MAGNETOSPHERIC AND SOLAR RADIATION FROM LOSS-CONE DRIVEN INSTABILITIES

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ABSTRACT

Loss-cone driven instabilities are thought to generate certain intense planetary and solar radio emissions. If the electron cyclotron frequency Ω_e and the plasma frequency ω_p satisfy $\Omega_e/\omega_p \gtrsim 1$, electron cyclotron maser emission produces intense radiation, predominantly in the magnetoionic x mode. However if $\Omega_e/\omega_p \ll 1$, plasma emission occurs, generating Langmuir waves which can interact with other waves or coalesce to produce weakly polarized escaping radiation. The theories of these two coherent emission mechanisms are reviewed. Extensions of the basic electron cyclotron maser theory to take into account inhomogeneities in the source region and to provide explanations for the fine structures in the emissions are also discussed.

Keywords: Loss-cone driven instabilities, Electron cyclotron maser emission, Inhomogeneities and fine structures, Plasma emission

1. INTRODUCTION

Nonthermal emission processes in astrophysical and geophysical plasmas may be classified as "incoherent" or "coherent". Incoherent emission is due to spontaneous emission by suprathermal particles. Self-absorption by the emitting particles limits the brightness temperature T_b of any escaping radiation to a value less than about the energy ε of the emitting particles (1 eV $\cong 10^4$ K). Thus, for example, radiation with $T_b \gtrsim 10^{10}$ K from nonrelativistic electrons ($E \ll 10^6$ eV) cannot be due to any incoherent emission process, and so must be due to a coherent mechanism. In any coherent emission at least one stage (if it is a multistage process) must involve an instability of some form. Instabilities can be classified as "maser" which are due to negative absorption, or as "reactive", which involve some form of phase bunching process. Maser instabilities are driven by gradients in momentum space of the distribution function $f(p)$ of the radiating electrons, and are thought to be-the more relevant of the two in most applications outside of the laboratory. Here we are concerned with coherent emission processes which involve a maser instability.

The most widely studied forms of coherent emission are "plasma emission" in solar radio bursts and electron cyclotron maser emission (ECME). These mechanisms are quite different. Plasma emission is a multistage process in which the instability produces Langmuir waves at frequencies close to the plasma frequency ω_p, and nonlinear processes in the plasma partially convert these Langmuir waves into escaping radiation at $\cong \omega_p$ and/or $\cong 2\omega_p$. ECME, however, is a single stage process in which cyclotron absorption is negative, resulting in direct maser emission of escaping radiation near the electron cyclotron frequency Ω_e, or perhaps its harmonics. Plasma emission is favoured in weakly magnetized plasmas ($\Omega_e \ll \omega_p$), and ECME is effectively possible only for $\Omega_e \gtrsim \omega_p$.

Plasma emission and related processes occur in the solar corona, the interplanetary plasma, especially near shock waves including planetary bow shocks, and within planetary magnetospheres, especially near the edges of denser plasma regions such as the plasmasphere and the Io plasma torus (Refs. 1a,2,3). In contrast the necessary condition $\Omega_e \gtrsim \omega_p$ for ECME is satisfied only in a few regions such as the auroral zones of the Earth, the inner regions of the magnetospheres of Jupiter and Saturn and in the solar corona above active regions. It is just these regions which lead to radiation with $T_b \gg 10^{10}$ K and seem to require that the ECME mechanism operates. The corresponding emissions are the Earth's auroral kilometric radiation (AKR, Ref. 4), the Jovian decametric radiation (DAM, Ref. 5). Saturn's kilometric radiation (SKR, Ref. 3) and solar decimetric spike bursts (Ref. 6). To these should be added emissions with $T_b \gtrsim 10^{10}$ K from certain types of flare stares (Ref. 7).

