ELECTRON ACCELERATION AND THE PRODUCTION OF NONTHERMAL ELECTRON DISTRIBUTIONS IN ACCRETION DISK CORONAE

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ABSTRACT

We consider electron acceleration by obliquely propagating fast-mode waves in magnetically dominated accretion disk coronae. For low coronal plasma densities, acceleration can exceed Coulomb drag at lower energies and energize electrons out of the thermal background, which results in a nonthermal tail. The extent of this tail is determined by the balance between acceleration and radiative cooling via inverse Compton scattering and synchrotron emission and usually goes out to tens of MeV. This will have direct applications for explaining the gamma rays from several Galactic black hole candidates, such as Cyg X-1 and GRO J0422, which show 0.5–5 MeV emissions in excess of what most thermal models predict. Detailed time evolutions of the particle distributions and wave spectra are also presented.

Subject headings: acceleration of particles — accretion, accretion disks — gamma rays: theory — waves

1. INTRODUCTION

Thermal Comptonization models have had much success in explaining the hard X-ray spectra from Galactic black hole candidates (GBHCs) with plasma temperature ~50–100 keV and Thomson depths of a few (see, e.g., Shapiro, Lightman, & Eardly 1976; Sunyaev & Titarchuk 1980; Harmon et al. 1994; Liang 1993). However, the most sensitive observations of GBHCs to date in the 0.5–5 MeV range by Compton Gamma-Ray Observatory have clearly revealed that persistent gamma rays (>1 MeV) are being produced in some GBHCs, notably Cyg X-1 (McConnell et al. 1997; Phlip et al. 1996; Ling et al. 1997) and GRO J0422 (van Dijk et al. 1995). These gamma-ray emissions are very difficult to accommodate with the pure thermal models (see, e.g., Sunyaev & Titarchuk 1980; Titarchuk 1994), which strongly suggests the need for modification (Skibo & Dermer 1993) or incorporation of some nonthermal processes.

Recently, Li, Kusunose, & Liang (1996) have proposed a gyroresonant stochastic electron acceleration model to account for the MeV emissions from GBHCs (see also Dermer, Miller, & Li 1996). In that model, they showed that high-frequency whistlers can accelerate electrons directly from the thermal background, after which Alfvén waves would continue the acceleration to higher energies. The result is an electron distribution with a hard non-Maxwellian tail. Compton scatterings from both thermal and nonthermal electrons produce a broadband X-ray to gamma-ray spectrum, in agreement with the observed soft photons (a few keV) in GBHCs (see, e.g., Sunyaev & Titarchuk 1973), from which most of the soft photons (a few keV) originate. In the inner part of the disk, we assume that the magnetic energy density $U_B = B_0^2/8\pi$ is in equipartition with the plasma thermal pressure $n_kT_k$, where $B_0$ is Boltzmann’s constant, $T_r = T_e + T_i$, and $T_i$ and $T_e$ (~50 keV) are the proton and electron temperatures, respectively. $T_i$ is, unfortunately, poorly determined and is chosen to be 10 MeV in this study. The plasma density $n_0$ is taken to be $\sim 1/(\sigma TR)$, where $\sigma_T$ is the Thomson cross section. Since $T_i \gg T_e$, it roughly corresponds to the two-temperature accretion disk situation. However, there have been discussions in the literature questioning whether a two-temperature plasma can occur at all in the accretion plasma (Phinney 1981; Rees et al. 1982; Begelman & Chiueh 1988). We emphasize that we use $T_i = 10$ MeV only to get a fiducial magnetic field value, and, in fact, our model works even better for an isothermal accretion disk because electrons have a much higher thermal speed than protons and are preferentially accelerated (see below).

