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Magnetically accelerated particles in an anisotropic disk radiation field

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Abstract. We consider the magnetic acceleration of charged particles in rotating magnetospheres of active galactic nuclei (AGNs). The accelerating particle loses its kinetic energy due to the inverse Compton scattering with the photons emitted from the disk. The disk radiation is anisotropic, so that the inverse Compton energy loss of the particle depends sensitively on the direction of motion of the particle. We find that the maximum Lorentz factor the accelerating electron can attain near the light cylinder is mainly determined by the direction of motion of the electron. In the cases of $L_{\rm disk}/L_{\rm Edd} \leq 10^{-2}$, the maximum Lorentz factor of a magnetically accelerated electron can be as high as a few thousand, if the electron is moving close to the normal direction to the disk. The maximum Lorentz factor becomes relatively low if $L_{\rm disk}/L_{\rm Edd} > 10^{-2}$.

Key words. galaxies: active, nuclei – acceleration of particles – accretion disks

1. Introduction

The origin of the nonthermal high energy emission in AGNs can be explained by Fermi-type particle acceleration mechanisms, if seed electrons with Lorentz factors higher than 100 are present. But the pre-acceleration of these seed electrons is still a problem to be solved. In pulsars, centrifugal acceleration of charged particles by rotating magnetic field is a usable acceleration mechanism (Gold 1968, 1969; Gangadhara 1996). The charged particle moves along the rapidly rotating magnetic field lines, and it will be accelerated due to the centrifugal force. The rotating kinetic energy of the pulsar is tapped to fuel the particle through the magnetic field. The particle can be accelerated to close to the light speed near the light cylinder. Such a mechanism is also adopted to explain jet acceleration in accreting black hole systems (Blandford & Payne 1982). This jet acceleration mechanism has been studied extensively by many workers (Pudritz & Norman 1986; Contopoulos & Lovelace 1994; Cao & Spruit 1994; Kudoh & Shibata 1995). The particle can be accelerated if the poloidal field direction is inclined at an angle less than 60° to the radial direction of the disk surface. For a rapidly rotating black hole system, the critical angle could be even as large as 90° (Cao 1997). Gangadhara & Lesch (1997) proposed that the charged particle can be accelerated to a high speed, if it moves along the magnetic field line in the rotating magnetosphere of the AGN. They found that the inverse Compton radiation and/or synchrotron radiation losses of the particle become quite significant close to the light cylinder, which set an upper limit on the Lorentz factor the magnetically accelerated electron can achieve.

Recently, Rieger & Mannheim (2000, 2001) explored this mechanism in detail and estimated the maximum Lorentz factor of the charged particles under the assumption of an idealized rotating magnetosphere. They showed that the maximum Lorentz factor of a charged particle cannot be larger than 1000. Their main conclusions are as follows:

1. The charged particles cannot be accelerated at all in the case of Eddington accretion, because the inverse Compton energy loss dominates the energy gain;
2. The charged particles can be accelerated to a maximum Lorentz factor in the sub-Eddington accretion case, when the energy gain is balanced by inverse Compton energy loss. The typical $\gamma_{\max}$ is in the range of 100 to 1000;
3. In the very low accretion case ($L_{\rm disk}/L_{\rm Edd} < 10^{-3}$), although the inverse-Compton energy loss is rather unimportant, the acceleration of the charged particles is nevertheless limited to Lorentz factors of a few thousand due to the breakdown of the bead-on-the-wire approximation.

