X-Ray Emissions from Three-Dimensional Magnetohydrodynamic Coronal Accretion Flows

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Abstract

We calculated the radiation spectrum and its time variability of a black-hole accretion disk–corona system based on a three-dimensional magnetohydrodynamic (MHD) simulation. In explaining the spectral properties of active galactic nuclei, it is often assumed that they consist of a geometrically thin, optically thick disk and hot, optically thin corona surrounding the thin disk. As for a model of the corona, we adopted the simulation data of three-dimensional, non-radiative MHD accretion flows calculated by Y. Kato and coworkers, while for a thin disk we assumed a standard-type disk. We performed Monte-Carlo radiative transfer simulations of the corona, while taking into account the Compton scattering of soft photons from the thin disk by hot thermal electrons and coronal irradiation heating of the thin disk, which emits blackbody radiation. By adjusting the density parameter of the MHD coronal flow, we could produce emergent spectra that are consistent with those of typical Seyfert galaxies. Moreover, we found rapid time variability in the X-ray emission spectra, originating from a density fluctuation produced by a magnetorotational instability in the MHD corona. The features of reflection component including iron fluorescent line emission are also briefly discussed.

Key words: accretion, accretion disks — black hole physics — radiative transfer — X-rays: general

1. Introduction

Thanks to rapid progress in observational studies in recent years, our understanding of the radiation properties of black-hole accretion flows has really deepened. Accreting black holes, such as active galactic nuclei (AGNs) and black hole binaries (BHBs) during their very high spectral state [state with luminosities around a few tenths of the Eddington limit ($L_{\text{Edd}}$)], show the radiation spectra to be dominated by two components: a thermal bump in the UV/soft X-ray band and power-law emission with a spectral index of $\alpha \approx 1$ in the X-ray band (possibly with a high-energy cutoff around MeV; for the multiwavelength spectrum of a typical AGN, see Reynolds et al. 1997). These components are often explained by a disk–corona model (Liang & Price 1977; Bisnovatyi-Kogan & Blinnikov 1977; Haardt & Maraschi 1991, 1993). In this model, the accretion flow consists of a geometrically thin, optically thick accretion disk whose structure has been studied by Shakura and Sunyaev (1973), and hot, optically thin corona surrounding the disk. The thermal bump is believed to be thermal emission from an optically thick disk (see Kishimoto et al. 2005 for its observational implication), and a power-law component is interpreted to be formed by photons that are emitted from the disk and Compton up-scattered by hot electrons in the corona (see Thorne & Price 1975; Shapiro et al. 1976, but in the context of the inner hot accretion flow model). Such a structure in which hot gas (i.e., the corona) coexists with cool gas (i.e., the disk) is justified by the existence of a reflection component observed in the X-ray spectra (Pounds et al. 1990; see also Guilbert & Rees 1988; Lightman & White 1988). This reflection component is accompanied by an iron fluorescent line emission broadened relativistically (for reviews, see Mushotzky et al. 1993; Fabian et al. 2000; Reynolds & Nowak 2003).

On the other hand, our theoretical understanding of black-hole accretion has also made rapid progress. The dynamics of accretion flows is presently understood in terms of magnetohydrodynamics (MHD), since Balbus and Hawley (1991) rediscovered magnetorotational instability (MRI) as being the fundamental mechanism of angular-momentum transfer in accretion disks. Detailed dynamical features of MHD accretion flows have been investigated via global three-dimensional numerical simulations (Matsumoto 1999; Stone & Pringle 2001; Machida et al. 2001; Hawley & Krolik 2001, 2002; Machida & Matsumoto 2003; Armitage & Reynolds 2003; De Villiers et al. 2003; Gammie et al. 2003; Igumenshchev et al. 2003). Such magnetically dominated accretion flows as those simulated in these calculations are radiatively inefficient accretion flows (RIAFs), whose mass accretion rates are much smaller than the critical value of $\sim L_{\text{Edd}}$ (see, e.g., Mineshige et al. 2002). In such flows, the dissipated energy would not be radiated away efficiently, because of their low density, and would be advected inward to the central black hole (Ichimaru 1977; Narayan & Yi 1994; Abramowicz et al. 1995; Kato et al. 1998).

