THE ASTROPHYSICAL JOURNAL, 230:621–626, 1979 June 1

A THEORY OF THE TERRESTRIAL KILOMETRIC RADIATION

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Received 1978 September 5; accepted 1978 December 15

ABSTRACT

During magnetospheric substorms, electrons with energies ~1 keV are injected from the plasma-sheet region into the auroral region. A fraction of these energetic electrons can precipitate into the upper atmosphere, and the rest are reflected because of the mirror effect of the convergent geomagnetic field. It is found that these reflected electrons can result in the amplification of electromagnetic waves via a relativistic normal cyclotron resonance. This process may explain the recently discovered terrestrial kilometric radiation.

Subject headings: Earth: atmosphere — Earth: aurorae — planets: radio radiation — radio sources: extended

I. INTRODUCTION

The phenomenon of terrestrial kilometric radiation (TKR) has recently attracted much attention. The nonthermal emission is very intense. The total power of TKR is estimated to be 10^7 W (Kaiser and Alexander 1977*a*) and sometimes 10^9 W, which is comparable to the power output of the decametric radiation from Jupiter. Observations concerning TKR have been reported and discussed by numerous authors (Gurnett 1974; Kurth, Baumback, and Gurnett 1975; Kaiser and Stone 1975; Alexander and Kaiser 1976, 1977; Kaiser and Alexander 1977*a*, *b*; Green, Gurnett, and Shawhan 1977; Voots, Gurnett, and Akasofu 1977; Gurnett and Green 1978; Alexander, Kaiser, and Rodriguez 1978; Kaiser *et al.* 1978). The main features may be briefly summarized as follows: (1) TKR is closely correlated with magnetospheric substorms (Gurnett 1974; Kaiser and Alexander 1977*a*, *b*; Voots *et al.*). Observational results indicate that many TKR events begin at low altitudes and high frequencies (~400–600 kHz) and spread to higher altitudes and lower frequencies as the substorm expands (Kaiser and Alexander 1977*b*). (2) TKR appears to originate from the nightside auroral regions (Gurnett 1974; Kaiser and Stone 1975) or dayside cusps (Alexander and Kaiser 1976; Kaiser and Alexander 1977*a*), with altitudes ranging from 1 to 3 Earth radii or higher, although effects of propagation and scattering may often create an impression that emissions can occur at a large distance from the Earth (Alexander *et al.*). (3) The typical frequencies are about 100–600 kHz (corresponding to wavelengths in the kilometric range). From the observed radiation frequencies and the projected locations of the source regions, it is inferred that the emission occur at frequencies close to the local electron gyrofrequency. (4) Observations also indicate, consequently, that the polarizations of TKR are largely in the fast extraordinary mode (or the X mode [Green *et al.*; Gurnett and Green 1978; Kaiser *et al.* 1978]) and that emissions usually occur in regions where the local cyclotron frequency is greater than the local electron plasma frequency. In an attempt to understand the radiation mechanism, several theories have been proposed in the literature (Benson 1975; Melrose 1976; Palmadesso *et al.* 1976; Jones 1977*a*, *b*; Boswell 1978). The details of the theories are quite different, but the emphasis is similar. The consensus is that, first, the observed radiation must be of induced nature rather than of spontaneous nature and, second, the precipitating electrons in auroral regions are responsible for the induced emission. The main issue is how these electrons give rise to TKR. Most of the aforementioned theories predict that TKR is in the ordinary mode (or the O mode) of polarization, which is in disagreement with most recent observations (Green *et al.*; Gurnett and Green 1978; Kaiser *et al.* 1978). Moreover, most of the theories require a conversion mechanism which usually reduces the efficiency substantially. In this regard, Melrose's (1976) theory is appealing. Also, it can explain the X-mode polarization. However, in the case of TKR his mechanism requires an exceedingly large energy anisotropy. In addition to those cited above, there are a number of recent works in which some highly sophisticated models are proposed (Gurnett 1974; Barbosa 1976; Galeev and Krasnoselsick 1977). Gurnett suggests that perhaps the $3/2 f_{ce}$ (where f_{ce} is the electron cyclotron frequency) electrostatic waves are excited first and that then these waves may produce phase bunching for the auroral electrons and eventually give rise to enhanced radiation. Without a detailed analysis, it is not clear whether this process is efficient enough to explain observational results. Moreover, experimentally, the $3/2 f_{ce}$ waves have not been measured in the polar region. Barbosa has considered the possibility that TKR is produced by the nonlinear coupling of two electrostatic waves with frequencies close to the upper hybrid frequency. In order to explain the observed TKR powers, the electrostatic waves are required to have a field strength of several volts per meter. This requirement may be difficult to fulfill. Galeev and Krasnoselsick (1977) argue that TKR may be generated by solitons which could be produced by the auroral electron beams. The idea is interesting. Observational evidence of the existence of Langmuir solitons in the auroral beams will be crucial. A review of most of these theories is given by Maggs (1977).