A simple isotropic disk radiation field has been adopted in their calculations, which simplifies the problem. In reality, the disk radiation field should be anisotropic, so that the inverse Compton energy loss will depend on the direction of motion of the electron. In this work, we extend

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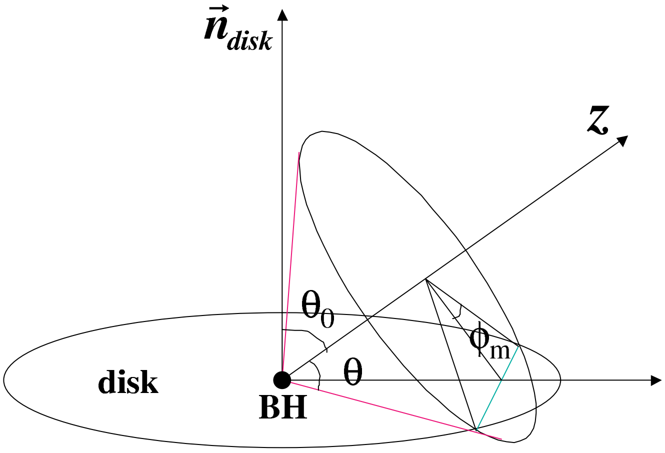

Fig. 1. The geometry used in the calculations.

their calculations on particle acceleration in an anisotropic disk radiation field. Equations describing the problem are given in Sect. 2. Sections 3 and 4 contain the results and discussion.

2. Basic equations

We consider an anisotropic radiation field from the accretion disk I_ν and assume that an electron with Lorentz factor γ moves along the magnetic field with an inclined angle θ_0 to the normal direction to the disk (see Fig. 1). In the electron rest frame $[S']$, the number of photons scattered by the electron to the direction $(\theta'_{\rm sc},\ \phi'_{\rm sc})$ is

$$\mathrm{d}N'_{\rm sc} = \frac{I'_{\nu'}}{h\nu'}\frac{\mathrm{d}\sigma(\delta')}{\mathrm{d}\Omega'_{\rm sc}}\mathrm{d}\Omega'_{\rm sc}\mathrm{d}\Omega'\mathrm{d}\nu'\mathrm{d}t', \tag{1}$$

where δ' is the angle between the directions of incoming photons and scattered photons, and $\frac{\mathrm{d}\sigma(\delta')}{\mathrm{d}\Omega'_{\rm sc}}$ is the differential scattering cross section in $[S']$:

$$\frac{\mathrm{d}\sigma(\delta')}{\mathrm{d}\Omega'_{\rm sc}} = \frac{r_0^2}{2}(1+\cos^2\delta') = \frac{3}{16\pi}\sigma_{\rm T}(1+\cos^2\delta'), \tag{2}$$

where r_0 is the classical electron radius, and $\sigma_{\rm T}$ is the Thomson cross-section.

Since the scattering number is a Lorentz invariant, we can obtain the scattering number measured in inertial frame $[S]$:

$$\mathrm{d}N_{\rm sc} = \frac{I_\nu}{h\nu}\frac{\mathrm{d}\sigma(\delta')}{\mathrm{d}\Omega'_{\rm sc}}\mathrm{d}\Omega'_{\rm sc}\mathrm{d}\Omega\mathrm{d}\nu\mathrm{d}t(1-\beta\cos\theta), \tag{3}$$

where β is the velocity of the electron, θ is the angle between the direction of the incoming photons and the direction of motion of the electron measured in frame $[S]$ (see Fig. 1, Rybicki & Lightman 1979; Jones 1965).

The power of a scattered photon is available:

$$\begin{aligned}P_{\rm sc} &= \int h\nu_{\rm sc}\frac{\mathrm{d}N_{\rm sc}}{\mathrm{d}t}\\ &= \int h\nu_{\rm sc}\frac{I_\nu}{h\nu}\frac{\mathrm{d}\sigma(\delta')}{\mathrm{d}\Omega'_{\rm sc}}\mathrm{d}\Omega'_{\rm sc}\mathrm{d}\Omega\mathrm{d}\nu(1-\beta\cos\theta),\end{aligned} \tag{4}$$

where $\nu_{\rm sc}$ is the frequency of scattered photons measured in $[S]$. In the cases we considered, $\gamma \gg 1$, i.e. $\beta \to 1$ is always satisfied. So we have

