbance of the fast magnetoacoustic wave, the Alfvén velocity, and the so-called ion inertia thickness, respectively (c, velocity of light; ω_{pi}, ion plasma frequency; ρ_0, undisturbed mass density). Equation (1) was derived by Kakutani *et al.* (1968) from the magnetohydrodynamic equations supplemented by the generalized Ohm's law under the assumption $c_s \ll v_A$, typical for the solar corona (c_s, sound speed). The velocity disturbance v_z is accompanied with a density variation ρ_1 and a perturbation $B_{1\parallel}$ of the magnetic field parallel to B_0 via the relation

$$\rho_1/\rho_0 = v_z/v_A = B_{1\parallel}/B_0 . \tag{2}$$

If L is the characteristic length of the wave packet, dispersion can be neglected against nonlinear effects in Equation (1) for $(v_z/v_A) \gg (d/L)^2$. According to our model the frequency of the escaping electromagnetic wave is given by $\omega_0 = 2\pi f_0 = \omega_{pe} + \omega$ (ω_{pe}, electron plasma frequency; ω, whistler frequency). Because the wave packet is connected with a density enhancement ρ_1 it radiates at a somewhat higher frequency $\omega_0 = \omega_{pe}(1 + \rho_1/\rho_0)^{1/2} + \omega$. This causes an absorption ridge at the low-frequency side of the slowly drifting envelope, as observed (Bernold, 1983; Mann, Karlický, and Motschmann, 1987). Thus, ρ_1/ρ_0 can be estimated from the frequency separation Δf between the emission ridge and the absorption edge

$$\Delta f/f_0 \approx 0.5\rho_1/\rho_0 . \tag{3}$$

With $\Delta f = 2$ MHz and $f_0 = 234$ MHz we get $\rho_1/\rho_0 \approx 2 \times 10^{-2}$. Adopting a coronal density height scale of 10^{10} cm, the length L of the wave packet can be estimated from the envelope bandwidth (2 MHz) to be $L = 2 \times 10^8$ cm. On the other hand, we find $d = 9 \times 10^2$ cm at the 234 MHz-plasma level which is much less than L. Thus, dispersion can be neglected against nonlinear effects for the case considered.

Neglecting the dispersion term in Equation (1) the solution has the form (Landau and Lifshitz, 1978)

$$v_z = F(\xi - 3/2 v_z t) . \tag{4}$$

The wave-breaking time t_b is given by

$$t_b = -(\tfrac{2}{3})\,(\partial F^*/\partial v_z)_{|\,v_z = v_z^*}\,, \tag{5}$$

with F^* being the inverse function of F. v_z^* is found from the relation $(\partial^2 F^*/\partial v_z^2) = 0$ (Landau and Lifshitz, 1978). Taking an initial Gaussian shape $v_z = v_{\max} \exp(-4\xi^2/L^2)$ the wave-breaking time t_b is found to be

$$t_b = L(2e)^{1/2}/6v_{\max}\,. \tag{6}$$

This relation gives $t_b = 80$ s for the above-mentioned values of L, $v_{\max}$, and for $v_A = 5 \times 10^7$ cm s^{-1} (see Section 4).

Thus, the deformation by dispersion and nonlinear effects of the wave packet is negligible for times fairly less than 80 s. However, it crosses the coronal loop of a typical radius of $R = 1.4 \times 10^8$ cm (Krüger, 1979) in a time of about $t = 2R/v_A \sin\varphi \approx 8$ s. Thus, the stability of the wave packet is expected to be destroyed after this time.

3. The Loss-Cone Instability

During the flare process energetic particles can be injected in a coronal loop where a part of them is trapped. Thus, a loss-cone distribution of these particles is established. Such a distribution superimposed on a Maxwellian distribution can excite whistler waves (Scharer, 1967; Ossakov, Ott, and Haber, 1972; Kuijpers, 1975; Berney and Benz, 1978).

In the present paper we investigate a loss-cone distribution of energetic electrons superimposed upon a Maxwellian distribution of a proton-electron plasma. The distribution functions are

$$f_e(v) = f_{e0}(v) + f_{lc}(v)\,, \tag{7.1}$$

$$f_i(v) = f_{i0}(v)\,, \tag{7.2}$$

with

$$f_{j0}(v) = n_0/(2\pi v_{j0}^2)^{3/2} \exp(-v^2/2v_{j0}^2) \tag{7.3}$$

($j = i$ protons, $j = e$ electrons) and

$$f_{lc}(v) = (n_{lc}/(((2\pi v_{lc}^2)^{3/2} \cos\alpha))\theta(v_\perp - v_\parallel \tan\alpha) \exp(-v^2/2v_{lc}^2) \tag{7.4}$$