In this paper we propose a different mechanism, one which can explain the X-mode polarization yet without the requirement of large anisotropy as demanded in Melrose's theory.

II. MODEL AND THEORY

Let us consider the energetic electrons inside the flux tube above the discrete auroral region. These electrons are injected from the plasma sheet and have energies about 1 keV. (In the region of dayside cusp, energetic electrons are injected from the magnetosheath.) It is conceivable that those injected electrons with sufficiently small pitch angles can precipitate into the upper atmosphere. However, because of the convergent magnetic field lines, descending electrons with pitch angles outside the atmospheric loss cone will be reflected owing to a magnetic mirror effect. Thus, at a given location in the flux tube, an observer should see more descending energetic electrons than ascending electrons (assuming that all ascending energetic electrons are due to the reflection of those injected from the plasma sheet with mirror points below the observer). Furthermore, we expect the distribution function of the ascending electrons, say, F_+, to possess a loss-cone feature whereas we do not expect the distribution function of the descending electrons, say, F_-, to have this feature. If F_+ and F_- are normalized to unity, the total energetic-electron distribution function may be written as $F_e = n_+F_+ + n_-F_-$, where n_+ and n_- denote the density of the two kinds of electrons. Evidently, within the context of our model (no other source), $n_+ < n_-$. Moreover, $|n_+ \int d^3v v_\parallel F_+| < |n_- \int d^3v v_\parallel F_-|$, implying that the net flux is downward. Hereafter, we are most interested in the radiation which can escape directly without refraction. These electromagnetic waves have wave vectors pointing away from the Earth. As a result, we find that the ascending electrons play a decisive role in the amplification of these waves. This will be shown later.

We next turn our attention to the low-energy ($\sim$1 ev) background electrons, which have a density n_b and a corresponding plasma frequency ω_b. For the altitude range 1–4 Earth radii (R_E) along the auroral field lines, the commonly accepted model suggests that $\omega_b < \Omega_e$, where Ω_e denotes the local electron frequency. We believe that during the inverted-V or other discrete auroral events, the background electrons may be greatly reduced in certain regions near or above 1 R_E. An explanation of this point is in order.

The implication of the inverted-V (Frank and Ackerson 1971; Gurnett and Frank 1973) event is discussed by Gurnett (1972), who suggests that, in the magnetic flux tube above the auroral region, a longitudinal component of electric field $E_\parallel$ can exist locally. Recent electric field measurements along auroral field lines (Mozer *et al.* 1977; Wescott *et al.* (1976) show that the particular type of potential structure suggested by Gurnett may indeed occur. Furthermore, the signature of parallel electric fields can be confirmed by the particle distributions measured by the *S33* satellite (Mizera and Fennell 1977; Croley, Mizera, and Fenell 1978). The potential drop along the auroral field lines is of the order of several kilovolts. In the present discussion we are not concerned with the generating mechanism of this electric field. Rather, we are interested in the consequence of this phenomenon. If we assume that the electric field can establish a quasi-stationary state, which seems to be indicated by observations, then in the region where the kilovolt potential drop takes place the background electrons should be depleted (at least locally) to a very low level. [For example, if T_e is the temperature of the background electrons (in units of electron volts), $-|\delta\phi|$ is the potential drop, and n_0 is the electron density in the region where $\delta\phi = 0$, then we expect that in the region $\delta\phi > 0$ the local electron density $n \approx n_0 \exp(-|e\delta\phi|/T_e)$. For $T_e \approx 1$ eV and $|e\delta\phi| \gg 1$ eV, we see $n \to 0$.] Thus we shall hereafter ignore the low-energy electrons completely. Let us also restrict our discussion to the case of nearly perpendicular propagation. Although the energetic electrons can give rise to a current along the field lines, the current is very weak and the extraordinary mode and the ordinary mode can be approximately decoupled while $k_\parallel^2 \ll k^2$. The condition of nearly perpendicular propagation is supported by data acquired by the radio experiment on the *ISIS 1* satellite (W. Calvert and R. F. Benson, private communication). Moreover, since we expect $\omega^2 \approx k^2c^2 \gg \omega_e^2$, where $\omega_e^2 \equiv 4\pi n_e e^2/m_e$ and $n_e \equiv n_+ + n_-$, we can write the dispersion equations of interest to us as

