GENERATION OF INTERMEDIATE DRIFT BURSTS
IN SOLAR TYPE IV RADIO CONTINUA THROUGH COUPLING
OF WHISTLERS AND LANGMUIR WAVES

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Abstract. The possible generation of intermediate drift bursts in type IV dm continua through
coupling between whistler waves, traveling along the magnetic field, and Langmuir waves, excited
by a loss-cone instability in the source region, is elaborated. We investigate the generation, propaga-
tion and coupling of whistlers.

It is shown that the superposition of an isotropic background plasma of $10^6$K and a loss-cone
distribution of fast electrons is unstable for whistler waves if the loss-cone aperture $2a$ is sufficiently
large ($\sec a \gtrsim 4$); a typical value of the excited frequencies is $0.1 \omega_{ce}$ ($\omega_{ce}$ is the angular
electron cyclotron frequency).

The whistlers can travel upwards through the source region of the continuum along the magnetic
field direction with velocities of $21.5-28 v_A$ ($v_A$ is the Alfvén velocity).

Coupling of the whistlers with Langmuir waves into escaping electromagnetic waves can lead to
the observed intermediate drift bursts, if the Langmuir waves have phase velocities around the velocity
of light.

In our model the instantaneous bandwidth of the fibers corresponds to a frequency of $0.1-0.5 \omega_{ce}$
and leads to estimates of the magnetic field strength in the source region. These estimates are in
good agreement with those derived from the observed drift rate, corresponding to $21.5-28 v_A$, if we
use a simple hydrostatic density model.

1. Introduction

Intermediate drift bursts or 'fiber' bursts (see Figure 1) have been observed in solar
type IV dm continua by Young et al. (1961), Slottje (1972a) and Elgarøy (1973); their drift rates are intermediate between type III and type II drift rates. Typically,
a fiber burst consists of two adjoining depression and enhancement ridges relative
to the surrounding continuum with the emission feature on the high-frequency side.
Both the drift rate and the ridge separation led to the suggestion (Kuijpers, 1973) that
fiber bursts are generated by whistler wave packets traveling through the source

![Figure 1](https://example.com/figure1.png)

Fig. 1. Example of intermediate drift bursts observed with the 60-channel Utrecht spectrograph
on March 6, 1972. The figure shows the flux variations (<3 s with respect to a floating zero level;
range $\pm 1.7$ dB). Each channel has a width of 0.9 MHz.

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region and coupling with the locally excited Langmuir waves, which presumably cause the observed continuum emission (Kuijpers, 1974).

In Section 2 the observations are reviewed. Section 3 deals with the nature and the propagation properties of whistler waves. The excitation of whistlers is the subject of Section 4. The possible coalescence of whistlers with Langmuir waves into transverse waves and the efficiency of this process are determined in Section 5. Finally, in Section 6 we argue that the model might apply to the situation in a type IV source region. In this paper we neglect the influence of collisions.

We assume that the source originates from fast electrons, injected into a stationary magnetic arch configuration and that the fast particles are, at each point, distributed outside a local loss-cone thus giving rise to an instability for Langmuir waves. Subsequently the transverse waves are produced by induced scattering on the thermal ions (Kuijpers, 1974).

2. Observations

The observational characteristics of intermediate drift bursts are summarized in Table I.

<table>
<thead>
<tr>
<th>TABLE I</th>
<th>Observed characteristics of intermediate drift bursts</th>
</tr>
</thead>
<tbody>
<tr>
<td>Frequency range (MHz)</td>
<td>950–500(^\text{a})</td>
</tr>
<tr>
<td>Instantaneous band-width (MHz)</td>
<td>10</td>
</tr>
<tr>
<td>Single frequency duration (s)</td>
<td>0.2–0.6</td>
</tr>
<tr>
<td>Frequency drift (MHz s(^{-1}))</td>
<td>(-10) to (\pm 50) ((-150))(^\text{d})</td>
</tr>
<tr>
<td></td>
<td>(+19) to (+25)</td>
</tr>
<tr>
<td>Frequency extent of single fibers (MHz)</td>
<td>((20))(^\text{a}) (50–150) ((300))(^\text{d})</td>
</tr>
</tbody>
</table>

\(^{a}\) Young et al. (1961).
\(^{b}\) Slottje (1972a, b); here the listed values refer to emission and absorption ridges separately.
\(^{c}\) Elgarøy (1973).
\(^{d}\) The values between brackets are rarely observed.

(1) In most cases the fibers appear as fine structures within a continuum burst which lasts 5–50 min, although the fibers are sometimes observed as emission ridges before or after the continuum (Young et al., 1961). Often the fibers cluster in time with (sometimes nearly constant) intervals of a few seconds or much less between successive fibers and are very similar within one group.

(2) The absorption edge is typically on the low frequency side; according to Young et al. (1961) and Elgarøy (1973) the absorption edge is at the high frequency side in some cases, but with the Utrecht 60-channel spectrograph no convincing example thereof has been recorded (Slottje, 1974).

(3) The absorption and emission ridges begin and end simultaneously (Young et al., 1961; Slottje, 1972b).

(4) The majority of fibers have a negative frequency drift rate. Generally the absolute value of the drift rate of an individual fiber decreases with decreasing fre-
frequency. Also the drift rates and instantaneous band-widths of fibers in different frequency bands tend to decrease with decreasing frequency; for the single frequency duration this is less clear (cf. Table I). From the observed drift rates Young et al. (1961) and Elgarøy (1973) deduced velocities for the exciting agent less than the electron thermal velocity under assumption that the emission originates at the local plasma frequency.

(5) The emission ridges and the surrounding continuum are strongly circularly polarized and in the same sense (Slottje, 1974).

3. Propagation Characteristics of Whistlers*

3.1. Whistler Waves in a Cold Plasma

Whistlers are transverse waves with frequencies below the electron cyclotron frequency. In a cold plasma and for frequencies well above the ion cyclotron frequency the dispersion relation for whistler waves propagating parallel to the magnetic field is (Kennel and Petschek, 1966):

\[ \frac{c^2k^2}{\omega^2} \approx 1 + \frac{\omega_{pe}^2}{\omega(\omega_{ce} - \omega)}, \]

where \( \omega \) is the angular wave frequency, \( k \) the wave number, \( \omega_{pe} \) the angular electron plasma frequency and \( \omega_{ce} \) the (positive valued) electron cyclotron frequency (see Figure 2). Since in the corona \( \omega_{pe}/\omega_{ce} \gg 1 \), the phase velocity \( v_p \) can be written as

\[ v_p = v_{Ae}\left[\frac{\omega}{\omega_{ce}}\right]^{1/2} \ll c, \]

where \( \omega/\omega_{ce}, v_{Ae} = B/(4\pi n_e m_e)^{1/2} \) is the electron Alfvén velocity, which is typically a velocity intermediate between the usual Alfvén velocity and the velocity of light. Consequently, it follows from Maxwell's law, that the wave character is dominantly magnetic

\[ E_{k,\omega} = \frac{v_p}{c} B_{k,\omega}, \]

where \( E_k \) and \( B_k \) are, respectively, the electric and magnetic wave amplitudes, perpendicular to the wave vector. (We use the Gaussian system of units in this paper).