($v^2 = v_\parallel^2 + v_\perp^2$). The distribution functions are normalized to the particle number densities. n_0, n_{lc} are the particle number densities of the Maxwellian background plasma and of the energetic particles, respectively; $v_{j0} = (k_B T/m_j)^{1/2}$ is the thermal velocity of the particle species j (standard notations) and v_{lc} is the mean velocity of the energetic electrons. α denotes the loss-cone angle and θ is the well-known θ-function. $v_\parallel$ and $v_\perp$ are the components of the particle velocities parallel and perpendicular to the external magnetic field B_0. Here, the external field is directed along the z-axis. We assume that the Maxwellian populations of electrons and protons have the same temperature, i.e., $v_{i0} = (m_e/m_i)^{1/2} v_{e0}$.

The dispersion relation of right-hand circularly polarized electromagnetic waves (field vector rotates in the sense of the electron gyration) propagating along the magnetic field is given by (see, for example, Krall and Trivelpiece, 1973)

$$\omega^2 = k^2c^2 + \pi \sum_j \frac{4\pi e^2}{m_j} \int_{-\infty}^{+\infty} \mathrm{d}v_\parallel \int_0^\infty \mathrm{d}v_\perp v_\perp^2 \frac{(\omega - kv_\parallel) \dfrac{\partial f_j}{\partial v_\perp} + kv_\perp \dfrac{\partial f_j}{\partial v_\parallel}}{kv_\parallel - \omega + \varepsilon_j \omega_{cj}} . \tag{8}$$

Here, $\omega_{cj} = |e| B_0/m_j c$ is the cyclotron frequency of the species j, $\varepsilon_e = 1$ and $\varepsilon_i = -1$. Inserting the distribution (7) into the expression (8) the integrals can be expressed in terms of the plasma dispersion function Z (Fried and Conte, 1961).

In order to discuss the properties of the instability the dispersion relation (8) is numerically solved for plasma parameters appropriate for the solar corona at the 234 MHz level ($n = 7 \times 10^8\ \mathrm{cm}^{-3}$, $T = 2 \times 10^6$ K). Furthermore, $v_{lc}/c = 0.33$ is assumed as being typical for energetic electrons in the solar corona (Krüger, 1979). $\omega_{pe}/\omega_{ce} = 15$ and $n_{lc}/n_0 = 10^{-2}$ are parameters which fit the values estimated from the observational data (see Section 4). The left-hand side of Figure 2 displays the region of

