SYNCHROTRON EMISSION DRIVEN BY THE CHERENKOV-DRIFT INSTABILITY IN ACTIVE GALACTIC NUCLEI

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ABSTRACT

In the present paper, we study the generation of synchrotron emission by means of the feedback of Cherenkov-drift waves on the particle distribution through the diffusion process. Despite the efficient synchrotron losses, it is demonstrated that the excited Cherenkov-drift instability leads to the quasi-linear diffusion (QLD), the effect of which is balanced by dissipation factors and, as a result, the pitch angles are prevented from damping, thus maintaining the corresponding synchrotron emission. We analyze the model for a wide range of physical parameters and determine that the mechanism of QLD guarantees the generation of electromagnetic radiation from soft X-rays up to soft $\gamma$-rays, which is strongly correlated with Cherenkov-drift emission ranging from IR up to UV energy domains.

Key words: galaxies: active – instabilities – plasmas – radiation mechanisms: non-thermal

1. INTRODUCTION

Interest in X-ray and $\gamma$-ray emission from active galactic nuclei (AGNs) has substantially increased recently due to new observational data from satellite telescopes. Investigations of cosmic objects have been aided by the 2007 and 2008 launches of the Astro-rivelatore Gamma a Immaginii Leggero satellite and the Fermi spacecraft. These new observational results are of fundamental importance in studying the emission properties of several X-ray and $\gamma$-ray sources, such as AGNs, pulsars, gamma-ray bursts, and others.

It is most commonly believed that highly relativistic electrons do not participate in the synchrotron emission. Energy losses are so efficient that particles loose their energies very rapidly and transit to the ground Landau states. In spite of this it was shown so efficient that particles lose their energies very rapidly and do not participate in the synchrotron emission. Energy losses are still might account for the radiation processes produced by very energetic particles in AGNs.

According to Kazbegi et al. (1991), in strong magnetic fields the plasma may induce unstable cyclotron waves. On the other hand, as it was shown by Lominadze et al. (1979), these waves via quasi-linear diffusion (QLD) perturb the particles both along and across the magnetic field lines. Such a feedback of cyclotron waves on particles will inevitably lead to the creation of pitch angles, restoring synchrotron emission. In addition to AGNs the QLD process was applied to pulsars as well and it was shown that the previously mentioned mechanism is very efficient for pulsar magnetospheres (Lominadze et al. 1979; Machabeli & Usov 1979; Malov & Machabeli 2001; Machabeli & Osmanov 2009, 2010).

This approach makes it possible for synchrotron radiation to be a working mechanism despite strong energy losses. A very interesting property of the QLD is that it enables us to produce highly correlated radiation in two different energy bands. In this process there are two major forces that influence the particle distribution function. On the one hand, the diffusion attempts to increase the values of the pitch angles of resonant particles and, on the other hand, the emitting particles are affected by dissipative forces intending to decrease their pitch angles. In the papers listed above it was shown that under certain conditions the dissipative and diffusion factors balance each other, the physical system reaches stationary state, and the pitch angles saturate. Consequently, the cyclotron emission redistributes resonant particles by means of diffusion and as a result synchrotron emission is produced. This mechanism guarantees strongly correlated radiation in two different energy bands.

In the framework of this emission model, the radiation is generated in two energy bands because of the feedback of the excited cyclotron waves on particles by means of the QLD, which start to radiate in the synchrotron regime. In general, a similar emission mechanism cannot be driven only via the cyclotron waves. Particularly, in the present paper, unlike the aforementioned articles, we consider the feedback of the Cherenkov-drift waves on the resonant particles via the QLD. This process should also switch the synchrotron radiation mechanism and must provide correlated emission in different energy domains.

The paper is arranged as follows. In Section 2, we introduce the QLD theory. In Section 3, we apply the model to AGNs and in Section 4 we summarize our results.

