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par
A. ROUX
R. PELLAT
Centre de Physique Théorique,
École Polytechnique
91128 Palaiseau Cedex
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PAR

A. ROUX
R. PELLAT *

Pour le Directeur
le Directeur Adjoint

M. PETIT

L'Ingénieur chargé du départemment ETE

Centre de Physique Théorique, Ecole Polytechnique
91128 Palaiseau Cedex
Coherent Generation of the Terrestrial Kilometric Radiation by Non-Linear Beatings Between Electrostatic Waves

by

A. ROUX
Centre de Recherches sur la Physique de l'Environnement
Centre National d'Etude des Télécommunications
92131 Issy les Moulineaux, France

and

R. PELLAT
Centre de Physique Théorique
Ecole Polytechnique
91128 Palaiseau Cedex, France

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ABSTRACT

The propagation of electrostatic plasma waves in an inhomogeneous and magnetized plasma is studied analytically. These waves, which are driven unstable by auroral beams of electrons, are shown to suffer a further geometrical amplification while they propagate towards resonances. Simultaneously their group velocities tend to be aligned with the geomagnetic field. Then it is shown that the electrostatic energy tends to accumulate at, or near $\omega_{LH}$ and $\omega_{UH}$, the local lower and upper hybrid frequencies. Due to this process, large amplitude electrostatic waves with very narrow spectra should be observed near these frequencies at any place along the auroral field lines where intense beam driven instability takes place. These intense quasi-monochromatic electrostatic waves are then shown to give rise by a coherent non-linear three wave process to an intense electromagnetic radiation. Depending upon the ratio $\omega_{pe}/\omega_{ce}$ between the electron plasma frequency and the electron gyrofrequency the electromagnetic wave can be radiated in the ordinary mode (at $\omega_{UH}$) or in the extraordinary mode (at $2\omega_{UH}$). As far as the ratio $\omega_{pe}/\omega_{ce}$ tends to be rather small it is shown that the most intense radiation should be observed at $2\omega_{UH}$ in the extraordinary mode.
INTRODUCTION

All magnetized planets seem to be intense radio sources, emitting in the decametric range for Jupiter (Warwick 1961, Carr and Desch 1976) in the hectometric range for Saturn (Brown 1975), and in the kilometric range for the Earth (Gurnett 1974, Alexander and Kaiser 1976). The temperature of brightness of these sources are so large that any classical incoherent radiation process is irrelevant. At least for the Earth and Jupiter it is widely admitted that this radio-emission process takes place in the regions of their magnetospheres where intense beams of electrons are known to exist. Thus it is clear that the electromagnetic radiation is directly or indirectly driven by a beam plasma instability. For non-relativistic beams of electrons most of the energy driven by this beam plasma instability is electrostatic. Thus the problem is to convert this electrostatic energy into electromagnetic radiation. The total peak power of Auroral Kilometric Radiation (AKR hereafter) was estimated by Gurnett to be $10^9$ W. which can be 1% of the total energy available during auroral phenomena. As far as the power of electrostatic noise should be less than the power involved in auroral phenomena, it is required that the conversion process have an efficiency of the order of at least 1%. It seems that this rules out any linear conversion process as those proposed by Oya (1974) and Benson (1975). Thus, what is required is a very efficient non-linear conversion process. Some non-linear theory were already proposed,
we will comment briefly on two of them. In the work of Palmadesso et al. (1976) electrostatic high frequency waves, driven unstable by an energetic electron beam, are coupled with ionic fluctuations. The required level of ionic fluctuation is very high ($\delta n_i/n_i = 0.3$). If such high level ionic fluctuations exist, they should be localized in the return current regions as shown by Kindel and Kennerly (1975), that is where low energy upgoing electrons are observed. These return currents are known to take place outside the region where downgoing primary electron beams are observed. Thus the coupling between high frequency electrostatic waves and ionic fluctuations should not be very strong. Furthermore the authors of the quoted works have used in an unstable situation the theory of Dawson and Oberman (1962) which was demonstrated to hold for damped plasma waves. The transposition of this theory to the unstable case seems to be irrelevant. Galeev and Krasnosel'skikh (1976) have proposed a different interpretation. They argued that due to the large level of electrostatic fluctuations at the plasma frequency there is a tendency to cavitation. Then they considered that each such caviton radiates a coherent electromagnetic wave at $2\nu_{pe}$. But, as far as they should be a large number of such cavitons inside an auroral electron beam, one cannot expect that the radiation of this set of cavitons is still coherent. Even so, the maximum efficiency of their process seems not to be large enough to account for the experimental requirements.
Melrose (1976) has discussed a generation mechanism for A.K.R. where escaping electromagnetic waves are directly driven unstable by beam plasma instability. In principle such a mechanism is tempting since it is a direct process, but the obtained threshold condition: \( \beta_{LO}^2 \geq \beta_{f0} \) would imply a very large thermal anisotropy \( T_L / T_e \geq 30 \) in the beam frame. In absence of experimental confirmation, at least theoretical arguments explaining how such a large anisotropy could be achieved over auroral region would be required.

We propose here a different interpretation where the electromagnetic radiation is generated through a coherent beatings between electrostatic waves. The required coherency being due to the large enhancement of the amplitude of the electrostatic waves when they propagate towards resonances.

