Astron. Astrophys. 328, 390–398 (1997)

ASTRONOMY AND ASTROPHYSICS

Plasma radiation of power-law electrons in magnetic loops: application to solar decimeter-wave continua

V.V. Zaitsev[1], **A. Krüger**[2], **J. Hildebrandt**[2], **and B. Kliem**[2]

[1] Institute of Applied Physics, Nizhny Novgorod 603600, Russia
[2] Astrophysikalisches Institut Potsdam, An der Sternwarte 16, D-14482 Potsdam, Germany

Received 24 April 1997 / Accepted 21 July 1997

Abstract. It is shown that flare-produced fast electrons with a power-law energy distribution trapped in magnetic loops are capable to produce plasma waves at the upper hybrid frequency due to a loss-cone instability. Conditions for instability are obtained in dependence of the parameters of the power-law distribution and the magnetic-loop parameters. The growth rate of the plasma waves has been calculated and their energy density has been estimated in the frame of quasi-linear saturation. This instability has been considered as the cause of the solar decimetric continuum which exhibits a strong temporal and spatial correlation with regions of flare-energy release and sources of microwave bursts. The strong absorption of the first harmonic in the decimeter range and also the peculiarities of the conversion of the plasma waves into electromagnetic waves yield a preference of the generation of the decimetric continuum at the second harmonic of the plasma frequency. In this case the polarization of the radiation corresponds to the ordinary wave mode in a wide cone of propagation angles around the direction perpendicular to the magnetic field.

Key words: solar radio emission – Sun: flares – Sun: activity – Sun: corona

1. Introduction

Hard X-ray observations demonstrate that energetic flare-electrons frequently exhibit an energy distribution which can be approximated by a power law (Lin 1974). The gyrosynchrotron radiation from these electrons trapped in magnetic fields of solar active regions can be regarded as the main source of microwave bursts at frequencies >1 GHz (cf. Kundu & Vlahos 1982). Another important burst component which is intimately connected with flare electrons is formed by the decimetric continuum comprising mainly type IVdm bursts, but also a number of substructures like pulsations, spikes, zebra patterns etc. in the frequency range between about 200 MHz and 1 or 2 GHz (Zheleznyakov 1970; Krüger 1979; Benz 1985, 1986, 1996; Isliker & Benz 1994). This connection is evident from the near-coincidence of the time profiles of microwave bursts and decimetric continua (with a small delay of the latter, cf., e.g., Aschwanden et al. 1990) and the high correlation between decimetric spikes and hard X-ray bursts (Aschwanden & Güdel 1992). The presence of peculiarities in the dynamical spectrum of the decimetric continuum such as zebra patterns and sudden reductions has lead to the assumption that this continuum is generated by plasma waves at the upper hybrid frequency $\omega_{\rm uh} = (\omega_{\rm p}^2 + \omega_{\rm g}^2)^{1/2}$, excited by fast electrons having a loss-cone anisotropy (Kuijpers 1974; Zaitsev & Stepanov 1975; Benz & Kuijpers 1976) while interaction with lower hybrid waves has been invoked to explain dm-pulsations (Benz 1980).

There arises the question whether the electrons generating the microwave bursts via gyrosynchrotron emission and those electrons generating the decimetric continuum by the plasma-wave mechanism represent two populations of power-law electrons with similar velocity distributions and the same origin.

In the present paper we consider the instability of a power-law distribution with a loss-cone anisotropy and show that such a distribution can generate plasma waves at all reasonable values of the mirror ratio and of the exponent of the power-law distribution function (Sect. 2). Furthermore, we calculate the energy density of the plasma waves excited by the loss-cone instability and relate it to the parameters of the power-law electron energy distribution and properties of the trapping magnetic loop (Sect. 3). In Sect. 4 we investigate the possibility of the generation of both, decimetric continua and microwave bursts by one original population of flare-generated electrons with a power-law energy spectrum and assume that the same power-law distribution (with different number densities) exists in both sources. In Sect. 5 we discuss the results and in Sect. 6 the conclusions are summarized.

Send offprint requests to: A. Krüger

2. Possibility of the generation of plasma waves by power-law electrons

2.1. Energy distribution function

We consider the generation of plasma waves in a system consisting of a background plasma with an electron density n and a minor part of fast electrons with a density $n_1 \ll n$ and a distribution function $f(v_\parallel, v_\perp)$, where $\parallel$ and $\perp$ denote the vector components parallel and perpendicular to the direction of the magnetic field $\boldsymbol{B}$, respectively. We suppose that the plasma is sufficiently dense ($\omega_\mathrm{p}^2 \gg \omega_\mathrm{g}^2$, where ω_p and ω_g are the plasma frequency and the electron gyrofrequency, respectively), and the plasma waves propagate nearly perpendicular to the magnetic field ($k_\perp^2 \gg k_\parallel^2$, where k is the wave number of the plasma waves at $\omega \approx \omega_\mathrm{p}$). The growth rate γ of the plasma waves is given by the following formula (Mikhailovskij 1974):

$$\gamma = \frac{\pi}{n}\frac{\omega_\mathrm{p}^4}{k^3}\int_{-\infty}^{\infty} \mathrm{d}v_\parallel \int_{\omega^2/k^2}^{\infty} \mathrm{d}v_\perp^2 \, \frac{\partial f/\partial v_\perp^2}{\sqrt{v_\perp^2 - \omega^2/k^2}}. \qquad (1)$$

It is seen from Eq. (1) that for instability ($\gamma > 0$) it is necessary that the derivative $\partial f/\partial v_\perp^2$ is positive at least at some part of the integration path. If we take the distribution function in the form used by Benz & Kuijpers (1976)

$$f(v_\parallel, v_\perp) = \frac{A}{(v_\parallel^2 + v_\perp^2)^\delta}\,\Theta\left(v_\perp^2 - \frac{v_\parallel^2}{\sigma - 1}\right), \qquad (2)$$

where

$$\Theta(x) = \begin{cases} 1 & \text{for} \quad x \geq 0 \\ 0 & \text{for} \quad x < 0 \end{cases} \qquad (3)$$

is the step function and $\sigma = B_\mathrm{max}/B_\mathrm{min}$ ("mirror ratio") characterizes the magnetic loop trapping the particles, we find that the derivative of the distribution function

$$\frac{\partial f}{\partial v_\perp^2} = -\frac{A\delta}{(v_\parallel^2 + v_\perp^2)^{\delta+1}}\,\Theta\left(v_\perp^2 - \frac{v_\parallel^2}{\sigma - 1}\right) + \frac{A}{(v_\parallel^2 + v_\perp^2)^\delta}\,\frac{\partial\Theta}{\partial v_\perp^2} \qquad (4)$$

is always negative except at the boundary of the loss cone where it tends to infinity, since $\partial\Theta/\partial v_\perp^2 > 0$ is the delta function. In particular this part of the derivative leads to instability and must be considered in the evaluation of the integral in Eq. (1).

