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AN EMISSION MECHANISM FOR EXTRAGALACTIC RADIO JETS

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ABSTRACT

The observed emissions of radio waves from extragalactic radio jets are customarily interpreted in terms of the incoherent synchrotron radiation produced by a nonthermal population of relativistic electrons (or positrons) gyrating in a magnetic field. Such an interpretation requires, by necessity, a mechanism (or mechanisms) to constantly reenergize the particles in directions transverse to the magnetic field in momentum space. This paper, on the other hand, presents a possible alternative mechanism for radiation emission that does not require transverse reacceleration, albeit the theory is only preliminary at this stage. The emission mechanism described in the present paper is an induced plasma process that results from a nonlinear coupling of the electrostatic beam mode with the transverse electromagnetic mode. The beam mode is excited when a stream of ultrarelativistic electron-positron pair plasma nonlinearly interacts with low-frequency intrinsic turbulence. The only source of free energy for the instability is the streaming motion of the pair plasma. Therefore, as opposed to the incoherent synchrotron model, this mechanism (if, indeed, it is operative) has an advantage of not having to rely on the concept of constant renenergization of particles in the transverse direction.

Subject headings: acceleration of particles — galaxies: jets — MHD — radiation mechanisms: nonthermal

1. INTRODUCTION

Many extragalactic radio sources are known to possess a pair of jets that emanates from the center of a galaxy, called the active galactic nucleus (AGN) (Perley, Dreher, & Cowan 1984; Serabym 1986; Kellerman & Pauliny-Toth 1981; Readhead & Pearson 1982; Bridle & Perley 1984; Begelman, Blandford, & Rees 1984; Blandford & Königl 1979; Scheuer & Readhead 1979; Cohen & Unwin 1982; Orr & Brown 1982; Fanti, Kellerman & Setti 1984; Porcas 1986). The jets are believed to be made of energetic plasma (the usual electron-proton plasma, or more likely, the electron-positron pair plasma). The radiation from these objects is observed in a radio-frequency band roughly corresponding to 100 MHz–10 GHz.

The standard interpretation of these radio sources is based upon the well-known incoherent synchrotron radiation by relativistic electrons gyrating in a magnetic field (Pacholyczyk 1970; Rybicki & Lightman 1979). The synchrotron emission is a spontaneous emission of electromagnetic waves by ultrarelativistic electrons (or positrons) undergoing gyration about the magnetic field. It is the transverse acceleration (i.e., the gyro motion) which is responsible for the spontaneous emission of photons.

In the radio jet environment, on the other hand, kinetic energy in the energetic plasma resides mainly in the streaming motion. Therefore, in order for these jets to radiate photons through synchrotron mechanism, a portion of streaming energy has to be converted to energy in the transverse direction. Once particles begin to emit photons via the synchrotron process, however, the transverse energy is transferred to radiation energy, and eventually, the individual particle motion becomes predominantly field aligned again. As these radio jets are observed to travel enormous distances (kiloparsecs, and even megaparsecs), it is obvious that a constant reacceleration of the particle in a transverse direction must be somehow maintained.

Meanwhile, from the standpoint of plasma physics, as the jet plasma plows its way through a surrounding intergalactic plasma, it is expected that certain collective interaction processes between the jet plasma and the background plasma will occur. In the past, a number of emission mechanisms based upon collective plasma processes have been proposed to explain various features in the radio jet (Baker et al. 1988; Yoon & Chang 1989; Weatherall & Benford 1991). This paper introduces yet another alternative radiation mechanism which relies on a nonlinear interaction between the jet plasma and the intrinsic density perturbation in the background medium. The nonlinear wave-particle interaction results in the coupling of the beam mode excited by the jet plasma and the transverse electromagnetic mode. Such a coupling leads to a direct excitation of high-frequency electromagnetic radiation.

The basic notion of the physical mechanism invoked in the present paper is conceptually very similar to the ion-ripple-laser theory in the laboratory experiment (Chen & Dawson 1992). Moreover, a similar notion was applied to the plasma emission phenomenon (Yoon & Wu 1994; Wu, Yoon, & Zhou 1994; Yoon 1995), which is widely accepted to be central to the solar radio bursts such as the type II and type III bursts (Kundu 1965; Melrose 1980; Goldman 1983).

It should be remarked at the outset that the present theory is only preliminary, and that detailed test and comparison of the theory against actual sources must wait until a plausible nonlinear saturation picture is developed. Such a task is not trivial, and at the present time, a nonlinear saturation model of the present mechanism does not exist. However, in view of the difficulties inherent in the standard synchrotron interpretation, it is timely to explore alternative radiation mechanisms.