$$1 - \frac{c^2k^2}{\omega^2} + \frac{\omega_e^2}{\omega^2 n_e}\int d^3v \left(\Omega_e \frac{\partial F_e}{\partial v_\perp} + k_\parallel v_\perp \frac{\partial F_e}{\partial v_\parallel}\right) \times \frac{v_\perp J_1'^2(b)}{(\omega - \Omega_e/\gamma - k_\parallel v_\parallel)} = 0 \quad (1)$$

(for the extraordinary mode when $c^2k^2 \gg \omega_e^2$) and

$$1 - \frac{c^2k^2_\perp}{\omega^2} + \frac{\omega_e^2}{\omega^2 n_e}\int d^3v \left(\Omega_e \frac{\partial F_e}{\partial v_\perp} + k_\parallel v_\perp \frac{\partial F_e}{\partial v_\parallel}\right) \times \frac{v_\parallel^2 J_1^2(b)}{v_\perp(\omega - \Omega_e/\gamma - k_\parallel v_\parallel)} = 0 \quad (2)$$

(for the ordinary mode), where $J_1'(b) = dJ_1/db$, $b = k_\perp v_\perp/\Omega_e$, and $\gamma = (1 - v^2/c^2)^{-1/2}$. Here a few remarks are necessary. First, in the original equations, a summation over n (n is a harmonic number such that $n = -\infty$ to $+\infty$) should be considered. In the present problem the argument of the Bessel function J_n is roughly of $O(v/c)$, which is small; therefore, the $n = 1$ term is most important in equation (1). On the other hand, in equation (2) the $n = 1$ term is again most important because Cerenkov resonance ($n = 0$) is impossible. Second, we are interested primarily in electrons with energies of a few keV, and the relativistic effect is important only in the denominator. Thus elsewhere we have set $\gamma = 1$. In the denominator, relativistic correction is of $O(\langle v^2\rangle/2c^2\Omega_e)$ and the Doppler-shift term is of $O(\Omega_e\langle v\rangle/c\cos\theta)$. Obviously, for $\cos\theta = k_\parallel/k \ll 1$, the two terms are of the same order. The term $\Omega_e v^2/2c^2$ provides a mechanism which couples the perpendicular electron motion $v_\perp^2$ with the waves and eventually allows the free energy available in the loss-cone distribution to be transferred to the electromagnetic waves. As a matter of fact, if the relativistic effect is ignored at the outset, no excitation is found in our theory. Third, in general, the dispersion equation for the extraordinary mode is represented by a 2×2 determinant. However, when $c^2k^2 \gg \omega_e^2$, electrostatic effect is negligible and, consequently, equation (1) represents a good approximation (Montgomery and Tidman 1964).