In our analysis we will apply a power-law distribution with a boundary of the loss cone that depends on energy:

$$f_1(v_\parallel, v_\perp) = \frac{A_1}{(v_\parallel^2 + v_\perp^2)^\delta}\,\left[1 - \exp(-y)\right] \qquad (5)$$

with

$$y = (\sigma - 1)\,\frac{v_\perp^2}{v_\parallel^2} \qquad (6)$$

and the condition $v > v_1$, where v_1 is the minimum value of the velocity of the power-law distribution. The distribution function (5) leads to a finite derivative $\partial f_1/\partial v_\perp^2$ at the boundary of the loss cone. In the opposite case we would leave the frame of the kinetic approximation for the description of the instability and Eq. (1) for the growth rate would no longer be valid (Mikhailovskij 1974).

The normalization coefficient A_1 in Eq. (5) follows from the condition

$$2\pi \int_{-\infty}^{\infty} \mathrm{d}v_\parallel \int_0^\infty f_1(v_\parallel, v_\perp)\, v_\perp \, \mathrm{d}v_\perp = n_1, \qquad (7)$$

where n_1 is the density of the power-law electrons. In the case that the mirror ratio is sufficiently large ($\sigma \gg 1$), the value of A_1 is given by the following equation

$$A_1 = \frac{n_1}{4\pi}\,\frac{(4\sigma - 4)}{(4\sigma - 5)}\,(2\delta - 3)\, v_1^{2\delta - 3}. \qquad (8)$$

The distribution functions (2) and also (5) are restricted to certain minimum values δ characterizing the hardness of the power-law spectrum. In order to keep n_1 finite, $\delta > 3/2$ is required, and in order to have a finite average velocity $\langle v \rangle = \int v\, f(\boldsymbol{v})\, \mathrm{d}^3\boldsymbol{v}$, it is necessary to take $\delta > 2$.

2.2. *Growth rate of instability*

Inserting the distribution function (5) into Eq. (1) and evaluating the integral, we obtain the following expression for the growth rate of the plasma waves at $\omega = (\omega_\mathrm{p}^2 + \omega_\mathrm{g}^2)^{1/2} \approx \omega_\mathrm{p}$:

$$\gamma = \frac{2\pi^{3/2}\,\omega^4\,A_1}{n\,k^3\,v_1^{2\delta}}\left\{\frac{(\delta + \sigma - 1)}{\sqrt{\sigma - 1}}\,\Gamma(\delta)\,\gamma^*(\delta, x) - \frac{1}{2\delta}\,\frac{\Gamma(\delta + \frac{1}{2})}{\Gamma(\delta)}\right\}, \qquad (9)$$

where γ^* and x are defined as

$$\gamma^*(\delta, x) = \frac{x^{-\delta}}{\Gamma(\delta)}\,\gamma(\delta, x), \qquad (10)$$

$$x = \frac{\omega^2(\sigma - 1)}{k^2 v_1^2}. \qquad (11)$$

$\Gamma(\delta)$ is the gamma function and $\gamma(\delta, x)$ is given by

$$\gamma(\delta, x) = \int_0^x \mathrm{e}^{-t}\, t^{\delta - 1}\, \mathrm{d}t. \qquad (12)$$

Eq. (9) is obtained under the conditions $(\sigma - 1) > 1$ and $\omega^2/k^2 < v_1^2$.

The function $\gamma^*(\delta, x)$ can be written in form of an expansion (Abramovitz & Stegun 1964)

$$\gamma^*(\delta, x) = \mathrm{e}^{-x}\left[\frac{1}{\Gamma(\delta + 1)} + \frac{x}{\Gamma(\delta + 2)} + \frac{x^2}{\Gamma(\delta + 3)} + \frac{x^3}{\Gamma(\delta + 4)} + \cdots\right], \qquad (13)$$

which is strongly decreasing for values $\delta > 2$ we are interested in. Therefore, we can restrict our analysis to the first term of the expansion. Then, Eq. (9) simplifies to

$$\gamma = \frac{2\pi^{3/2}\,\omega^4\,A_1}{n\,k^3\,v_1^{2\delta}}\left[\frac{\delta + \sigma - 1}{\sqrt{\sigma - 1}}\,\mathrm{e}^{-x} - \frac{\Gamma(\delta + \frac{1}{2})}{2\Gamma(\delta)}\right]. \qquad (14)$$

From Eq. (14) follows that the instability is maximal for $\omega/k \leq v_1/\sqrt{\sigma-1}$, i.e., for plasma waves with a sufficiently small phase velocity. For $\omega/k \gg v_1/\sqrt{\sigma-1}$ the first term in the brackets of Eq. (14) is exponentially small and the instability vanishes. The instability condition for plasma waves with small phase velocities ($\omega/k \ll v_1/\sqrt{\sigma-1}$) can be written as follows

$$\frac{(\delta+\sigma-1)}{\sqrt{\sigma-1}}\,\frac{\Gamma(\delta)}{\Gamma(\delta+\frac{1}{2})} > \frac{1}{2}\,. \tag{15}$$

From Eq. (15) it can easily be seen that instability occurs for all relevant mirror ratios σ and exponents δ of the power-law spectrum. However, the maximum growth rate

$$\gamma_{\max} = 0.36\,\frac{n_1}{n}\,\omega\,\frac{(4\sigma-4)}{(4\sigma-5)}\,\frac{(2\delta-3)}{\delta}\,\frac{(\delta+\sigma-1)}{(\sigma-1)^2} \tag{16}$$

decreases with increasing mirror ratio σ which is related to a diminishing of the number of particles outside the loss cone contributing to the instability.

2.3. Influence of the background plasma

In the ambient plasma the excited plasma waves are subject to Landau damping. The damping rate γ_L for the most unstable waves with $\omega^2/k^2 = (3/2)v^2/(\sigma-1)$ is given by the formula

$$\gamma_L = \sqrt{\frac{\pi}{2}}\left(\frac{3}{2}\right)^{3/2}\left(\frac{v_1}{v_T}\right)^3\frac{\omega_p}{(\sigma-1)^{3/2}}\times \exp\left[-\frac{3}{4}\left(\frac{v_1}{v_T}\right)^2\frac{1}{(\sigma-1)}\right], \tag{17}$$

where $v_T = (\kappa T/m)^{1/2}$ is the electron thermal velocity of the background plasma (T – temperature, κ – Boltzmann's constant). The condition $\gamma_{\max} > \gamma_L$ determines the values of the density of fast electrons for which the excited plasma waves are not seriously damped:

$$\frac{n_1}{n} > 6.4\left(\frac{v_1}{v_T}\right)^3\frac{(4\sigma-5)}{(4\sigma-4)}\,\frac{\delta}{(2\delta-3)}\,\frac{\sqrt{\sigma-1}}{(\delta+\sigma-1)}\times \exp\left[-\frac{3}{4}\left(\frac{v_1}{v_T}\right)^2\frac{1}{(\sigma-1)}\right]. \tag{18}$$