$$\delta' \simeq \pi - \theta'_{\rm sc}. \tag{5}$$

For a typical disk radiation field in AGNs, $\gamma h\nu \ll m_{\rm e}c^2$. So we can also approximate the frequency of scattered photon as:

$$\nu_{\rm sc} \simeq \nu\gamma^2(1-\beta\cos\theta)(1+\beta\cos\theta'_{\rm sc}), \tag{6}$$

where ν is the frequency of the incoming photon measured in $[S]$. Substituting Eqs. (5) and (6) into Eq. (4), we can rewrite Eq. (4) as

$$\begin{aligned}P_{\rm sc} = \frac{3}{8}\sigma_{\rm T}\gamma^2 \int (1+\cos^2\theta'_{\rm sc})(1+\beta\cos\theta'_{\rm sc})\mathrm{d}\cos\theta'_{\rm sc}\\ \int(1-\beta\cos\theta)^2\mathrm{d}\Omega\int I_\nu\mathrm{d}\nu.\end{aligned} \tag{7}$$

For simplicity, we approximate the anisotropic disk radiation field as

$$I_\nu = I_{0\nu}\cos\alpha, \tag{8}$$

where $I_{0\nu}$ is constant, α is the angle between the disk normal direction and the direction of motion of the incoming photon measured in $[S]$. It is given by

$$\cos\alpha = -\sin\theta\cos\phi\sin\theta_0 + \cos\theta\cos\theta_0, \tag{9}$$

where $(\theta,\ \phi)$ refers to the direction of motion of the incoming photon measured in $[S]$.

Assuming the accretion disk has an infinite parallel plane geometry, we can integrate Eq. (7) over $\theta'_{\rm sc}$:

$$\begin{aligned}P_{\rm sc} &= \gamma^2\sigma_{\rm T}I_0\int(-\sin\theta\sin\theta_0\cos\phi+\cos\theta\cos\theta_0)\\ &\quad\times(1-\beta\cos\theta)^2\mathrm{d}\Omega\\ &= \gamma^2\sigma_{\rm T}I_0\int_0^{\frac{\pi}{2}+\theta_0}(1-\beta\cos\theta)^2\sin\theta\mathrm{d}\theta\\ &\quad\int_{\phi_{\rm m}}^{2\pi-\phi_{\rm m}}(-\sin\theta\sin\theta_0\cos\phi+\cos\theta\cos\theta_0)\mathrm{d}\phi,\end{aligned} \tag{10}$$

where $I_0 = \int I_{0\nu}\mathrm{d}\nu$, $\phi_{\rm m}$ is related to θ and θ_0 by (Ghisellini et al. 1991)

$$\phi_{\rm m} = \begin{cases} 0, & 0 \le \theta \le \frac{\pi}{2}-\theta_0\\ \arccos\left(\frac{\cos\theta\cos\theta_0}{\sin\theta\sin\theta_0}\right), & \frac{\pi}{2}-\theta_0 < \theta < \frac{\pi}{2}+\theta_0.\end{cases} \tag{11}$$

Integrating Eq. (10) over ϕ, we obtain

$$P_{\rm sc} = \gamma^2\sigma_{\rm T}I_0 f(\theta_0,\beta), \tag{12}$$

where

$$\begin{aligned}f(\theta_0,\ \beta) = \int_0^{\frac{\pi}{2}+\theta_0}&[2\sin^2\theta\sin\theta_0\sin\phi_{\rm m}\\ &+(2\pi-2\phi_{\rm m})\sin\theta\cos\theta_0\cos\theta](1-\beta\cos\theta)^2\mathrm{d}\theta.\end{aligned} \tag{13}$$

Finally, we have the net energy loss due to the inverse Compton scattering with the soft photons emitted from the disk:

$$P_{\rm sc}^{\rm IC} = (\gamma^2 - 1)\sigma_{\rm T} I_0 f(\theta_0, \beta)\cdot \quad (14)$$