The organization of this paper is as follows: In § 2, we present the general formulation of the problem, in which we derive the formal nonlinear dispersion relation (3), with linear dispersion tensor defined by equation (4), the second-order susceptibility being zero, and the third-order susceptibility given by equation (7). Then, in § 3, a simplified dispersion relation (15) is derived from the formal result (14), from which numerical calculation of the complex roots are determined. Finally, § 4 presents the discussion and conclusions.

2. NONLINEAR DISPERSION EQUATION

In what follows, essential assumptions made in the present analysis are explained. First, we assume that the jet material is made of an ultrarelativistic electron-positron (e^+-e^-) pair plasma. Such a plasma is assumed to maintain charge and current neutrality. The intergalactic medium is assumed to be made of the usual electron-proton plasma. The temperature of the intergalactic plasma is considered to be cold. Let us assume that the background medium, through which the jet travels, supports a fluctuation in the ion number density. The ion density fluctuation generates an electrostatic field, $\boldsymbol{E}(\boldsymbol{r}) = -\nabla\phi(\boldsymbol{r})$. In spectral form, the electric field generated by the ion density modulation (i.e., the turbulence) is written as

$$\boldsymbol{E}(\boldsymbol{r}) = -i \sum_{\boldsymbol{k}} \boldsymbol{k}\phi_{\boldsymbol{k}} \exp(i\boldsymbol{k}\cdot\boldsymbol{r}) . \tag{1}$$

Here $\phi_{\boldsymbol{k}}$ is the spectral component of the electrostatic potential, $\phi(\boldsymbol{r})$, associated with the turbulence.

In what follows, we model the momentum distributions for ultrarelativistic e^+-e^- pair jet plasma in such a way that these particles are distributed with zero momentum spread in the transverse direction. That is, the momentum distribution functions are one-dimensional. Furthermore, for the sake of simplicity, let us ignore the spread in the parallel momentum. Taking the z axis as the direction parallel to the jet propagation, we adopt the following model distribution function for the jet plasma:

$$f_\pm(\boldsymbol{p}) = \frac{n_\pm}{2\pi p_\perp}\,\delta(p_\perp)\delta(p_\parallel - p_\pm) , \tag{2}$$

where $f_\pm$ is normalized according to $\sum_{+,-} 2\pi \int_0^\infty dp_\perp p_\perp \int_{-\infty}^\infty dp_\parallel f_\pm(\boldsymbol{p}) = \sum_{+,-} n_\pm = n_0$, and $n_\pm$ is the number density for $e^\pm$, respectively. For the sake of simplicity, we will limit ourselves to the situation where $n_+ = n_- = n_0/2$ and $p_+ = p_- = p_0$. In this case, the charge and current neutrality conditions can be satisfied.

In the stability analysis that follows, we consider the following physical situation. If the jet plasma is sufficiently dense in comparison with the intergalactic plasma, it can be conceived that the jet can excite a beam electrostatic mode, in addition to the usual electrostatic Langmuir oscillation. We are interested in the coupling of the beam mode and the transverse electromagnetic mode, through the nonlinearity provided by the ion density modulation. Therefore, for the sake of mathematical clarity, we shall ignore the contribution from the background electrons altogether in the stability analysis. Also, owing to the high-frequency wave, which is of interest to us, the ions play no dynamical role, except to provide the density modulation. Under these assumptions, a nonlinear dispersion equation has been derived, following the method outlined in Yoon & Wu (1994), Wu et al. (1994), and Yoon (1995). Here, we present the formal result

$$0 = \Lambda_{ij}(\boldsymbol{k}, \omega)E^j_{\boldsymbol{k}\omega} - \sum_{\boldsymbol{k}'} \langle\phi^2\rangle_{\boldsymbol{k}'}[\chi^{(2)}_{il}(\boldsymbol{k}'\,|\,\boldsymbol{k}-\boldsymbol{k}', \omega)\Lambda^{-1}_{lm}(\boldsymbol{k}-\boldsymbol{k}', \omega)\chi^{(2)}_{mj}(-\boldsymbol{k}'\,|\,\boldsymbol{k}, \omega) - \chi^{(3)}_{ij}(\boldsymbol{k}'\,|\,-\boldsymbol{k}'\,|\,\boldsymbol{k}, \omega)]E^j_{\boldsymbol{k},\omega} . \tag{3}$$