In what follows we will focus on the A.K.R. The relevance of the present theory to the explanation of other planetary radio-emissions will be examined elsewhere.
PROPAGATION CHARACTERISTICS OF THE PRIMARY ELECTROSTATIC WAVES

According to Green et al (1977), AKR is generated along auroral field lines in the regions where the plasma frequency \( w_{pe} \) is less or of the order of the electron gyrofrequency \( w_{ce} \). Then provided that we have still \( \omega_i \gg \omega_c \) where the subscript \( i \) stands for ions, the dispersion relation for electrostatic waves writes in the zero temperature limit:

\[
D(\omega,k) = 1 - \frac{\omega_{pe}^2}{k^2} \left[ \frac{k^2}{\omega^2} + \frac{k^2}{\omega_c^2} \right] - \frac{\omega_i^2}{\omega^2} = 0
\]  

(1)

As long as \( \omega/k || \gg V_t \) where \( V_t \) is the thermal velocity of the so-called "cold plasma", and \( n_c \gg n_b \), \( n_c \) and \( n_b \) being respectively the cold plasma and beam densities, the expression (1) is valid. Of course very far from the earth where \( n_w \) the warm plasma density can be larger than \( n_c \), thermal effects would be essential giving rise to Bernstein waves having \( \omega \simeq (n + 1/2)\omega_c < \omega_{UH} \). The coherent non-linear process that we will discuss later can also be pertinent in such situations but we will not examine this in the present paper.

Let us discuss now the dispersion relation (1). There are two propagation windows: \( \omega_{LH} < \omega < \min(\omega_{pe}, \omega_{ce}) \) and \( \max(\omega_{pe}, \omega_{ce}) < \omega < \omega_{UH} \). Near from \( \omega_{pe} \) and \( \omega_{ce} \), \( k || \gg k_\perp \). The modulus of the wave number does not intervene in (1) thus \( \mathbf{k} \cdot \mathbf{V_Z} = 0 \); since \( k_\parallel \) is fixed by the resonance condition \( |V_g| \) is well determined, even in the zero temperature limit. Then the introduction of a static magnetic field avoid the usual degeneracy of Langumir waves. From the previous relation we deduce that the group velocity is aligned with the static magnetic field \( \mathbf{B}_0 \) when \( \omega = \omega_{LH} \) or
\( \omega_{\text{UH}} \). On the other hand, \( \mathbf{V}_g \perp \mathbf{B}_o \) when \( \omega = \omega_{\text{pe}} \) or \( \omega_{\text{ce}} \).

In a situation where beam-plasma driven instability takes place only those electrostatic waves which have their parallel phase velocities in the same direction as the beam are amplified. As stated earlier, waves having \( \omega = \omega_{\text{pe}} \) or \( \omega_{\text{ce}} \) essentially propagate across the magnetic field which implies that they rapidly leave the auroral tube of force, that is the beam region. Thus what we will discuss now is the behavior of waves having frequencies far enough from \( \omega_{\text{pe}} \) and \( \omega_{\text{ce}} \). An unstable wave, in the lower frequency window, has its parallel group velocity directed towards the earth, consequently it propagates towards regions where \( \omega_{\text{DH}} \) increases, thus the ratio \( \omega/\omega_{\text{LH}}(Z) \) (\( Z \) being the distance along field lines) decreases. Simultaneously its group velocity becomes more and more aligned with \( \mathbf{B}_o \). On the other hand, an unstable wave having \( \max(\omega_{\text{pe}}, \omega_{\text{ce}}) < \omega < \omega_{\text{UH}} \) propagates outwards since its parallel group velocity is directed away from the earth. Then the ratio \( \omega/\omega_{\text{UH}}(Z) \) decreases while \( \mathbf{V}_g \) becomes more and more aligned with \( \mathbf{B}_o \). Then in both cases the energy of plasma waves propagates towards resonances, while their group velocities tend to be aligned with \( \mathbf{B}_o \).

During this propagation, as long as the WKB approximation holds, \( k_\perp \) is constant and consequently \( k_\parallel(Z) \) has to decrease in order to satisfy locally (1). The medium being essentially stationary it is clear that \( \omega \) is constant. Then, as the plasma waves propagate towards resonance their parallel phase velocities increase while their parallel group velocities decrease (since \( \mathbf{V}_g \perp k_\parallel \)) and simultaneously \( |\mathbf{V}_g \parallel / \mathbf{V}_g \parallel | \) decreases.
Depending upon the ratio \( k_{||}/k_{\perp} \) between their parallel and perpendicular wave numbers, the electrostatic waves can remain inside the beam region or rapidly leave it. Then it is important to evaluate the critical value \( (k_{||}/k_{\perp})_{cr} \) corresponding to the transition between these two regimes. Let us first discuss this problem for the upper propagation window where the ratio \( k_{||}/k_{\perp} \) can change over distances much shorter than it does in the lower one. Starting with a wave having \( \omega, k_{||} = k_{||}(Z_0), k_{\perp} = k_{\perp}(Z_0) \) it will reach resonance at \( Z = Z_0 + \Delta Z \) where \( w_{UH}(Z_0 + \Delta Z) = \omega \). Let \( R_b \) be the radial size of the beam, if \( \Delta Z/V_{||} > R_b/V_{\perp} \) the wave will still be inside the beam when it reaches resonance. This condition also writes \( k_{||} \Delta Z = k_{\perp} R_b \). Expanding \( w_{UH}(Z_0 + \Delta Z) = w_{UH}(Z_0) [1 + \beta \Delta Z] \) where \( \beta \approx 3/R \), and using (1) we find:

\[
\Delta Z = (k_{||}/k_{\perp})^2 \left( \frac{w_{pe}^2 w_{ce}^2 R}{6 w_{UH}^4} \right) \tag{2}
\]

Combining this expression and the relation \( k_{||} \Delta Z = k_{\perp} R_b \) we get:

\[
(k_{\perp}/k_{||}) \left[ (k_{\perp}/k_{||})^2 + 1 \right] = \left( \frac{w_{pe}^2 w_{ce}^2 R}{6 R_b w_{UH}^4} \right) \tag{3}
\]

Setting \( R_b \approx 10^5 \) m and hence \( R_E/R_b \approx 10 \) we obtain for \( w_{ce} = \frac{1}{4} w_{pe} \): \( R = 2 R_E \):

\( (k_{||}/k_{\perp})_{cr} \approx 1 \), which corresponds to \( \omega \approx (w_{ce} + w_{pe})/2 \). Similar calculations performed near \( w_{LH} \) give \( (k_{||}/k_{\perp})_{cr} \approx 2.4 \times 10^{-2} \) which corresponds to a wave starting with \( \omega/w_{LHO} \approx 1.4 \).
LINEAR INSTABILITY OF BEAM DRIVEN PLASMA WAVES

We will now discuss the convective amplification of the electro-static waves which are excited by the auroral electron beam. The problem under discussion in this section was worked out numerically by Maggs (1976). However, as far as he was interested in the production of the auroral hiss he has not discussed the convective amplification in the upper propagation window. Furthermore, for the purpose of our future discussion we need analytical (even though less precise) expressions.