Let us consider waves with frequencies above the cutoff frequency and denote $\omega = \omega_r + i\omega_i$. For the extraordinary mode we find that, if we assume $\omega_r \gg \omega_i$,

$$\omega_i = \frac{\pi^2\omega_e^2}{4\omega_r}\frac{n_+}{n_e}\int_0^\infty dv_\parallel \int_0^\infty dv_\perp v_\perp^2 \times \delta\left[\omega_r - \Omega_e\left(1 - \frac{v^2}{2c^2}\right) - k_\parallel v_\parallel\right] \times \left(\Omega_e \frac{\partial F_+}{\partial v_\perp} + k_\parallel v_\perp \frac{\partial F_+}{\partial v_\perp}\right), \quad (3)$$

where we may ignore the effect of the energetic electrons on the real frequency and write $\omega_r \approx kc > \Omega_e (1 + \omega_e^2/\Omega_e^2)$. On the other hand, for the ordinary mode

$$\omega_i = \frac{\pi^2\omega_e^2}{4\omega_r}\frac{n_+}{n_e}\int_0^\infty dv_\parallel \int_0^\infty dv_\perp \frac{v_\perp^2}{c^2} v_\parallel^2 \times \delta\left[\omega_r - \Omega_e\left(1 - \frac{v^2}{2c^2}\right) - k_\parallel v_\parallel\right] \times \left(\Omega_e \frac{\partial F_+}{\partial v_\perp} + k_\parallel v_\perp \frac{\partial F_+}{\partial v_\parallel}\right), \quad (4)$$

where $\omega_r \approx ck_\perp > \Omega_e$. If $\omega_i > 0$, waves are amplified. Since we consider $\omega_r > \Omega_e$, resonance $\omega_r - \Omega_e(1 - v^2/2c^2) - k_\parallel v_\parallel = 0$ is possible only for particles with $k_\parallel v_\parallel > 0$. For waves propagating upward ($k_\parallel > 0$), only the ascending electrons ($v_\parallel > 0$) can resonate with the waves. Thus only F_+ appears in equations (3) and (4). To demonstrate that ω_i can be positive, we consider a model distribution function F_+ which should contain the essential feature $F_+(v_\parallel, v_\perp = 0) = 0$ (loss-cone effect). This feature implies $\partial F_+/\partial v_\perp > 0$ for a certain range of $v_\perp$. It is apparent from equations (3) and (4) that ω_i is essentially determined by the integral

$$\int dv_\parallel \partial F_+/\partial v_\perp|_{v_\perp = v_R(v_\parallel, w, k_\parallel)}$$

(integrating over a certain range of $v_\parallel$), where $v_R = c[2c_\parallel/c\cos\theta - v_\parallel^2/c^2 - 2(\omega_r/\Omega_e - 1)]^{1/2}$, and $\cos\theta = k_\parallel/k$. Since v_R must be positive and real, one can readily see that $v_\parallel$ must be within the range $v_1 < v_\parallel < v_2$, where $v_1 = c[\cos\theta - (\cos^2\theta - 2\Delta\omega)^{1/2}]$, $v_2 = c[\cos\theta + (\cos^2\theta - 2\Delta\omega)^{1/2}]$, and $\Delta\omega \equiv (\omega_r/\Omega_e - 1)$. Here it has been postulated that the condition $\cos^2\theta \geq 2\Delta\omega$ is satisfied. In other words for a given frequency ω_r, there exists a minimum value of $\cos\theta$ such that the resonance condition is meaningful. It is instructive to point out that for given θ and ω_r such that $\cos^2\theta > 2\Delta\omega$, there exists a $(v_R)_{\max} = c(\cos^2\theta - 2\Delta\omega)^{1/2}$. Thus a sufficient condition for positive ω_i is $\partial F_+/\partial v_\perp \gtrsim 0$ at $v_\perp = (v_R)_{\max}$. It is also obvious that if $F_+ = F_\perp(v_\perp)F_\parallel(v_\parallel)$ and if $F_\parallel$ peaks at $v_\parallel = U$, maximum growth rates occur when $\cos\theta \approx U/c$. To discuss the problem further we choose

$$F_+(v_\parallel, v_\perp) = 2(v_\perp^2/\alpha_\perp^2)\pi^{-3/2}\alpha_\perp^{-2}\alpha_1^{-1}(1 - \alpha_2/\alpha_1)^{-1} \times \exp(-v_\perp^2/\alpha_\perp^2) \times [\exp(-v_\parallel^2/\alpha_1^2) - \exp(-v_\parallel^2/\alpha_2^2)],$$