In the above, $\langle\phi^2\rangle_{\boldsymbol{k}'}$ is the square of the spectral electrostatic potential ensemble averaged over the random phase; $\boldsymbol{E}_{\boldsymbol{k}\omega}$ is the self-consistently generated electric field, whose wave frequency $\omega = \omega_{\boldsymbol{k}} + i\gamma_{\boldsymbol{k}}$ is to be determined from the above nonlinear dispersion equation ($\gamma_{\boldsymbol{k}} > 0$ corresponds to the instability); and the quantity $\Lambda_{ij}(\boldsymbol{k}, \omega)$ is the linear dispersion tensor given by

$$\Lambda_{ij}(\boldsymbol{k}, \omega) = \delta_{ij} + \chi_{ij}(\boldsymbol{k}, \omega) - \frac{c^2k^2}{\omega^2}\left(\delta_{ij} - \frac{k_i k_j}{k^2}\right) . \tag{4}$$

Here $\chi_{ij}(\boldsymbol{k}, \omega)$ is the total (e^- and e^+) linear susceptibility tensor defined by

$$\chi_{ij}(\boldsymbol{k}, \omega) = \sum_{+,-} \frac{4\pi e^2}{\omega} \int d\boldsymbol{p} v_i\, G_{\boldsymbol{k}\omega} L^j_{\boldsymbol{k}\omega} f_\pm(\boldsymbol{p}) . \tag{5}$$

In equation (3), $\chi^{(2)}_{ij}(\boldsymbol{k}'\,|\,\boldsymbol{k}-\boldsymbol{k}', \omega)$ is the second-order nonlinear susceptibility tensor defined by

$$\chi^{(2)}_{ij}(\boldsymbol{k}'\,|\,\boldsymbol{k}-\boldsymbol{k}', \omega) = \sum_{+,-} \frac{\pm 4\pi e^3}{\omega} \int d\boldsymbol{p} v_i\, G_{\boldsymbol{k}\omega} L_{\boldsymbol{k}'} G_{\boldsymbol{k}-\boldsymbol{k}'\omega} L^j_{\boldsymbol{k}-\boldsymbol{k}'\omega} f_\pm(\boldsymbol{p}) , \tag{6}$$

while $\chi^{(3)}_{ij}(\boldsymbol{k}'\,|\,\boldsymbol{k}''\,|\,\boldsymbol{k}-\boldsymbol{k}'-\boldsymbol{k}'', \omega)$ is the third-order nonlinear susceptibility tensor defined by

$$\chi^{(3)}_{ij}(\boldsymbol{k}'\,|\,\boldsymbol{k}''\,|\,\boldsymbol{k}-\boldsymbol{k}'-\boldsymbol{k}'', \omega) = \sum_{+,-} \frac{4\pi e^4}{\omega} \int d\boldsymbol{p} v_i\, G_{\boldsymbol{k},\omega} L_{\boldsymbol{k}'} G_{\boldsymbol{k}-\boldsymbol{k}'\omega} L_{\boldsymbol{k}''} G_{\boldsymbol{k}-\boldsymbol{k}'-\boldsymbol{k}''\omega} L_{\boldsymbol{k}-\boldsymbol{k}'-\boldsymbol{k}''\omega} f_\pm(\boldsymbol{p}) , \tag{7}$$

where

$$G_{\boldsymbol{k}\omega}(\boldsymbol{v}) = \frac{1}{\omega - \boldsymbol{k}\cdot\boldsymbol{v}} , \qquad L^i_{\boldsymbol{k}\omega}(\boldsymbol{p}) = \left(1 - \frac{\boldsymbol{k}\cdot\boldsymbol{v}}{\omega}\right)\frac{\partial}{\partial p_i} + \frac{v_i}{\omega}\,\boldsymbol{k}\cdot\frac{\partial}{\partial \boldsymbol{p}} , \qquad L_{\boldsymbol{k}}(\boldsymbol{p}) = \boldsymbol{k}\cdot\frac{\partial}{\partial \boldsymbol{p}} . \tag{8}$$

In the above, the usual relationship between the relativistic momentum and velocity, $\boldsymbol{p} = \gamma m \boldsymbol{v}$, is implicit; $\gamma = (1 + p^2/m^2c^2)^{1/2} = (1 - \beta^2)^{-1/2}$ being the Lorentz mass factor, m being the electron/positron mass, c being the speed of light in vacuo, $\beta = v/c$ being the usual relativistic beta.