3.2. Interaction with Particles in a Finite Temperature Plasma

For parallel propagation the whistler is purely transverse and circularly polarized in the same sense as a gyrating electron. Therefore, the wave can interact with a particle of given kind, if in the reference frame of the gyration centre of the particle the wave frequency equals the corresponding cyclotron frequency and if in the same frame the wave has the correct polarization (cyclotron resonance). Thus, for electrons the resonance velocity \( v_\parallel \) parallel to the ambient field direction is determined by

\[ \omega - k_\parallel v_\parallel = \omega_{ce}, \]

* See also the appendix for the meaning of the symbols.
Fig. 2. Relation of the real part of the frequency with the wave number for whistlers in an isotropic Maxwellian plasma with $T = 10^6$ K, $\omega_{pe}/2\pi = 300$ MHz and $\omega_{pe}/\omega_{ce} = 30$ (curve 1), resp. 10 (curve 2). The curves end at $-\gamma/\omega_{ce} \approx 0.1$. The dashed curve indicates the cold dispersion relation. The values of $k$ are in cm$^{-1}$.

and for ions

$$\omega - k_v v_\parallel = - \omega_{ci},$$

(4)

where the velocity component $v_\parallel = (v \cdot B)/B$ and $k_v = (k \cdot B)/B$.

Since, according to Equations (3) and (4) the values of the particle resonant energies at frequencies above 0.1 $\omega_{ce}$ are much smaller for electrons than for ions, the damping in a Maxwellian plasma will be determined mainly by the electrons. In this case the dispersion relation for parallel propagation is (Scharer and Trivelpiece, 1967)

$$\frac{c^2 k^2}{\omega^2} = 1 + \frac{\omega_{pe}^2}{\omega k_v \sqrt{2 \omega_{ce}}} Z_t,$$

(5)

where $v_{te}^2 = K_B T_e/m_e$ and $Z_t = Z(\psi)$ is the plasma dispersion function (Fried and Conte, 1961) with argument $\psi = (\omega - \omega_{ce})/(kv e \sqrt{2})$. For large values of the real part of the argument $\psi$, relation (5) reduces to the cold plasma dispersion relation (1). In general, a solution of Equation (5) with real $k$ is complex: $\omega = \omega_r + i\gamma$, with $\gamma < 0$ corresponding to damping. If the damping is sufficiently weak, that is if $|\gamma| \ll \omega_r$ and $|\gamma| \ll \omega_{ce} - \omega_r$ and $|\omega, \text{Re} Z| \gg |\gamma \text{ Im} Z|$, one can easily solve relation (5) for $\omega_r$ and $\gamma$. 
Fig. 3. Cyclotron damping of parallel whistlers in an isotropic Maxwellian plasma with $T=10^6\,\text{K}$, $\omega_{pe}/2\pi=300\,\text{MHz}$ and $\omega_{pe}/\omega_{ce}=30$ (curve 1), resp. 10 (curve 2). The effect of an extra isotropic hot component ($n_h=10^{-3}, \sqrt{2} v_h=10^{10}\,\text{cm}\,\text{s}^{-1}$) is shown by the dashed curves.

From Figure 3 it can be seen, that under coronal circumstances ($T_e=10^6\,\text{K}, \omega_{pe}/2\pi=300\,\text{MHz}, \omega_{pe}/\omega_{ce}=30-10$) and in the absence of a loss-cone distribution the role of the cyclotron damping is so important that the propagating whistlers will not occur above $0.5\,\omega_{ce}$. On the other hand Figure 2 shows that the relation between the real part of the frequency and the wave number is hardly affected in the accessible frequency regime ($-\gamma/\omega_{ce}<0.1$).

3.3. **Oblique wave vectors**

In the case of propagation at an angle $\theta$ to the magnetic field direction the 'cold' dispersion relation is (Kennel and Scarf, 1969)

$$v_p = v_{Ac} \left[ x (\cos \theta - x) \right]^{1/2}. \quad (6)$$

From this relation one can derive the well-known property of whistlers that the direction of the group velocity $((d\omega/dk)=\hat{k}(\partial\omega/\partial k)+(\hat{\theta}/k)(\partial\omega/\partial \theta); \hat{\theta}$ denotes a unit vector) is confined to a small cone with aperture $2\alpha_{max}$ around the field direction at the lower frequencies. (See Table II; cf. Stix, 1962).

<table>
<thead>
<tr>
<th>$x$</th>
<th>$\alpha_{max}^a$</th>
<th>$\alpha_{max}^b$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>19.5°</td>
<td>0°</td>
</tr>
<tr>
<td>0.5</td>
<td>0°</td>
<td>30°</td>
</tr>
</tbody>
</table>

$^a$ local maximum for lower $k$-values

$^b$ value for $k \to \infty (\cos \theta \to x)$.  

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Since in the oblique case whistlers have an electric field component parallel to the magnetic field direction (Allis et al., 1963), such waves can also interact with particles through Cherenkov-resonance, provided that $\omega - k v = m\omega_{ce}$. They are therefore subject to Landau damping if the velocity distribution is monotonically decreasing when integrated over the velocity components perpendicular to the magnetic field (see for a numerical expression p. 21 in Kaplan and Tsytovich (1973)). Moreover, now higher harmonic cyclotron interactions are possible at

$$\omega - k v = m\omega_{ce}$$

for any entire positive number $m$, apart from the already mentioned fundamental cyclotron damping at $m=1$ (cf. Equation (3)). Therefore in practice the occurring wave vector directions will certainly fall well within the cone given by $\cos^2 \theta = x$ and the alignment of the group velocity with the magnetic field direction will be even better than suggested by Table II.

Due to refraction the whistlers change their direction with respect to the magnetic field. Therefore their damping rate changes. Whistlers refract towards the magnetic field direction during their passage through the inhomogeneous corona if the phase velocity at a given frequency or (from Equation (6)) the local Alfvén velocity decreases along their path. On the other hand, in the reverse situation, a whistler essentially refracts away from the field direction over a distance $\Delta r = \pi (2 \sin \theta d \ln v_A/dr)^{-1}$ and will be Landau damped.

3.4. Propagation velocity

The value of the group velocity of parallel whistlers in the allowed frequency range ($x < 0.5$; cf. Section 3.2) follows from Equation (6) as

$$v_g = 2v_A [x(1-x)^3]^{1/2}.$$  

The maximum group velocity is $0.65 v_A$ at $x = 0.25$. At frequencies in between 0.1–0.5 $\omega_{ce}$ the group velocity takes a value in between 0.5–0.65 $v_A$ (21.5–28 $v_A$).