The dependence of the minimum density of the power-law electrons, beginning with that being able to generate plasma waves, on the mirror ratio σ is shown in Fig. 1 for $\delta = 4$ using the ratio v_1/v_T as a parameter. The curves were derived from Eq. (18) under the assumption that Landau damping acts as the main attenuation mechanism of plasma waves in the background plasma. They were found to depend only weakly on δ in the range $2 < \delta \leq 5$. It can be concluded from Fig. 1 that Landau damping suppresses the generation of plasma waves only in the immediate source region of the microwave bursts, i.e., in those parts of the flaring loops where a sufficiently hot plasma with temperatures of the order $(0.5–1)\times 10^7$ K can be assumed. Hence for the excitation of plasma waves either an extremely

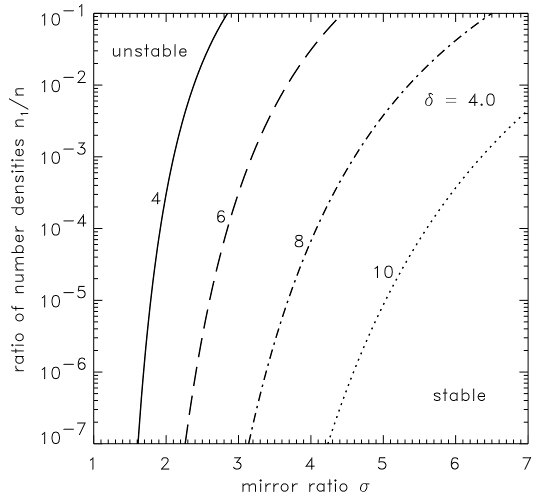

Fig. 1. Dependence of the minimum density of energetic power-law electrons allowing for the generation of plasma waves in a sufficiently hot plasma on the mirror ratio σ. The numbers on the curves are the values of the ratio v_1/v_T according to the temperature of the background plasma and the minimum energy ε_1 of the power-law energy distribution. Assuming $\varepsilon_1 = 10$ keV, the corresponding temperatures would be 2.8×10^7 K (4), 1.25×10^7 K (6), 7×10^6 K (8), and 4.5×10^6 K (10). The curves were calculated on the basis of inequality (18) under the assumption that the main damping mechanism of the plasma waves is Landau damping.

high concentration of energetic electrons ($n_1/n \gtrsim 10^{-3}$) or a sufficiently large minimum energy value $\varepsilon_1 > 20$ keV of the power-law electron energy distribution is necessary. The possibility of generation of plasma waves within the source region of the microwave bursts must be considered as a rather extreme case, although this possibility cannot be fully excluded.

On the other hand, the plasma in the immediate vicinity of a flaring magnetic loop can be sufficiently cold with a temperature $T \approx 7\times10^5$ K (Benz et al. 1992). In this case even a relatively small number of power-law electrons escaping from the flare volume can generate plasma emission and the threshold of instability is determined by the damping of plasma waves due to electron-ion collisions in the plasma. Then the damping rate is given by

$$\gamma_c = \frac{\nu_{ei}}{2}, \tag{19}$$

where

$$\nu_{ei} = \frac{5.5\,n}{T^{3/2}}\ln\left(\frac{10^4\,T^{2/3}}{n^{1/3}}\right) \qquad \text{for } T > 4\times10^5 \text{ K} \tag{20}$$

is the effective collision frequency of electron-ion encounters in the plasma (Zheleznyakov 1996, p. 261).

From Eqs. (17) and (19) follows that for $n = 3\times10^9$ cm^{-3} and $v_1 = 5.84\times10^9$ cm s^{-1} ($\varepsilon_1 = 10$ keV) the collisional damping of the plasma waves begins to dominate over the Landau damping if the temperature of the background plasma satisfies the condition

$$T < T^* = \frac{3}{\sigma-1}\times10^6 \text{ K}. \tag{21}$$

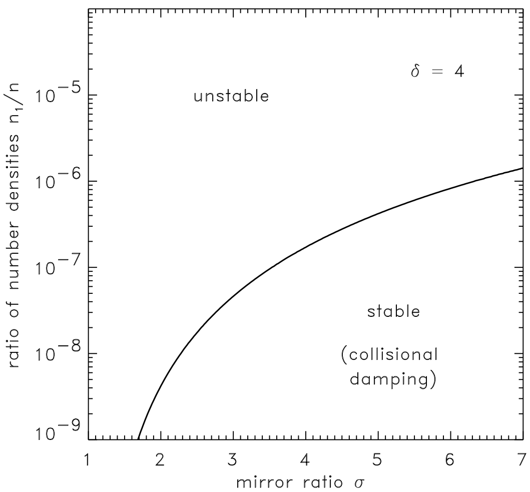

Fig. 2. Dependence of the threshold of the excitation of plasma waves on the mirror ratio σ for the case that the damping of the plasma waves is due to collisions of electrons with ions in the background plasma, i. e., if the temperature of the plasma satisfies condition (21). Parameter values are $n = 3 \times 10^9$ cm^{-3}, $\delta = 4$, and $v_1 = 5.84 \times 10^9$ cm s^{-1} ($\varepsilon_1 = 10$ keV).

In this case the maximum growth rate of the instability given by Eq. (16) exceeds the collisional damping if

$$\frac{n_1}{n} > 2.3 \times 10^{-8} \, \frac{(4\sigma - 5)}{(4\sigma - 4)} \, \frac{(\sigma - 1)^{7/2}}{(\sigma + 3)} . \tag{22}$$

Inequality (22) corresponds to the condition $\gamma_{max} > \nu_{\rm ei}/2$ under the assumption that $T = T^*$, where T^* is determined by Eq. (21). Fig. 2 shows the dependence of the ratio n_1/n on the mirror ratio σ for $\delta = 4$ and values of the temperature determined by condition (21). We see from Fig. 2 that in the given case the threshold of the excitation of plasma waves is sufficiently small ($(n_1/n)_{\rm min} \lesssim 10^{-6}$, depending on the mirror ratio).

3. The energy density of the plasma waves

The plasma waves generated by fast electrons lead to a diffusion of these electrons inside the loss cone and to their escape from the magnetic trap. The characteristic diffusion time $T_{\rm d}$ depends on the size of the source region of the fast particles. In the case of a sufficiently large loop size the diffusion time is smaller than the mean life time of the fast particles in the magnetic trap (Bespalov et al. 1991):

$$T_{\rm d} < \sigma \, \frac{L_{\|}}{2\langle v \rangle} \qquad \text{or} \qquad T_{\rm d} < \frac{L_{\|}}{2\langle v \rangle} . \tag{23}$$

Here $L_{\|}$ denotes the length of the magnetic trap and $\langle v \rangle$ the average velocity of the fast particles. For characteristic values of $L_{\|}$ and $\langle v \rangle$ the diffusion time is of the order of some fraction of a second to a few seconds which is much less than the flaring time of the generation of fast particles (which is about 1 minute for an impulsive flare). Thus, for an estimation of the energy density of the plasma waves we can apply the quasi-linear theory under stationary conditions since the quasi-linear diffusion of the particles into the loss cone is compensated by the source of fast particles, and the growth rate of plasma waves under quasi-linear conditions compensates its damping by electron-ion collisions or Landau damping.