In the present discussion, as we have stated before, we take $n_+ = n_- = n_0/2$ and $p_+ = p_- = p_0$ in equation (2) (this preserves the charge and current neutrality in the jet). For such a choice and the distribution function (2), one can readily derive the following explicit expression for the linear susceptibility tensor:

$$\chi_{ij}(\boldsymbol{k}, \omega) = -\frac{\omega_p^2}{\gamma_0 \omega^2}\left[\delta_{ij} + \frac{n_i \beta_{0j} + n_j \beta_{0i}}{1 - \boldsymbol{n} \cdot \boldsymbol{\beta}_0} - \frac{(1 - n^2)\beta_{0i}\beta_{0j}}{(1 - \boldsymbol{n} \cdot \boldsymbol{\beta}_0)^2}\right], \tag{9}$$

where

$$\boldsymbol{\beta}_0 = \frac{p_0}{\gamma_0 mc}\,\boldsymbol{e}_z = \frac{v_0}{c}\,\boldsymbol{e}_z\,, \qquad \omega_p = \left(\frac{4\pi n_0 e^2}{m}\right)^{1/2}, \qquad \gamma_0 = \left(1 + \frac{p_0^2}{m^2c^2}\right)^{1/2} = \frac{1}{(1 - \beta_0^2)^{1/2}}\,, \qquad \boldsymbol{n} = \frac{c\boldsymbol{k}}{\omega}\,. \tag{10}$$

In the above, ω_p is the total (electron plus positron) plasma frequency, and $\boldsymbol{n}$ is the vector index of refraction.

It is interesting to note that

$$\Lambda(\boldsymbol{k}, \omega) = \det \Lambda_{ij}(\boldsymbol{k}, \omega) = \left[1 - \frac{\omega_p^2/\gamma_0^3}{\omega^2(1 - \boldsymbol{n} \cdot \boldsymbol{\beta}_0)^2}\right]\left(1 - n^2 - \frac{\omega_p^2}{\gamma_0 \omega^2}\right)^2. \tag{11}$$

Linear dispersion relations for normal modes in a plasma can be obtained by simply letting $\Lambda(\boldsymbol{k}, \omega)$ equal to zero, and solving for the roots, $\omega = \omega_k$.

For the present situation, as a result of the assumption $n_+ = n_- = n_0/2$ and $p_+ = p_- = p_0$, it can be seen that the second-order susceptibility tensor (6) vanishes identically by virtue of exact cancellation:

$$\chi_{ij}^{(2)}(\boldsymbol{k}' \,|\, \boldsymbol{k} - \boldsymbol{k}', \omega) = 0\,.$$

The third-order nonlinear susceptibility tensor (7), meanwhile, can be evaluated formally as follows:

$$\chi_{ij}^{(3)}(\boldsymbol{k}' \,|\, -\boldsymbol{k}' \,|\, \boldsymbol{k}, \omega) = \frac{e^2}{m^2c^4}\frac{\omega_p^2}{\gamma_0\omega^2}\left\{\frac{\partial}{\partial \beta_{0j}} - \beta_{0j}\boldsymbol{\beta}_0 \cdot \frac{\partial}{\partial \boldsymbol{\beta}_0} + [\boldsymbol{n} - (\boldsymbol{n} \cdot \boldsymbol{\beta}_0)\boldsymbol{\beta}_0] \cdot \frac{\partial}{\partial \boldsymbol{\beta}_0}\frac{\beta_{0j}}{1 - \boldsymbol{n} \cdot \boldsymbol{\beta}_0}\right\}$$
$$\times \left\{\frac{1}{\gamma_0}[\boldsymbol{n}' - (\boldsymbol{n}' \cdot \boldsymbol{\beta}_0)\boldsymbol{\beta}_0] \cdot \frac{\partial}{\partial \boldsymbol{\beta}_0}\left[\frac{1}{\gamma_0}\frac{\boldsymbol{n}' - (\boldsymbol{n}' \cdot \boldsymbol{\beta}_0)\boldsymbol{\beta}_0}{1 - (\boldsymbol{n} - \boldsymbol{n}') \cdot \boldsymbol{\beta}_0} \cdot \frac{\partial}{\partial \boldsymbol{\beta}_0}\frac{\beta_{0i}}{1 - \boldsymbol{n} \cdot \boldsymbol{\beta}_0}\right]\right\}. \tag{12}$$

In the above, a generic form for the various derivatives is very cumbersome to perform explicitly. For this reason, in the present analysis, let us concentrate on the situation in which the density modulation associated with the background medium is directed predominantly along the direction of the jet propagation. More specifically, we are interested in those spectral components that has an opposite sign with respect to the beam propagation direction:

$$\boldsymbol{k}' = -k'\boldsymbol{e}_z\,.$$

This is to ensure that nonlinear wave-particle interaction can occur. Spectral component with positive k values cannot interact with the particles. In this case, the reader can easily verify that only the zz-component of the third-order nonlinear susceptibility tensor (12) remains nonzero. Thus, by explicitly writing $\beta_{0i} = \beta_0\,\delta_{iz}$, and $\boldsymbol{n}' = -n'\boldsymbol{e}_z$, we obtain

$$\chi_{zz}^{(3)}(-k' \,|\, k' \,|\, \boldsymbol{k}, \omega) = \frac{e^2 n'^2}{m^2c^4}\frac{\omega_p^2}{\gamma_0^3\omega^2}\frac{\partial}{\partial \beta_0}\left\{\frac{1}{\gamma_0^3}\frac{1}{1 - n\beta_0\cos\theta}\frac{\partial}{\partial \beta_0}\left[\frac{1}{\gamma_0^3}\frac{1}{(1 - n\beta_0\cos\theta)^2}\frac{1}{1 - (n\cos\theta + n')\beta_0}\right]\right\}. \tag{13}$$