In the absence of dissipation isolated whistler wave packets can occur in the form of solitons if the nonlinear effects are balanced by dispersion (Kakutani et al., 1967). Typical frequencies of the involved whistlers are 0.25–0.5 $\omega_{ce}$ (Tidman and Krall, 1971). The propagation velocity of whistler solitons moving along the magnetic field direction is 21.5 $v_A$ if the plasma is cold and if the wave magnetic field amplitude is sufficiently small (Kakutani, 1966). In our case the maximum wave field amplitude does not exceed the field strength of the background magnetic field (see Section 6.2) and the latter condition is fulfilled. Since an initial disturbance of a system described by the Korteweg-de Vries equation leads to the formation of solitons (Davidson, 1972), we expect whistler solitons to evolve in our case if the whistler waves are generated during a short time (cf. Section 4.2) and at sufficiently large wave frequencies (cf. Section 4.1). However, henceforth we will not use any property of whistler solitons since they are not essential for the explanation of the radio observations.
Summarizing briefly the last section, we have found that whistlers can propagate slightly damped under coronal circumstances at frequencies below 0.5 $\omega_{ce}$. The group velocities are directed around the field within a small cone. Whistlers focus towards the magnetic field direction if the Alfvén velocity decreases along the path but refract away from it in the reverse situation. Finally, their group velocity is in between 21.5–28 $v_A$ at frequencies in between 0.1–0.5 $\omega_{ce}$.

4. Excitation

4.1. Linear Instability

As has been shown by Brice (1964) the energy change of a particle upon emission or absorption of whistler waves satisfies

$$\frac{\Delta W_\perp}{\Delta W} = 1 - \frac{v_\parallel}{v_p} = \pm \frac{\omega_{ce_i}}{\omega},$$

(9)

where $\Delta W$ is the change in total particle kinetic energy and $\Delta W_\perp = mv_\perp \Delta v_\perp$ is the change in transverse energy only. For the last part of Equation (9) we have used $k_\parallel = k$, Equation (3) for the electrons and Equation (4) for the ions. From Equation (9) it follows that the transverse energy of an emitting electron ($\Delta W < 0$) decreases. Consequently, anisotropic electron distributions with an excess of average transverse energy in comparison with the average parallel energy are expected to be unstable for the generation of whistler waves through cyclotron resonance. We therefore investigate the eventual generation of whistlers from a loss-cone type distribution of fast electrons superposed on the thermal coronal plasma.

For simplicity we make the approximation that the fast electron distribution $f_h$ is Maxwellian outside the loss-cone with half-aperture $\alpha$

$$f_h(v) = n_h (2\pi v_e^2)^{-3/2} \exp\left(-v^2/2v_h^2\right) \quad \text{for} \quad \left|\arctan(v_\perp/v_{\parallel})\right| > \alpha$$

(10)

and $f_h(v) = 0$ otherwise.

It should be noted that this distribution function is normalized to $n_h \cos \alpha$, where $n_h$ is the fractional fast particle density for a zero degree loss-cone. Then the dispersion relation for parallel propagating whistlers is (Scharer, 1967; Ossakow et al., 1972)

$$\frac{c^2k^2}{\omega^2} = 1 + \frac{\omega_{pe}^2}{\omega k v_e \sqrt{2}} Z_h + n_h \omega_{pe} \left[ \frac{\sec \alpha}{\omega k v_h \sqrt{2}} + \frac{\tan^2 \alpha}{\omega^2} \left\{ 1 + \psi_h^2 (1 + \psi_h Z_h^{-1}) \right\} \right].$$

(11)

Here $Z_h \equiv Z(\psi_h)$ with $\psi_h = (\omega - \omega_{ce})/(kv_h \sqrt{2} \cos \alpha)$, cf. Equation (5), and $n_h$ is assumed to be much smaller than unity. From Equation (11) the growth rate $\gamma$ is derived for the region where $|\gamma| \ll \omega_r$ and $|\gamma| \ll |\omega_{ce} - \omega_r|$ as in Section 3.2 and is

$$\gamma = \frac{-\text{Im} Z_h - n_h \text{Im} Z_h}{v_e} + \frac{n_h \tan^2 \alpha}{v_h} \psi_h \left( \frac{\omega_{ce}}{\omega_r} - 1 \right) \text{Im} Z_h,$$

$$\frac{\omega_r}{v_e} = \frac{2 \sqrt{3} k \omega_r}{v_h} + \frac{n_h \text{Re} Z_h}{v_h} + \frac{2 n_h \tan^2 \alpha}{v_h} \psi_h + \frac{3 n_h \tan^2 \alpha \psi_h^2}{v_h} \text{Re} Z_h.$$

(12)
In deriving Equation (12) we used the cold dispersion relation Equation (1) for the real part of the frequency. Examples of the growth (damping) rate are shown in Figure 4 as a function of frequency for a number of different cases. Their global behaviour is readily understood from the following simple reasoning. With a characteristic velocity of the emitting particles $v_{\parallel} = v_h \cos \alpha$, one can derive a relation between the energy of the 'hot' particles, the loss-cone angle $\alpha$ and the unstable frequencies from Equations (2) and (3)

$$v_h \cos \alpha = v_{\parallel} = c \frac{\omega_{ce} (1 - \chi)}{\omega_{pe} x^{1/2}}.$$  

(13)

First, since the interaction at the higher frequencies ($\chi \to 1$) takes place with particles of small parallel velocities (Equation (13)), where the distribution is dominated by the (stable) background Maxwellian, the whistlers are expected to be unstable below a certain frequency in agreement with Figure 4. Secondly, from Equation (13) we expect that the growth rate becomes larger and the unstable frequencies shift to higher values when the loss cone widens, in agreement with the behaviour of curves 1 to 6 in Figure 4. Finally, when the ratio $\omega_{pe}/\omega_{ce}$ decreases, Equation (13) indicates that the instability shifts to higher frequencies relative to the cyclotron frequency, again in agreement with Figure 4 (curves 5 and 7).

We conclude, that the configuration of a closed magnetic arc, where fast particles have loss-cone distributions ($n_h \approx 10^{-3}$) in an otherwise isotropic background plasma, can lead to the generation of whistlers. The value of $\sec \alpha$ varies much more along the source tube than the ratio $\omega_{pe}/\omega_{ce}$ or the background density. Therefore the change in growth rates at different heights is mainly due to the change in degree of anisotropy. From Figure 4 we infer that the excitation of whistler waves is concentrated in the lower parts where the fast particle distribution is most anisotropic. With neglect of collisions typical values of the loss-cone in the whistler unstable region correspond to $\sec \alpha \approx 4$ and the growing waves have frequencies around 0.1 $\omega_{ce}$. This situation is in contrast with the magnetospheric case where whistlers are excited predominantly near the top of a magnetic arch (Kennel and Petschek, 1966). The cause of this difference is that in the magnetosphere the background density and the ratio $\omega_{pe}/\omega_{ce}$ vary by three orders and, respectively, one order of magnitude over the source. As a consequence the fraction of emitting particles near the footpoints of an arch is, even for moderate values of the loss-cone apertures, minute in comparison with the fraction at the top. The resulting effect on the growth rates at different heights counts more heavily than the opposite effect of the change in anisotropy along the arch.