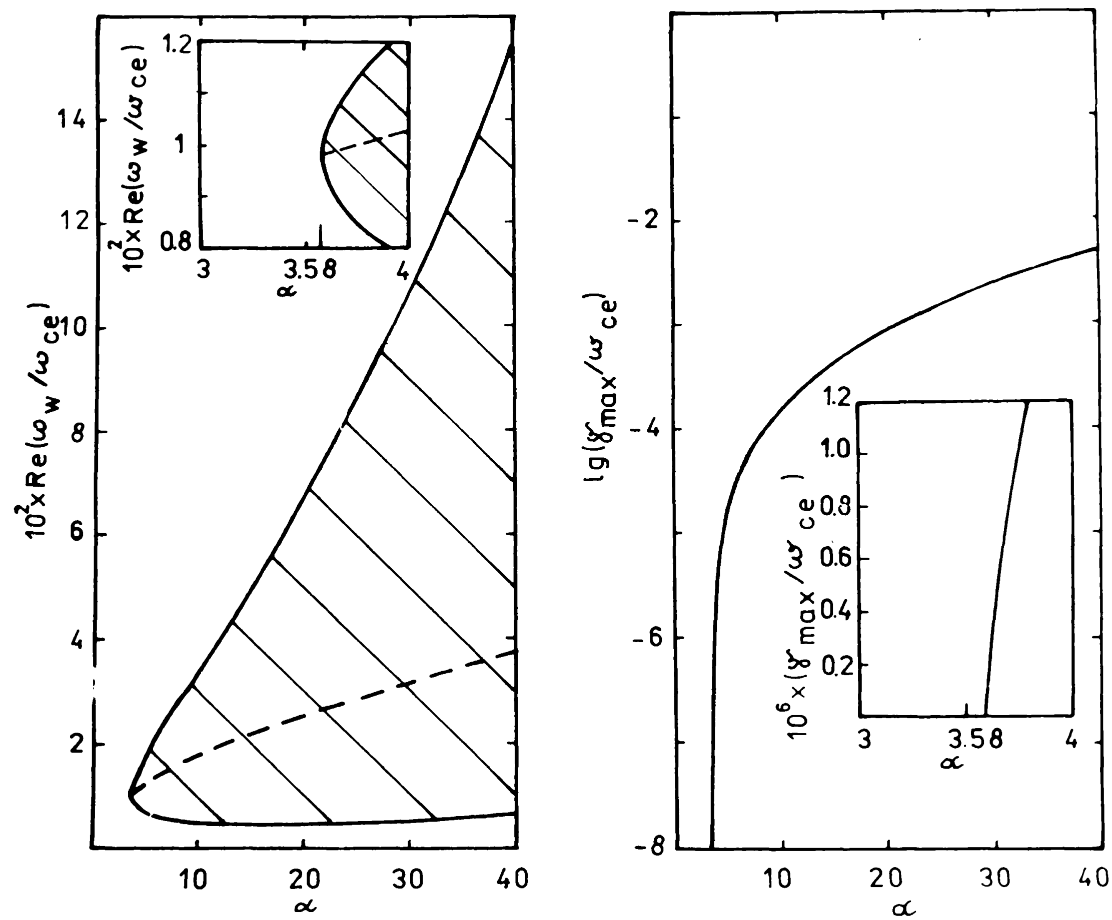

Fig. 2. *Left*: Domain of whistler loss-cone instability in the $\omega - \alpha$ plane (hatched). The dashed line indicates the location of maximum growth rate. *Right*: Maximum growth rate as function of the loss-cone angle.

instability in the $\omega - \alpha$ plane. An instability only appears for loss-cone angles exceeding a critical value determined to be 3.58°. We emphasize that this threshold behaviour of the loss-cone instability is crucial for our model (see Section 1). The threshold angle proves to be insensitive to the variation of all relevant parameters, except the temperature of the loss-cone electrons. With decreasing temperature of the energetic electrons, the threshold angle is shifted to higher values (e.g., 8 for $v_{lc}/c = 0.05$), accompanied by both increasing threshold frequency and maximum growth rate. It should be emphasized, however, that the loss-cone whistler threshold behaviour only appears when the proton population is taken into consideration. In contrast to the thermal electrons which are responsible for the high-frequency damping of whistlers near the electron-cyclotron frequency, the thermal protons give rise to damping at low frequencies ($\omega \gtrsim \omega_{ci}$). Thus, it is not surprising that the ions contribute to the dispersion in the frequency region considered (some tens of the ion-cyclotron frequency).

4. Estimate of Plasma Parameters

In this section some plasma parameters are deduced from the observational data on the basis of the presented model.

Assuming radio emission near the local electron plasma frequency from a localized emitter moving with the velocity V through the inhomogeneous corona, the drift rate is defined by

$$D = \mathrm{d}f/\mathrm{d}t = (fV/2n)\,(\mathrm{d}n/\mathrm{d}s) \tag{9}$$

(n, particle number density; f, emitted radio frequency). d/ds represents the spatial derivative along the magnetic field. According to our model, the slowly drifting envelope is assumed to be produced by the fast magnetosonic wave packet which propagates with local Alfvén velocity v_{A} at an angle φ to the magnetic field. The whistlers, which are suggested to be responsible for the fibers, move with their group velocity $v_G = 2c(\omega_{ce}/\omega_{pe})\,(x(1-x)^3)^{1/2}$, $x = \mathrm{Re}\,\omega/\omega_{ce}$, along the magnetic field. Thus, the ratio of the drift rates of fibers and the envelope is

$$D_{\mathrm{fiber}}/D_{\mathrm{env.}} = (v_G/v_{\mathrm{A}})\cos\varphi = 2(m_i/m_e)^{1/2}\,(x(1-x)^3)^{1/2}\cos\varphi\,. \tag{10}$$

Because the whistlers are preferentially excited near the frequency which maximizes the growth rate, i.e., at $x = 0.01$, an angle $\varphi = 54°$ is obtained from Equation (10) inserting the observed drift rates. Then, relation (2) gives a relative magnetic field enhancement $B_{1\,\parallel}/B_0 = 2 \times 10^{-2}$ within the fast magnetosonic wave packet.

Adopting a coronal density height scale of 10^{10} cm the group velocity of whistlers is estimated to be $v_G = 4.3 \times 10^8$ *cm s*$^{-1}$ ($M_G = v_G/c_s = 34$ with a coronal temperature of $T = 2 \times 10^6$ K). Then, $\omega_{pe}/\omega_{ce} = 15$ is reproduced using $x = 0.01$. Thus, we are able to calculate the Alfvén velocity, the basic magnetic field and the plasma-beta to be $v_{\mathrm{A}} = 5 \times 10^7$ cm s^{-1}, $B_0 = 6$ G and $\beta = (8\pi n k_B T)/B_0^2 = 0.13$, respectively, in the source region. Such values are typical for the solar corona (Dulk and McLean, 1978).