Although the detailed structures and behavior of MHD accretion flows are shown by numerical methods, it is still unclear whether these simulational results can be applied to realistic situations in the universe. In recent years, research concerning the observational properties, like the spectra and
their time variabilities expected from numerically simulated accretion flows, has been extensively performed. Since the RIAF model is believed to fit the emergent spectrum of Sagittarius A* (Hawley & Balbus 2002; Goldstone et al. 2005; Ohshiga et al. 2005; Moscibrodzka et al. 2007). However, as for accretion flows with a moderately high mass accretion rate, which are considered to be a good model for AGNs and BHs showing a thermal bump and power-law emission in their spectra, no attempts have been made so far to calculate the spectra predicted from MHD simulations that take into account realistic radiation processes. As noted above, in order to reproduce such spectra, the simulated accretion flows should have two components: a geometrically thin cool disk and optically thin hot corona. However, no simulation data that incorporate such two-component effects are available.

In this study we calculated for the first time the emergent spectra of two-component accretion flows based on a three-dimensional MHD simulation by Kato et al. (2004, hereafter KMS04). Instead of fully solving the dynamics of two-component accretion flows, we adopted the simulation result of RIAF-like MHD accretion flow for an optically thin corona, and assumed that the optically thick, geometrically thin disk is embedded in the corona with mutual interaction through radiation. Other interactions, such as mass evaporation/condensation or heat conduction, were neglected, for simplicity (Meyer & Meyer-Hofmeister 1994; Meyer et al. 2000; Liu et al. 2002a, 2007). The disk is emitting soft photons with a thermal spectrum, and those photons are up-scattered by hot electrons in the corona. After being scattered in the corona, a part of the upscattered photons return to the disk and are thermalized there, thereby heating the accretion disk. The disk emission is, hence, enhanced by this returning process. In the present work we performed three-dimensional Monte-Carlo radiative transfer simulations to properly calculate such radiation processes and the emergent spectra.

The plan of this paper is as follows. We show the details of the model and method used in our calculation in section 2. The results are presented in section 3, and are compared with the spectra observed in typical AGNs. We also discuss the similar points and discrepancies between them in section 4. In section 5 we summarize our study.

2. Model and Calculation Methods

2.1. Overview of Adopted MHD Simulations

KMS04 investigated the evolution of a torus threaded by weak localized poloidal magnetic fields by performing a three-dimensional MHD simulation. They solved the following basic equations of the resistive MHD in cylindrical coordinates, \( r, \phi, z \):

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0, \tag{1}
\]

\[
\frac{\partial (\rho \mathbf{v})}{\partial t} + \nabla \cdot \left( \rho \mathbf{v} \otimes \mathbf{v} - \frac{\mathbf{B} \otimes \mathbf{B}}{4\pi} \right) = -\nabla \left( \rho \mathbf{v} \right) - \rho \nabla \psi, \tag{2}
\]

\[
\frac{\partial \mathbf{B}}{\partial t} = -\mathbf{v} \times \mathbf{E}, \tag{3}
\]

\[
\frac{\partial \mathbf{E}}{\partial t} = -\mathbf{B} \times \mathbf{v}, \tag{4}
\]