When the whistlers travel upwards into regions with decreasing field strength, their frequency rises relative to the local cyclotron frequency. Therefore, eventually their energy is transferred to the thermal background plasma when they enter regions with small loss-cone angles where the growth rate becomes negative. We expect this damping to be more important than the extra damping caused by refraction (cf. Section 3.3) since the magnetic field strength varies more rapidly over a certain height than the Alfvén velocity.
Fig. 4. The vertical axis indicates the whistler growth rate and the horizontal axis the wave frequency, both in units of the electron cyclotron frequency. The numbered curves correspond to the parameter values $n_h = 10^{-3}$, $\sqrt{2} \, v_{ce} = 5.5 \times 10^8 \, \text{cm s}^{-1}$, $\sqrt{2} \, v_h = 10^{10} \, \text{cm s}^{-1}$ and

<table>
<thead>
<tr>
<th>curve</th>
<th>$\sec \alpha$</th>
<th>$\omega_{pe}/\omega_{ce}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>1.01</td>
<td>30</td>
</tr>
<tr>
<td>2</td>
<td>1.1</td>
<td>30</td>
</tr>
<tr>
<td>3</td>
<td>2</td>
<td>30</td>
</tr>
<tr>
<td>4</td>
<td>4</td>
<td>30</td>
</tr>
<tr>
<td>5</td>
<td>8</td>
<td>30</td>
</tr>
<tr>
<td>6</td>
<td>15</td>
<td>30</td>
</tr>
<tr>
<td>7</td>
<td>8</td>
<td>10</td>
</tr>
</tbody>
</table>

The dashed curves present part of the damping rates (absolute value of $\gamma/\omega_{ce}$); the dashed curve at the left is common to cases 1–6, the other corresponds to case 7.
4.2. QUASILINEAR DEVELOPMENT

We assume that non-linear wave-wave interactions are not of importance for the actual development of the instability in comparison with the effect of the excited whistlers on the fast particles themselves (quasilinear diffusion). Then the main effect for these low-frequency waves is scattering of the particles into the loss-cone with a characteristic isotropization time \( \tau_{qi} \) (Kennel and Petschek, 1966)

\[
\tau_{qi} \approx \frac{(\Delta \alpha)^2}{D_\alpha} \approx \frac{\pi^2 v B^2}{32 \omega_{ce}^2 B_k^2},
\]

where \( D_\alpha \) is the pitch angle diffusion coefficient and \( B_k^2/8\pi \) the wave magnetic energy density per unit wave number. We have taken \( \Delta \alpha \) equal to \( \pi/4 \) and used a pitch angle of \( \pi/3 \) for the original particle with velocity \( v \). Now we expect that after the onset of the instability the waves continue to grow until their energy density is so high that the quasilinear relaxation time \( \tau_{qi} \) for the electrons becomes comparable with the growth time of the waves \( (2\gamma)^{-1} \). For the values \( \gamma = 10^{-3} \omega_{ce} \), \( v = 10^{10} \text{ cm s}^{-1} \), \( \omega_{ce}/2\pi = 10 \text{ MHz} \) we find a value of 0.05 erg cm\(^{-2} \) as an upper limit for the maximum wave energy density per unit wave number. With the above growth rate taken to be constant the wave will reach this level from its thermal equilibrium level given by \( B_k^2/8\pi \approx k^2 K_B T/2\pi^2 \approx 1.8 \times 10^{-14} \text{ erg cm}^{-2} \) \( (k \approx 0.05 \text{ cm}^{-1} \); cf. Figure 2) after a time

\[
\tau_{\max} \approx 29/(2\gamma) \approx 2.3 \times 10^{-4} \text{ s}.
\]

When a whistler of fixed frequency travels upwards, the ratio \( \omega/\omega_{ce} \) increases and, provided the ratio \( \omega_{ce}/\omega_{pe} \) does not increase, the parallel energy of the resonant electrons decreases (Equation (13)). On the other hand, the parallel energy of a particle that resonates with the wave at a given altitude increases when it travels upwards due to the adiabatic invariance of the flux through the particle’s orbit (Jackson, 1962). Consequently, the wave and the particle can interact only over a very limited distance. We consider the simplified picture in which both the wave growth and the scattering of particles are confined to one and the same region over which the loss-cone parameter \( \sec \alpha \) varies from infinity to \( \sqrt{10} \) (cf. Figure 4). Within this distance \( L \) the magnetic field strength varies \( 10\% \) and we take \( L = 10^3 \text{ km} \). The passage time of a whistler through the unstable region is

\[
\tau_{\text{pass}} = L/v_G \approx 0.2 \text{ s},
\]

where we have used \( v_G = 21.5 \text{ v}_\Lambda \) (cf. Section 3.4) and \( \omega_{pe}/\omega_{ce} = 30 \).

Finally we estimate the bounce time of the fast particles in the magnetic loop to be

\[
\tau_{\text{bounce}} \approx 1 \text{ s}.
\]

Then since \( \tau_{\max} \ll \tau_{\text{pass}} \ll \tau_{\text{bounce}} \) the instability evolves in the following way. First the whistlers reach their maximum amplitudes and scatter the fast particles in the unstable region into the loss-cones. Subsequently these particles escape within 0.01 s towards
the lower atmosphere where they lose their energy. The excited whistlers travel upwards and leave the unstable region after a time $\tau_{\text{pass}} \approx 0.2$ s. Meanwhile the unstable region is refilled by particles from above. However, as long as a major part of the whistlers is present in the unstable region, these particles are scattered immediately into the loss-cones and enhance the wave amplitudes slightly. Only after the disappearance of the whistlers from the scattering region renewed generation of whistlers can take place. This will happen if $\tau_{\text{pass}} \ll \tau_{\text{bounce}}$.

We therefore expect the whistlers to be generated in a periodic manner with time intervals exceeding $\tau_{\text{pass}} \approx 0.2$ s.

5. Coupling

5.1. Admissibility

We are interested in a three-wave coupling process: a whistler wave couples with a Langmuir wave to produce a transverse wave above the plasma frequency. Both in the random phase approximation and in the case of fixed phases, the necessary conditions for coupling are (Tsytovich, 1970)

$$\omega_w + \omega_l = \omega_t$$

and

$$k_w(\omega_w) + k_l(\omega_l) = k_t(\omega_t)$$

which can be considered as conservation of the total energy and momentum of the coalescing plasmons. Here the index $w$ indicates the whistler, $l$ the Langmuir wave, $t$ the escaping transverse wave and $\omega_j$ is the frequency related to the wave vector $k_j$ through the linear dispersion relation for wave mode $j$.