The local loss-cone angle is defined by

$$\alpha(s) = \arcsin((B(s)/B_{\max})^{1/2}) . \quad (11)$$

Here, $B(s)$ and $B_{\max}$ denote the magnetic field strengths at the height s in a coronal loop and the magnetic field strength, where all particles with $v_\perp \neq 0$ are reflected, respectively. Due to the magnetic disturbance $B_{1\parallel}$ associated with the fast magnetoacoustic wave packet, the loss-cone angle is locally shifted to a value

$$\alpha = \arcsin((\sin\alpha_0)(1 + B_{1\parallel}/B_0)^{1/2}) \quad (12)$$

with α_0 as the undisturbed loss-cone angle. This shift is sufficient for the whistler loss-cone instability threshold to be exceeded. Assuming $\alpha_0 = 3.56°$ and $B_{1\parallel}/B_0 = 2 \times 10^{-2}$ we get a shifted angle $\alpha = 3.60°$. From Figure 2(a) it is seen that indeed this shift changes the whistler wave propagation to become unstable. The corresponding maximum growth rate is $\gamma_{\max}/\omega_{ce} = 3 \times 10^{-7}$ at $x = 0.01$, i.e., $\gamma_{\max}/\omega = 3 \times 10^5$ or $\gamma_{\max} = 30\ \mathrm{s}^{-1}$, where the estimated electron-ion collision frequency is fairly less than this value.

With $\alpha = 3.6°$ and $B_0 = 6$ G we get $B_{\max} = 1520$ G. This might appear to be a too high value for the low corona. Note, however, that a maximum magnetic field strength of 3000 G was measured in the active region connected with the event on April 24, 1985 at the photospheric level during this event (Borovik, Watrushin, and Korjavin, 1987).

The following estimates should finally demonstrate that even a very low maximum growth rate as derived above is able to produce quickly enough the whistler wave energy density level indicated by the measurements.

Owing to the action of the ponderomotive force (Karpman and Washimi, 1977) the generated whistler waves induce a relative density enhancement

$$\rho_2/\rho_0 = \varepsilon W^w/(M_G^2 - 1) \quad (13)$$

with $W^w = |E_w|^2/16\pi n k_B T$ ($M_G = v_G/c$, E_w electric field of the whistler wave) and $\varepsilon = (\omega_{pe}/\omega_{ce})^2/x(1 - x)$. Thus, the whistlers radiate (via coalescence with Langmuir waves) at a somewhat higher frequency $\omega_0 = 2\pi f_0 = \omega_{pe}(1 + \rho_2/\rho_0)^{1/2} + \omega$ forming an absorption edge at the low-frequency side of the emission stripe of the fiber bursts (Bernold, 1983; Mann, Karlický, and Motschmann, 1987). The ponderomotively induced nonlinear frequency shift of the whistlers is several orders of magnitude smaller than that due to influence of the density enhancement. Using Equation (3) we find $W^w/W^w_{\mathrm{therm}} = 1.5 \times 10^4$ for $\Delta f = 0.5$ MHz and $f_0 = 234$ MHz from the above estimated parameters. Employing the above-mentioned value of the maximum growth rate, this whistler wave level is reached in a time of 0.16 s according to $W^w/W^w_{\mathrm{therm}} = \exp(2\gamma_{\max} t)$, where the thermal level of the whistler waves is assumed to be $W^w_{\mathrm{therm}} \approx (2n\lambda_0^3)^{-1}$ with $\lambda_0 = v_{e0}/\omega_{pe}$ (Krall and Trivelpiece, 1973). This is consistent with the observations because the individual fibers appear in a time up to 0.2 s (see Figure 1).

5. Summary

Slowly drifting chains of narrow band fiber bursts are suggested to reflect the radio signature of whistler waves excited within a localized wave packet of the fast magnetosonic mode. A loss-cone distribution of energetic electrons superimposed upon a Maxwellian distribution, unstable only for loss-cone angles exceeding a critical value, is assumed to be the driver of the whistlers. The magnetic field variation associated with this wave packet shifts the local loss-cone angle beyond the critical value and triggers the excitation of the whistlers finally coalescing with Langmuir waves into escaping electromagnetic waves. The high-energetic particles are scattered into the loss-cone by quasi-linear pitch angle diffusion, which suppresses the whistler excitation. The bounce movement of the fast electrons leads to a renewal of the loss-cone distribution giving rise to a new whistler generation. This periodic whistler excitations form the chain of the narrow band fiber bursts.

The plasma parameters derived on the basis of this model are consistent with the generally accepted picture of the solar corona (Krüger, 1979; Dulk and McLean, 1978).

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