with $\alpha_2 < \alpha_1$ and $v_\parallel > 0$. Here we stress that the specific functional form is chosen for purposes of illustration only. The loss-cone feature is described by a form similar to that first considered by Dory, Guest, and Harris (1965) in plasma stability theory. Now we may define $U = \int d^3v v_\parallel F_+$, $\mathscr{E}_\perp = \frac{1}{2}m_e \int d^3v v_\perp^2 F_+$, and $\mathscr{E}_\parallel = \frac{1}{2}m_e \int d^2v (v_\parallel - U)^2 F_+$. From these definitions we can relate the two sets of parameters $(\alpha_\perp, \alpha_1, \alpha_2)$ and $(\mathscr{E}_s = \frac{1}{2}m_e U^2, \mathscr{E}_\perp, \mathscr{E}_\parallel)$. We have computed ω_i for two slightly different situations: (1) $\mathscr{E}_s = 1$ keV, $\mathscr{E}_\perp = 1$ keV, and $\mathscr{E}_\parallel = 0.5$ keV; and (2) $\mathscr{E}_s = 5$ keV, $\mathscr{E}_\perp = 5$ keV, and $\mathscr{E}_\parallel = 1$ keV. Evidently, for a given frequency ω_r, the growth rate ω_i for the extraordinary mode is significantly larger than that for the ordinary mode. The maximum growth rates $\omega_{i,\max}$ are plotted against ω_r/Ω_e in Figure 1. The corresponding values of $k_\parallel/k$ are also shown.

III. DISCUSSION AND SUMMARY

In a discussion of amplification of radiation, it is convenient to model a discrete aurora (inverted-V or auroral arc) by a slab geometry as shown in Figure 2. According to observations, the typical width in the x-direction varies from 100 m to 50 km. However, the length along the y-direction may exceed 10^3 km. Inside this slab the electron density is essentially n_e, which is significantly lower than that of the background electrons exterior to this slab. (Note that background electrons are assumed to have been depleted by an electric field along the magnetic field lines.) Obviously, the cutoff frequency of the X mode outside the slab $\omega_x(\mathrm{O})$ is higher than that inside the slab $\omega_x(\mathrm{I})$. If the radiation frequency ω_r satisfies the condition $\omega_x(\mathrm{O}) > \omega_r > \omega_x(\mathrm{I})$, the wave is locally trapped until it reaches a region where $\omega_r > \omega_x(\mathrm{O}) > \omega_x(\mathrm{I})$. Thus, even when the discrete auroral region is thin, the extraordinary-mode radiation can be trapped and sufficiently amplified for a long time. This trapping