To solve Equation (18) for a range of relevant values of $\omega_t$, $\omega_w$ and suitable directions of the different wave vectors, we proceed along the line described by Höijer and Wilhelmsson (1970). Using the notation $\omega_+ \equiv \omega_w + \omega_l$, $k_+ \equiv k_w + k_l$, and adopting a certain value for $k_l$ and $\theta_w$ ($\theta_j$ is the angle of the wave vector of kind $j$ with the magnetic field direction), we examine graphically whether the curve formed by the pairs $(k_+, \omega_+)$ intersects the appropriate dispersion curves $k_j(\omega_+)$ for a sufficiently large interval of $\theta$, so that we can be sure that $\theta_+$ is equal to $\theta$ at the point of intersection for some $\theta$. We use the whistler dispersion curve for parallel propagation as drawn in Figure 2. To obtain the dispersion curves for the Langmuir waves and for the two escaping transverse modes we solve the well-known dispersion equation derived in the continuum approximation and for infinitely heavy ions (Denisse and Delcroix, 1961)

$$\begin{align*}
(1 - a^2 n^2 - X) \left\{ (1 - n^2) (1 + Y_L) - X \right\} &= (1 - n^2) (1 - Y_L) - X \right\} = \\
Y_T^2 (1 - n^2) (1 - n^2 - X)
\end{align*}$$

(19)

for wave numbers at given frequencies and a given direction $\theta$ with respect to the magnetic field (see Figure 5). Here $n = ck/\omega$ is the index of refraction, $Y = \omega_{ce}/\omega$, $X = \omega_{pe}/\omega^2$, $Y_L = Y \cos \theta$, $Y_T = Y \sin \theta$ and $a^2 = 3 v_{te}^2/c^2$. For the coronal conditions

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T = 10^6 K, ω_{pe}/ω_{ce} = 30, the dispersion curves so derived are plotted in Figure 5 for various angles φ. Since for ω_{pe}/ω_{ce} = 30 the whistlers are strongly damped by the thermal plasma above ω/ω_{ce} = 0.37 (cf. Section 3.2), the maximum allowable wave number is 0.056 cm^{-1}. Since moreover k_+ > k_1 - k_w, it can be seen in Figure 5 that coupling can only take place if the longitudinal wave number is small enough, that is if the frequency of the Langmuir wave is very close to the plasma frequency. For the values k_1 ≈ 0.06 and φ_1 = 0°, 10° and 20° the three curves k_+ (ω_+) for various whistler frequencies up to ω/ω_{ce} = 0.37 are drawn in Figure 5. Because the whistler frequencies cannot exceed a value of 0.5 ω_{ce} under coronal conditions, clearly only the ordinary electromagnetic wave can be produced and this only then if the angle between longitudinal and whistler waves is rather small (≤10°) (see Figure 5).

The amount of energy radiated in a particular mode equals the work j_k^{(2)}. E_k f done by the electric field of the outgoing electromagnetic wave E_k f on the nonlinearly (by
coalescing whistler and longitudinal wave) generated current, \(j_k^{(2)}\). Considering only the response of the electrons in a cold and magnetized plasma, we arrive, after Fourier transformation and expansion of the fluid variables in the wave electric fields in a way analogous to Tsytovich (1970) for an isotropic plasma, at the following result for the nonlinear current:}

\[
\begin{aligned}
j_k^{(2)} &= - \frac{\omega_{pe}^2 e}{8\pi m_e} \int \frac{d^3k_1 \ d\omega_1 \ d^3k_2 \ d\omega_2}{\omega_1 \omega_2} \delta (\mathbf{k} - \mathbf{k}_1 - \mathbf{k}_2) \\
&\quad \times \left[ \frac{1}{\omega_2} (\mathbf{I} + \mathbf{A}_1) : \mathbf{E}_{k_1} \cdot (\mathbf{k}_2 \cdot (\mathbf{I} + \mathbf{A}_2) : \mathbf{E}_{k_2}) + \\
&\quad + \frac{1}{2\omega} (\mathbf{I} + \mathbf{A}) : \mathbf{k} (\mathbf{E}_{k_1} \cdot \mathbf{E}_{k_2}) + \frac{1}{\omega} (\mathbf{I} + \mathbf{A}) : \mathbf{k}_1 (\mathbf{E}_{k_1} \cdot \mathbf{A}_2 : \mathbf{E}_{k_2}) + \\
&\quad + \frac{1}{\omega} \mathbf{B}_1 : \mathbf{E}_{k_1} ( (\mathbf{I} + \mathbf{A}_2) : \mathbf{E}_{k_2} \cdot \mathbf{k}_1 ) + 1 \equiv 2 \right].
\end{aligned}
\tag{20}
\]

Here the indices 1 and 2 refer to the coalescing modes while the outgoing wave is indicated without indices; the index \(k_j\) is an abbreviation of \((\mathbf{k}_j, \omega_j)\), \(\mathbf{I}\) is the identity and the operators \(\mathbf{A}\) and \(\mathbf{B}\) take the following form on the basis \(((1/\sqrt{2}) (\hat{x} + i\hat{y}), (1/\sqrt{2}) (\hat{x} - i\hat{y}), \hat{z})\) with \(\hat{z}\) parallel to the constant external magnetic field:

\[
\mathbf{A}_j = \begin{pmatrix}
(x_j - 1)^{-1} & 0 & 0 \\
0 & -(x_j + 1)^{-1} & 0 \\
0 & 0 & 0
\end{pmatrix},
\]

\[
\mathbf{B}_j = \begin{pmatrix}
x (x - 1)^{-1} (x_j - 1)^{-1} & 0 & 0 \\
0 & -x (x + 1)^{-1} (x_j + 1)^{-1} & 0 \\
0 & 0 & 0
\end{pmatrix},
\]

where as before \(x = \omega / \omega_c\). If the external magnetic field vanishes, the operators \(\mathbf{A}\) and \(\mathbf{B}\) vanish also, and Equation (20) reduces to Equation (2.41) in Tsytovich (1970).

In the case of parallel propagating whistlers \((\mathbf{k}_w = (0, 0, k_w)\) and \(\mathbf{E}_{k_w} = (0, E_{k_w}, 0))\) and parallel Langmuir waves \((\mathbf{k}_l = (0, 0, k_l)\) and \(\mathbf{E}_{k_l} = (0, 0, E_{k_l}))\) no transverse waves of the ordinary mode \((\mathbf{k}_i = (0, 0, k_i)\) and \(\mathbf{E}_{k_i} = (E_{k_i}, 0, 0)\), that is, with circular polarization opposite to that of the whistler) are generated, as follows from inspection of the explicit form of \((\mathbf{j}_k^{(2)} \cdot \mathbf{E}_{k})\) (using Equation (20)). At the same time, it follows that waves in the ordinary mode can be produced from parallel whistlers and oblique Langmuir waves, provided the conservation relations (18) are satisfied.