Let us let $E \propto \exp \left( \int k_{\parallel} dz + \hat{k}_{\perp} \cdot \hat{r} - \omega t \right)$, $k_{\parallel} = Jm(k_{\parallel})$ and $G = - \int k_{\parallel} dz$, $k_{\parallel}$ is the spatial growth rate and $G$ the effective gain along a ray path. $f_{b}$ is the electron beam distribution function, it is isotropic and centered around $V_{b}$, its bulk velocity; its thermal velocity is $V_{b} \cdot f_{th}$; $0$, $U_{th}$ are the corresponding quantities for thermal electrons. Provided that $f_{th}(V_{b} - U_{b}) < f_{b}(V_{b} - U_{b})$ one can safely describe the thermal plasma by a Dirac "function". A discussion of the intermediate case where the beam is not completely detached from thermal electrons can be found in Maggs'. Then, assuming a cold plasma background with $n_{c} >> n_{b}$ we obtain by using standard methods the contribution of the Cerenkov pole to the growth rate:

$$k_{\parallel} \sim + \frac{\pi}{n_{c} k_{\parallel}} \frac{w_{pe}^{2}}{w_{\perp}^{2}} \frac{w^{2}}{w_{\perp}^{2}} \left\{ \int dV_{\perp} \frac{2}{J_{0}^{2}} \frac{k_{\perp} V_{\perp}}{w_{ce}} \frac{\delta f_{b}}{\delta V_{\parallel}} \right\} V_{\parallel} = \frac{w}{k_{\parallel}} \tag{4}$$
where $J_0$ is a Bessel function. When $k_\perp \gg \omega_{ce}/U_b$ the contribution to the growth is very small, so that we have restricted ourselves to the case $k_\perp \ll \omega_{ce}/U_b$ which will be checked a posteriori. Then $J_0^2 \approx 1$, and the expression (2) is very simple. What we have to do now is to evaluate $\int |k_\parallel| \, dz$. Let us look first at the upper propagation window and assume that $\omega_{ce} > \omega_{pe}$. For the other cases: $\omega_{ce} < \omega_{pe} < \omega < \omega_{UH}$ and $\omega_{UH} < \omega \ll \min(\omega_{pe}, \omega_{ce})$ the calculations are very similar and we will solely mention the results. From (1) we get:

$$\frac{w^2}{k_\parallel(Z)} = \frac{w^2 - \omega_{ce}^2(Z)}{w_{UH}^2(Z) - w^2}$$  \hspace{1cm} (5)

If $\omega_{pe}$ is significantly less than $\omega_{pe}$, then, $(w^2 - \omega_{pe}^2)/\omega_{ce}^2 \approx 1 \gg (w^2 - \omega_{ce}^2)/\omega_{ce}^2$ or $(w^2 - \omega_{UH}^2)/\omega_{ce}^2$ these last two quantities being less than $w_{pe}^2/2\omega_{ce}^2$. Then the variation of $w_{pe}^2$ can be neglected. Expanding the gyrofrequency $\omega_{ce}(Z) = \omega_{ceo} [1 + \beta(Z-Z_0)]$ we get from (5):

$$\frac{dx}{dV_{\parallel}} = \frac{V_{\parallel}}{2\beta} \frac{w_{pe}^2}{\omega_{ce}^2} \frac{w^2}{[w^2 + V_{\parallel}^2]^2}$$  \hspace{1cm} (6)

Combining (2) and (4) and setting $F_b(V_{\parallel}) = \int f_b(V_{\parallel}, V_{\perp}) \, dV_{\perp}^2$ we obtain:

$$G = -\int_{Z_0}^{Z_1} k_\parallel \, dz = \frac{n^2 w_{pe}^4}{2\pi e \omega^2 - \omega_{ce}^2} \omega \int_{-\infty}^{\infty} g(V_{\parallel}, W) \frac{dF_b}{dV_{\parallel}}$$  \hspace{1cm} (7)

where $g(V_{\parallel}, W) = g(V_{\parallel}, W) = w^2 \cdot V_{\parallel}^2 \cdot [w^2 + V_{\parallel}^2]^{-2}$ and $Z_1$ corresponds to $k_\parallel \to 0$, that is to $\omega_{UH}(Z_1) = \omega$. Consequently, due to the inhomogeneity along field lines
the parallel phase velocity explores the beam from \( V_{q_0} \) to infinity. If 
\( V_{q_0} > V_b \) the integral will be negative since \( \delta F_b / \delta V < 0 \), then \( G < 0 \).
Thus it is needed that (for a beam symmetric with respect to \( V_b \))
\( V_{q_0} < V_b - U \). But this is not sufficient, if \( g \) was constant we would then
get \( g = 0 \). It is easy to see that as long as \( W < V_b \) the growth will
be larger than further damping. In order to evaluate the integral in
(7) we have to specify \( F_b \), which is a difficult task since it is poorly
known. The simplest way is to use a square-ended beam distribution:
\[ F_b = \left( n_b / 2U \right) \left\{ H [V_{q}] - (V_b - U)] - H [V_{q}] - (V_b + U) \right\} \]
where \( H \) is
the Heaviside "function". Then we obtain to the lowest order in
\( U/V_b \) (which must be a small parameter):
\[ G = \pi^2 \frac{n_b}{n_c} \frac{w_{pe}}{\left( w^2 - w_{pe}^2 \right)} \frac{w}{w_{pe}} \frac{(1-w^2/V_b^2)}{V_b \left( 1+w^2/V_b^2 \right)^3} \]
(8)