A stationary model of the generation of plasma waves at the upper hybrid frequency $\omega = (\omega_{\rm p}^2 + \omega_{\rm g}^2)^{1/2}$ was considered by Shaposhnikov (1988). He found that under the assumption of a one-dimensional diffusion along the line

$$v_{\|}^2 = \frac{\omega - s\,\omega_{\rm g}}{\omega_{\rm g}} \, v_{\perp}^2 \tag{24}$$

in velocity space with conservation of the pitch angle of the fast electrons the energy density of plasma waves can be expressed by the formula

$$W_{\rm L} = \frac{J m \langle v^2 \rangle}{4\,\nu_{\rm d}} \ln \sigma , \tag{25}$$

where J[cm^{-3} s^{-1}] is the source function which determines the amount of fast particles released in the flare source per volume unit and time unit, $\nu_{\rm d}$ denotes the dissipation rate of the energy of the plasma waves which is, in our case, determined either by Landau damping or by electron-ion collisions:

$$\nu_{\rm d} = \begin{cases} 2\gamma_{\rm L} & \text{for} \quad T > T^* \\ \nu_{\rm ei} & \text{for} \quad T < T^* \end{cases} \tag{26}$$

In the stationary state the source function J can be expressed in terms of the fast-particle density n_1 and the mean lifetime T_1 of particles in the magnetic trap:

$$J = \frac{n_1}{T_1} . \tag{27}$$

T_1 can be estimated in the following way (Bespalov et al. 1991):

$$T_1 = \frac{N}{2S} , \tag{28}$$

where

$$N = \iint \frac{B_{\rm max}}{B} \, f(\boldsymbol{v}) {\rm d}^3\boldsymbol{v} \, {\rm d}z \tag{29}$$

is the number of energetic particles in the magnetic trap with unit cross-section at the footpoint where $B = B_{\rm max}$, and

$$S = 2\pi \int_0^\infty \int_{-1}^{1} f(v,\sigma)\, v^3 \, \cos\vartheta \, {\rm d}v \, {\rm d}\cos\vartheta \tag{30}$$

is the particle flux at the footpoint of the magnetic tube.

For moderate diffusion, where $L_{\|}/(2\langle v \rangle) < T_{\rm d} < \sigma L_{\|}/(2\langle v \rangle)$, the velocity distribution function is quasi-isotropic and the loss cone is practically filled with fast particles, although there is no turbulent mirror at the footpoints of the flux tube yet which could impede the escape of particles from the loss cone.

Hence one obtains the following relations for N, S, and T_1 (Bespalov & Trakhtengerts 1986):

$$N \approx \sigma\, n_1\, L_\parallel, \qquad S \approx n_1 \langle v \rangle, \qquad T_1 \approx \sigma\, \frac{L_\parallel}{2\langle v \rangle}. \tag{31}$$

Then Eqs. (25), (27), and (31) allow to calculate the energy density of the plasma waves from the density of the fast electrons n_1 and the main parameters of the magnetic flux tube:

$$W_{\mathrm{L}} \approx \frac{n_1\, m \langle v^2 \rangle}{2} \left(\frac{\langle v \rangle}{\nu \sigma L_\parallel} \right) \ln \sigma. \tag{32}$$

For strong quasi-linear diffusion, i. e., if $T_{\mathrm{d}} < L_\parallel/(2\langle v \rangle)$, at the footpoints of the magnetic traps turbulent mirrors develop which reduce the escape of particles from the loss cone. This fact diminishes the stream of particles at the footpoints of the trap ($S < n_1 \langle v \rangle$). In this case Eq. (32) yields, for a given particle density n_1, an upper limit of the energy density of the plasma waves.

Usually, under conditions of the lower solar atmosphere, the parameter $\langle v \rangle/(\nu \sigma L_\parallel)$ is small; therefore the energy density of the plasma waves reaches only a fraction of the energy density of the fast particles. Let, e. g., n = (1–3)$\times 10^9$ cm^{-3}, T = (7–10)$\times 10^5$ K, $L_\parallel = 10^{10}$ cm, $\sigma = 3$, and $\langle v \rangle = 5 \times 10^9$ cm s^{-1}, we obtain from Eq. (32) $W_{\mathrm{L}} \approx$ (2.4–6.5) $\times 10^{-4} n_1 m \langle v^2 \rangle$. Consequently, the ratio of the energy density of the plasma waves to the thermal energy density of the background plasma, which determines the efficiency of the conversion of the plasma waves into electromagnetic waves, is of order

$$w_{\mathrm{L}} = \frac{W_{\mathrm{L}}}{n \kappa T} \approx (6\text{–}16) \times 10^{-2}\, \frac{n_1}{n}. \tag{33}$$

4. Decimetric continuum

As already mentioned in the Introduction, there are several arguments which lead to the assumption that the solar decimetric continuum is generated by plasma waves at the upper hybrid frequency $\omega = (\omega_{\mathrm{p}}^2 + \omega_{\mathrm{g}}^2)^{1/2}$ (Kuijpers 1974; Zaitsev & Stepanov 1975; Benz & Kuijpers 1976). These waves are generated by trapped energetic electrons having an anisotropic velocity distribution of the loss-cone type.

It was shown in Sect. 3 that the flare-related magnetically trapped fast electrons with a power-law energy distribution are able to generate plasma waves at the upper hybrid frequency, which in turn may lead to the solar decimetric continuum. The close association of the sources of the decimetric continuum with the region of the primary flare energy release follows not only from the near-coincidence of the time profiles of the decimetric continuum and microwave bursts already mentioned in Sect. 1, but also from the following observations:

- The decimetric continuum is generated in source regions with small sizes comparable with those of microwave bursts (Zheleznyakov 1970).
- The sources of the decimetric continuum do not show a strong displacement across the solar disk similar as the sources of microwave bursts (Fleisher & Oshima 1961; Krishnan & Mullaly 1961).

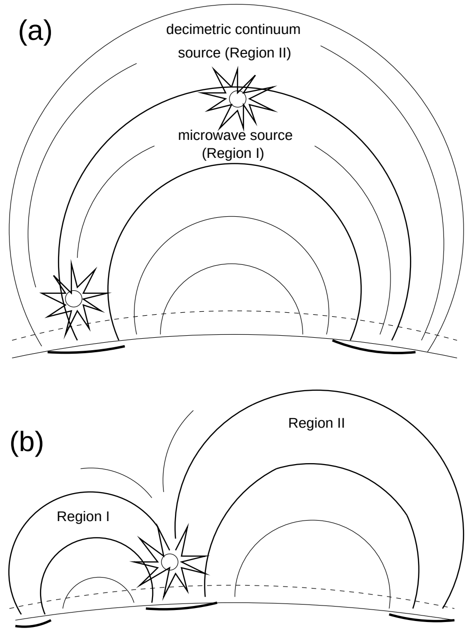

Fig. 3a and b. Schematic view of the relative positions of the source regions of the microwave burst (Region I) and the decimetric continuum burst (Region II) invoked in the present paper. The flare occurs either at the top of Region I or at its footpoints and acts as primary source of energetic power-law electrons for both, the microwave burst and the decimetric continuum. The source of the decimetric continuum may be a high-lying relatively cool loop (**a**) or a loop adjacent to the microwave burst source (**b**).