2. MODEL

In this section we develop the model of the QLD of particles driven by means of the feedback of the Cherenkov-drift modes in the magnetospheres of AGNs. Normally, in the region mentioned, values of the Lorentz factors of electrons may vary from $\sim 1$ to $\sim 10^7$ with typical lengthscales $l \sim 10^{13}–10^{14}$ cm. This range for the Lorentz factors was implied in Osmanov et al. (2007) and Rieger & Aharonian (2008), where it was shown that acceleration of particles becomes extremely efficient in the light cylinder zone (a hypothetical area where the linear velocity of rotation exactly equals the speed of light) due to the relativistic effects of rotation. We apply the methods developed in those papers in the framework of the present article. For simplicity one can assume that the magnetosphere of AGNs is composed of a low energy plasma component with the Lorentz factors $\gamma_p$ and a highly relativistic part—the beam component with
As we have already mentioned, the synchrotron radiation is suppressed for strongly relativistic electrons. In particular, one can see that the cooling timescale of beam electrons is given by $t_{\text{syn}} \approx \gamma_b m c^2 / P_{\text{syn}}$, where $m$ is the electron’s mass, $c$ is the speed of light, $P_{\text{syn}} \approx 2e^4 \gamma_b^5 B^2 / 3m^3 c^3$ is the single particle synchrotron emission power, $B$ is the magnetic induction, and $\alpha$ is the electron’s charge. By taking into account typical values of the magnetic field in an ambient close to the supermassive black hole (SMBH), $B \approx 10^7 \text{–} 10^4 \text{ G}$, one can see that the cooling timescale varies in the following interval $5 \times (10^{-7} \text{–} 1) \text{ s}$. On the other hand, close to the light cylinder zone, due to the efficient curvature drift instability, the magnetic field lines significantly twist (Osmanov 2008); therefore, the lengthscale is still of the order of $l$ and the corresponding plasma escape timescale, $l/c$, is of the order of $3 \times 10^{3} \text{ s}$. As we can see, the synchrotron cooling timescale is much less than the kinematic timescale. Therefore, in AGN magnetospheres, due to the strong energy losses, particles very rapidly transit to the ground Landau states resulting in the termination of the emission process.

Generally speaking, as was explained by Shapakidze et al. (2002), the necessary condition for the development of the Cherenkov-drift instability (ChDI) is the presence of the beam component in plasmas. Since the magnetic field lines are always curved, it is evident that the electrons will drift along a direction perpendicular to the plane of the curved magnetic field lines with the following velocity:

$$u_\psi = \gamma \beta c e / \omega_x \rho,$$  

where $\omega_x = e B / (m c)$ is the cyclotron frequency of the particle and $\rho$ is the curvature radius of the magnetic field line. It is clear that for highly relativistic particles ($\gamma_B > 1$) the drift velocity becomes significant and the ChDI with the following resonance condition arises (Shapakidze et al. 2002):

$$k_x v_x = \xi v_\psi = 0,$$  

where $k_x$ and $v_x$ are the longitudinal components (along the magnetic field line), respectively, of the wave vector and the particle velocity, and $k_x$ is the wave vector’s component along the drift direction.

During the process of ChDI, the transverse $(t)$ waves might be excited, having the growth rate (Shapakidze et al. 2002)

$$\Gamma_k = \frac{\pi \omega_0^2 \gamma_B}{2 \omega} \omega_B \rho, $$

where $\omega_B = \sqrt{4 \pi n_b e^2 / m}$ is the Langmuir frequency, $n_b$ is the beam number density, $A_k = \left(k_x / k_x^2 \right)$, and $k_x$ is the component of the wave vector along the curvature radius of magnetic field lines. Throughout the paper we assume $A_k = 1$ (Kazbegi et al. 1991).

During their motion in a nonuniform magnetic field, the charged relativistic particles undergo two major dissipative forces: $\mathbf{H}$—responsible for conservation of the adiabatic invariant (Landau & Lifshitz 1971)

$$H_\perp = - \frac{c}{\rho} v_\perp, \quad H_1 = \frac{c}{\rho p_\parallel} p_\perp^2,$$

and the synchrotron radiation reaction force (Landau & Lifshitz 1971)

$$F_\perp = -\frac{\alpha}{m c^2} \frac{p_\perp}{p_\parallel} \left(1 + \frac{p_\perp^2}{m^2 c^4} \right), \quad F_\parallel = -\frac{\alpha}{m c^2} \frac{p_\parallel^2}{p_\perp},$$  

where $\alpha = e^2 \omega_B / 3c^2$, $p_\perp$ and $p_\parallel$ are transversal and longitudinal components of the momentum, respectively.