It is clear from (8) that we must have \( w^2 < V_b^2 \) in order to get an
amplification. If we start at some point \( Z_0 \) with \( w/k_{lo}^2 \approx V_b - U_b \),
the preceding condition can be written: \( W^2/V_b^2 \approx k_{||}^2 w^2/w_{pe}^2 \approx \)
\( k_{||}^2 \approx (w^2 - w_{pe}^2)/w^2 \). On the other hand, as mentioned above, it is
necessary that \( (k_{||}^2/k_{lo}^2) \approx \left( U_b/V_b \right)^2 \) otherwise finite larmor radius
effect would begin to decrease the size of the gain factor. Combining
these conditions we get \( (U_b/V_b)^2 < (k_{||}^2/k_{lo}^2) \approx (w^2) \approx (w^2 - w_{pe}^2) \) which
seems reasonable since \( w^2/(w^2 - w_{pe}^2) > 1 \) and \( (U_b/V_b)^2 \approx 0.2 \). See the observations
resumed in figure 2 of Maggs(1976). It is worthwhile to notice that the above
condition is also compatible with (3) which insures that most of the
wave energy remains inside the beam and hence can be convectively ampli-
ified by it.
In order to have an order of magnitude for $G$ we have set

$$(W/V_b)^2 \approx 1/2, \quad \omega_{pe} \ll \omega_{ce}, \quad \text{so that } \omega^2 - \omega_{pe}^2 \approx \omega_{ce}^2, \quad V_b = 10^7 \text{ m/s}, \quad \beta = 3/R$$

and $\omega = 2\pi \times 10^5 \text{ Hz}, \quad R = 2.5R_E$. Then

$$G = 2 \times 10^5 \frac{n_b}{n_c} \left( \frac{\omega_{pe}}{\omega_{ce}} \right)^4$$

Then with $n_b/n_c = 10^{-3}$ and $\omega_{pe}/\omega_{ce} \approx 2$ we get $G = 12.5$ which is a rather large factor. It should be said at this point that the choice of a square ended beam has increased the size of $G$ as compared as it would have been for a more smooth beam distribution.

The discussion of the case $\omega_{ce} < \omega < \omega_{uh}$ is essentially similar one gets with the same square-ended beam distribution and $W_2^2 = (\omega^2 - \omega_{ce}^2)/\gamma^2 \approx \omega_{ce}$:

$$G_2 = \pi^2 \frac{n_b}{n_c} \frac{\omega}{\beta V_b} \frac{W_2^2 / V_b}{(1 + W_2^2 / V_b^2)^2}$$

Now $G$ is always positive, the maximum corresponds to $W_2^2 \sim V_b^2$; for the same reasons than those discussed above we will require that $(V_b/\gamma)^2 < \gamma^2 / \omega_{ce}^2 < 1$ the last inequality now comes from (3) in the case $\omega_{pe} \sim 2\omega_{ce}$.

Finally, when $\omega_{lh} \ll \omega \ll \min(\omega_{pe}, \omega_{ce})$ we get also:

$$G_3 = \pi^2 \frac{n_b}{n_c} \frac{\omega}{\beta V_b} \frac{W_3^2 / V_b}{(1 + W_3^2 / V_b^2)^2}$$
But now \( \omega_3^2 = \omega_0^2 (\omega^2 - \omega_{\text{LH}0}) / \kappa_{\text{LH}}^2 \omega_{\text{LO}}^2 \). \( \omega_3 \) is smaller than \( \omega_2 \) since now, in \( \omega_3 \) the frequency of interest is: \( \omega \sim \sqrt{\omega_{\text{LH}}} \), instead of being \( \omega \sim \omega_{\text{UH}} \). Consequently, the convective amplification is less than in the upper frequency window.
GEOMETRICAL AMPLIFICATION OF ELECTROSTATIC WAVES NEAR THE RESONANCES

When an electrostatic wave propagates towards resonance, as said earlier, \( k_{\parallel} \) and hence \( V_{\parallel} \) decrease in order to satisfy (1) for increasing ratios \( \omega/\omega_{\text{UH}}(Z) \) or decreasing ratios \( \omega/\omega_{\text{LH}}(Z) \). If we assume that \( V_{\parallel} \) is conserved then \( B^2 \) increases and would even become infinite when \( k_{\parallel} \to 0 \) that is when \( \omega \approx \omega_{\text{UH}} \) or \( \omega_{\text{LH}} \). Of course then the WKB procedure fails, we have a turning point. Let us evaluate this in a more appropriate way. Near a turning point (1) becomes in the upper propagation window:

\[
\frac{d^2 \omega}{dz^2} + k^2 \varphi \frac{[\omega_{\text{UH}}(Z) - \omega]^2}{[\omega^2 - \omega_{\text{ce}}(Z)][\omega^2 - \omega_{\text{pe}}(Z)]} = 0 \tag{12}
\]

As we are interested in the behavior of the electrostatic potential \( \varphi \) near from resonance, we can assume that: \( \omega^2 - \omega_{\text{pe}}^2 \approx \omega_{\text{ce}}^2 \) and that \( \omega^2 - \omega_{\text{pe}}^2 \approx \omega_{\text{ce}}^2 \) and keep them constant. Then expanding the upper hybrid frequency we find: \( \omega_{\text{UH}}(Z) = \omega_{\text{UHO}} [1 + \beta (Z-Z_0)] \) where \( Z_0 \) is our starting point. Then \( \omega_{\text{UH}}^2 - \omega^2 = \omega_{\text{UHO}}^2 - \omega^2 + 2 \beta \omega_{\text{UHO}}^2 (Z-Z_0) \) setting that \( \omega_{\text{UH}}(Z_1) = \omega \), we get:

\[
Z_0 - Z_1 = (\omega_{\text{UHO}}^2 - \omega^2)/2\beta \omega_{\text{UHO}}^2 \tag{13}
\]