Combining the results of the previous paragraphs we conclude that the coalescence of parallel whistlers with Langmuir waves can only lead to radiation in the ordinary mode and in directions non-parallel to the magnetic field.

5.2. Efficiency

To estimate the coupling efficiency of a Langmuir wave and a whistler with the appropriate wave vectors (cf. previous section) into an electromagnetic wave, we use the growth relation derived for isotropic spectra of the coalescing modes and averaged
over the angles of the outgoing waves (Kaplan and Tsytovich, 1973, p. 284),
\[
\frac{\partial W_{lt}^i}{\partial t} = \sqrt{2} \frac{\omega_{pe}^{9/2}}{\omega_{ce}^{3/2} n_e m_e c^4} \frac{W_{lt}^i W_{kw}^w}{k} \approx 3.1 \times 10^7 W_{lt}^i W_{kw}^w \text{erg cm}^{-2} \text{s}^{-1},
\]
(21)
with \( k_w \approx k_l \equiv k \). \( W_k^j \) is the total wave energy density per unit wave number for mode \( j \). We have used in the right hand part of the equality \( \omega_{pe}/2\pi = 300 \text{ MHz}, \ \omega_{pe}/\omega_{ce} = 30 \) and \( k = \omega_{pe}/c \).

In the present case we have to compare Equation (21) with the production rate of electromagnetic waves from induced scattering of the Langmuir waves on the thermal ions (Kuijpers, 1974). This scattering process is described by Tsytovich (1970, Equation (8.61)):
\[
\frac{\partial W_{kt}^l}{\partial t} \sim W_{kt}^l W^l.
\]
(22)
Owing to the explicit dependence on \( W_{kt}^l \) of the right hand part the total conversion efficiency depends sensitively on the detailed geometry of the source region. Therefore a comparison with Equation (21) becomes rather arbitrary. On the other hand the amount of radiation produced according to Equation (21) can be compared directly with the observed flux values. This will be done in Section 6.2.

If the continuum radiation is the result of induced scattering of the Langmuir waves on particles, harmonic generation by coalescing Langmuir waves is of minor importance in comparison with it and therefore also with the whistler coupling. The angle-averaged generation rate of first harmonic radiation from isotropic longitudinal waves with phase velocities smaller than the velocity of light is (Kaplan and Tsytovich, 1973, p. 284)
\[
\frac{\partial W_{kt}^l}{\partial t} = \frac{4 \sqrt{3} \pi}{5} \frac{\omega_{pe}^4}{n_e m_e c^5} \frac{1}{\Delta k_l} \int \left( \frac{W_{lt}^i}{k_l} \right)^2 dk_l.
\]
(23)
Comparing Equation (21) with Equation (23) the ratio between the growth rates is found to be of the order
\[
\frac{1}{3} \left( \frac{\omega_{pe}}{\omega_{ce}} \right)^{3/2} \frac{W_{kw}^w}{W_{kt}^l}
\]
(24)
for \( \Delta k_l = k_l \), \( \Delta k_i = \omega_{pe}/c \) and on the assumption that the Langmuir waves have the same wave numbers in both reactions. Apparently, the coalescence of Langmuir waves with whistler waves is more efficient than the mutual coalescence of Langmuir waves into electromagnetic waves under coronal conditions where \( \omega_{pe}/\omega_{ce} \) is typically of order 30, if \( W_{kw}^w/W_{kt}^l > 0.5 \omega_{ce}/\omega_{pe} \).

6. Application to the Type IV dm Source Region

6.1. Langmuir Wave Spectrum

As we have stressed earlier (Section 5.1), coupling of longitudinal waves with whistlers
can only lead to an observable structure in dynamic radiospectra if a substantial fraction of the Langmuir waves is concentrated around small wave numbers (i.e. high phase velocities). Since the Langmuir waves are excited by fast particles, initially their phase velocities are smaller than the velocity of light. But we know that in a Maxwellian plasma induced scattering of longitudinal waves into longitudinal waves leads to a degradation of the Langmuir wave spectrum towards lower frequencies and higher phase velocities (Kaplan and Tsytovich, 1973). Moreover, in a magnetic field this induced scattering results in a 'condensation' of the Langmuir waves parallel to the magnetic field (Kaplan and Tsytovich, 1973) due to the angle dependence of the approximate dispersion relation for the weak field case

$$\omega^2 = \omega_{pe}^2 + 3k^2v_{te}^2 + \omega_{ce}^2 \sin^2 \theta.$$  \hspace{1cm} (25)

This alignment of the waves along the field direction occurs for phase velocities exceeding $\sqrt{3} v_{te} \omega_{pe}/\omega_{ce} \approx 2 \times 10^{10} \text{ cm s}^{-1}$, and is therefore important in our case where the coupling takes place with Langmuir waves with phase velocities of the order of the velocity of light.

To find the actual shape of the longitudinal wave spectrum we compare the induced scattering rate of longitudinal into longitudinal waves on one hand with that of longitudinal into transverse waves on the other hand. These processes are described by (Kaplan and Tsytovich, 1973)

$$\frac{\partial W_{k_1}^{\parallel}}{\partial t} = \frac{\pi}{108} \frac{\omega_{pe}^3}{n_e m_i v_{te}^4} W_{k_1}^{\parallel} \frac{\partial W_{k_1}^{\parallel}}{\partial k_1}$$ \hspace{1cm} (26)

and

$$\frac{\partial W_{k_1}^{\perp}}{\partial t} = \frac{\pi}{108} \frac{\omega_{pe}^3}{n_e m_i v_{te}^4} W_{k_1}^{\perp} \frac{k_1}{\partial k_1} (k_1 W_{k_1}^{\parallel}).$$ \hspace{1cm} (27)

Consequently the ratio of the transfer rates is $W_{k_1}^{\parallel}/W_{k_1}^{\perp}$, if we neglect $W_{k_1}^{\parallel}$ in comparison with $k_1 \partial W_{k_1}^{\parallel}/\partial k_1$. Since $W_{k_1}^{\parallel} \gg W_{k_1}^{\perp}$, the Langmuir waves will be concentrated at high phase velocities and, from the above discussion, along the magnetic field direction.

6.2. Dimensions and energetics of the whistler wave packet

The whistler wave packet has to be sufficiently extended in space and of sufficiently large amplitude in order to give rise to an observable structure in the radio signals.