2. SUCCESSES AND LIMITATIONS OF THE ECME THEORY

The first ECME theory was proposed by Twiss in 1958 (Ref. 8) and applied by him to the interpretation of solar radio bursts (Ref. 9). Neither this nor later ECME theories (e.g. Refs. 10,11) for metre-wave bursts have become widely accepted. It has long been thought that DAM is due to ECME (Ref. 12) and a detailed ECME theory was proposed for it and for AKR by Melrose in 1976 (Ref. 13). This theory had two notable successes. First it requires $\Omega_e > \omega_p$ in the source region, and soon after it was proposed, Benson and Calvert (Ref. 14) showed that auroral cavities with $\Omega_e > \omega_p$ exist and that AKR originates in such cavities. Second the theory implies emission in the x mode of magnetoionic

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theory at frequencies close to Ω_e and it is found that AKR is indeed predominantly x mode at the appropriate frequencies (e.g. Ref. 15). The theory requires, however, an extreme type of pitch angle anisotropy and the so-called "inverted V" electrons, identified as the generators of AKR, do not have an anisotropy of the required form.

In 1979 Wu and Lee (Ref. 16) proposed an alternative form of ECME which could be driven by a loss-cone anisotropy. The observed distributions of reflected inverted V electrons do have loss-cone features and detailed numerical calculations (Refs. 17,18), which imply that these are capable of producing growth of the maser instability, provide strong support for this theory. The same theory applied to DAM provides a natural explanation (Ref. 19) for the inferred angular pattern of emission - on the surface of a wide-angled hollow cone (e.g. Ref. 20). The Wu and Lee form of ECME has become the accepted one in astrophysical and geophysical applications.

One of the strengths of a maser theory is that it is independent of detailed phase relations, making it possible to develop a semi-quantitative treatment using quasilinear theory. However this is also a weakness in that both DAM and AKR exhibit fine structures (e.g. Refs. 21,22) which cannot be explained by the ECME theory, at least in its simplest form. The present ECME theory should be regarded as one for the averaged (over the fine structures) properties of the radiation, to be supplemented with a detailed phase-coherent theory for the fine structures. In the following sections the phase-averaged theory is presented in detail and several qualitative comments are made on the possible form of the phase-coherent interactions which may be involved.

A related difficulty (Ref. 18) is that the timescale over which the development of the instability would cause the electron distribution to be significantly modified is much less than the time required to measure the distribution function. This implies that the features which actually drive the instability can not be observed directly. Moreover whatever these features are they must be generated on a comparably short timescale, and at present there is no theory for this.

3. ELECTRON CYCLOTRON MASER THEORY

The theory and applications of ECME have recently been reviewed by Wu (Ref. 23). The following presentation is based on a discussion by Melrose (Ref. 24).

The absorption coefficient $\gamma_M(\underset{\sim}{k})$ for waves in a mode M with wavevector $\underset{\sim}{k}$ is of the form

$$\gamma_M(\underset{\sim}{k}) = -\sum_{s=-\infty}^{\infty}\int d^3\underset{\sim}{p}\; w_M(\underset{\sim}{k},\underset{\sim}{p},s)\; \hbar\left(\frac{s\Omega}{v_\perp}\frac{\partial}{\partial p_\perp} + k_\parallel\frac{\partial}{\partial p_\parallel}\right) f(\underset{\sim}{p}) \qquad (1)$$

where the electrons are described by their parallel and perpendicular momentum components $p_\parallel = \gamma m v_\parallel$ and $p_\perp = \gamma m v_\perp$ with distribution function $f(\underset{\sim}{p})$. The (positive definite) probability function $w_M(\underset{\sim}{k},p,s)$ contains a δ-function requiring that the resonance condition

$$\omega - s\Omega - k_\parallel v_\parallel = 0 \qquad (2)$$

be satisfied with $\omega = \omega_M(\underset{\sim}{k})$ determined by the dispersion relation for the mode M and with $\Omega = \Omega_e/\gamma$ where Ω_e is the non-relativistic gyrofrequency and $\gamma = (1 - v_\parallel^2/c^2 - v_\perp^2/c^2)^{-\frac{1}{2}}$ is the Lorentz factor. In practice, only one of the terms in the infinite sum over s in Eq. 1 (usually s = 1 or 2) makes a significant contribution to $\gamma_M(\underset{\sim}{k})$.