The effect of the aforementioned dissipative forces decreases the pitch angles, which inevitably causes attenuation of the corresponding synchrotron emission process. The situation drastically changes due to the influence of diffusion on the longitudinal and transversal components of momentum, since this process tries to increase the pitch angles. Under certain conditions diffusion and dissipation forces might balance each other and, as a result, the synchrotron radiation process will be maintained. The QLD influences the particle distribution function and the corresponding kinetic equation can be written as (Chkheidze et al. 2011; Machabeli & Usov 1979; Malov & Machabeli 2001)

$$\frac{df^0(p)}{dt} + \frac{\partial}{\partial p_x} \left\{ \left[ F_x + H_1 \right] f^0(p) \right\} + \frac{\partial}{\partial p_\perp} \left\{ p_\perp \left[ F_\perp + H_\perp \right] f^0(p) \right\}$$

$$= \frac{\partial}{\partial p_x} \left\{ p_x D_{x,x} \frac{df^0(p)}{dp_x} \right\} + \frac{\partial}{\partial p_\perp} \left\{ p_\perp D_{\perp,\perp} \frac{df^0(p)}{dp_\perp} \right\}, $$

where $f^0(p)$ is the distribution function, and

$$D_{x,x} = 8 \pi \left( \frac{u_x}{c} \right)^2 \left( \frac{m e}{c} \right)^2 W F_x,$$

$$D_{\perp,\perp} = 8 \pi \left( \frac{u_\perp}{c} \right)^2 \left( \frac{m e}{c} \right)^2 W F_\perp,$$

are the diffusion coefficients. Here $W$ is the energy density of the excited waves and we have taken into account that $D_{x,x} = D_{\perp,\perp} = 0$ (Shapakidze et al. 2003).

It should be noted that pitch angles $\psi$ acquired by particles during the QLD are small enough to apply the condition, $\partial / \partial p_\perp \gg \partial / \partial p_\parallel$, which reduces the kinetic equation (7) to the following form:

$$\frac{df^0(p)}{dt} + \frac{\partial}{\partial p_\perp} \left\{ p_\perp \left[ F_\perp + H_\perp \right] f^0(p) \right\} = \frac{\partial}{\partial p_\perp} \left\{ p_\perp D_{\perp,\perp} \frac{df^0(p)}{dp_\perp} \right\}. $$

On the light cylinder lengthscales one can simplify Equation (9) by estimating the ratio $F_\perp / H_\perp$. Taking into account parameters typical for the light cylinder area, one obtains

$$\frac{H_\perp}{F_\perp} \approx 350 \times \frac{M_{8}}{A_{0.1} \times L_{40}} \times \left( \frac{10^{-3} \text{ rad}}{\psi} \right)^2 \times \frac{10^7}{\gamma_B} \frac{R_{lc}}{\rho},$$

where $A_{0.1} = \Omega / (0.1 \Omega_{\text{max}})$, $\Omega$, and $\Omega_{\text{max}} \approx c^3 / (GM_{\text{BH}})$ are the dimensionless angular velocity, the actual angular velocity, and the maximum angular velocity of the SMBH, $M_{\text{BH}}$ is its mass, $G \approx 6.7 \times 10^{-8} \text{ dyne cm}^2 \text{ g}^{-2}$ is the gravitational constant, $L_{40} = L / 10^{40} \text{ erg s}^{-1}$ and $L$ are the dimensionless and actual values of luminosity of AGN, $M_{8} = M_{\text{BH}} / (10^8 M_{\odot})$ is its dimensionless mass, $M_{\odot}$ is the solar mass, and $R_{lc} = c / \Omega$ is the light cylinder radius.

In general, it is evident from Equation (10) that depending on physical parameters, one has two major regimes: (I) $F_\perp \ll H_\perp$ and (II) $H_\perp \ll F_\perp$. 