- According to limb observations the source heights of the decimetric continuum are about (4–5) $\times 10^9$ cm which is close to the heights of magnetic loops where the gyrosynchrotron radiation of flare electrons, i. e., solar microwave bursts are generated (Kundu & Firor 1961).
- The high correlation of decimetric bursts with hard X-ray bursts (Aschwanden & Güdel 1992).

We will assume that the decimetric continuum originates in relatively weak, closed magnetic fields directly adjoining the flare loops. The flare originates either at the top or the feet of a magnetic loop and acts as the general source of the energetic electrons for both the microwave burst and the decimetric continuum (Fig. 3).

In the following the microwave-burst and decimeter-continuum source will be denoted as Region I and Region II, respectively, both strongly differing in their parameters.

On the average, the solar microwave emission has a maximum flux density at, say, the frequency $\nu_{\mathrm{m}} \approx 5$ GHz. Inferring gyrosynchrotron radiation of energetic electrons with a power-law energy spectrum this corresponds roughly to magnetic fields of $B \approx 450$ G in the source region (Kundu & Vlahos 1982) and a value of the gyrofrequency $\nu_{\mathrm{g}} \approx 1.25$ GHz. In Region I, on the average, the plasma frequency should obey the relation $\nu_{\mathrm{p}} \leq \nu_{\mathrm{g}}$. In the case $\nu_{\mathrm{p}} \gg \nu_{\mathrm{g}}$ a strong suppression of the gyrosynchrotron radiation occurs (Razin-Tsytovich effect). For example, for ν_{p} = (5–6) ν_{g} the brightness temperature of the

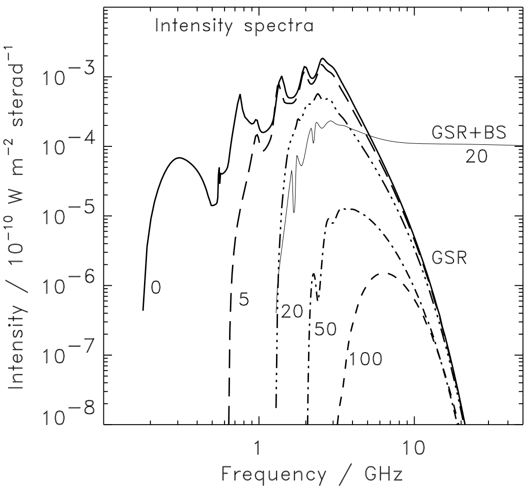

Fig. 4. Theoretical spectra of gyrosynchrotron radiation (GSR) generated by power-law electrons assuming various densities of the thermal background plasma in the (homogeneous) source volume. The numbers on the curves are density values in units of 10^9 cm^{-3}. The low-frequency suppression (Razin-Tsytovich Effect) can clearly be recognized for increasing density and will suppress the gyrosynchrotron emission below the level of thermal bremsstrahlung (BS) for $n > 2 \times 10^{10}$ cm^{-3}. For comparison, the resulting spectrum including bremsstrahlung for $n = 2 \times 10^{10}$ cm^{-3} is also shown in the plot (thin solid line). The other parameters are: B = 200 G, n_1 = 10^6 cm^{-3}, $L = 10^9$ cm, $T = 3 \times 10^6$ K, $\delta = 4$, ε_1 = 10 keV. In these calculations the power-law distribution was cut at the energy ε_2 = 500 keV.

source of the microwave burst emission decreases by 2–3 orders of magnitude to the level of Coulomb bremsstrahlung of the plasma filling the source volume (cf. Fig. 4). The condition $\nu_p \leq \nu_g$ gives a restriction on the density of the background plasma in Region I to values $n < 2 \times 10^{10}$ cm^{-3}.

On the average the spectral maximum of the decimetric continuum lies in the frequency range $\nu \approx$ (0.5–1) GHz (Isliker & Benz 1994). Hence we obtain a plasma density in the source region of $n \approx$ (3–6) $\times 10^9$ cm^{-3} for $\nu_m = (\nu_p^2 + \nu_g^2)^{1/2}$ or $n \approx$ (0.8–1.5) $\times 10^9$ cm^{-3} for $\nu_m = 2\,(\nu_p^2 + \nu_g^2)^{1/2}$. Here we assume that in Region II the gyrofrequency is much smaller than the plasma frequency. This condition $\nu_g^2 \ll \nu_p^2$ is necessary in order to exclude strong gyroresonance absorption at the levels $\nu = 2\nu_g$ and $\nu = 3\,\nu_g$ at the escape of the radiation from Region II outwards.

The temperature of the plasma inside (hot) flare loops is typically $T \sim 10^7$ K (Tsuneta 1996). We will assume that outside these flare loops the temperature is lower. For Region II (as sketched in Fig. 3 b) we assume $T \approx$ (0.7–2) $\times 10^6$ K although a heating of this region during the flare cannot be excluded (Aschwanden & Benz 1995; Tsuneta 1996)). Hence, summarizing the above arguments we use the following estimates for the characteristic plasma parameters in the Regions I and II taking into account that they may slightly differ towards higher or lower values:

- Region I (microwave burst source):
  $B \approx 500$ G
  $T \approx$ (0.5–1) $\times 10^7$ K
  $n < 2 \times 10^{10}$ cm^{-3}.
- Region II (source of decimetric continuum):
  $B \ll 2 \times 10^2$ G
  $T \approx$ (0.7–2) $\times 10^6$ K
  $n \approx$ (1–5) $\times 10^9$ cm^{-3}.

4.1. Brightness temperature of the radiation in the source region of the decimetric continuum

Analyzing the flare-energy support of the source of the decimetric continuum, the absorption of the radiation by electron-ion collisions at the wave propagation from the source through the solar atmosphere is important. Since this absorption is rather high, the effective temperatures of the decimetric continuum up to 10^{10} K (Kundu 1961; Krishnan & Mullaly 1961) can be ensured only as a result of a sufficiently high number of fast electrons injected from the flare region into the source region of the decimetric continuum comparable with the number of fast electrons in the microwave burst source.