In the above, $\boldsymbol{k} = k\sin\theta\boldsymbol{e}_x + k\cos\theta\boldsymbol{e}_z$ has been used, which entails the relation, $\boldsymbol{n} \cdot \boldsymbol{\beta}_0 = n\beta_0\cos\theta$.

For the sake of completeness, in what follows we present equation (13) with the derivatives actually carried out:

$$\chi_{zz}^{(3)}(-k' \,|\, k' \,|\, \boldsymbol{k}, \omega) = \frac{e^2 n'^2}{m^2c^4}\frac{\omega_p^2/\gamma_0^5}{\omega^2(1 - n\beta_0\cos\theta)^3[1 - (n\cos\theta + n')\beta_0]}\left[\!\left[12\beta_0^2 - \frac{3}{\gamma_0^2}\left[1 + \frac{7n\beta_0\cos\theta}{1 - n\beta_0\cos\theta} + \frac{3(n\cos\theta + n')\beta_0}{1 - (n\cos\theta + n')\beta_0}\right]\right.\right.$$
$$\left.\left. + \frac{1}{\gamma_0^4}\left\{\frac{8n^2\cos^2\theta}{(1 - n\beta_0\cos\theta)^2} + \frac{2(n\cos\theta + n')^2}{[1 - (n\cos\theta + n')\beta_0]^2} + \frac{5n\cos\theta(n\cos\theta + n')}{(1 - n\beta_0\cos\theta)[1 - (n\cos\theta + n')\beta_0]}\right\}\right]\!\right].$$

Using the above result, one can show that the formal nonlinear dispersion equation (3) is now rewritten as

$$0 = \left(1 - n^2 - \frac{\omega_p^2}{\gamma_0\omega^2}\right)\left[1 - \frac{\omega_p^2/\gamma_0^3}{\omega^2(1 - n\beta_0\cos\theta)^2}\right] + \left(1 - n^2\cos^2\theta - \frac{\omega_p^2}{\gamma_0\omega^2}\right)\sum_{k'}\langle\phi^2\rangle_{k'}\,\chi_{zz}^{(3)}(-k' \,|\, k' \,|\, \boldsymbol{k}, \omega)\,. \tag{14}$$

Note again, that setting $\Lambda(\boldsymbol{k}, \omega)$ equal to zero and solving for the solutions gives the linear dispersion relation (see eq. [11]). Specifically, one observes that $\Lambda(\boldsymbol{k}, \omega) = 0$ supports the transverse wave dispersion relation, $\omega^2 = c^2k^2 + \omega_p^2/\gamma_0$, as well as the Doppler-shifted longitudinal Langmuir mode dispersion relation, $(\omega - \boldsymbol{k} \cdot \boldsymbol{v}_0)^2 = \omega_p^2/\gamma_0^3$.

As we have noted earlier, the present analysis ignores the contribution from the background electrons to the dielectric tensor for reasons we have explained before. Had we kept their contribution, in addition to the Doppler-shifted Langmuir mode solution, there should also exist the usual solution describing the plasma oscillation associated with the background (i.e., intergalctic) medium. However, such a result would only make the analysis cumbersome, and essential physical notions can easily get lost in the complexity of the algebra. To elucidate the basic idea, we thus have chosen to consider only the e^+-e^- pair plasma as the only active component.

In the absence of the background ion density fluctuation (i.e., $\langle\phi^2\rangle_{k'} = 0$), it can be shown that the (Doppler-shifted) Langmuir mode and the transverse electromagnetic mode do not couple. To put it another way, regardless of the propagation angle θ and the relativistic beam speed β_0, the two modes defined by

$$n^2 = 1 - \frac{\omega_p^2}{\gamma_0\omega^2} ,$$

and

$$(1 - n\beta_0 \cos\theta)^2 = \frac{\omega_p^2}{\gamma_0^3\omega^2}$$

have no common solution. However, as we shall demonstrate later, the background ion fluctuation ($\langle\phi^2\rangle_{k'} \neq 0$) gives rise to another Doppler-shifted beam mode, $\omega \simeq (ck\cos\theta + ck')\beta_0$, which can couple to the radiation mode, under certain circumstances, in such a way that at the intersection of the two modes in $\omega - k$ space, there exists an instability.