The kinematic restrictions on ECME are determined by Eq. 2 and the relevant dispersion relation. For given ω, $k_\parallel$ and Ω_e the equation describes a "resonance ellipse" in $v_\parallel - v_\perp$ space, whose centre lies on the line $v_\perp = 0$; three examples of these are shown in Figure 1. Only the portions of the ellipses lying in a region $v \lesssim v_0 \ll c$, where there are significant numbers of electrons in the distribution $f(p)$, contribute appreciably to $\gamma_M(\underset{\sim}{k})$. It is relevant to distinguish between two cases, the non-relativistic and the semi-relativistic. In examples a and b of Figure 1, it is reasonable to use the non-relativistic approximation $\gamma = 1$ in Eq. 2 which then describes a straight line parallel to the $v_\perp$ axis at $v_\parallel = (\omega - s\Omega_e)/k_\parallel$. In example a where $k_\parallel^2c^2 \gg \omega^2$ the resonance ellipse has a large eccentricity and becomes unphysical after touching the boundary $v = c$ at $v_\parallel = \omega/k_\parallel$. In example b where $k_\parallel^2c^2 \cong \omega^2$ the resonance ellipse is nearly semi-circular and has a large radius and a centre located well away from the origin. In example c, however, it is not sensible to use the non-relativistic approximation. In this semi-relativistic case where $k_\parallel^2c^2 \ll \omega^2$ the resonance ellipse is nearly semi-circular with radius less than v_0, and has a centre located within v_0 of the origin. The importance of the semi-relativistic case in this context was first recognized by Wu and Lee in 1979 (Ref. 16).

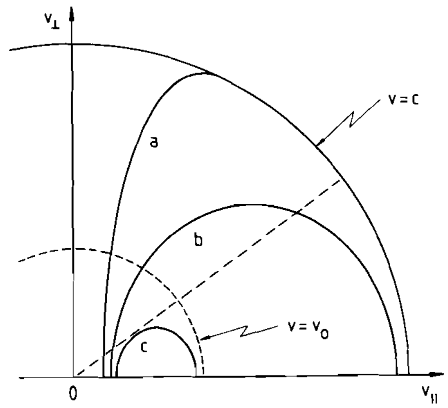

Figure 1. Resonance ellipses in $v_\parallel - v_\perp$ space for (a) $k_\parallel^2c^2 \gg \omega^2$, (b) $k_\parallel^2c^2 \cong \omega^2$ and (c) $k_\parallel^2c^2 \ll \omega^2$. The portion of the ellipse which is not shown in case (a) is unphysical. The dashed line indicates the edge of a loss-cone feature in the electron distribution function. With this source of free energy for the instability, growth is possible only in case (c) where the resonance ellipse lies entirely within the loss cone. The ellipse in this case is approximately semi-circular with radius v_R and centre located at $v_\parallel = v_c$, $v_\perp = 0$.

It may readily be shown (e.g. Ref. (18) that in the non-relativistic case the term involving the $p_\perp$-derivative in Eq. 1 leads to a positive semidefinite contribution to $\gamma_M(\underset{\sim}{k})$. Thus in this case any growth (i.e. negative absorption) can be attributed entirely to the term containing the $p_\parallel$-derivative and may be called parallel driven. This is the case discussed in 1976 by Melrose (Ref. 13), who found that growth requires streaming electrons with average values $\langle (v_\parallel - \langle v_\parallel \rangle)^2 \rangle \ll \langle v_\perp^2 \rangle$. It is impossible in principle for a loss-cone feature to cause such a parallel-driven instability. (In Refs. 25 and 26 it was erroneously predicted that such an instability would occur; this error was due to an inconsistency between the signs of s in the authors' versions of Eqs. 1 and 2.) There are however features in the observed distributions of inverted V electrons which could drive this instability - "bumps" due to trapped electrons and downgoing "holes" (Ref. 27).