where \( \psi = -GM/(R - r_s) \) is the pseudo-Newtonian potential (Paczynski & Wiita 1980), \( \varepsilon = \rho v^2/2 + \rho \mathbf{v} \cdot \mathbf{B} / (\gamma - 1) \) is the energy of the gas (here \( \gamma \) is fixed to be 5/3), and \( \mathbf{E} = -(\varepsilon/c) \times \mathbf{B} + (4\pi \eta/c^2) \mathbf{J} \) is Ohm’s law. Here, \( R = (r^2 + z^2)^{1/2} \) is the distance from the origin, \( r_s \equiv 2GM/c^2 \) is the Schwarzschild radius (with \( G, M, \) and \( c \) being the gravitational constant, the mass of a black hole, and the speed of light, respectively), and \( \mathbf{J} = (c/4\pi) \mathbf{V} \times \mathbf{E} \) is the electric current. As for \( \eta \), we adopted the anomalous resistivity model that is used in many solar-flare simulations (for the detail, see Yokoyama & Shibata 1994). The calculation was started with a rotating torus in hydrostatic balance located around \( r = r_0 = 40r_s \). The initial magnetic fields were confined to within a torus and purely poloidal (see KMS04 for the detail). In the simulation, they used \( 300 \times 32 \times 400 \) nonuniform mesh points. The grid spacing was uniform (\( \Delta r = \Delta z = 0.16r_s \)) within the inner calculation box of \( 0 \leq z \leq 10r_s \), and increased by 1.5% from one mesh to the adjacent outer mesh outside this box up to \( r \leq 20r_s \) and \( z \leq 20r_s \), and increased by 3% beyond that. The entire computational box size was \( 0 \leq r \leq 200r_s, 0 \leq \phi \leq 2\pi, \) and \( -50r_s \leq z \leq 50r_s \), and they simulated a full 360° domain (see Kato 2004 for more detail).

The simulated MHD flow is slightly oscillating because of turbulence driven by MRI, and a geometrically thick density distribution is produced. In this quasi-steady accretion flow, the density profile is \( \rho \propto r \) in the inner part \( (r < 20r_s) \), while being \( \rho \propto r^{-1} \) in the outer part \( (r > 20r_s) \) (see figure 4 in KMS04). We used this quasi-steady density distribution as well as the ion temperature distribution in modeling the corona in which hot thermal electrons up-scatter soft photons emerging from a cold disk virtually located in the equatorial plane.

2.2. Physical Quantities of the Disk and Corona

As mentioned in section 1 briefly, we adopted the following assumptions in constructing the disk–corona model:

1. In the equatorial plane, we assume a standard accretion disk (Shakura & Sunyaev 1973) with an infinitesimal height for a given mass-accretion rate, \( \dot{M}_{\text{disk}} \).

2. The RIAF-like accretion flow, whose structure is provided by a three-dimensional MHD simulation (KMS04), surrounds the standard thin disk as a hot corona. That is, we assume that coronal flow is created at large radius by evaporation of the disk material, which moves freely inward.\(^1\) Mass evaporation and condensation occurring between them and heat conduction from the corona to the disk were neglected in the calculation. Radiative cooling of coronal plasma was also neglected.

3. The photons radiated from the thin disk are partly up-scattered by hot electrons in the corona, and the remaining

\(^1\) Note that evaporation dominates over condensation at large radii, whereas the opposite is the case at small radii (e.g., Liu et al. 2007).
portion penetrate through the corona, creating a power-law hard emission in the spectrum.

4. As the seed photon field from the underlying disk, we include the reprocessed radiation, which comes from the coronal irradiation on the disk, as well as the intrinsic disk radiation.

We took three-dimensional data of the density and proton temperature distributions in the accretion flows calculated in Kato (2004). This has the same initial condition as Model B in KMS04. The data of the physical properties were given at each point in the simulation box associated with the Cartesian coordinates \((x, y, z)\), in which the black hole was located at the origin of the coordinate axes, the \(z\)-axis was set to be the rotation axis of the accretion flow, and the \(x\)-\(y\) plane corresponded to the equatorial plane. We employed Cartesian grids with numbers \((N_x, N_y, N_z) = (101, 101, 101)\) of cells. The size of the calculating box was \(2X \times 2Y \times 2Z\), where we set \((X, Y, Z) = (99.9r_s, 99.9r_s, 99.9r_s)\).