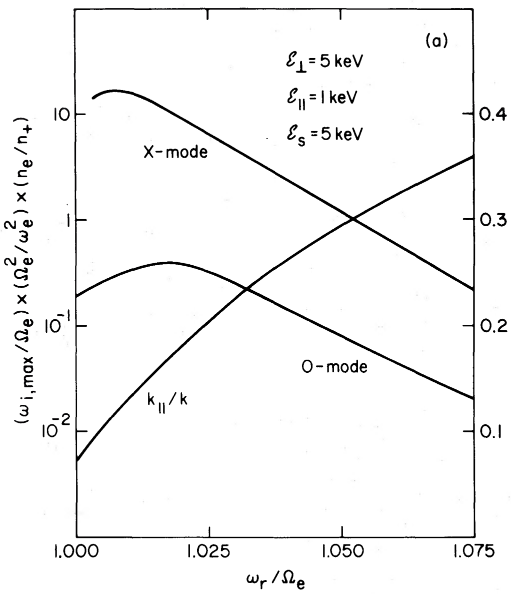


FIG. 1*a*

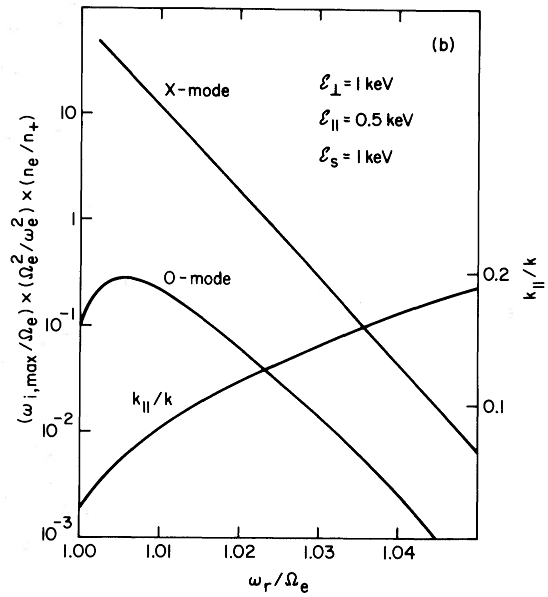


FIG. 1*b*

FIG. 1.—Maximum growth rate and the corresponding value of $k_{\parallel}/k$ versus frequency ω_r

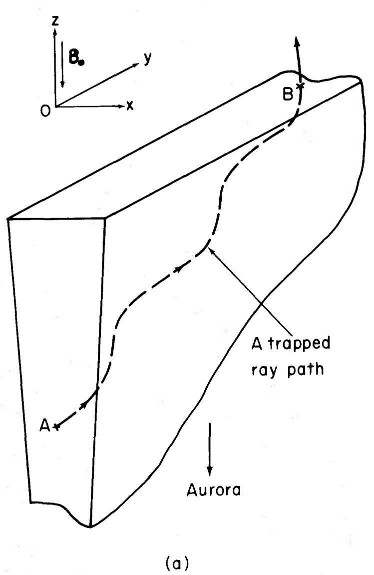
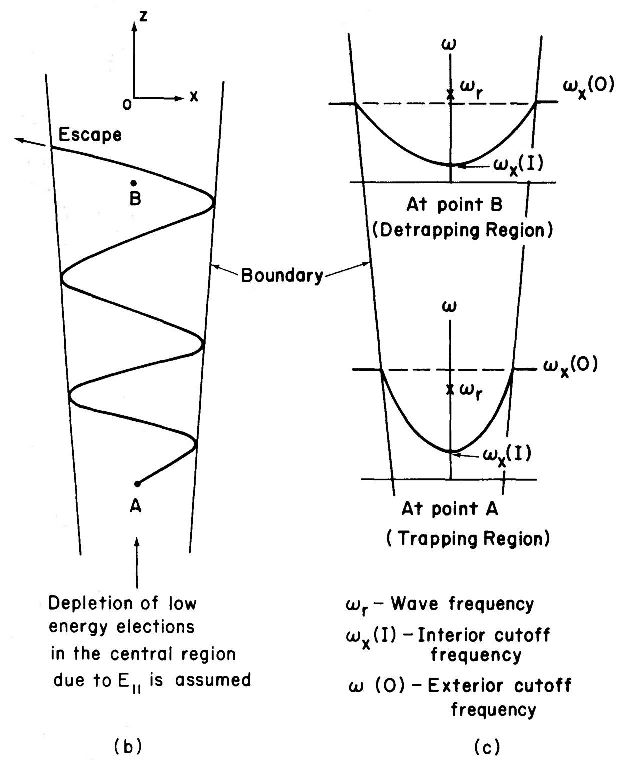


FIG. 2*a*



FIGS. 2*b–c*

FIG. 2.—Trapping of a ray path in a simple slab geometry. (*a*) Schematic drawing of a trapped ray path in a portion of a discrete aurora (three-dimensional view). (*b*) Projection into the *xz*-plane. (*c*) Cutoff frequency ω_x versus position. At point A, $\omega_x(\mathrm{I}) < \omega_r < \omega_x(\mathrm{O})$ (trapping). At point B, $\omega_r > \omega_x(\mathrm{O})$ (detrapping).

does not occur for the ordinary-mode radiation, and the amplification is therefore much less effective.