Then, working with \( \xi = k_{\parallel}^{2/3}(Z-Z_1/Z_0-Z_1)^{-1/3} \) instead of \( Z \) we can rewrite (12) in the form:
\[
d^2\varphi/d\xi^2 - \xi \varphi = 0 \quad (14)
\]

whose solution is an Airy function \( \text{Ai}(\xi) \). According to this solution the electrostatic potential \( \varphi \) increases when \( Z \to Z_1 \), that is when the wave propagates towards resonance. Using the results of White and Chen (1974) we get:

\[
(\varphi_{\text{max}}/\varphi_0) \approx 1.93 \left[ k_{||0} (Z_0 - Z_1) \right]^{1/6} \quad (15)
\]

\( \varphi_0 \) being the potential of the incident wave far from \( Z_1 \). This result is valid as long as the reflection coefficient is one. Then starting with \( k_{||0} \approx k_{10} \) and assuming that \( \beta \approx 3/R, V_b \approx 10^7 \text{ m/s}, \omega_{ce} \approx 2\omega_{pe} \approx 2\pi \times 10^5 \text{ Hz}, R = 2.5 R_E \), we get \( (\omega_{lh0}^2 - \omega^2)/\omega_{lh0}^2 \approx \omega_{pe}^2/2\omega_{ce}^2 \approx 1/8 \) and then from (15):

\[
(\varphi_{\text{max}}/\varphi_0) \approx 10 \text{ and } (E_{\text{max}}/E_0) \approx (k_{10}/k_{1||0}^2)^{1/2} (\varphi_{\text{max}}/\varphi_0) \approx 7.
\]

This is a large increase since \( (E_{\text{max}}/E_0) \approx 50 \).

For lower frequencies, that is for waves which propagate towards increasing lower hybrid frequencies the situation is very similar to the previous case. By similar calculations we get: \( (\varphi_{\text{max}}/\varphi_0) = 1.93 \left[ k_{||0} \right] (\omega_{lh0}^2 - \omega^2)/(2\beta \omega_{lh0})^{1/6} \). As mentioned above, only waves starting at very small values of \( (k_{||0}/k_{10}) \) can reach resonance inside the beam then \( (E_{\text{max}}/E_0) \) is quite as large as previously \( \sim 7 \).

Consequently, electrostatic waves having \( k_{||0}/k_{10} \) small enough to reach resonance inside the beam are linearly amplified by beam driven instability and also geometrically amplified when they propagate towards resonances. Thus, the electrostatic waves spectrum should peak.
near $\omega_{\text{UH}}$ and $\omega_{\text{LH}}$ at a given place along auroral field lines. At a given place the frequency width of these spectra are essentially fixed by the Airy function. A wave having $\omega + \delta \omega$ will reach resonance at $\omega_{\text{UH}} + \delta \omega_{\text{UH}}$ that is at a point $Z_1 + \delta Z_1$ where $\delta \omega_{\text{UH}}/\omega_{\text{UH}} = 36Z_1/R_E$. We will fix $\delta Z_1$ by setting that the argument of the Airy function varies from $-2$ to $+2$. Then $\delta Z_1 \approx 4(Z_0 - Z_1)^{1/3} k_{||}^{2/3}$. Consequently we get:

$$\frac{\delta \omega_{\text{UH}}}{\omega_{\text{UH}}} \approx \frac{12}{R} \left[ \frac{\omega_{\text{UHO}}^2 - \omega_{\text{UH}}^2}{\omega_{\text{UHO}}^2} \right]^{1/3}$$

(16)

where $k_{||} \sim \omega/V_b$. Setting as above $\omega \approx \omega_{ce} \approx 2 \omega_{pe} \approx 2\pi \times 10^5 \text{ Hz}$, $R = 2.5 \text{ R}_E$, $V_b = 10^7 \text{ m/s}$. We get $(\omega_{\text{UHO}}^2 - \omega_{\text{UH}}^2)/\omega_{\text{UHO}}^2 \approx 1/8$ and $\delta \omega_{\text{UH}}/\omega_{\text{UH}} \approx 1.2 \times 10^{-3}$. Thus $\delta \omega_{\text{UH}} = 2\pi \times 130 \text{ Hz}$, which is a very narrow spectrum. Of course this only corresponds to the largest peak of the Airy function.
NON-LINEAR INTERACTION BETWEEN PLASMA WAVES

We want to generate, by non-linear beating between electrostatic waves an electromagnetic wave which can escape from the source region. Then, if the subscript 1 and 2 refer to the primary electrostatic wave while 3 stands for the electromagnetic one, we have to insure that:

\[ \mathbf{k}_3 = \mathbf{k}_1 + \mathbf{k}_2, \]
\[ \omega_3 = \omega_1 + \omega_2 \]
\[ D(\omega_i, \mathbf{k}_i) = 0 \]

where \( D \) is the dispersion relation, \( i = 1, 2, 3 \). In order to produce an electromagnetic wave it is necessary that \( |k| \sim \omega/c \ll |k_{1,2}| \sim \omega/V_b \). Consequently we must have \( |k_{1,2}| \sim |k_{1,2}| \). The first of these conditions is not easy to fulfill, for instance if we want to couple \( \omega_1 \) and \( \omega_2 \) and get \( \omega_3 = \omega_1 + \omega_2 \), then, since amplified waves have their parallel phase velocities in the same direction as the beam, only one of these waves can be amplified. Of course, the second condition is easy to satisfy, provided that the radial size of the beam is large and the distribution of \( k_{1,2} \) is isotropic around \( B_0 \). As we know from previous discussion that the electrostatic waves have large amplitudes near \( \omega_{\text{HE}} \) and \( \omega_{\text{UR}} \) any strong coupling between such waves is
a priori likely to occur. Furthermore, near from resonances $|k_{11}|$ and $|k_{12}|$ are already very small, hence the matching condition $k_{13} = k_{11} + k_{12}$ will be easily fulfilled. Then, in all what follows we will assume that primary electrostatic waves are near resonance.

An important question must now be discussed. Depending upon the width of the spectrum of the primary waves the non-linear decay or mode coupling can be coherent, or incoherent.