In the semi-relativistic case where $k_\parallel$ is required to be small, the term involving the $p_\parallel$-derivative in Eq. 1 is small and only the term with the $p_\perp$-derivative need be considered to a first approximation. A perpendicular-driven instability can arise due to a loss-cone feature: if the resonance ellipse lies entirely within the loss cone, so that $\partial f/\partial p_\perp$ is positive at every point on the ellipse, then Eq. 1 necessarily implies negative absorption (see Fig. 1). A loss-cone feature arises naturally when precipitating electrons are partially reflected by a magnetic field whose strength increases with decreasing height; electrons with small pitch angles are absorbed by the lower atmosphere while those with large pitch angles are reflected to form an upward directed loss cone. Most discussions of ECME are based on the approximation that the dispersive properties of the waves are determined by a cold, relatively dense background plasma in the source region. If the values of ω_p and Ω_e are fixed, $k_\parallel$ and hence the resonance ellipse for a given wave mode and harmonic s are functions only of ω and 8. The range of resonant ellipses which lie within the loss cone therefore determines the ranges of frequencies and angles at which waves can grow.

Figure 2 shows the region in ω/Ω_e - θ space for which growth is possible for the s = 1 x mode when $\omega_p/\Omega_e = 0.1$. The right hand boundary of this region is determined by the condition for a resonance ellipse to exist, i.e. $\omega^2 < c^2k_\parallel^2 + s^2\Omega_e^2$, which follows directly from Eq. 2. The resonance ellipses vanish on this boundary curve and no resonance is possible to the right of it for $\theta < 90^\circ$, i.e. for upward propagating waves. The left hand boundary is determined by the upward directed loss-cone source of free energy. The ellipses for values of ω/Ω_e and θ to the left of this boundary are too large to fit into the loss cone and so damping rather than growth occurs. The growth region near the x-mode cutoff is extremely narrow since the refractive index of the mode varies rapidly with frequency in the neighbourhood of the cu-toff. The region of maximum growth is localized near $\omega/\Omega_e \sim 1.04$ and $\theta \sim 70^\circ$ implying that emission takes place predominantly within a narrow range of w and 8 (Ref. 28). This provides a plausible explanation for the beaming of DAM into a small angular range about the surface of a hollow cone (Refs. 19,20).

Figure 2 also shows the trajectories of three rays (each with fixed ω) in ω/Ω_e - θ space, calculated under the assumptions that the magnetic field is vertical and that ω_p and Ω_e decrease linearly with altitude. The highest frequency ray does not enter the resonance region so no growth is possible. The lowest frequency ray rapidly passes through the region of growth and is then strongly damped before emerging from the source region. The intermediate frequency ray however passes through the region of maximum growth and does not undergo damping. The emergence angle of this ray ($\sim 62^\circ$) differs only slightly from the value of θ where the wave is growing most rapidly ($\sim 70^\circ$).

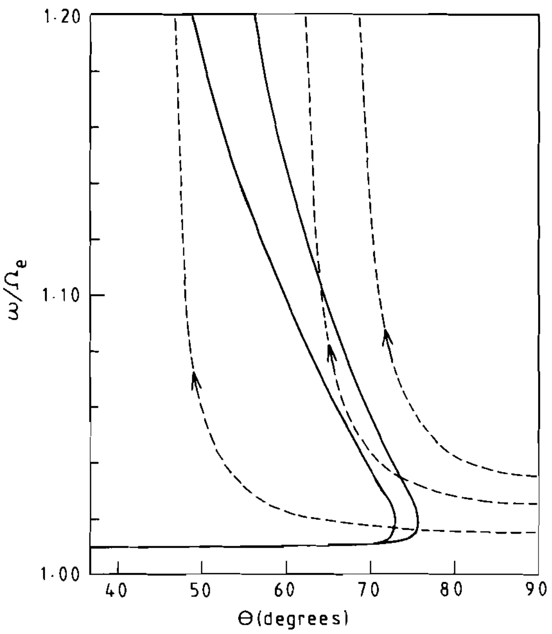

Figure 2. Boundary curves (solid) plotted in ω/Ω_e - θ space for the growth of upgoing waves in the s = 1 x mode for $\omega_p/\Omega_e = 0.1$; the instability is driven by an upgoing loss-cone distribution of energetic electrons. Resonance cannot occur in the region between the right hand boundary and its complement (i.e. $180^\circ - \theta$) for downgoing waves. Damping rather than growth occurs to the left of the left hand boundary curve and in the remaining part of the space for $\theta > 90^\circ$. The dashed curves show the trajectories in ω/Ω_e - 8 space of three waves each with a fixed ω, propagating in a stratified source region. The magnetic field is assumed to be vertical and both ω_p and Ω_e to decrease linearly with altitude.