In MHD simulations with no radiative loss, the density was given as a non-dimensional number, \(\rho_0\), with the normalization factor, \(\rho_0\), which was treated a free parameter in our calculation. Basically, the coronal density is determined by evaporation of the disk gas, but here we determined \(\rho_0\) so that the power-law indices of the evaluated spectra would agree with the observations. The proton temperature does not depend on the density parameter, and is given by the MHD simulation as \((\rho m_p c^2/k_B)\), where \(\rho\) is the mean molecular weight \((=0.5)\), \(m_p\) is the proton mass, \(k_B\) is the Boltzmann constant, and \(c\) is the normalized sound velocity obtained by the simulation.

Here, we should note that ions (protons) and electrons in the plasma simulated in KMS04 have the same temperature, though electrons would be radiatively cooled, and have a lower temperature in the realistic situation. Assuming that ion temperature coincides with the simulated one, and that the electrons have a Maxwellian distribution, we evaluated the electron temperature, \(T_e\), through the energy balance of the electrons between Coulomb collisions with ions and radiative cooling,

\[
\int_{-\infty}^{\infty} \int_{r_{\text{in}}}^{r_{\text{out}}} \lambda_{\text{ele}} 2\pi r dr dz = \int L_{\text{cor}}(v; r_{\text{in}} \leq r \leq r_{\text{out}}) dv.
\]

(5)

Here \(\lambda_{\text{ele}}\) is the energy transfer rate from ions to electrons (Stepney & Guilbert 1983) and \(L_{\text{cor}}(v)\) is the coronal luminosity at frequency \(v\). In the present study, we divided the corona into three regions \([0 < r \leq 10r_s, 10r_s < r \leq 30r_s, \text{ and } 30r_s < r; (r_{\text{in}}, r_{\text{out}}) = (0, 10r_s), (10r_s, 30r_s) \text{ and } (30r_s, \infty)]\) for simplicity, and supposed that the electrons in the accretion flux in each region have a temperature that is independent of the radius \(r\) and altitude \(z\). The coronal luminosity, \(L_{\text{cor}}\), was obtained by Monte-Carlo simulations (see next subsection) for a guess value of \(T_e\). Since this \(T_e\) does not always satisfy equation (5), we should do some iterations to calculate the appropriate electron temperature and the emergent spectrum.

2.3. Radiative Transfer Simulations

Our model consists of a cold disk that produces blackbody radiation at each radius with the temperature being determined by the standard model (Shakura & Sunyaev 1973):

\[
T_{\text{disk}}^0 = \left[ \frac{3GM \dot{M}_{\text{disk}}}{8\pi r^3 \sigma} \left(1 - \frac{r}{r_{\infty}} \right) \right]^{1/4},
\]

(6)

(as long as the reprocessed radiation is unimportant), where \(\sigma\) is the Stephan–Boltzmann constant. Note that when irradiation flux, \(F_{\text{irr}}\), by the corona is available, the disk emits blackbody radiation with an enhanced temperature,

\[
T_{\text{disk}} = \left[ (T_{\text{disk}}^0)^4 + F_{\text{irr}}/\sigma \right]^{1/4}.
\]

(7)

In the following calculation we set \(r_{\infty} = 3r_s\), \(\dot{M}_{\text{disk}} = 10^{-3}M_{\text{edd}}\) (with \(M_{\text{edd}} = 10L_{\text{edd}}/c^2\)), and the mass of a central black hole to be \(M = 10^8M_\odot\), which is believed to be the typical value for AGNs. Thus, the normalized time corresponds to \(t \gg r_s/c = 10^3\) s.