For the anisotropic disk radiation field we adopted, the energy flux F is

$$F = \int I_\nu {\rm d}\nu {\rm d}\Omega = \int I_0 \cos\alpha {\rm d}\Omega = \pi I_0. \quad (15)$$

Assuming the luminosity of the disk to be $L_{\rm disk}$, we have:

$$\pi I_0 \simeq \frac{L_{\rm disk}}{4\pi r_{\rm L}^2} \quad (16)$$

where $r_{\rm L}$ is the light cylinder radius of the rotating magnetosphere.

So the cooling time scale for the inverse Compton energy loss can be expressed as:

$$t_{\rm cool}^{\rm IC} = \frac{\gamma m_{\rm e} c^2}{P_{\rm sc}^{\rm IC}}\cdot \quad (17)$$

The accelerating time scale is given by Rieger & Mannheim (2000):

$$t_{\rm acc} = \frac{\gamma}{\dot\gamma} = \frac{\sqrt{1-\Omega^2 r^2}}{2\Omega^2 r\sqrt{1-\tilde m(1-\Omega^2 r^2)}}\cdot \quad (18)$$

where $\tilde m = (1-\Omega^2 r_0^2 - v_0^2)/(1-\Omega^2 r_0^2)^2$, r_0, v_0 are the initial position and velocity of the injected particle, and Ω is the angular velocity of the rotating magnetic field lines.

Thus, we can obtain the maximum Lorentz factor of the electron $\gamma_{\rm max}$ by equating these two time scales given by Eqs. (17) and (18).

3. Results

In order to show the influence of the anisotropic disk radiation field on the inverse Compton energy loss, we plot $f(\theta_0,\ \beta)$ as functions of θ_0 in Fig. 2. We note that value of $f(\theta_0,\ \beta)$ is insensitive to the Lorentz factor γ, if $\gamma > 10$. (For a large γ, $\beta \to 1$ in Eq. (13).) $f(\theta_0,\ \beta)$ increases rapidly with θ_0, i.e., the inverse Compton energy loss is strongly reduced in small θ_0 case (close to the normal direction to the disk).

In our calculations of the maximum electron Lorentz factor $\gamma_{\rm max}$, we use the same parameters as those used by Rieger & Mannheim (2000) for comparison: a light cylinder radius $r_{\rm L} \simeq 10^{15}$ cm, an initial position $r_0 \simeq 0.4 r_{\rm L}$ and an initial velocity $v_0 \simeq 0.6c$ of the injected electron. We assume the mass of the central black hole in the AGN $M_{\rm BH} = 10^8\ M_\odot$, where $M_\odot$ denotes the solar mass. We also express the disk luminosity as $L_{\rm disk} = l_{\rm e} \times L_{\rm Edd}$, and $L_{\rm Edd} \simeq 10^{46}$ ergs s^{-1} is the maximum luminosity corresponding to the mass of the black hole.

The two time scales $t_{\rm cool}^{\rm IC}$ and $t_{\rm acc}$ as functions of the Lorentz factor γ are shown in Fig. 3. We use $l_{\rm e} = 5\times 10^{-3}$ and $\theta_0 = 30°$, $60°$, $90°$ in the calculations. In this figure, $t_{\rm acc}$ is denoted by the dotted line, and the other three solid lines correspond to $t_{\rm cool}^{\rm IC}$ for the different incoming directions $\theta_0 = 30°$, $60°$ and $90°$. It can be seen that the electron can attain a high $\gamma_{\rm max}$, when the electron is accelerated close to the normal direction to the disk.