Here we remark that if we define an amplification distance $L \equiv v_g/\omega_i \approx c/\omega_i$, where v_g is the group velocity, we find that typically L varies from a few kilometers to several hundred kilometers depending upon the polarization and various parameters. Obviously, if L is of a few kilometers, the above-mentioned trapping process is unnecessary. In short, for the X mode propagating in the y-direction, as shown in Figure 2, sufficient amplification can be easily obtained. The power output per unit area may be estimated. If we assume $n_+ \approx 10\ \mathrm{cm}^{-3}$ (estimated from *ISIS 1* measurements), typical energy of the reflected electron $\mathscr{E} \approx 5$ keV, and only 0.1% of the electron energy to be converted into the radiation energy, we find $P = d\overline{W}/dt = 10^{-3} n_+ \mathscr{E} L \omega_i \approx 3 \times 10^{-3}\ \mathrm{W\ m^{-2}}$. The typical TKR bandwidth, according to Kaiser and Alexander (1977a), is 2×10^5 Hz. Thus the power flux in the source region is roughly $10^{-8}\ \mathrm{W\ m^{-2}\ Hz^{-1}}$. We now compare this with the observational result reported by Gurnett (1974) of a peak power flux $10^{-14}\ \mathrm{W\ m^{-2}\ Hz^{-1}}$ at 25 R_{E}. (This perhaps represents one of the extreme cases since in most cases the flux density is lower than $10^{-15}\ \mathrm{W\ m^{-2}\ Hz^{-1}}$ [Kaiser and Alexander 1977*a*].) Assuming that the radiation is beamed into a cone-shaped region with a typical solid angle, say, 4 sr, we find that the flux density at 25 R_{E}, $S_{2^a R_{\mathrm{E}}}$, is related to that in the source region, S_s, by

$$S_{25 R_{\mathrm{E}}} \approx \frac{A}{4(25\ R_{\mathrm{E}})^2} S_s\,,$$

where A denotes an effective area of the source region. If $S_s \approx 10^{-8}\ \mathrm{W\ m^{-2}\ Hz^{-1}}$, $R_{\mathrm{E}} \approx 6 \times 10^6$ m, and $S_{25 R_{\mathrm{E}}} \approx 10^{-14}\ \mathrm{W\ m^{-2}\ Hz^{-1}}$, it is found that

$$A \approx 4 \times 2.3 \times 10^{12}\ \mathrm{m}^2 = 9.2 \times 10^4\ \mathrm{km}^2\,.$$

If we assume that the radiation frequency is close to the local gyrofrequency, the bandwidth (2×10^5 Hz) implies that the dimension along the magnetic field flux tube is roughly 1 R_{E} (or 6×10^3 km). On the other hand, the typical thickness of an inverted-V region is 50–100 km in the ionosphere. If this thickness is mapped to the source region, say, 3 R_{E}, the thickness is roughly 300–600 km. Thus an effective area $A \approx 10^5$ km is consistent with observations. It should be remarked that in the above calculation we have assumed the flux density at 25 R_{E} to be uniform over the entire solid angle of 4 sr. This assumption is justified by observational results (Green *et al.*). Here we note that in the above calculation we have considered $S_{25 R_{\mathrm{E}}} = 10^{-14}\ \mathrm{W\ m^{-2}\ Hz^{-1}}$, which perhaps represents the most extreme value ever observed. If we had adopted a lower value for $S_{25 R_{\mathrm{E}}}$, the calculated effective area would be smaller.

In summary, we suggest in this paper a new mechanism which is basically different from the models proposed by previous authors. In the present theory the emission mechanism of TKR is attributed to a maser effect associated with the trapped energetic electrons which originate in the plasma sheet during a substorm. The free energy of the electrons is transferred to electromagnetic waves via a relativistic normal Doppler resonance process. The theory has a few appealing features: (1) no large thermal anisotropy is required; (2) the model explains why the X-mode radiation prevails over the O-mode radiation; (3) the amplification process is direct; and (4) the model does not need to implicitly assume that in the source region electrons are much hotter than ions. Most important of all is its simplicity. The present theory does not have to rely on the complicated and efficiency-limiting processes suggested by some of the previous models. Of course, the theory is preliminary. Evidently it is desirable to examine the effects of the density inhomogeneity and magnetic field gradient in the source region on the radiation process. Furthermore, the auroral electrons can also excite Langmuir waves and very low frequency waves. Whether these waves can in turn significantly modulate TKR is another area that deserves study. The recent results acquired with ISEE *1/2* satellites may provide us with many important clues and will be valuable for future research on the subject.