The optical depth τ_c of the corona due to electron-ion collisions for radiation propagating from the source region of the decimetric continuum to the observer is given by

$$\tau_c\,(\nu) \approx \frac{\nu_p^2}{\nu^2}\,\frac{\nu_{ei}}{c}\,\frac{h_0}{\cos\vartheta}, \tag{34}$$

where ν_{ei} is determined by Eq. (20) and ϑ is the angle between the line of sight and the direction of the density gradient. Furthermore,

$$h_0 = \frac{\kappa T}{m\,g_\odot} \tag{35}$$

is the scale height of plasma density variation in the source region of the decimetric continuum. Using Eqs. (20) and (35), Eq. (34) can be rewritten:

$$\tau_c\,(\nu) \approx \frac{74.6}{\cos\vartheta}\,\nu_p^2\left(\frac{\nu_p}{\nu}\right)^2\left(\frac{10^6}{T}\right)^{1/2}, \tag{36}$$

assuming that ν is given in GHz. Taking ν_p = 0.5 GHz, $T = 7\times10^5$ K, $\cos\vartheta = 0.7$, we obtain from Eq. (36) $\tau_c \approx 32\,(\nu_p/\nu)^2$. We can see that even in the most favourable case ($\nu \approx 2\nu_p$) the optical depth becomes $\tau_c(\nu = 2\nu_p) = 8$ and the radiation at the second harmonic is weakened by 3×10^3 at the transit through the source of the decimetric continuum while the radiation of the first harmonic is weakened by a factor $\sim 10^{14}$!

The necessity of a consideration of a significant absorption in the corona for the study of the decimetric continuum was already early mentioned by Kuijpers (1974). Our estimations show that, due to the rather strong absorption of the first harmonic, the most favourable mechanism for the decimetric continuum is the generation of the second harmonic as a result of coupling of strong plasma waves excited by the loss-cone instability. Here the observed brightness temperatures $T_b^{(obs)} \approx 10^{10}$ K should correspond to brightness temperatures inside the source of the decimetric continuum which are about three orders higher, i. e.,

$$T_{\rm b}^{\rm (source)}\,(\nu \approx 2\nu_{\rm p}) \approx 10^3\,T_{\rm b}^{\rm (obs)} \approx 10^{13}\ {\rm K}. \tag{37}$$

This brightness temperature is less than the temperature of the upper hybrid Langmuir waves

$$T_{\rm L}^{\rm eff} \approx \frac{W_{\rm L}}{2\pi^2\,\kappa\,k^3} \approx 5\times 10^{11}\,T\,\frac{W_{\rm L}}{n\,\kappa\,T} \tag{38}$$

which for the case given here is of the order $T_{\rm L}^{\rm eff} \approx 10^{14}$–$10^{15}$ K.

4.2. Generation of the second harmonic

The emission at the second harmonic of the plasma frequency $\omega_{\rm t} = 2\omega_{\rm p}$ is generated by nonlinear coalescence of two plama waves (combination scattering) if the resonance condition

$$\omega + \omega' = \omega_{\rm t}, \qquad \boldsymbol{k} + \boldsymbol{k}' = \boldsymbol{k}_{\rm t} \tag{39}$$

is fulfilled. The transfer equation for the brightness temperature of the emission has the form

$$\frac{{\rm d}T_{\rm b}}{{\rm d}l} = \alpha - (\mu_N + \mu_{\rm c})\,T_{\rm b}. \tag{40}$$

Here α is the emission coefficient, μ_N is the absorption coefficient related to the decay of an electromagnetic wave of the frequency $2\omega_{\rm p}$ into two plasma waves, $\mu_{\rm c}$ is the absorption coefficient due to the absorption of electromagnetic waves by electron-ion collisions inside the source of radiation, and l is the coordinate along the ray propagation.

If the source has a steady inhomogeneous distribution of the plasma density n with a characteristic scale height $L_n = |n/({\rm d}n/{\rm d}l)|$, the integration of Eq. (40) should be carried out through a thin layer $\Delta l \ll L$ in which the frequency of the elctromagnetic wave is approximately constant, i. e., $\omega[l, k(l)]+\omega'[l, k'(l)]$ = const. The depth of this layer is (Zaitsev & Stepanov 1983):

$$\Delta\,l \approx 3\,L_n\,\frac{v_T^2}{\omega_{\rm p}^2}\,(k_{\rm max}^2 - k_{\rm min}^2) \approx 6\,L_n\,\frac{k^2\,v_T^2}{\omega_{\rm p}^2}, \tag{41}$$

where $k_{\rm max}$ and $k_{\rm min}$ are the maximum and minimum values of the wave number in the excited wave spectrum, respectively. In the second part of of Eq. (41) we took $k_{\rm max} - k_{\rm min} \approx k$, where k is the average wave number of the plasma waves.

The coefficients α and μ_N for the process obeying Eq. (39) have the following form (Zheleznyakov 1996):

$$\alpha \approx \frac{(2\pi)^5}{15\sqrt{3}}\,\frac{c^3}{\omega_{\rm p}^2\langle v_{\rm ph}\rangle}\,\frac{w_{\rm L}^2}{\xi^2}\,nT, \tag{42}$$

$$\mu_N \approx \frac{(2\pi)^2}{15\sqrt{3}}\,\frac{\omega_{\rm p}}{\langle v_{\rm ph}\rangle}\,\frac{w_{\rm L}}{\xi}, \tag{43}$$

where ξ characterizes the width of the spectrum of the excited plasma waves, $(\Delta\boldsymbol{k})^3 = \xi\,\omega_{\rm p}^3/c^3$, and

$$w_{\rm L} = \frac{W_{\rm L}}{n\kappa T}. \tag{44}$$

In our case we have $(\Delta\boldsymbol{k})^3 = 2\pi\int k_\perp dk_\perp dk_\parallel = \pi k_\perp^2\Delta k_\parallel$. Taking $\Delta k_\parallel \approx \frac{1}{3}\,k_\perp$, we obtain $(\Delta\boldsymbol{k})^3 \approx k_\perp^3$. Maximum growth of the plasma waves occurs at the wave number $k_\perp = (2/3)^{1/2}\,(\sigma-1)^{1/2}\,\omega_{\rm p}/v_1$; therefore we have

$$\xi = \frac{c^3}{v_1^3}\,\left(\frac{2\,(\sigma-1)}{3}\right)^{3/2} \tag{45}$$

which yields $\xi = 7\times 10^2$ for $v_1 = 5.84\times 10^9$ cm s^{-1} (ε_1 = 10 keV) and $\sigma = 3$.

It can easily be seen that collisional absorption at the second harmonic with the absorption coefficient $\mu_{\rm c} = \nu_{\rm ei}/(2\sqrt{3}c)$ within the limit of a layer $\Delta\,l$ as given by formula (41) does not give a significant contribution to the general absorption coefficient. Therefore, integrating Eq. (40) inside a layer Δl under the assumption $\mu_{\rm c} \ll \mu_N$, we obtain at $\nu \approx 2\nu_{\rm p}$

$$T_{\rm b}^{\rm (source)} \approx (2\pi)^3\,\left(\frac{c}{\omega_{\rm p}}\right)^3\,\frac{w_{\rm L}nT}{\xi}\,\left[1 - \exp(-\tau_N)\right], \tag{46}$$

where

$$\tau_N \approx \mu_N\,\Delta l \approx 10^5\,w_{\rm L} \tag{47}$$

is the optical depth of the layer Δl for the decay of electromagnetic waves of the frequeny $\omega = 2\omega_{\rm p}$ into two plasma waves under the assumption that the scale height is $L_n = 10^9$ cm. Eq. (46) yields, together with (45) and (47), for $\nu_{\rm p}$ = 0.5 GHz, $n = 10^9$ cm^{-3}, and $T = 7\times 10^5$ K the required brightness temperature in the source region $T_{\rm b}^{\rm (source)}(\nu \approx 2\nu_{\rm p}) \approx 10^{13}$ K if the energy density of the waves is $w_{\rm L} \approx 10^{-5}$.