3. NONLINEAR BEAM INSTABILITY

To further simplify the general nonlinear dispersion equation (14), let us consider the simple limiting case of monochromatic ion density modulation. Suppose that the ion fluctuation has a characteristic wave vector

$$k' = k_0 ,$$

so that the spectral energy density for the electric potential associated with the density fluctuation is given by

$$\langle\phi^2\rangle_{k'} = \phi_0^2 .$$

Note that in the formal nonlinear dispersion equation (14), there are two types of resonances associated with the density fluctuation. One of them is defined by the denominator $1 - n\beta_0\cos\theta$, and the other is associated with the denominator $1 - (n\cos\theta + n_0)\beta_0$. Of these, the first resonance (which characterizes the Doppler-shifted Langmuir oscillation) is of no interest to us, since it cannot interact with the transverse mode. The second resonance, which comes from the nonlinear wave-particle interaction, leads to the Doppler-shifted beam electrostatic mode,

$$1 - (n\cos\theta + n_0)\beta_0 \sim 0 .$$

Unlike the case with the Langmuir mode, it can be shown that the above beam electrostatic mode $[1 = (n\cos\theta + n_0)\beta_0]$, and the transverse electromagnetic mode $[n = (1 - \omega_p^2/\gamma_0\omega^2)^{1/2}]$ do have a common solution, provided the condition

$$(ck_0\beta_0)^2 > \frac{\omega_p^2}{\gamma_0}(1 - \beta_0^2\cos^2\theta)$$

is satisfied. Since such an intersection between the beam electrostatic mode and the transverse mode is essential for an unstable mode to exist, we find that the necessary and sufficient condition for the instability is

$$ck_0\beta_0 > \frac{\omega_p}{\gamma_0^{1/2}} . \tag{15}$$

Henceforth, let us assume that the condition (15) is maintained. Let us turn to equation (14). We are interested in solutions to the nonlinear dispersion equation in the vicinity of both $\omega^2 = \omega_p^2/\gamma_0 + c^2k^2$ and $\omega = c(k\cos\theta + k_0)\beta_0$. Therefore, we let

$$1 - n^2 - \frac{\omega_p^2}{\gamma_0\omega^2} \simeq 2\frac{\delta\omega}{\omega} \qquad 1 - (n\cos\theta + n_0)\beta_0 \simeq \frac{\delta\omega}{\omega} . \tag{16}$$

Substituting equation (16) into (14), we readily obtain the following algebraic equation for $\delta\omega/\omega$:

$$0 = \frac{\delta\omega^4}{\omega^4} + \frac{1}{2\beta_0^2}\sin^2\theta\left(\frac{e\phi_0}{\gamma_0 mc^2}\right)^2 \frac{n^2(1-n^2)(1-n\beta_0\cos\theta)}{\gamma_0^2(1-n\beta_0\cos\theta)^2 - (1-n^2)}$$
$$\times\left\{\left[12 + \frac{6}{\gamma_0^2} - \frac{21}{\gamma_0^2}\frac{1}{1-n\beta_0\cos\theta} + \frac{8}{\gamma_0^4}\frac{n^2\cos^2\theta}{(1-n\beta_0\cos\theta)^2}\right]\frac{\delta\omega^2}{\omega^2} - \frac{1}{\gamma_0^2}\left(9 - \frac{5}{\beta_0\gamma_0^2}\frac{n\cos\theta}{1-n\beta_0\cos\theta}\right)\frac{\delta\omega}{\omega} + \frac{2}{\beta_0^2\gamma_0^4}\right\} , \tag{17}$$

where $n = (1 - \omega_p^2/\gamma_0\omega^2)^{1/2}$. In deriving the above expression, we have replaced $n_0\beta_0$ by $1 - n\beta_0\cos\theta$, and replaced the factor $1 - n^2\cos^2\theta - \omega_p^2/\gamma_0\omega^2$ by $n^2\sin^2\theta$. Solving equation (17) for $\delta\omega$ as a function of ω gives the desired nonlinear growth rate $\gamma = \mathrm{Im}(\delta\omega)$ for the nonlinear beam instablity. However, one should keep in mind that in equation (17), not all ω is of interest to us. Remember that ω that appears in equation (17) is a point where the beam mode and the electromagnetic mode intersects. In other words, ω must satisfy both equations, $\omega^2 = c^2k^2 + \omega_p^2/\gamma_0$ and $\omega = ck\beta_0\cos\theta + ck_0\beta_0$ simultaneously. Solving these two equations, we therefore obtain

$$\omega = \frac{ck_0 \beta_0}{1 - \beta_0^2 \cos^2 \theta} \left\{ 1 + \beta_0 \cos \theta \left[1 - \frac{\omega_p^2/\gamma_0}{c^2 k_0^2 \beta_0^2} (1 - \beta_0^2 \cos^2 \theta) \right]^{1/2} \right\} . \tag{18}$$

As a consequence, in determining the growth rate $\gamma = \mathrm{Im}\,(\delta\omega)$, we can now substitute equation (18) into equation (17), and solve for $\delta\omega$.