As for the radiation process, we took into account both of the intrinsic disk radiation [equation (6)] and the thermal reprocessing from the irradiated disk as a seed photon field [see equation (7)], and the Compton/inverse Compton scattering in the corona. We neglected synchrotron emission/absorption and free-free emission/absorption. Some studies about the disk–corona model have shown that the power-law component of the emergent spectra can be explained as Comptonized emission (Haardt & Maraschi 1991, 1993), so for our present purpose this approximation is justified, at least in the X-ray energy bands. According to the simulation by KMS04, the magnetic field in coronal flow with \(B_0 = 1.6 \times 10^{-14}\) g cm\(^{-2}\) is \(B \sim 10^7\) G, for which the Compton scattering is the most efficient cooling mechanism. This fact justifies the method of deriving the coronal temperature described in the previous section.

The method of the Monte-Carlo simulation is based on Pozdnyakov et al. (1977). In order to efficiently calculate the emergent spectra, we introduced a photon weight, \(w\). When emerged from the disk we set that each photon would have a weight of \(w_0 = 1\), and then calculated the escape probability, \(P_0\). The escape probability of a photon after \(i\)-th scattering (for \(i \geq 1\)), \(P_i\), is evaluated as

\[
P_i = \exp \left[ - \int \frac{\sigma_{\text{KN}}(x_i, y_i, z_i)}{m_p} dl \right],
\]

(8)

where \((x_i, y_i, z_i) = (x_0, y_0, 0)\) corresponds to the point on the equatorial plane (i.e., the disk plane), in which the thermal soft photons are generated, \((x_i, y_i, z_i)\) is the point where a photon is subject to the \(i\)-th scattering, \(m_p\) is the proton mass, \(\sigma_{\text{KN}}\) is the Klein–Nishina cross section (Rybicki & Lightman 1979), and the integral of \(dl\) should be done along the photon direction there from the point \((x_i, y_i, z_i)\) to the boundary of the calculating box. The quantity \(w_0 P_0\) represents the transmitted portion of photons, which is recorded to calculate the penetrated spectrum if the path of the photon does not cross with the equatorial plane, and will no longer continue to be counted and will be regarded as being absorbed by the disk. As for the remaining portion of a photon, its weight becomes \(w_1 = w_0(1 - P_0)\). The transmitted portion of photons after the \(i\)-th scattering, \(w_i P_i\), is recorded to calculate the emergent spectrum, and the remaining portion, \(w_i(1 - P_i)\), undergoes \((i + 1)\)-th scattering. This calculation is continued until the
weight, $w_i$, becomes sufficiently small ($w_i \ll 1$), or the path of the remaining photon crosses the equatorial plane, and is considered to be absorbed by the disk. The whole process was simulated by the Monte-Carlo method. Finally, we supposed that the inner boundary of the underlying standard disk is $r = 3r_S$, and general-relativistic effects, like the light bending and energy shift, were neglected in our study.

3. Results

3.1. Spectral Features

First, we show the emergent spectra from the accretion disk with MHD coronal flow with various density parameters for the corona in figure 1. In this calculation, the mass-accretion rate of the underlying disk is set to be $M_{\text{disk}} = 10^{-3}M_{\text{Edd}}$. This value is relatively lower than that assumed in the standard picture ($M_{\text{disk}} \sim 0.1-0.01M_{\text{Edd}}$). However, according to an idea of Haardt and Maraschi (1991), the liberated gravitational energy accompanying the mass accretion, $L_G$, should be mostly dissipated in the corona ($fL_G$, where $f \simeq 1$), while only a small fraction of the energy $[(1 - f)L_G]$ is dissipated in the underlying disk. Then, the temperature of the intrinsic disk radiation would be reduced by a factor of $(1 - f)^{1/4}$. In our calculation the mass-accretion rate of the disk only appears when giving the intrinsic disk temperature equation (6), and so the small mass-accretion rate corresponds to the reduced disk temperature. Due to this fact, the relatively small mass-accretion rate adopted in our calculation is justified.