As it was shown in Sect. 3 [cf. Eq. (33)], the generation of plasma waves by fast electrons with a power-law energy spectrum leads to an energy level $w_{\rm L} \approx 6\times 10^{-2} n_1/n$. Hence the required brightness temperature in the source can be explained if $n_1/n \approx 1.5\times 10^{-4}$. For $n = 3\times 10^9$ cm^{-3} this corresponds to a density of fast electrons of $n_1 \approx 4.5\times 10^5$ cm^{-3}. If the source volume is $V \approx L_n^2 L_\parallel \approx 10^{28}$ cm^3 (by taking $L_n \approx 10^9$ cm and $L_\parallel \approx 10^{10}$ cm) the total number of fast electrons injected into the source region of the decimetric continuum amounts to $N_1 = n_1 V \approx 4.5\times 10^{33}$.

4.3. Polarization of the decimetric continuum

Commonly the polarization of the decimetric continuum is believed to correspond to the ordinary wave mode (cf., e.g., Zheleznyakov 1970, Kuijpers 1980). However, broadband decimetric pulsations were found to be polarized in the extraordinary sense according to the leading spot hypothesis (Aschwanden 1986).

We will suppose that the plasma waves are generated within a certain cone of angles with a half-width ψ_0, where the direction of the cone differs slightly from the direction perpendicular to the magnetic field:

$$W_{\boldsymbol{k}} = \begin{cases} W_1/\sin\psi_0 & \text{for} \quad \psi < \psi_0 \\ 0 & \text{for} \quad \psi > \psi_0 \end{cases} \tag{48}$$

Then, transforming the results of Zlotnik (1981) to our case, we obtain the following expression for the degree of polarization:

$$\rho = \frac{F_x - F_o}{F_x + F_o} = \frac{\omega_g}{\omega_p}\left[P(\theta,\psi) - \frac{5}{16}\cos\theta\right], \quad (49)$$

where F_x and F_o are the observed fluxes of the extraordinary and ordinary wave modes, respectively, and θ is the angle between the direction of wave propagation and the magnetic field. Thus, the cases $\rho > 0$ and $\rho < 0$ correspond to the polarization of the extraordinary and ordinary waves, respectively. For our selected plasma-wave spectrum given in Eq. (48), the function $P(\theta,\psi_0)$ in Eq. (49) has the following form:

$$P(\theta,\psi_0) = \frac{1}{\cos\theta}\frac{P_1}{P_2} \quad (50)$$

with

$$\begin{aligned} P_1 &= \frac{8}{15}\cos^2\theta - \cos^4\psi_0\sin\psi_0 \\ &\times\left(\frac{19}{12}\sin^4\theta - \frac{37}{15}\sin^2\theta + \frac{4}{5}\right) \\ &-\frac{4}{16}\cos^2\theta\,(1-\sin\psi_0)^2\,(2+\sin\psi_0), \end{aligned} \quad (51)$$

$$\begin{aligned} P_2 &= \frac{4}{15} - \sin\psi_0\cos^2\psi_0 \\ &\times\left(\sin^2\theta\cos^2\theta - \frac{1}{4}\sin^2\theta\cos^2\psi_0\,(1+7\cos^2\theta)\right. \\ &\left.+\frac{2}{3}\cos^2\psi_0\right) - \frac{2}{15}(1-\sin\psi_0)^2(2+\sin\psi_0)\,. \end{aligned} \quad (52)$$

In the simplest case of a narrow plasma-wave spectrum, where $\psi_0 \to 0$, the degree of polarization becomes

$$\begin{aligned} \rho(\theta,\psi_0\to 0) &= \frac{\omega_g}{\omega_p}\left[-\frac{3.96}{\cos\theta}(\sin^4\theta - 0.88\cos^2\theta)\right. \\ &\left.-\frac{5}{16}\cos\theta\right]. \end{aligned} \quad (53)$$

We find that the sense of polarization corresponds to the ordinary wave ($\rho < 0$) for angles $\theta > 50°$. Here, for $\theta = 60°$, the degree of polarization becomes $\rho = -2.75\,\omega_g/\omega_p$. For angles $\theta < 50°$ the polarization corresponds to the extraordinary mode. However, for sources of radio emission concentrated in magnetic traps, as in our case, the observation of a major part of the source under angles θ near $\pi/2$ is very likely. This can mean that, regardless of a part of the source (at the roots of magnetic loops) seen at angles $\theta < 50°$ and hence yielding extraordinary polarization, the main polarization of the decimetric continuum corresponds to the ordinary wave mode. Significant extaordinary polarization requires strongly asymmetric source locations inside the trapping loop close to one footpoint (Aschwanden 1986).

5. Discussion

We have shown that fast flare electrons with a power-law energy distribution trapped in magnetic loops are generating plasma waves at the upper hybrid frequency which is here considered as the source of the solar decimetric continuum. Evidently this source is located near a flaring loop and its magnetic field should be sufficiently weak satisfying the condition $\omega_p^2 \gg \omega_g^2$. This condition is necessary in order to prevent a strong gyroresonance absorption at the layers $2\omega_g$ and $3\omega_g$ at the escape of the radiation from the source region.

The absorption of the decimetric continuum due to free-free transitions in the corona at $\omega \approx \omega_p$ is rather high, which favours emission at the second harmonic of the plasma frequency $\omega \approx 2\omega_p$. This emission turns out to be polarized in the ordinary sense within a wide cone of angles θ between the (perpendicular) magnetic field and the direction to the observer.

The direct vicinity of the source of the decimetric continuum to the flare loop (or system of flare loops, cf. Fig. 3) allows to explain the good temporal and spatial correlation between the decimetric continuum and the microwave bursts. This circumstance allows to conclude that both components are feeded by one and the same source of fast electrons originating during the flare process.

The main difference between our results and the conclusion by Benz & Kuijpers (1974) is connected with the problem of the use of a distribution function with a sharp boundary of the loss cone. A loss cone with a smooth boundary appears physically more realistic and gives the main contribution to the instability since it provides a part of a positive derivative $\partial f/\partial v_\perp^2 > 0$ in Eq. (1) for the growth rate of the instability.