A tenuous, quasi-spherical extended corona surrounding the hot inner disk is postulated (see, e.g., Liang & Price 1977;
Haardt & Maraschi 1991; Haardt, Maraschi, & Ghisellini 1994). The coronal electron temperature is again taken to be 50 keV initially, and the coronal optical depth \( \tau = n_e \sigma_T \sim 1 \) is varied, where \( n_e \) is the coronal electron density. We further assume that the magnetic field in the corona is the same as that in the disk \( B_0 \), so that the corona is magnetically dominated and the plasma \( \beta = n_e k_B T_e / U_B \ll 1 \). Notice that \( \beta = \tau \), with these conditions. Adopting the black hole mass \( M \) to be 10 \( M_\odot \), and the size \( R \) of the system to be \( \approx 30 \) \( M_\odot \), within which most of the high-energy emissions are produced, we find that the coronal magnetic field \( B_0 \) and dimensionless Alfvén speed \( v_\alpha / c \) are \( \sim 3.7 \times 10^{-3} \) G and \( B_0/(4\pi m_e c^2)^{1/2} = 0.15 \beta_b^{-1/2} \), respectively. The \( \beta < 1 \) condition also implies that the proton thermal speed is always less than \( v_\alpha \).

Here we make a key assumption that a fraction of the total available energy goes into generating MHD turbulence. This wave turbulence consists of fast-mode and shear Alfvén waves (the slow mode will be heavily Landau damped and will probably not be excited), but it is just the fast-mode waves that are relevant for electron acceleration. Fast-mode waves propagating obliquely with respect to the ambient magnetic field have a parallel magnetic field component, which can couple strongly (or resonate) with a particle when the parallel phase speed of the wave \( w_k / \beta_b \) is equal to the parallel component of particle velocity \( v_p \); i.e., when \( v_p = v_\alpha / \gamma \), where we have used the dispersion relation for the fast-mode waves and \( \gamma \) is the cosine of the wave propagation angle. This process is referred to as transit-time damping or magnetic Landau damping (see, e.g., Stix 1992, pp. 273–292; Lee & Völk 1975; Achterberg 1981), and a broadband wave spectrum leads to both rapid wave damping and particle acceleration (MLM). This process is essentially the resonant form of Fermi acceleration (Fermi 1949). Particle interactions with the parallel magnetic field variations (caused by the waves) can be viewed as either head-on (gaining energy) or trailing (losing energy) “collisions.” Head-on collisions occur more often than trailing ones, so that net acceleration results. The resonant condition implies that the acceleration threshold is \( v_\alpha \). For the coronal plasma, \( v_\alpha \) is comparable to the electron thermal speed but greater than the proton thermal speed if \( \beta < 1 \); so that electrons are preferentially accelerated.

2.2. The Electron and Wave Diffusion Equations and Energy Transfer Rates

To simplify the calculations, we consider an isotropic fully ionized hydrogen plasma permeated by a homogeneous background magnetic field \( B_0 \). At some large-scale \( \lambda_{\text{ij}} \) an unspecified mechanism generates fast-mode waves. The evolution of the electron distribution \( N(E) \) is given by the Fokker-Planck equation

\[
\frac{\partial N}{\partial t} = -\frac{\partial}{\partial E} \left[ \left( \frac{dE}{dt} \right)_{\text{acc}} + \left( \frac{dE}{dt} \right)_{\text{loss}} \right] N + \frac{1}{2} \frac{\partial^2}{\partial E^2} (D + D_p) N,
\]