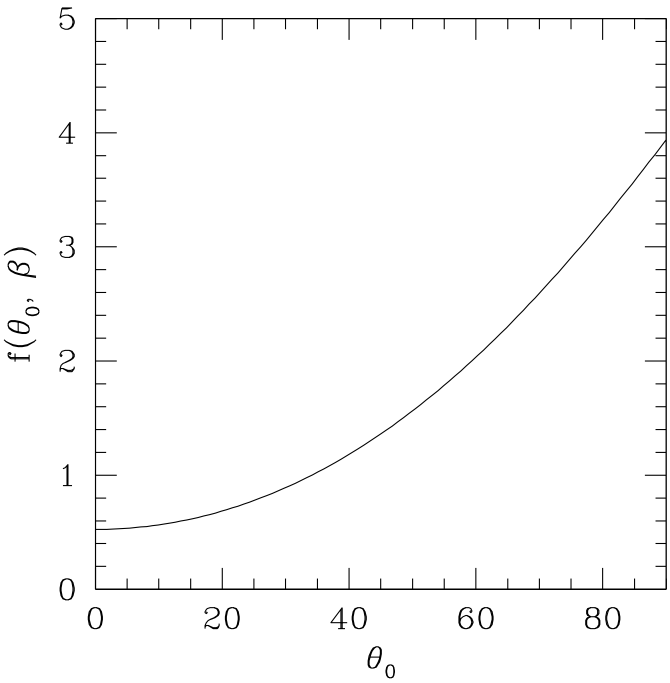

Fig. 2. The dependence of $f(\theta_0, \beta)$ given by Eq. (13) on θ_0, where $\gamma = 1000$ is adopted.

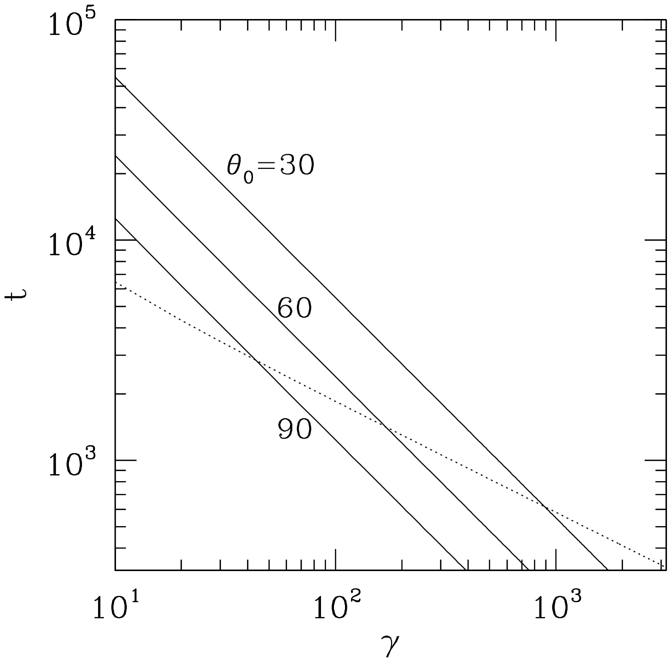

Fig. 3. Acceleration time scale $t_{\rm acc}$(dotted) and inverse Compton cooling time scale $t_{\rm cool}^{\rm IC}$ (Eq. (17)) for $\theta_0 = 30°$, $60°$ and $90°$, as functions of the electron Lorentz factor γ, where $l_{\rm e} = 5\times 10^{-3}$ is adopted.

The maximum Lorentz factor $\gamma_{\rm max}$ as a function of $l_{\rm e}$ for the different electron incoming directions is plotted