At each instant the wave packet must extend all over that part of the source where the excited Langmuir waves of one frequency are concentrated. Assuming that the Langmuir waves lie within a frequency band $(\omega_{pe}, \omega_{pe} + \Delta \omega)$ with

$$\Delta \omega = \omega_{pe} \left( \frac{\omega_{ce}^2}{2 \omega_{pe}^2} + \frac{3}{2} \frac{v_{te}^2}{v_p^2} \right) \approx 2.2 \times 10^{-3} \omega_{pe}$$ \hspace{1cm} (28)

corresponding to phase velocities exceeding $30 v_{te}$ (cf. Equation (25)), the whistler packet should cover the same range of plasma frequencies in the solar corona. For an isothermal hydrogen corona in hydrostatic equilibrium the dependence of the density
$n$ on the distance from the solar centre $R$ (in solar radii) is

$$n = n_0 \exp \left[ \frac{11.5}{(T_6 R)} \right],$$

(29)

where $T_6$ is the temperature in units of $10^6$ K (cf. Van de Hulst, 1950). From this relation we then find a value of order 100 km for the spatial extent of the whistler wave packet at $R \approx 1$. Similarly, the lateral extent of the wave packet must be of the order of the lateral source dimension, say $10^4$ km.

We can now compare the amount of radiation produced by the coalescing whistlers and Langmuir waves with the observed flux values by means of Equation (21).

If the continuum radiation is the result of induced scattering of Langmuir waves on ions, the Langmuir wave energy density must exceed $10^{-5}$ erg cm$^{-3}$ (Kuijpers, 1974). Using $\Delta k_i \approx \omega_{pe}/(30 v_{te})$, $\omega_{pe}/2\pi = 300$ MHz and $T = 10^6$ K we find $W^{l}_{k_i} \approx 6.2 \times 10^{-5}$ erg cm$^{-2}$. For the whistler energy density per unit wave number we take the quantity $3(\omega_{ce}/\omega_{pe})^{3/2} W^{l}_{k_i}$, which is derived on the assumption that the coalescence of Langmuir waves with whistler waves is at least as efficient as the mutual coalescence of Langmuir waves (cf. (24)). Then $W^{w}_{k_i} \approx 1.1 \times 10^{-6}$ erg cm$^{-2}$ for $\omega_{pe}/\omega_{ce} = 30$. If we assume that the total production rate of $W^{l}_{k_i}$ is simply given by multiplication of the right hand part of Equation (21) with the whistler volume $\pi 2.5 \times 10^{24}$ cm$^3$ and that, moreover, the transverse radiation is not weakened by collisions, the observed flux at a distance of 1 AU is $7.6 \times 10^{-5}$ erg cm$^{-2}$ s$^{-1}$ cm, or, in frequency measure, $2.5 \times 10^{-15}$ erg cm$^{-2}$ s$^{-1}$ Hz$^{-1}$. This value well exceeds the flux variations of $6.5 \times 10^{-18}$ erg cm$^{-2}$ s$^{-1}$ Hz$^{-1}$ that have been observed in the type IV outburst of March 6, 1972 during the appearance of ‘sudden reductions’ (Krüger, 1972; cf. Kuijpers, 1974). We have assumed that the source is optically thin for the coupling process, that is, we have neglected the decay of a transverse wave into a whistler and a Langmuir wave.

Taking account of this decay one can show that the effective temperature of the transverse waves that are produced by the coalescence of whistlers and Langmuir waves has an upper bound of (Melrose, 1970)

$$T_{k_i} \leq \frac{\omega_i T_{k_i} T_{k_w}}{\omega_w T_{k_i} + \omega_i T_{k_w}}.$$  

(30)

Here the effective temperature for wave mode $j$ is defined according to

$$K_B T_{k_j} = W^{l}_{k_j} (2\pi)^3,$$

(31)

where $W^{l}_{k_j}$ is the wave energy density per unit volume in wave vector space. For the above situation one can derive $T_{k_i} \approx 1.0 \times 10^{15}$ K and $T_{k_w} \approx 1.2 \times 10^{14}$ K. Then from Equation (30) we find $T_{k_i} \approx 9 \times 10^{14}$ K for $\omega_c = 0.5 \omega_{ce}$ and $\omega_{pe}/\omega_{ce} = 30$. The above calculated flux corresponds to $T_{k_i} \approx 3.3 \times 10^{14}$ K. The observed flux corresponds to a brightness temperature of $8.4 \times 10^{11}$ K and leads to a lower bound (due to collisional damping) of the effective temperature in the source $T_{k_i} \approx 8.4 \times 10^{11}$ K. Consequently
we conclude that the adopted source model is optically thick for the coupling process and can give rise to the observed flux variations.

6.3. EXPLANATION OF THE OBSERVATIONS

(1) The association of intermediate drift bursts with continua naturally results from the suggested mechanism, in which Langmuir waves are an essential ingredient, since Langmuir waves can produce a structureless continuum on their own (Kuijpers, 1974; 1975).

(2) If at each instant the whistler wave packet extends over that part of the source where the Langmuir waves of one frequency are present (cf. Section 6.2) and if the Langmuir waves are concentrated around wave vectors such that the conservation relations (18) can be fulfilled (cf. Section 6.1), coupling with the Langmuir waves can lead to an observable flux enhancement at a frequency $\omega_{pe} + \omega_w$ relative to the ambient continuum (Section 6.2).

If, moreover, the coupling with whistlers leads to a temporal decrease of the energy density of Langmuir waves by a few per cent, the net flux of radiation at $\omega_{pe}$ can decrease temporally by a considerable fraction depending on the actual amplification length in the source for induced scattering at $\omega_{pe}$ (cf. Equation (22)). Then this leads to a depression of the observed flux at $\omega_{pe}$ relative to the ambient continuum.

Then the result of the coupling of whistler waves, that travel along the magnetic field in the form of a discrete wave packet, with the Langmuir waves clearly is a structure that drifts in the frequency-time plane and having an instantaneous frequency profile consisting of a depression relative to the ambient continuum on the low frequency side and an enhancement shifted to higher frequencies over a distance equal to the dominant whistler frequency. Moreover, the depression-enhancement structure should begin (and end) suddenly, in agreement with the observations.

(3) Since the coupling mechanism is independent of the induced scattering mechanism, in the sense that the latter needs a minimum amplification length in the source (Kuijpers, 1974) in contrast with the former, one expects to observe sometimes only emission ridges without an ambient continuum, in conformity with a few observations by Young et al. (1961).

(4) Since the whistler instability is faster at the lower parts of the source where the loss-cones are large (Section 4.1), whereas the Langmuir waves are unstable for a wide range of loss-cones with half-apertures $7-80^\circ$ (Kuijpers, 1974), one expects to observe predominantly fibers with negative drift rates, in agreement with the observations.

(5) The drift rate of an individual fiber is expected to decrease with decreasing frequency, since the projected velocity of the wave packet along the density gradient decreases as the packet reaches the top of the magnetic arch. Occasionally one expects to see fibers bending from negative to positive drift rates (cf. Figure 1).

(6) Since the magnetic field strength decreases with increasing height, the ridge separation should decrease for successive fibers centered around decreasing frequencies, in agreement with the observations (Table I).
(7) The occurrence of intermediate drift bursts in clusters of similar fibers can be explained by the effect of quasilinear diffusion in pitch angle and the finite extent of the whistler unstable region. (Section 4.2).