where \( E \) is the kinetic energy. Here we have neglected the escape and (possible) e⁺e⁻ pair production. Those coefficients associated with wave-particle interactions are the systematic mean acceleration rate \( \left( \frac{dE}{dt} \right)_{\text{acc}} = \frac{p^2}{m_e} \frac{dD}{dp} / \delta p \) and the diffusion coefficient \( D = 2v^2 D(p) \), where \( v \) and \( p \) are the electron speed and momentum, and \( D(p) \) is the momentum diffusion coefficient (given below). The other convection term \( \left( \frac{dE}{dt} \right)_{\text{loss}} \) represents the sum of electron energy change rates from inverse Compton scattering and synchrotron (ICS) losses \( \left( \frac{dE}{dt} \right)_{\text{ics}} \) and e⁺e⁻ Coulomb collisions \( \left( \frac{dE}{dt} \right)_{\text{coll}} \), (Dermer & Liang 1989), which also give rise to diffusion \( D \) (Spitzer 1962; Dermer & Liang 1989). These are the most important processes for electrons in our parameter regime. The processes that are neglected include the electron-proton Coulomb interaction since it is much slower than e⁺e⁻ (when \( T_e < 100 \) MeV), the diffusion due to Compton scatterings since the electron recoil is typically small and wave-particle diffusion will dominate at high energies, and the energy gain due to the convergence of the flow since the accretion time is much longer than dynamic timescale \( \sim R/c \). The momentum diffusion coefficient \( D(p) \) for transit-time damping is (MLM)

\[
D(p) = (m_e c^2) \frac{\pi}{16} \left( \frac{v_\alpha^2}{c} \right) \frac{c(k) \xi}{m_e c} \frac{c}{v} F(\mu_0),
\]

where \( m_e \) is the electron mass, \( \xi = U_z / U_B \), \( U_z \) is the fast-mode wave energy density, \( (k) \) is the mean wavenumber of the fast-mode waves, \( \mu_0 = v_\alpha / v \), and \( F(\mu_0) \) is basically an efficiency factor, which equals zero when \( v < v_\alpha \) (i.e., when resonance is impossible) and equals \( -5/4 - (1 + 2\mu_0^2) \ln \mu_0 + \mu_0^2 + 3\mu_0^4 \) otherwise. The evolution of the isotropic wave spectral density \( W \) can be approximated by (see Zhou & Matthaeus 1990)

\[
\frac{\partial W}{\partial t} = \frac{\partial}{\partial k} \left[ k^2 D w \frac{\partial}{\partial k} (k^2 W_k) \right] - \gamma W + Q_w \delta(k - k_0),
\]

where we have included a term \( \gamma(k) \) for the wave damping by the electrons and a term \( Q_w \) for the injection. The wave-damping rate \( \gamma(k) \) can be obtained from the relation \( f_\delta d \kappa \gamma(k) W_k(k) = f_\delta d E \gamma(E) \langle dE/dt \rangle_{\text{acc}} \) (see eq. [2]). Since \( \langle dE/dt \rangle_{\text{acc}} \propto \langle k \rangle \), this implies that large-scale waves (small \( k \)) will suffer little loss and that rapid wave dissipation occurs only when \( k \) is sufficiently large.

At steady state, the volumetric wave energy injection rate \( Q_w \) (ergs cm⁻³ s⁻¹) at \( k_0 = 2\pi / \lambda_{\text{ij}} \) must equal the rate at which energy is transferred to smaller scales \( 2^{1/2} v_\alpha k_0^{3/2} U_B \) (MLM). This energy flow is eventually dissipated at higher wavenumbers by electrons, which in turn produce emissions in the X-ray to gamma-ray range. This implies that the volumetric gamma-ray energy production rate \( Q_w \approx L_\gamma / (4\pi R^2) \), where \( L_\gamma \) is the gamma-luminosity, must be smaller than or equal to (steady state) \( Q_w \). This gives

\[
\frac{U_z}{U_B} \leq \frac{3 v_\alpha}{c} \frac{R}{\lambda_{\text{ij}}} \varepsilon^{1/2},
\]

where \( U_z = L_\gamma / (4\pi R^2 c) \) is the energy density of gamma-ray photons. Letting \( \lambda_{\text{ij}} / R \sim 0.1 \), we obtain \( U_z / U_B \approx 4.5 \beta_b^{-1/2} \varepsilon^{1/2} \). This implies that in order to get \( U_z / U_B \approx 0.1 \) as suggested by the 60–1000 keV luminosity from Cyg X-1 (Phillips et al. 1996), \( \varepsilon \) must be \( \approx 0.08 \beta_b^{-1/2} \), well within the weak turbulence limit.