The magnitude of the growth rate (for either parallel or perpendicular-driven cases) can be quite large. The maximum value for a loss-cone driven maser in the s = 1 x mode is

$$|\gamma_{max}| \simeq \pi \frac{\omega_p^2}{\Omega_e} \frac{n_H}{n_C} \frac{c^2}{v_R v_c} \qquad (3)$$

where v_R and v_c are defined in the caption to Figure 1 and where n_H/n_C is the ratio of the number densities of the energetic and background electrons (Ref. 18). The ECME mechanism can also produce radiation in the s = 1 o mode, in the z mode, and at higher harmonics. At low plasma densities however the s = 1 x mode is dominant; its growth rate is typically an order of magnitude larger than that of the s = 1 o mode. As ω_p/Ω_e is increased, the growth region shown in Figure 2 moves rapidly to higher frequencies and smaller angles and the s = 1 x mode is strongly suppressed for $\omega_p/\Omega_e \gtrsim 0.3$ (Ref. 28).

4. INHOMOGENEITIES AND FINE STRUCTURES IN ECME

Two areas of continuing theoretical interest in the ECME theory for AKR and DAM concern the effects of inhomogeneities in the source regions and the interpretation of fine structures in the emissions.

Due to the gradient in B with height, the resonance condition in Eq. 2 changes due to the change in $\Omega_e \propto B$ with height, and also to the changes in $k_\parallel$ (as a result of refraction) and in $v_\parallel$ (which varies such that the quantity $v_\perp^2/B$ is constant and the quantity $v_\parallel^2 + v_\perp^2$ is 2/m times the energy). Resonance is particularly favourable when these changes balance such that the condition

$$d(\omega - s\Omega_e/\gamma - k_\parallel v_\parallel) = 0 \quad (4)$$

is satisfied. The waves and particles which satisfy Eq. 4 remain in resonance for an exceptionally long time and the limiting effect of the inhomogeneity is minimized. However Eq. 4 is not readily satisfied because the rate of change of Ω_e is by far the largest term in most circumstances. Recently Le Quéau, Pellat and Roux (Ref. 29) proposed a model in which the emission occurs just above the cutoff where the refractive index rises rapidly from zero towards unity, and so the gradient in $k_\parallel$ is anomalously large. At any given height the waves which can satisfy Eq. 4 in this manner are in a very restricted frequency range, leading to a natural explanation for fine structures in the emission.

The fine structures in DAM (Ref. 21) and in AKR (Refs. 22,30) are rapidly drifting in frequency. Melrose (Ref. 31) argues for an analogy with Helliwell's phase bunching mechanism for discrete VLF emissions (Ref. 32) and proposed that Eq. 4 can be satisfied by invoking a parallel electric field in the source region. The electric field $E_\parallel$ accelerates the particles causing both $v_\parallel$ and γ to have gradients so that a large term involving $E_\parallel$ appears explicitly when Eq. 4 is evaluated. Moreover in Helliwell's theory ω changes in Eq. 4, due to the wave-particle interaction occurring in a moving interaction region. The characteristic frequency drift from high to low frequency in DAM is then interpreted in terms of an upward drift of the interaction region. There is also observational evidence for a correlation between AKR and $E_\parallel$ (Ref. 30) and this theory provides a possible explanation for this correlation. The observed triggering of AKR by solar radio emissions (Ref. 33) might also be explained naturally by this theory, using the analogy with the well-documented triggering of discrete VLF emissions.