The adopted density parameters, $\rho_0$, for the corona are $\rho_0 = 5.1 \times 10^{-15}$, $1.6 \times 10^{-14}$, and $5.1 \times 10^{-14}$ g cm$^{-3}$. These values correspond to the number density of $\sim 10^9$ cm$^{-3}$. Such a density is consistent with some corona models (see Liu et al. 2002b, 2003). Due to inverse Compton scattering of the soft thermal photons from the underlying disk in the corona, the power-law component with a spectral index of $\alpha \sim 1-2$ ($F_\nu \propto \nu^{-\alpha}$) appears in the higher energy band. In the calculated spectra, we cannot clearly find a bump-like structure in the UV/soft X-ray band, which is usually seen in typical spectra of AGNs. As $\rho_0$ increases, the total luminosity increases, while the power-law index decreases. The corona with a higher density (especially with a higher scattering optical depth) would irradiate the underlying disk with a higher energy flux, because a larger number of photons that originated from the disk would gain energy and be backscattered. Because the disk is heated by the corona, the energy flux of the seed photon field would increase. This is why the total luminosity increases with the coronal density.

In table 1 we summarize the scattering optical depths, $\tau$ (evaluated by integrating the scattering opacity over the z-direction from the equatorial plane), the Compton $y$ parameters of the corona (averaged in each region), and the spectral indices of the power-law component, $\alpha$, estimated from the calculation results for 3 density parameters. A MHD simulation of coronal flow was started from the initial condition of a magnetized torus located around $r \sim 40r_S$. The radius of the maximum gas density (and of maximum $\tau$) decreased inward with time, and remained at around $r \sim 20r_S$. According to the theory of unsaturated inverse Compton scattering, the power-law index of the Comptonized emission component depends on the plasma density and the temperature through the following equation:

$$\alpha = -\frac{3}{2} + \frac{9}{4} + \frac{4}{y}. \quad (9)$$

Table 1. Coronal properties.

<table>
<thead>
<tr>
<th>$\rho_0$ (g cm$^{-3}$)</th>
<th>Region</th>
<th>$T_{\text{cor}}$ (K)</th>
<th>$\tau$</th>
<th>$y$</th>
<th>$\alpha$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$5.1 \times 10^{-14}$</td>
<td>$0 &lt; r &lt; 10r_S$</td>
<td>$\sim 5.2 \times 10^9$</td>
<td>$\sim 0.2$</td>
<td>$\sim 0.5$</td>
<td>$\sim 1.30$</td>
</tr>
<tr>
<td></td>
<td>$10r_S &lt; r &lt; 30r_S$</td>
<td>$\sim 3.4 \times 10^9$</td>
<td>$\sim 0.9$</td>
<td>$\sim 3$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$30r_S &lt; r$</td>
<td>$\sim 7.6 \times 10^8$</td>
<td>$\sim 0.7$</td>
<td>$\sim 0.5$</td>
<td></td>
</tr>
<tr>
<td>$1.6 \times 10^{-14}$</td>
<td>$0 &lt; r &lt; 10r_S$</td>
<td>$\sim 5.2 \times 10^9$</td>
<td>$\sim 0.05$</td>
<td>$\sim 0.2$</td>
<td>$\sim 1.60$</td>
</tr>
<tr>
<td></td>
<td>$10r_S &lt; r &lt; 30r_S$</td>
<td>$\sim 3.4 \times 10^9$</td>
<td>$\sim 0.25$</td>
<td>$\sim 0.55$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$30r_S &lt; r$</td>
<td>$\sim 7.5 \times 10^8$</td>
<td>$\sim 0.2$</td>
<td>$\sim 0.1$</td>
<td></td>
</tr>
<tr>
<td>$5.1 \times 10^{-15}$</td>
<td>$0 &lt; r &lt; 10r_S$</td>
<td>$\sim 5.2 \times 10^9$</td>
<td>$\sim 0.02$</td>
<td>$\sim 0.05$</td>
<td>$\sim 2.00$</td>
</tr>
<tr>
<td></td>
<td>$10r_S &lt; r &lt; 30r_S$</td>
<td>$\sim 3.4 \times 10^9$</td>
<td>$\sim 0.09$</td>
<td>$\sim 0.15$</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$30r_S &lt; r$</td>
<td>$\sim 7.6 \times 10^8$</td>
<td>$\sim 0.07$</td>
<td>$\sim 0.03$</td>
<td></td>
</tr>
</tbody>
</table>
Here, \( y = [4k_B T_{\text{cor}}/(m_c e^2)] \) is the Compton \( y \)-parameter (Rybicki \& Lightman 1973), where \( T_{\text{cor}} \) and \( y \) are the temperature and Thompson optical depth of the corona, respectively. The spectral indices derived from this equation using the \( y \)-parameters in table 1 do not always agree with those estimated from the spectra. This is because \( y \) and \( T_{\text{cor}} \) have spatial distributions, and we cannot evaluate the \( y \)-parameter of the corona uniquely. Even with such a situation, however, we can see the tendency that the spectra become flatter with a higher coronal density (and then higher \( y \)-parameter), which is consistent with the theory mentioned above.