Next, we look at the electron energy gain and loss rates. The mean acceleration rate of electrons by the waves is given as (MLM)

\[
\left\langle \frac{d\gamma}{dt} \right\rangle = \frac{\pi}{4} \left( \frac{v_\alpha^2}{c} \right) c(k) \xi \left( \frac{p}{m_e c} \right) G(\mu_0),
\]
where $G(\mu) = F(\mu) + (4\mu_0^2 \ln \mu_0 - \mu_0 + 1)/(4\gamma^2)$ when $\mu_0 < 1$ and equals 0 otherwise. The product $(k)\zeta$ has to be obtained from the simulations (see next section). The ICS losses can be written as $(d\gamma/dt)_\text{ics} = \gamma (4/3) \left( \Theta_j / \tau_{\text{dyn}} \right) (1 + U_{\text{ph}}/U_p) (p/m,c)^2$, where $\tau_{\text{dyn}} = R/c$ and $\Theta_j = k_0 T_j/m,c^2$. The soft photon energy density $U_{\text{ph}}$ is fixed to be the same as $U_p$ in this study. Note that this loss rate may be an overestimate for mildly relativistic particles; thus, more careful treatment of the losses will help the acceleration.

When the Coulomb loss timescale is longer than both the acceleration and the ICS cooling timescales, we can define a critical energy $\gamma_c$ at which acceleration balances radiative cooling, $(d\gamma/dt)_\text{ics} = (d\gamma/dt)_\text{rad}$, which gives

$$\gamma_c \approx 1.3 \times 10^4 \left( \frac{R}{4.5 \times 10^2} \right) (k)\zeta, \quad (6)$$

where we have utilized the fact that when $p/(m,c) \sim \gamma_c \gg 1$, $G(\mu_0)/\tau_c$ is almost constant (ranging from 0.76 to 1.1) for $0.1 \leq \gamma \leq 1$. In order to get substantial acceleration—say $\gamma > 10-(k)\zeta$ must be greater than $8 \times 10^{-4} \text{ cm}^{-2}$. This offers a direct test of our simulations.

3. RESULTS

We solve equations (1) and (3) using the Crank-Nicholson method (MLM) and concentrate on the time evolution of particles and waves from the start of wave injection until the steady state is reached, during which waves are constantly injected at $k_0 = 2 \pi/(0.1 R)$. This period turns out to be always less than or comparable to the dynamic timescale $\tau_{\text{dyn}} \approx 1.5 \times 10^{-5} \text{s}$. This validates our assumption of neglecting escape. The Coulomb loss is fixed at the rate that corresponds to the initial Maxwellian. Even though this treatment is not self-consistent, our particle acceleration results should not be affected much since Coulomb loss plays a negligible role compared to the ICS cooling at relativistic energies. The Kolmogorov phenomenology for the wave evolution is assumed so that $W_r = W_{\text{IC}} k^{-5/3}$, with $W_0 = (Q_{\text{in}}/\nu_0) (k_0^3 U_p)^{2/3}$. The injection rate $Q_{\text{in}}$ is chosen as $\approx 2.8 \times 10^4 (\nu_0/c) \text{ ergs cm}^{-3} \text{s}^{-1}$, which corresponds to $\zeta = 0.2$ at steady state, and $U_r = f_0 W k^{-5/3} dk$ should be a constant for all the cases considered here; these are all confirmed by the simulations to within the numerical error.