A specific feature of some observed fine structures in AKR is a characteristic frequency separation. Grabbe (Ref. 34) and later Le Quéau et al (Ref. 29) argued for an interpretation involving ion cyclotron waves (ICW) in the source region. In principle, an ICW modulation could be due to a direct role of the ICW in the emission process (Ref. 34) or to an ICW modulation of the electron distribution. The latter idea seems favourable in that the $E_\parallel$ structures in the auroral zones involve an ICW oscillation (e.g. Ref. 35).

5. PLASMA EMISSION

The most familiar form of plasma emission was developed originally for type III solar radio bursts by Ginzburg and Zheleznyakov in 1958 (Ref. 36): Langmuir waves are generated by an electron stream and are converted partially into fundamental ($\cong \omega_p$) radiation by scattering off thermal ions and partially into second harmonic ($\cong 2\omega_p$) radiation by coalescence with other Langmuir waves. The details of the original treatment are now outdated and modern versions invoke a variety of nonlinear plasma phenomena (Refs. 37, 1b). The application to type III bursts continues to be emphasized. However there are several other types of solar radio bursts for which there is no evidence for streaming of electrons, notably type IV emission and more specifically the flare continua (Ref. 1c). A plausible type of theory for such cases involves Langmuir waves generated by trapped energetic electrons with a loss-cone anisotropy.

The growth of Langmuir waves due to specific types of loss cone distribution has been discussed by Benz and Kuijpers (Ref. 38) and by Zaitsev and Stepanov (Ref. 39). A systematic discussion of the unmagnetized case has been given by Hewitt and Melrose (Ref. 40) who included relativistic effects and noted that the resonance condition $\omega - \underset{\sim}{k}.\underset{\sim}{v} = 0$ corresponds to a resonance hyperboloid in momentum space. In addition to the loss cone anisotropy, an important requirement for growth is that the electrons have a "gap" distribution, i.e. that the distribution function $f(p)$ has a maximum at $v = v_0$ and a minimum at $v \ll v_0$. Such a gap can form when an initial trapped distribution with $\partial f(p)/\partial p < 0$ for all p experiences collisional losses. Hewitt and Melrose presented numerical results for idealized but plausible loss-cone distributions $f(p)$, separable in momentum p and pitch angle α. The results showed that the instability grows only when (i) the pitch angle distribution falls off sufficiently rapidly for $\alpha \lesssim \alpha_0$, the loss-cone boundary angle, (ii) the "gap" is well defined, (iii) the initial electron momentum spectrum is relatively hard (e.g. a power law distribution $\propto p^{-\delta}$ with $\delta \lesssim 4$) and (iv) the background electrons are relatively cool so that Landau damping does not dominate. Growth is considerably stronger for double-sided loss-cone distributions, i.e. those with loss-cone boundaries at α_0 and $180^\circ - \alpha_0$. Maximum growth then occurs when k is close to perpendicular to the axis of the loss cones.

The distribution of Langmuir waves generated by electrons with a loss-cone anisotropy therefore contains waves with nearly anti-parallel $\underset{\sim}{k}$'s and these can coalesce directly to produce radiation at $\cong 2\omega_p$ (with $|\underset{\sim}{k}| \cong \sqrt{3}\,\omega_p/c \ll$ the wavenumber of the Langmuir waves). This is unlike the case of a streaming instability where the k's for the Langmuir waves are predominantly parallel to the direction of the stream and an intermediate scattering of some of these waves into the backward (antistreaming) direction is required before coalescence can produce $2\omega_p$ radiation. Thus the loss-cone generated Langmuir waves are more favourable for second harmonic plasma emission. The flare continua plausibly result from trapped electrons radiating in this manner.

This type of plasma emission from a loss-cone distribution is expected to occur for $\Omega_e \ll \omega_p$ whereas ECME occurs for $\Omega_e > \omega_p$. In the intermediate range $\Omega_e \lesssim \omega_p$ a loss-cone distribution is likely to produce z-mode waves and Bernstein waves which can then play the same role as the Langmuir waves in plasma emission.

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