In the incoherent case, the random phase approximation (RPA hereafter) must be used. Due to the averaging over random phases the coupling between incoherent waves is usually rather small. In a recent work Barbosa has studied in the RPA, the generation of the AKR by non-linear coupling between two electrostatic waves having $\omega \approx \omega_{\text{UH}}$. Then he found that the primary electrostatic waves should have at least $3V/m$ in order to generate the observed large powers of electromagnetic waves. Which is probably too large by at least one order of magnitude, (which means four orders of magnitude as far as the electromagnetic energy is concerned).

In the two previous sections we have shown that near from $\omega_{\text{LH}}$ and $\omega_{\text{UH}}$ the spectrum of electrostatic waves is very peaked. Then a coherent three wave process is likely to occur.
Such a conclusion is valid if the phase velocity variation remains small over a non-linear growth time. This condition can be written \( \gamma_{NL} \geq \Delta \omega \).

\( \gamma_{NL} \) is the non-linear growth rate that one obtains assuming that the non-linear wave-wave interaction process is coherent. \( \Delta \omega \), the phase variation is linked to the spreading of the linear growth rate of plasma waves; expanding, up to the second order, the phase of the electrostatic wave we get

\[
\Delta \omega/2\pi = \left( \frac{\gamma_{L}^2}{2\omega_{e}^2} \int (\gamma_{L}/\gamma_{e}) \, dz \right)^{-1/2}
\]

which corresponds to a phase variation of \( 2\pi \).

We have shown in the previous part that, in the upper propagation window

\[
\int (\gamma_{L}/\gamma_{e}) \, dz = 12,
\]

the characteristic width of the convectively unstable waves (\( \delta \omega \)) is much smaller than the upper propagation band itself, say 1/5 of it. Then we get \( \delta \omega = (\omega_{ce}/5)(\omega_{pe}/2\omega_{e}^2) \); with \( \omega_{ce} = 3\omega_{pe} \) and \( \omega_{ce} = 2\pi \times 100 \text{kHz} \) we get \( \delta \omega/2\pi = 1 \text{kHz} \) and \( \Delta \omega/2\pi = 300 \text{Hz} \).

Consequently we must have \( \gamma_{NL} \geq 2 \times 10^3 \).

Let us now discuss the coherent three wave process, keeping in mind that the previous condition must hold. Two cases are possible:

\( \omega_1 \ll \omega_2 \ll \omega_{UH} \) and \( \omega_1 \ll \omega_{UH} \), \( \omega_2 \ll \omega_{LH} \), the frequencies of the resulting electromagnetic radiation being respectively \( \omega_3 \simeq 2\omega_{UH} \) and \( \omega_3 \simeq \omega_{UH} \).
RADIATION AT $\omega_3 \approx \omega_{\text{UH}}$

The non-linear currents which produce the electromagnetic radiation can be obtained by combining the Maxwell and fluid equations. In a situation where $\omega_{\text{pe}} \approx \omega_{\text{ce}}$, such a calculation is a priori very lengthy. Nevertheless, since $\omega_1 \approx \omega_{\text{UH}}$ and $\omega_2 \approx \omega_{\text{LH}}$, their wave numbers are nearly perpendicular to the static magnetic field; furthermore, as said above $|k_3| \ll |k_1|, |k_2|$. Then the problem is greatly simplified. Parallel and perpendicular non-linear currents have similar order of magnitude. But, since we need a freely escaping electromagnetic radiation, only the parallel current, which generates an $O$-wave, is of interest here. We get three coupled equations, but will only write the one which describes the production of the electromagnetic wave. Retaining only the largest contribution to the non-linear current we obtain:

\[
\frac{\delta D_3}{\delta k_{i3}} \frac{\delta E_3}{\delta \mathbf{r}} \approx \frac{\mu_0}{m} \frac{k_{i2} k_{i3}}{k_1} \frac{E(w, k_1) E(\omega_2, k_2)}{\omega_3} \]

\[
D_3 = \frac{k_3^2 c^2}{\omega_3^2} - 1 + \frac{\omega_{\text{pe}}^2}{\omega_3^2} \tag{18}
\]

where $\omega_1 \approx \omega_{\text{UH}}, \omega_2 \text{ near from } \omega_{\text{LH}}, \omega_3 \approx \omega_{\text{UH}}$. Since the $O$-mode propagates essentially across the magnetic field we have expanded $D_3 E_3 \approx 1 \frac{\delta D_3}{\delta k_1} \frac{\delta E_3}{\delta \mathbf{r}}$. Then using the following relation: $\omega_2/k_{i2} \approx \omega_{\text{pe}}/k_{i2} \approx \omega_{\text{pe}}/k_{11}$, which holds even when $\omega_2$ is near (but not at) $\omega_{\text{LH}}$, (18) becomes:
\[ \frac{\delta E_3}{\delta r} \approx \frac{e}{2mc^2} \frac{\omega_3^2}{\omega_{pe}^2} E_1 E_2^* \]  

(19)

Let us assume that \( |E_1|, |E_2| \gg |E_3| \), then (19) describes the production of the O-radiation by non-linear coupling of electrostatic waves. As long as \( |E_3| \) remains small one can neglect the two other equations which describe the reaction of the electromagnetic wave on the production of the electrostatic ones. Then setting \( \frac{\delta E_3}{\delta r} \approx \frac{E_3}{\delta r} \), \( r \) being the distance measured in a direction perpendicular to \( B_0 \), and \( \Delta r \) the characteristic length over which \( E_1 \) and \( E_2 \) have large amplitudes, which should be of the order of the radial extension of an inverted V event, that is, for \( L \sim 2-3 R_E \), \( \Delta r \approx R_0 \approx 10^5 \) m. Setting \( E_1 \approx E_2 \approx 100 \) mV/m, \( \omega_3 \approx 3 \omega_{pe} \) we get \( E_3 \approx 3 \) mV/m, which seems to be an accurate value for large AKR events in the source region. The largest AKR events as observed by Hawkeye, in the 178 kHz channel have \( E_3 \approx 10-20 \) mV/m if one takes care of the size of the propagation cone and of the dependance in \((R_E/R)^2\) of the power flux (see Green et al., 1977). Of course we also need the amplitudes of the electrostatic waves which is difficult to know since measurements in the presumed source region are lacking. Only recently, during low altitude (\( R \sim 1.7 R_E \)) over the southern hemisphere in auroral regions electric (from 1.78 Hz to 178 kHz) and magnetic (from 1.78 to 562 Hz) dc fields were measured on board Hawkeye 1. When there is some indication of a field aligned current in the night sector on auroral field lines, most of the electric field channels measured strong noise, (see fig.8 of Gurnett and Green 1978).