Figure 1 summarizes the time evolution of the particle density distribution $N(E)$ (upper panels) as a function of kinetic energy $E$ and the corresponding wave spectral density $W_r$ (lower panels) as a function of wavenumber $k$. Three different densities are considered. Each plot has 16 curves in it, corresponding to five evenly spaced time intervals in each of three periods, $t = 0 - 0.04 \tau_{\text{dyn}}$, $0.04 - 0.16 \tau_{\text{dyn}}$, and $0.16 - 0.6 \tau_{\text{dyn}}$, respectively. The particle distributions soften first (e.g., curves 1–8 in the $\tau_c = 0.1$ case) to the fact that acceleration is very inefficient initially since waves have not fully cascaded (i.e., small $k$), as shown in the lower panel) and Coulomb and ICS losses dominate at high energies. As waves cascade over the inertial range, $k$ quickly grows to a level at which acceleration overcomes both Coulomb and ICS losses. Electrons are then energized out of the thermal background, and the nonthermal tail forms. This is indicated, for example, by curves 6–15 in the $\tau_c = 0.1$ case. After that, both the particle and wave spectra gradually reach steady states.

The $\tau_c = 0.1$ case clearly illustrates several other points. From simulations, $(k)\zeta \approx 10^{-3} \text{ cm}^{-1}$; thus, $\gamma_c \sim 10$ using equation (6), which is in perfect agreement with curve 15. Furthermore, the nonthermal tails start to develop only at $E/m,c^2 \sim 0.13$ (corresponding to $\nu_0/c = 0.46$), complying with the acceleration threshold. The threshold energies for $\tau_c = 1$ and 0.5 are buried in the thermal distributions.

That the particle’s high-energy cutoff gets larger as $\tau_c$ decreases can be understood from equation (6). As $\tau_c$ decreases, fewer particles are available to absorb the wave energy, which results in larger $(k)\zeta$. Smaller $\tau_c$ also reduces the number of particles with $\nu > \nu_{\text{ta}}$, as is evident from the upper panels in Figure 1.

Figure 2 shows the fraction of electrons with $E \geq 511 \text{ keV}$ out of the total electron population as a function of coronal
optical depth $\tau_c = 0.1-1$, both in energy content $f_{E511} = \int f(E) dE$ ([upper panel]) and in number density $f_{N511} = \int f(E) dE$ ([lower panel]). The horizontal dashed line indicates the initial values for a 50 keV Maxwellian. When $\tau_c$ is high, electrons are mostly nonrelativistic owing to cooling, but a good fraction of electrons becomes relativistic and nonthermal from the wave acceleration when $\tau_c \approx 0.5$.

4. CONCLUSIONS AND DISCUSSION

We have studied particle acceleration in Galactic black hole accretion disk coronae via interactions between electrons and fast-mode waves—specifically, via the transit-time damping process. Including Coulomb collisions, inverse Compton scattering, and synchrotron losses, we show that particles with speeds higher than the Alfvén speed can be accelerated out of the thermal background, and we obtain steady state particle distributions composed of a Maxwellian plus a nonthermal high-energy tail extending into several tens of MeV. Detailed radiation modeling will be presented in a forthcoming work, and we expect that the Maxwellian and the nonthermal tail of the particle distribution will be responsible for the power-law spectra in the tens of keV and the high-energy gamma rays observed from several GBHCs such as Cyg X-1 and GRO J0422, respectively.

The generation of plasma wave turbulence in accretion disk environments is a fairly unexplored topic, but it is reasonable to suppose that the fast-mode waves will be excited since it is an intrinsic long-wavelength mode of a magnetized plasma. We emphasize that the coronal plasma $\beta$ must be less than 1 for electrons to get most of the wave energy; otherwise, proton acceleration becomes possible, which reduces the energy flow to the electrons.

The particle acceleration mechanism discussed here also has direct implications on the high-energy radiation from accretion disks in AGNs, notably Seyfert galaxies. Preliminary analyses have indicated that most of our results are insensitive to the size of the system; thus, we expect greater than MeV emissions are also being produced in Seyfert galaxies as well, although high-quality spectra above 200 keV are clearly needed to settle this issue firmly.

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