In order to obtain an instability for plasma waves from fast flare-electrons with a power-law energy spectrum, Benz & Kuijpers (1974) investigated the deformation of the initial distribution function by collisions of the fast electrons with particles of the background plasma. According to their estimations the time necessary for producing the instability is about 37 s for plasma parameters corresponding to the source region of the decimetric continuum. In our case, collisions of fast electrons with particles of the background plasma should not have an essential influence on the instability because the characteristic time of the formation of the loss cone for the distribution function of the fast electrons after their injection into the trap is of the order $L_\parallel/v \approx$ (1–2) s. This caracteristic time was found using the following values for the characteristic length of the magnetic trap $L_\parallel$ and the mean velocity of the fast electrons v: $L_\parallel \approx 10^{10}$ cm and $v \approx 5\times10^9$–10^{10} cm s^{-1}. The characteristic time of the development of the instability is $\gamma_{\max}^{-1} \approx 5\times10^{-6}$ s [cf. Eq. (16)]. After the generation of the instability a further evolution of the distribution function takes place due to the interaction of the electrons with the plasma waves as a quasilinear effect, which is a more rapid process than the collisions of the fast electrons with the particles of the background plasma.

We estimated the total number of fast electrons needed for the generation of the decimetric continuum with an observed brightness temperature $T_b^{(obs)} \approx 10^{10}$ K to be of order $N_L \approx$

4.5×10^{33}. In order to compare this quantity with the number of electrons necessary for the generation of the related microwave burst, one can use the formula for the maximum frequency of the spectrum of gyrosynchrotron radiation from electrons with a power-law energy distribution (Dulk & Marsh 1982):

$$\nu_{\rm peak} \approx 2.72 \times 10^3 \left[10^{0.27\delta} (\sin\theta)^{0.41+0.03\delta} \times (n_1 L)^{0.32-0.03\delta} B^{0.68+0.03\delta}\right] . \quad (54)$$

From Eq. (54) follows that for $\nu_{\rm peak}$ = 5 GHz, B = 500 G, and δ = 5 we obtain $n_1 L \approx 2.4 \times 10^{16}$ cm^{-2}. Taking a characteristic size of the microwave-burst source $L \approx 10^9$ cm we obtain $n_1 \approx 2.4 \times 10^7$ cm^{-3} and the total number of fast electrons $N_1 \approx n_1 L^3 \approx 2.4 \times 10^{34}$. The estimation of the number N_1 made in Sect. 4.2 for the solar decimetric continuum is of the order of 20 % of this quantity. Hence a non-negligible part of the fast flare electrons should be injected into the source region of the decimetric continuum to provide brightness temperatures up to 10^{10} K, as observed.

The formation of an instability of plasma waves inside the source of the microwave burst at a temperature of the background plasma $T \approx 5 \times 10^6$–10^7 K appears less likely because this would require a relatively high density of fast electrons $n_1/n > 10^{-3}$, although for sufficiently large events this possibility cannot be excluded.

6. Conclusions

We have considered the possibility of an instability produced by fast flare-electrons with a power-law energy spectrum generating plasma waves at the upper hybrid frequency due to an anisotropy resulting from a loss-cone in a magnetic trap.

The resulting radiation can reach the observed brightness temperatures of the solar decimetric continuum under the conditions that the continuum corresponds to the second harmonic of the plasma frequency and that a sufficiently large number of fast electrons is injected into the source region which is estimated to be ~ 20 % of the number of fast electrons inside the source of the related microwave burst.

The polarization of the emission of the second harmonic corresponds to the ordinary wave mode for a wide cone of propagation angles around the direction perpendicular to the magnetic field.

Acknowledgements. The authors are indebted to Drs. L. Vlahos, H. Auraß, G. Mann, and A.V. Stepanov for useful discussions and to the unknown referee for helpful comments. One of the authors (V.V.Z.) would like to thank for the hospitality received during his stay at the Astrophysical Institute Potsdam. The work was supported by the Russian Foundation of Basic Research, by INTAS, and by the Deutsche Agentur für Raumfahrtangelegenheiten (DARA grants 50 QL 9208, 50 QL 9301, 50 QL 9602).

References

Abramowitz, M., Stegun, I.A. (eds.), 1964, Handbook of Mathematical Functions. National Bureau of Standards, p. 88
Aschwanden, M.T., 1986, Solar Phys. 104, 57
Aschwanden, M.T., Benz A.O., Kane S.R., 1990, A&A 229, 206
Aschwanden, M.T., Güdel, M., 1992, ApJ 401, 736
Aschwanden, M.T., Benz, A.O., 1995, ApJ 438, 997
Benz, A.O., 1980, ApJ 240, 892
Benz, A.O., 1985, Solar Phys. 96, 357
Benz, A.O., 1986, Solar Phys. 104, 99
Benz, A.O., 1996, in: Taylor A.R., Paredes J.M. (eds.) Radio Emission from the Stars and the Sun. ASP Conference Series 93, 347
Benz, A.O., Kuijpers, J., 1976, Solar Phys. 46, 275
Benz, A.O., Magun, A., Stehling, W., Su, H., 1992, Solar Phys. 141, 335
Bespalov, P.A., Trakhtengerts, V.Yu., 1986, Rev. Plasma Physics 10, 155
Bespalov, P.A., Zaitsev, V.V., Stepanov, A.V., 1991, ApJ 374, 269
Dulk, G.A., Marsh, K.A., 1982, ApJ 259, 350
Fleischer, R., Oshima, M., 1961, AJ 66, 43
Isliker, H., Benz, A.O., 1994, A&AS 104, 145
Kuijpers, J., 1974, Solar Phys. 36, 157
Kuijpers, J., 1980, in: Kundu M.R., Gergely T.E. (eds.) Radio Physics of the Sun. 93, 347
Krishnan, T., Mullaly, R.F., 1961, Nat 1922, 58
Krüger, A., 1979, Introduction to Solar Radio Astronomy and Radio Physics. D. Reidel Publishing Company. Dordrecht, Boston, London
Kundu, M., 1961, ApJ 134, 96
Kundu, M., 1965, Solar Radio Astronomy. Interscience Publishers. New York, London, Sydney
Kundu M.R., Firor, J.W., 1961, ApJ 134, 389
Kundu, M.R., Vlahos, L., 1982, Space Sci. Rev. 32, 405
Lin, R.P., 1974, Space Sci Rev. 16, 189
Mikhailovskii, A.B., 1974, Theory of Plasma Instabilities. Consultant Bureau, New York
Shaposhnikov, V.E., 1988, Pisma v Astronomicheskii Zhurnal 14, 644 (in russian)
Tsuneta, S., 1996, ApJ 456, 840
Zaitsev, V.V., Stepanov, A.V., 1975, A&A 45, 135
Zaitsev, V.V., Stepanov, A.V., 1983, Solar Phys. 88, 297
Zheleznyakov V.V., 1970, Radio Emission of the Sun and Planets. Pergamon Press, Oxford
Zheleznyakov, V.V., 1996, Radiation in Astrophysical Plasmas. Kluwer Academic Publishers, Dordrecht, Boston, London
Zlotnik, E.Ya. 1981, A&A 101, 250

This article was processed by the author using Springer-Verlag LaTeX A&A style file *L-AA* version 3.