Let us introduce the following dimensionless quantities:

$$x \equiv \gamma_0^{1/2} \frac{\omega}{\omega_p} , \qquad \Gamma = \gamma_0^{1/2} \frac{\mathrm{im}\ \delta\omega}{\omega_p} \left(\frac{e\phi_0}{2mc^2} \right)^{-1/2} , \qquad \kappa = \gamma_0^{1/2} \frac{ck_0 \beta_0}{\omega_p} \tag{19}$$

Here x represents the normalized wave frequency, Γ is the normalized growth rate, and κ is the normalized characteristic wave number associated with the low-frequency ion density fluctuations. In equation (17), let us keep only the leading term, i.e., the third term within the curly bracket on the right-hand side of equation (17). We then obtain an analytical expression for the normalized growth rate Γ:

$$\Gamma^2 = |\sin \theta| \frac{x^2}{\beta_0^2 \gamma_0^3} \frac{n(1 - n^2)^{1/2}(1 - n\beta_0 \cos \theta)^{1/2}}{[\gamma_0^2(1 - n\beta_0 \cos \theta)^2 - (1 - n^2)]^{1/2}} , \qquad n = \left(1 - \frac{1}{x^2} \right)^{1/2} ,$$
$$x = \frac{\kappa}{1 - \beta_0^2 \cos^2 \theta} \left[1 + \beta_0 \cos \theta \left(1 - \frac{1 - \beta_0^2 \cos^2 \theta}{\kappa^2} \right)^{1/2} \right] . \tag{20}$$

Note that the condition for instability (15) dictates that κ must be greater than 1.

To illustrate the general behavior of the solution (20), we have plotted the normalized radiation frequency x (Fig. 1) and the square of the normalized growth rate $\Gamma^2(\beta_0^2 \gamma_0^3)$ (Fig. 2) versus the angle of radiation propagation, θ (in degree), for $\kappa = 2$, and for various relativistic beam speed, $\beta_0 = 0.5$, 0.6, 0.7, 0.8, and 0.9. According to Figure 1 the normalized growth rate $x = \gamma_0^{1/2}(\omega/\omega_p)$ generally decreases for increasing propagation angle θ, and for increasing β_0 value, x increases as well. Note that for $\theta = 90°$, $x = \kappa$. Moreover, the higher the β_0 value, the higher the maximum value of x attains. The normalized growth rate $\Gamma^2(\beta_0^2 \gamma_0^3)$ vanishes for $\theta = 0$, but it reaches a finite value for $\theta = 90°$. The peak growth rate shifts to smaller propagation angle as β_0 increases.

Finally, in Figure 3, which is in the same format as Figures 1 and 2, we vary κ from 2, to 4, to 6, while holding $\beta_0 = 0.9$ constant. For larger κ values, one can see that both the normalized radiation frequency x and the normalized wave growth rate $\Gamma^2(\beta_0^2 \gamma_0^3)$ increases. To aid the readers, the growth rate curves are indicated by arrows.

4. CONCLUSIONS AND DISCUSSION

In this paper, we have presented a possible radio wave emission mechanism appropriate for extragalactic radio jet environment, which is based upon the concept of collective plasma instability. The specific instability mechanism is the nonlinear beam instability due to a nonlinear interaction of electron-positron (e^+-e^-) stream (eq. [2]) and the low-frequency density fluctuation (eq. [1]) that may naturally exist in a background (intergalactic) plasma. The e^+-e^- plasma stream gives rise to a beam mode, which nonlinearly couples to the transverse electromagnetic wave as the stream interacts with the low-frequency density fluctuation. At the point of mode-coupling, we have found that there exists an instability. An approximate analytical expression for the wave growth (eq. [20]) is obtained, and it is shown that the wave amplification occurs over a broad range of wave frequency (Figs. 1–3).