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In the framework of the first approximation, one can neglect its contribution of $F_{\perp}$ in Equation (9), which, for the stationary case ($\partial/\partial t = 0$), has the following solution:

$$
 f_I(p_{\perp}) = C \exp \left( \int \frac{H_{\perp}}{D_{\perp}} \, dp_{\perp} \right) = Ce^{-\left(\frac{p_{\perp}}{p_{\perp0}}\right)^4},
$$

(11)

where

$$
 p_{\perp0} = \left( \frac{2 \rho D_{\perp}}{c} \right)^{1/2}.
$$

(12)

As we see from Equation (11), particles are distributed differently for different values of transverse momentum, and therefore, it is natural to examine the average value of $p_{\perp}$, which in turn defines the mean value of the pitch angles. A straightforward calculation leads to the following expression:

$$
 \langle p_{\perp} \rangle_I = \frac{\int_0^\infty p_{\perp} f_I(p_{\perp}) \, dp_{\perp}}{\int_0^\infty f_I(p_{\perp}) \, dp_{\perp}} = \frac{p_{\perp0}^I}{\sqrt{\pi}},
$$

(13)

and the corresponding average value of the pitch angle

$$
 \langle \psi_I \rangle = \frac{\langle p_{\perp} \rangle_I}{p_{\parallel}} = \frac{p_{\perp0}^I}{\sqrt{\pi} p_{\parallel}}.
$$

(14)

For the second case ($H_{\perp} \ll F_{\perp}$), the distribution function reduces to

$$
 f_{II}(p_{\perp}) = C \exp \left( \int \frac{F_{\perp}}{D_{\perp}} \, dp_{\perp} \right) = Ce^{-\left(\frac{p_{\perp}}{p_{\perp0}^II} \right)^4},
$$

(15)

where

$$
 p_{\perp0}^{II} = \left( \frac{4 \gamma p_m c^3 D_{\perp}}{\alpha} \right)^{1/4},
$$

(16)

and the corresponding mean values of the transverse momentum and the pitch angle, respectively, are given by

$$
 \langle p_{\perp} \rangle_{II} = \frac{\int_0^\infty p_{\perp} f_{II}(p_{\perp}) \, dp_{\perp}}{\int_0^\infty f_{II}(p_{\perp}) \, dp_{\perp}} = \frac{\sqrt{\pi}}{4 \Gamma(\frac{3}{4})} p_{\perp0}^{II},
$$

(17)

$$
 \langle \psi_I \rangle_{II} = \frac{\langle p_{\perp} \rangle_{II}}{p_{\parallel}} = \frac{\sqrt{\pi}}{4 \Gamma(\frac{3}{4})} \frac{p_{\perp0}^{II}}{p_{\parallel}},
$$

(18)

where $\Gamma(x)$ is the gamma function.

3. DISCUSSION

As we see from the results of the previous section, the ChDI strongly influences the particle distribution by means of the QLD and prevents the synchrotron emission. In this section we apply the developed model to AGNs and study the production of the nonthermal emission.

For this purpose we consider the light cylinder zone of the magnetosphere. In this area the value of the magnetic induction might be estimated in the framework of the equipartition approximation. Thus, we assume that the magnetic field energy density is of the order of plasma energy density, which defines the magnetic induction

$$
 B_{lc} \approx 5.5 \times A_{0.1} \times L_{40}^{1/2} \text{G}.
$$

(19)

As we see, the magnetic field is quite strong and since in this area particles normally achieve very high Lorentz factors, $\sim 10^6 - 10^7$, the resulting synchrotron emission should be quite efficient. In particular, relativistic electrons moving in a strong magnetic field and having the pitch angles expressed by Equations (14) and (18) will emit photons with energies (Rybicki & Lightman 1979):}

$$
 \epsilon_w^{I} \approx 1.2 \times 10^{-11} \gamma_b^2 B_{lc} \frac{1}{\sqrt{\pi}} \frac{p_{\perp0}^I}{p_{\parallel}},
$$

(20)

$$
 \epsilon_w^{II} \approx 3 \times 10^{-12} \gamma_b^2 B_{lc} \frac{\sqrt{\pi}}{\Gamma(\frac{3}{4})} \frac{p_{\perp0}^{II}}{p_{\parallel}},
$$

(21)

In the framework of the QLD, the problem is usually treated by means of the method of iteration. It is evident that the instability must be saturated by means of nonlinear effects (the corresponding study is beyond the scope of the paper) and since there is some energy budget, it is clear that after the process is relaxed, the energy of waves must be of the same order as that of plasmas associated with the energy budget (Malov & Machabeli 2001). Therefore, in the framework of the paper, we assume $W \sim \gamma_b n_b$.