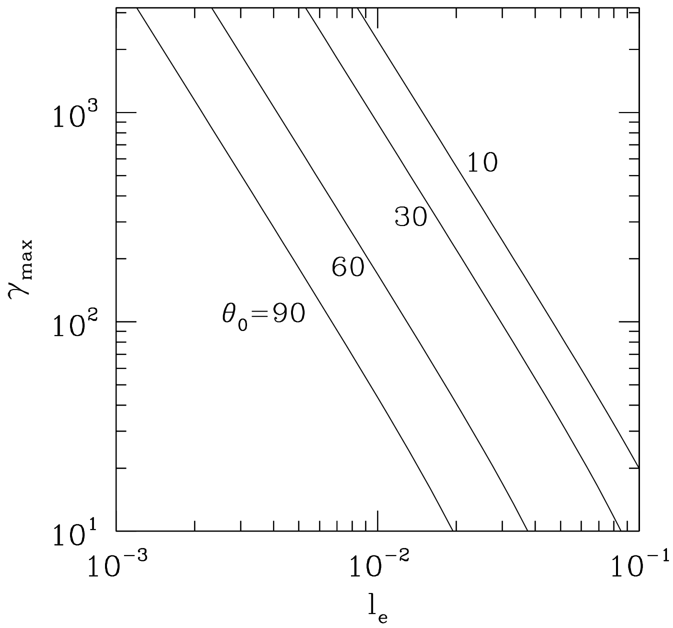

Fig. 4. Maximum electron Lorentz factor as a function of the disk luminosity $l_e = L_{disk}/L_{Edd}$, for $\theta_0 = 90°, 60°, 30°$ and $10°$, respectively.

in Fig. 4. From right to left, each line corresponds to $\theta_0 =$ 10°, 30°, 60° and 90°, respectively.

4. Discussion

We considered the centrifugal particle acceleration mechanism for an electron moving along a rotating magnetic field line in AGNs. The electron gains energy from centrifugal acceleration, while it also loses energy due to the interaction with the photons emitted from the disk (inverse Compton scattering). The maximum Lorentz factor is determined by equating the acceleration time scale and the cooling time scale.

The radiation field of the disk is anisotropic, so the inverse Compton energy loss will depend on the direction of motion of the electron. We therefore calculate the inverse Compton energy loss in the case of an anisotropic radiation field from the accretion disk. We found that the inverse Compton energy loss sensitively depends on the angle θ_0, i.e., the angle between the axis of the accretion disk and the direction of motion of the electron (see Fig. 1). The maximum Lorentz factor of the electron is therefore significantly influenced by this angle. We can see in Fig. 2, that the maximum Lorentz factor γ_{max} varies with the direction of motion of the electron for fixed disk luminosity. γ_{max} is less than 200, when the electron is moving parallel to the disk plane. This is similar to the results given by Rieger & Mannheim (2000). The value of γ_{max} could be higher than 1000, if the angle $\theta_0 < 45°$. It would be very high if the electron moves close to the normal direction of the disk ($\theta_0 \rightarrow 0$). It indicates that the electron can be magnetically accelerated more efficiently if the direction of motion is close to the normal direction to the disk. This is consistent with the well known phenomenon that jets are always perpendicular to the disk plane.

In Fig. 4, we found that the electron can be magnetically accelerated to $\gamma_{max} > 1000$ in sub-Eddington cases ($l_e < 10^{-2}$), if the electron is moving close to the normal direction to the disk ($\theta_0 < 10°$). In the cases of $l_e > 0.1$, the inverse Compton energy loss dominates over the energy gain from the magnetic field line, even if the electron is moving close to the disk axis. It implies that the centrifugal acceleration of the electron is mainly determined by the disk luminosity for high Eddington cases. It is therefore worthwhile to explore the relation between the radio properties and the accretion type in AGNs, which would be a useful test on this acceleration mechanism of electrons. This is beyond the scope of the present work, but it will be considered in future work.

Besides the inverse Compton scattering with the soft photons emitted from the disk, the emission from the broad-line region (BLR) will interact with the electron. The total broad-line luminosity in the AGN is about 10 percent of its optical continuum luminosity (Cao & Jiang 1999, 2001). The radius of the BLR is usually much more larger than light cylinder r_L (Kaspi et al. 2000). So, the inverse Compton energy loss due to the interaction with the soft photons from the BLR can be neglected compared with that caused by the disk radiation, even if the electron is moving close to the disk axis.

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