The present research was supported by the National Aeronautics and Space Administration under grant NGL 21-002-005. We thank J. K. Alexander and R. F. Benson of NASA Goddard Space Flight Center for making their works and other information available to us before publication.

REFERENCES

Alexander, J. K., and Kaiser, M. L. 1976, *J. Geophys. Res.*, **81**, 5948.

———. 1977, *J. Geophys. Res.*, **82**, 98.

Alexander, J. K., Kaiser, M. L., and Rodriguez, P. 1978, preprint.

Barbosa, D. D. 1976, Ph.D thesis, UCLA.

Benson, R. F. 1975, *Geophys. Res. Letters*, **2**, 52.

Boswell, R. W. 1978, *Geophys. Res. Letters*, **5**, 395.

Croley, D. R., Jr., Mizera, P. F., and Fennell, J. F. 1978, Aerospace Corp. Rept., No. SAMSO-TR-78-38.

Dory, R. A., Guest, G. E., and Harris, E. G. 1965, *Phys. Rev. Letters*, **14**, 131.

Frank, L. A., and Ackerson, K. L. 1971, *J. Geophys. Res.*, **76**, 3612.

Galeev, A. A., and Krasnoselsick, V. V. 1977, preprint.

Green, J. L., Gurnett, D. A., and Shawhan, S. D. 1977, *J. Geophys. Res.*, **82**, 1825.

Gurnett, D. A. 1972, in *Critical Problems of Magnetospheric Physics*, ed. E. R. Dyer (National Academy of Sciences), p. 123.

———. 1974, *J. Geophys. Res.*, **79**, 4227.

Gurnett, D. A., and Frank, L. A. 1973, *J. Geophys. Res.*, **78**, 145.

Gurnett, D. A., and Green, J. L. 1978, *J. Geophys. Res.*, **83**, 689.

Jones D. 1977*a*, *Astr. Ap.*, **55**, 245.

———. 1977*b*, *Geophys. Res. Letters*, **4**, 121.

Kaiser, M. L., and Alexander, J. K. 1977*a*, *J. Geophys. Res.*, **82**, 3273.

———. 1977*b*, *J. Geophys, Res.*, **82**, 5283.

Kaiser, M. L., Alexander, J. K., Riddle, A. C., Pearce, J. B., and Warwick, J. W. 1978, *Geophys. Res. Letters*, **5**, 857.

Kaiser, M. L., and Stone, R. G. 1975, *Science*, **189**, 285.

Kurth, W. S., Baumback, M. M., and Gurnett, D. A. 1975, *J. Geophys. Res.*, **80**, 2764.
Maggs, J. E. 1977, review paper presented at IAGA meeting, Seattle.
Melrose, D. B. 1976, *Ap. J.*, **207**, 651.
Mizera, P. F., and Fennell, J. F. 1977, Aerospace Corp. Rept., No. SAMSO-TR-77-195.
Montgomery, D. C., and Tidman, D. A. 1964, *Plasma Kinetic Theory* (New York: McGraw-Hill), p. 151.
Mozer, F. S., Carlson, C. W., Hudson, M. K., Torbert, R. B., Parady, B., Yatteau, J., and Kelly, M. C. 1977, *Phys. Rev. Letters*, **38**, 292.
Palmadesso, P., Coffey, T. P., Ossakow, S. L., and Papadopoulos, K. 1976, *J. Geophys. Res.*, **81**, 1762.
Voots, G. R., Gurnett, D. A., and Akasofu, S. I. 1977, *J. Geophys. Res.*, **82**, 2259.
Wescott, E. M., Stenback-Nielsen, H. C., Hallinan, T. J., and Davis, T. N. 1976, *J. Geophys. Res.*, **81**, 4495.

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