As a first example, we consider the case when $H_{\perp}$ exceeds the corresponding component of the radiation reaction force. By combining Equations (3), (8), (12), and (20), one can study the emission characteristics. In Figure 1 we show the dependence of energy of radiated synchrotron photons on the beam electron Lorentz factors for different values of plasma component Lorentz factors in the case of $H_{\perp} \gg F_{\perp}$. The set of parameters is $A_k = 1, M_b = 1, A_{0.1} = 1, L = 10^{40} \text{erg s}^{-1}, \gamma_p \in [1; 3; 10], n_b = n_{g3}$.

![Figure 1](image.png)

Figure 1. Behavior of energy of radiated synchrotron photons with the beam electron Lorentz factors for different values of plasma component Lorentz factors in the case of $H_{\perp} \gg F_{\perp}$. The set of parameters is $A_k = 1, M_b = 1, A_{0.1} = 1, L = 10^{40} \text{erg s}^{-1}, \gamma_p \in [1; 3; 10], n_b = n_{g3}$.
QLD guarantees emission in the keV energy domain (X-rays), which in turn must be strongly correlated with the lower energy emission provided by ChDI. From Equations (1) and (4) one can obtain the energy of photons corresponding to the Cherenkov-drift waves:

\[
\epsilon_{\text{Ch}}^b \approx 0.06 A_{0.1} \gamma_p^{-3/2} \left( \frac{n_b}{n_{GJ}} \right)^{1/2} \left( \frac{L}{10^{40} \text{ erg s}^{-1}} \right)^{1/2} \quad (22)
\]

We see that for the mentioned physical parameters the Cherenkov instability produces the IR photons with energies \((2 \times 10^{-3} - 6 \times 10^{-2})\) eV, and since the QLD is achieved by means of the feedback of these waves on particles, it is evident that the IR emission must be strongly connected to the soft X-ray radiation (1–12) keV.

Generally speaking, the instability is supposed to be efficient if the corresponding growth rate is high enough. From Equation (3) we obtain

\[
\Gamma_k \approx 6.5 \times 10^{-3} \times \gamma_p^{-1/2} \times \gamma_p \left( \frac{n_b}{n_{GJ}} \right)^{1/2} \times \left( \frac{L}{10^{40} \text{ erg s}^{-1}} \right)^{1/2} \text{ s}^{-1}.
\]

It is clear that the instability timescale, \(t_{\text{inst}} \sim 1/\Gamma_k\), is of the order of 500 s. On the other hand, particles are moving inside the magnetosphere and the kinematic timescale (escape timescale) \(t_{\text{kin}} \sim R/c \approx 5 \times 10^3 \text{ s}\) exceeds the instability timescale, which indicates high efficiency of the Cherenkov-drift mechanism.

From Equation (10) it is clear that \(H_0/F_0\) is sensitive to the Lorentz factor of beam components and the pitch angles (note that \(\psi\) itself nontrivially depends on \(\gamma_p\); see Equations (14) and (18)) and hence, for other parameters, the relation between forces may be different. Therefore, it is reasonable to consider another limit \((H_0 \ll F_0)\) and see what happens in this particular situation.