Nevertheless it is difficult to know what are the electrostatic wave amplitudes for \( \omega = \omega_{UH} \). As said earlier the peak at \( \omega_{UH} \) should be very narrow; since the 178 kHz channel of Hawkeye is also very narrow there is very little chance that
the 178 kHz channel be measuring the UHR noise at the right time and place. Furthermore, since there are no plasma frequency measurements, the plasma frequency (and thus the upper hybrid frequency) cannot be determined precisely. For the above reasons one cannot deduce from present Hawkeye measurements the electrostatic waves amplitude and spectrum for \( \omega \gtrsim \omega_{\text{UH}} \) in regions where inverted V events take place.

Nevertheless in the lower frequency channel, there is a strong indication that we have a large electrostatic noise peaked around \( \omega_{\text{LH}} \). The 1.78 kHz and 5.62 kHz channels are often at or near saturation when there is some indication of a field aligned current over auroral regions. This means that near from \( \omega_{\text{LH}} \) the amplitude should be of the order or larger than 40 mV/m which is not far from what was required by the previous calculation.

We must insure that the condition \( \gamma_{\text{NL}} \gtrsim \Delta \omega \) holds. Since \( \gamma_{\text{NL}} \approx C/R_b \approx 3000 \) and \( \Delta \omega \approx 2 \times 10^3 \) (see the previous part) the previous condition is marginally satisfied.

It is also possible to couple low and high frequency electrostatic waves having \( \omega_1 = \omega_{\text{LH}} \) and \( \omega_2 = \omega_{\text{UH}} \) in order to produce an extraordinary wave with \( \omega_3 = \omega_{\text{UH}} \). The coupling coefficient is even stronger than the previous one at least when \( \omega_{\text{ce}} > \omega_{\text{pe}} \). Nevertheless, such an electrostatic wave cannot escape from the source because its frequency is below the R-X cut-off except for very small ratios \( \omega_{\text{pe}}/\omega_{\text{ce}} \) which are unlikely on the earth but might be reasonable in Jupiter's magnetosphere.

For the above mentioned reason we will delete here this possibility for AKR generation.
RADIATION AT $2\omega_{\text{EH}}$

Setting $\omega_1 \approx \omega_2 \approx \omega_{\text{EH}}$ for the primary electrostatic waves we get by non-linear beatings a large non-linear current perpendicular to both the electrostatic wave numbers and $\hat{B}_0$. The corresponding parallel non-linear current is completely negligible. Hence the radiated wave is polarized in the X mode. Since the interaction involved here is a kind of self-interaction between the electrostatic wave which approaches resonance and a reflected wave, the matching conditions $k_{11} + k_{12} \approx 0$, $k_{11} + k_{12} \approx 0$ are easily satisfied. Then we get:

$$\frac{\partial D_3}{\partial k_1} \approx \frac{e}{m} k_1 \frac{B \cdot E}{w_{\text{pe}}} \left\{ 1 + \frac{w_{\text{pe}}^2}{w_{\text{ce}}^2} \left[ 1 + \frac{w_{\text{pe}}^2}{w_{\text{ce}}^2} \right] \right\}$$

and

$$D_3 = \left\{ \frac{k_1^2 e^2}{w_{\text{ce}}} - 1 + \frac{w_{\text{pe}}^2}{w_{\text{ce}}^2} \left[ 1 + \frac{w_{\text{pe}}^2}{w_{\text{ce}}^2} \right] \right\}_{w_{\text{pe}}, w_{\text{ce}}}

Provided that $w_{\text{pe}} < w_{\text{ce}}$ the last term in the RHS of the first equation is small and this equation rewrites:

$$\frac{\partial E_3}{\partial r} \approx E_3/\Delta r \approx \frac{e}{mc^2} \left( \frac{2w_{\text{EH}}}{w_{\text{pe}}} \right)^2 |E_1|^2 \quad \text{(21)}$$

Setting $\Delta r \approx R_b \approx 10^5$ m, $\omega_{\text{EH}} = 3w_{\text{pe}}$, $E_1 = 50$ mV/m we get $E_3 \approx 18$ mV/m which is fairly typical of strong TKR events. According to the
gyrofrequency and plasma frequency model of Green et al. (1977): $w_{ce} \gg \omega_{ce}$ over a wide range of altitude along auroral field lines; which increases significantly the coupling between electromagnetic and electrostatic waves as can be seen in (21). It should be noted that if one uses the same parameters as above namely $\omega_{UU} \approx 3\omega_{pe} E_1 \approx E_2 \approx 50$ mV/m, $\Delta r \approx 10^5$ m for both O-wave and X-wave generation. The latter is 25 more efficient as far as the energy is concerned.