(8) The field strength is most directly determined from the instantaneous ridge separation. From the observed values (Section 2) we have determined values of the magnetic field strength under the assumption that the whistler frequency is in between 0.1–0.5 \( \omega_{ce} \) (see Table III).

TABLE III

Magnetic field strength from observed intermediate drift bursts

<table>
<thead>
<tr>
<th>Plasma frequency (MHz)</th>
<th>Magnetic field strength (G)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>(1)</td>
</tr>
<tr>
<td>900</td>
<td>7.2 –36</td>
</tr>
<tr>
<td>320</td>
<td>1.1 –11</td>
</tr>
<tr>
<td>160</td>
<td>0.36– 1.8</td>
</tr>
</tbody>
</table>

(1) Determined from instantaneous ridge separation assuming that \( \Delta \omega = 0.1–0.5 \omega_{ce} \).

(2) Determined from observed drift rates using relation (32) with \( T_0 = 1 \) and \( R = 1 \). For the drift rates at 900, 320 and 160 MHz we have used, respectively, \(-50\), \(-10\) and \(-3.5\) MHz s\(^{-1}\).

On the other hand the observed drift rates allow to find a value for the field strength by means of the coronal density model (29) and the relationship \( v_G \approx 21.5–28 \, v_A \). Neglecting projection effects one finds the magnetic field strength in gauss to be

\[
B = \left| \frac{d\omega_{pe}}{dt} \right| \frac{2\pi}{T_0 R^2 \epsilon},
\]

(32)

where \( (2\pi)^{-1} \frac{d\omega_{pe}}{dt} \) is the plasma-frequency drift rate in MHz s\(^{-1}\), which is about equal to the observed drift rate and \( \epsilon = 0.29 \), respectively 0.22 for \( v = 21.5 \, v_A \) and 28 \( v_A \). From Table III it follows that the values of the field strengths determined in both ways are in good agreement with each other.

(9) If the Langmuir waves are concentrated in a narrowband of width \(< 10^{-2} \, \omega_{pe} \) above the local plasma frequency (Section 6.1), the escaping radiation is purely in the ordinary mode both in the case of coupling with whistlers (Section 5.1) and in the case of induced scattering by the thermal particles, since the frequency of the outgoing wave does not exceed a value of \( \omega_{pe} + 0.5 \omega_{ce} \).

We mention that the result for the dominant sense of polarization is the reverse of that given by Kuijpers (1974) concerning the generation of the continuum, where the bandwidth of the longitudinal waves was assumed to be much larger (\( \approx 0.1 \, \omega_{pe} \)).

The result that the radiation produced by induced scattering as well as that produced
by coalescing waves has the same sense of polarization is in agreement with the observations with the Utrecht spectrograph.

7. Conclusion

We have found that the fiber phenomenon can be explained by coupling of whistler wave packets with Langmuir waves in the source region of the continuum under the condition that the Langmuir turbulence is concentrated at phase velocities approaching the velocity of light. The whistlers, like the Langmuir waves, can be excited by loss-cone distributions of fast electrons superimposed upon a thermal background for the assumed source conditions, but the whistlers originate preferably at lower heights than the Langmuir waves. Their pulse-like way of emission can be explained by quasi-linear pitch angle scattering of the anisotropic fast particle distribution into the loss-cone with a resultant quenching of the instability. Only after the whistlers have left the unstable region, anisotropic loss-cone distributions can be built up again by the fast particles that travel downwards into the unstable region.

The observed frequency separation of the absorption and emission ridges corresponds to a fraction of 0.1–0.5 of the local electron cyclotron frequency and leads to estimates of the magnetic field strength in the source region. On the other hand the observed drift rate should correspond to the component of the propagation velocity along the density gradient and also leads to estimates of the field strength. The latter estimates are in good agreement with the former if one uses a simple hydrostatic equilibrium model.

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References


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Appendix

A 1. LIST OF SYMBOLS

\[ a = 3^{1/2} v_{te} / c \]  
loss-cone half-aperture, particle pitch angle

\[ \alpha \]  
magnetic field vector

\[ c \]  
velocity of light

\[ D_\alpha \]  
pitch angle diffusion coefficient

\[ E \]  
electric field vector

\[ e \]  
electron charge in absolute value

\[ f_{c, h} \]  
cold and, resp., hot electron distribution function

\[ j \]  
electric current

\[ K_B \]  
Boltzmann constant

\[ L \]  
length of the whistler generating region

\[ m_j \]  
particle mass

\[ n = ck / \omega \]  
index of refraction

\[ n_j \]  
particle density

\[ n_h \]  
fractional fast particle density

\[ \omega_j, k_j \]  
angular frequency, resp., wave vector

\[ \theta_j \]  
angle between wave vector of kind \( j \) and magnetic field

\[ k_+ = k_w + k_l \]  
\[ \omega_+ = \omega_w + \omega_l \]  
\[ \omega_r, \gamma \]  
real, resp., imaginary frequency part (\( \gamma > 0 \): growth)

\[ \omega_{pj} = (4\pi n_j e^2 / m_j)^{1/2} \]  
plasma frequency
\[ \omega_{ej} = eB/m_jc \]  
\( r \)  
\( R \)  
\( T_j \)  
\( T_{k_j} \)  
\( \tau_{qi} \)  
\( \tau_{\text{max}} \)  
\( \tau_{\text{pass}} \)  
\( \tau_{\text{bounce}} \)  
\( \mathbf{v} \)  
\[ v_{ij} = (K_BT_j/m_j)^{1/2} \]  
\( v_h \)  
\[ v_A = B/(4\pi n_e m_i)^{1/2} \]  
\[ v_{Ac} = B/(4\pi n_e m_e)^{1/2} \]  
\( v_p = \omega/k \)  
\[ v_G = \partial \omega/\partial k \]  
\( W \)  
\( W^j = \int_0^\infty W^j_{k_j} dk_j \)  
\( W^j_{k_j} \)  
\( W^j_{k_j} \)  
\( \mathbf{x} = \omega/\omega_{ce} \)  
\( X = \omega_{pe}^2/\omega^2 \)  
\( Y = \omega_{ce}/\omega \)  
\( Y_L = Y \cos \vartheta \)  
\( Y_T = Y \sin \vartheta \)  
\[ \psi_r = (\omega - \omega_{ce})/(kv_{te} 2^{1/2}) \]  
\[ \psi_p = (\omega - \omega_{ce})/(kv_{he} 2^{1/2} \cos \alpha) \]  
\( Z_r = Z(\psi_r) \)  
\( Z_h = Z(\psi_h) \)  

**A2. Indices**

- \( l, \perp \): component parallel, resp., perpendicular to the magnetic field
- \( e \): electrons
- \( i \): ions
- \( w \): whistlers
- \( l \): Langmuir waves
- \( t \): electromagnetic waves above the electron plasma frequency
- \( k \): \((\mathbf{k}, \omega)\)