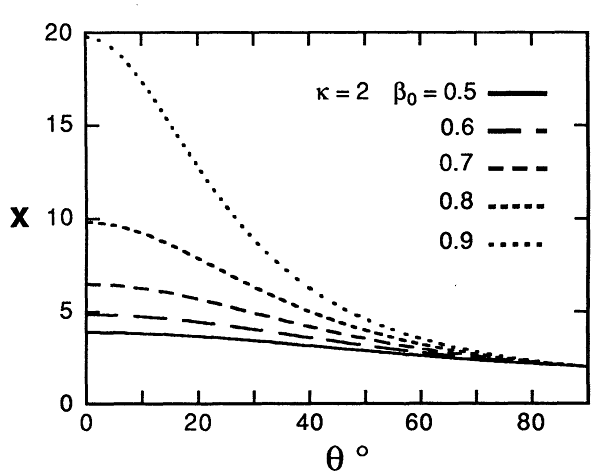

FIG. 1

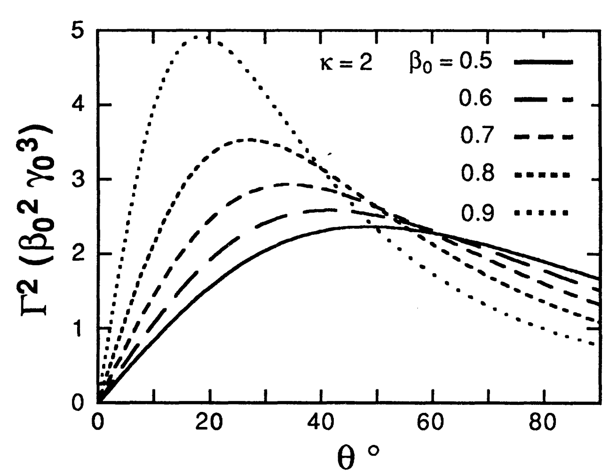

FIG. 2

FIG. 1.—Normalized radiation frequency x vs. the wave propagation angle θ. Note that the smaller the angle of propagation, the higher the excited wave frequency, and the higher the β_0, the higher the radiation frequency.

FIG. 2.—Normalized growth rate vs. the wave propagation angle. The peak growth rate shifts to smaller angles as β_0 increases.

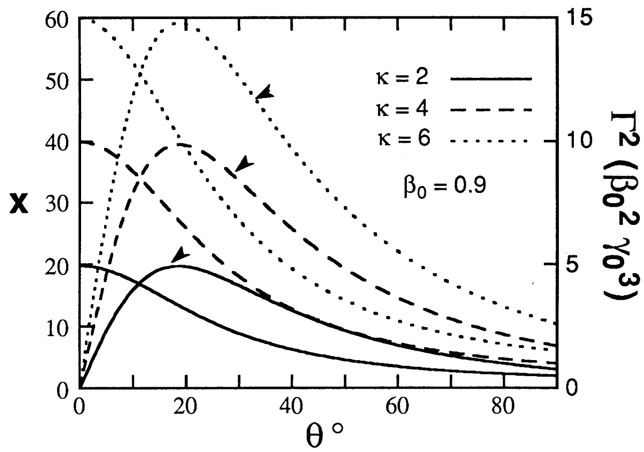

FIG. 3.—The same as Figs. 1 and 2, except that the dependence of the radiation frequency and the growth rate on the low-frequency fluctuation wave length κ is shown. In general, the shorter the wavelength (larger κ), the higher the radiation frequency and growth rate.

In the present paper, we have concentrated on a relatively simple physical model. Thermal effects, effects due to a spectrum of low-frequency turbulence, and the effect of background electron population have all been ignored. Also, this paper considers only the basic physical property of the proposed radio emission mechanism. Therefore, we have made no attempt to compare our theoretical result with any specific observations, as our theory is only preliminary.

The advantage of the present theory over the conventional spontaneous synchrotron model is that the present mechanism does not rely on constant reenergization of charged particles in transverse direction with respect to the magnetic field, which the synchrotron model requires. Another advantage of the present model (as is the case with other induced emission mechanisms) is that this model may resolve the high apparent brightness temperature issue. The disadvantage of the present model (as it is with many other induced emission theories) is, of course, that the present theory only establishes the condition for the radiation emission to occur (that is, the nonlinear instability), but the actual application cannot be made until a fully nonlinear saturation model is worked out. The detailed saturation picture, which must be properly worked out in the future, involves the dynamical evolution of particles and various wave modes. Such a task is not trivial and is beyond the scope of the present paper.

Because the synchrotron theory produces readily applicable formulae, it is widely used for the interpretation of observed features. However, in view of the many difficulties the synchrotron model suffers, especially in the extragalactic jet environment, we feel that it is timely and worthwhile to explore and investigate other radiation emission models, such as ours, although admittedly the model introduced in this article is only preliminary. It is our hope that the present discussion will stimulate the interest of other researchers in the community.

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