In Figure 2 we show the dependence of energy of synchrotron emission on beam Lorentz factors for different values of \(\gamma_p\). The set of parameters is \(A_0 = 1, M_b = 1, A_{0.1} = 1, L = 10^{40} \text{ erg s}^{-1}, \gamma_p \in [1; 3; 10], \ n_b = n_{GJ}\). The difference from the previous case is \(\gamma_p\), which varies from \(6 \times 10^6\) to \(10^7\). From Equation (10) it is clear that for the mentioned parameters the following approximation \(H_0 \ll F_0\) is satisfied. From the plots we see that \(\epsilon_{\text{keV}}\) lies in the interval \(\sim (100\text{–}600)\) keV, which is much higher than the photon energies shown in Figure 1, being a result of a steeper dependence, \(\epsilon_{\text{keV}} \propto \gamma_p^{1.7/3}\) (see Equations (3), (8), (16), and (21)). Since the \(\epsilon_{\text{Ch}}^b\) expression does not depend on the beam Lorentz factors, the hard X-ray radiation is strongly connected with the corresponding Cherenkov-drift emission generated in the same energy domain as in the previous case: IR band.

The emission pattern strongly depends on the values of the AGN luminosity. In particular, in the framework of the model, we use the equipartition magnetic field and the induction becomes dependent on \(L\). Such a dependence is motivated by the fact that the diffusion coefficient behaves as \(1/B^2\) (see Equation (8)), which leads to \(\epsilon_{\text{keV}} \propto L^{-9/8}\) (see Equations (14), (18), (20), and (21)). On the other hand, unlike the results shown on the previous two graphs, in this case \(\epsilon_{\text{Ch}}^b\) also depends on the AGN luminosity, and since synchrotron and Cherenkov-drift emission are generated simultaneously, it is interesting to investigate both.

In Figure 3 we present the behavior of the synchrotron photon energy on the Cherenkov-drift emission for different values of angular momentum. Unlike the previous cases it is worthwhile to consider higher luminosity values: \([0.5\text{–}1] \times 10^{42} \text{ erg s}^{-1}\). We do not examine extreme luminous sources, because in this case the Lorentz factors of relativistic electrons might be less than considered in the present paper. The set of parameters is \(A_0 = 1, M_b = 1, A_{0.1} \in [1; 5; 10], \gamma_p = 10^7, \gamma_p = 1, n_b = n_{GJ}\). It is straightforward to show that the synchrotron radiation reaction force is less than the force responsible for conservation of the adiabatic invariant. It is clear from the plots that \(\epsilon_{\text{keV}}(\epsilon_{\text{Ch}}^b)\) is a continuously decreasing function, which is a direct result of the fact that with bolometric luminosity \(\epsilon_{\text{keV}}\) decreases (\(\propto L^{-9/8}\)), whereas the energies of Cherenkov-drift photons behave as \(L^{1/2}\) (see Equation (22)). We also show the plots for different values of angular momentum of the SMBH. As can be clearly seen, the bigger the angular momentum, the bigger the resulting emission energy for both mechanisms. In
particular, from Equations (22), (8), (16), and (21)) one can see that $\epsilon_{\text{keV}}$ and $\epsilon^{\text{Ch}}_{\text{eV}}$ both are increasing functions of $A$. According to the present results, we see that for the mentioned parameters, the QLD provides the generation of synchrotron emission from the soft up to hard X-rays: $\sim(10-250)\text{ keV}$. These energies are strongly connected with the Cherenkov-drift emission in the energy interval $\sim(1-14)\text{ eV}$, corresponding to emission from IR up to extreme UV.

4. SUMMARY

The main aspects of the present work can be summarized as follows.

1. In the present paper we studied the role of Cherenkov-drift emission in maintaining the synchrotron regime despite the efficient energy losses.
2. It is shown that in the magnetospheres of SMBHs the excited ChDI is efficient enough to influence the particle distribution by means of the QLD.
3. In the framework of the model, we examine equations governing the process of the QLD. It is shown that for physically reasonable parameters, the synchrotron radiation reaction force prevails over other dissipative factors. By taking into account this fact, the corresponding kinetic equation is analytically solved and the emission characteristics are estimated.
4. The emission pattern is investigated for a variety of physical parameters. We found that for typical AGNs the QLD might guarantee the excitation of strongly coupled Cherenkov-drift emission in the eV domain and the synchrotron radiation in the keV energy band.

As we have already seen, the ChDI might guarantee the maintenance of the synchrotron emission process for a variety of AGNs. Another important issue we would like to address is the problem of radiation spectral index, which, in some sense, will complete the problem. It is necessary to investigate this task as well; therefore, we plan to examine it in the future.

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