It should be noted that the convective amplification which runs like $n_b/n_c (\omega_{pe}/\omega_{ce})^4$ is greatly reduced when $\omega_{pe}$ decreases furthermore the geometrical amplification which runs like $(\omega_{pe}/\omega_{ce})^{1/3}$ also decreases when $\omega_{pe}/\omega_{ce}$ decreases. Then, too small ratios $\omega_{pe}/\omega_{ce}$ would finally decrease the energy of the electromagnetic radiation. The optimum value for $\omega_{pe}/\omega_{ce}$ greatly depends upon the value of $n_b/n_c$ which is very poorly known. In any case, it seems that this optimum value is below the expected value of $\omega_{pe}/\omega_{ce}$ hence the conclusion that a decrease of $\omega_{pe}/\omega_{ce}$ increases the efficiency of the radiation process seems to be valid. It should be said also that too large amplitude for the electromagnetic wave will tend to saturate the emission process by pumping too much electrostatic energy.
CONCLUSION

We have studied analytically the convective amplification of electrostatic plasma waves by beam-driven instability. In both the lower and upper propagation windows, amplified waves propagate towards resonances ($\omega_{LH}$ and $\omega_{UH}$); simultaneously the modulus of their group velocities and their angle with the geomagnetic field tend to decrease. Thus the convectively amplified electrostatic waves suffer a further "geometrical" amplification while they propagate towards resonances. Consequently the spectrum of electrostatic waves in the regions where inverted-V events take place should peak near $\omega_{LH}$ or $\omega_{UH}$.

Due to the narrowness of the spectra of these primary waves the non-linear three wave interaction process which leads to electromagnetic wave generation is coherent. This is essentially the reason why the efficiency of the transformation of electrostatic waves into electromagnetic is so large. In principle the escaping electromagnetic radiation could be either in the O-mode with $\omega \approx \omega_{UH}$ or in the X-mode with $\omega \approx 2\omega_{UH}$. The respective efficiencies of the two preceding processes are essentially controlled by the ratio $\omega_{pe}/\omega_{ce}$ which is smaller than unity over auroral regions for $R \approx 2 - 4 R_E$ (see Green et al 1977). When it is so, the generation of the X-mode at $2\omega_{UH}$ becomes very efficient and largely overcomes the generation of the O-mode at $\omega_{UH}$. The strong A.K.R. events (with amplitudes $E \sim 20$ mV/m) can easily be generated by beatings between two electrostatic waves having narrow spectra and $E \sim 50$ mV/m for $\omega \approx \omega_{UH}$. Such an amplitude does not seem unrealistic especially if one takes care of the geometrical amplification effect; 50mV/m at $\omega_{UH}$ can correspond to
10mV/m just after the linear amplification stage, that is for $\omega$ not too near from $\omega_{U}\iota$.

The energy balance is discussed in the appendix.

One should keep in mind that the generation of electromagnetic radiation in the X-mode at or near $\omega_{U}\iota$, which was not considered here since it doesn't escape from the source region, could be an important component of the Jupiter Decametric Radiation (J.D.R.). This difference between A.K.R. and J.D.R. is essentially due to the very small ratios $\omega_{pe}/\omega_{ce}$ that presumably exist in Jupiter's magnetosphere. Then, as shown by Smith (1976), tunneling through the "stop zones" is very efficient. For very small ratios $\omega_{pe}/\omega_{ce}$ one could even have $\omega_{LH}+\omega_{U}\iota>\omega_{RX}$ the cutoff of the R-X mode, in such a case one could directly produce an escaping X-mode by beatings between electrostatic waves having $\omega_1 \simeq \omega_{U}\iota$ and $\omega_2 \simeq \omega_{LH}$.
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APPENDIX

Let us assume that the auroral region can be modeled by a parallelepiped $L \times H \times d$ where $d$ is the latitudinal extension of the auroral region at a distance $R \approx 2.5 R_E$, $L$ is the longitudinal extension of the A. K. R. source at the same distance and $h$ is vertical size. We have used $L \approx H \approx 10^7$ m, $d \approx 10^5$ m.

We want to evaluate the efficiency of the conversion of electrostatic waves into electromagnetic ones. We have seen that the generation of the $X$ mode at $2\omega_{UP}$ is the more likely to occur, thus we will only consider it here, which simplify the matter since only the electrostatic wave energy near from $\omega_{UP}$ is to be considered. The power flux of electrostatic waves is $\varepsilon V L d E_{e.s.}$. Near from resonance the optical geometry fails and $E_{e.s.}$ greatly increases; thus we will evaluate the power flux far from resonance, that is before that the focussing effect associated with this resonance has increased $E_{e.s.}$ by a large factor.

The production of the electromagnetic wave at a given frequency takes over a small region $L \times \delta z_1 \times d$, where $\delta z_1$ its vertical size is defined by the spatial extension of the Airy peak (see formula 16 and discussion above). We get $\delta z_1 \approx 4 (z_o - z_1)^{1/3} k_{0^2/3}$, for $\omega_{ce} = 2\omega_{pe} = 2 \chi 2\pi \times 10^5$ Hz, $V_b \approx 10^7$ m/s and $R \approx 2.5 R_E$ we get $z_o - z_1 \approx R_E/12$ and $\delta z_1 \approx 1.3 \times 10^3$ m.
Then power flux conservation writes:

(A.1) \[(1 - R) E_{e.m}^2 = \alpha E_{e.s}^2\]

where \(\alpha = Vg\pi d/\delta z_1 C\), \(R\) being the energy reflection coefficient. The reflected part of the electrostatic energy is proportional to \(R E_{e.s}^2\).

hence eq. 21 rewrites:

(A.2) \[E_{e.m}^2 = M^2 R E_{e.s}^4\]

where \(M\) is the coupling coefficient. Then combining (A.1) and (A.2) we get:

(A.3) \[E_{e.m}^2 = \frac{E_{e.s}^2}{\alpha + 1 / M^2 E_{e.s}^2}\]

For large enough values of \(E_{e.s}\), the asymptotic solution of (A.3) is \(E_{e.m}^2 = \alpha E_{e.s}^2\) with the previous values of the parameters and \(Vg\pi /C \ll 1/30\) we get \(\alpha \ll 3\). The electromagnetic wave energy is comparable to the electrostatic wave energy (as taken far from resonance).

The same calculation can be made at any distance from the earth within the source, so that the real width of the A. K. R. spectrum, as observed far from the earth, results from the integration over the vertical size of the source of the locally monochromatic radiation processes.
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