In our calculation, the coronal density parameter \( \rho_0 \) was determined so as to reproduce the observation. By tuning \( \rho_0 \), we could reproduce the luminosity of the power-law component, which is as intense as that of the thermal component originated from an optically thick disk. Such a feature is typical in Seyfert galaxies. However, the coronal temperature cannot be chosen freely, but should be determined by imposing the energy balance of the electrons between Coulomb collisions and the cooling via inverse Compton scattering, as we have done. Nevertheless, the resulting coronal temperature is substantially reduced from the plasma temperature derived by the simulation (\( \sim 10^{13} \) K) to \( \sim 10^7 \) K, which makes the high-energy cutoff of computed spectra consistent with the observations.

### 3.2. Time Variation

The spectral variation caused by the time variation of the MHD coronal flow structure is shown in figure 2. In the highest energy range (\( \gtrsim 10^{13} \) Hz) the spectra show fluctuations because of poor photon statistics. As for the soft X-ray band (with \( \log \nu = 17–18 \)), where the spectra show a smooth power-law shape, the spectral index slightly changes with time, and then the X-ray flux fluctuates a little (see also figure 3). According to the MHD simulation, on which our radiative transfer calculations are based, the three-dimensional structure of the coronal accretion flow is fluctuating everywhere in each timestep. On the other hand, the spectral index depends on the distribution of the \( y \)-parameter of the corona, as we note in the last subsection. We can thus conclude that the fluctuations of the spectral indices of the computed spectra in figure 2 reflect the fluctuation of \( y \), which comes from the density fluctuations (which is supposed to be due to MRI) in the coronal flow.

Figure 3 shows the X-ray lightcurve derived from our simulations. From this plot we can see that the X-ray luminosity from our disk–corona system can change by factors of a few tens of percent on timescales of the orbital period near the last stable orbit \( r = 3r_S \), i.e., about \( 10^{3}(M/10^8M_\odot) \) s. In the whole calculation, we did not vary the properties of the soft photon source (i.e., the underlying cold disk). This variation that we obtained is due to the density fluctuation (and accompanying temperature fluctuation) of the coronal flow.

### 4. Discussion and Conclusion

We calculated the emergent spectra and their time variabilities predicted based on the disk–corona model, in which a cold standard disk at the equatorial plane is sandwiched by a hot coronal flow. As for the structure and dynamics of the corona, we used the three-dimensional MHD simulation data by Kato (2004). In this section we discuss our results, especially in the context of AGNs.

We have shown that our black-hole disk–corona system can reproduce the power-law X-ray emission with the photon index \( \alpha \sim 1–2 \) by properly adjusting the density parameter. The power-law indices and the cutoff energy scales of the spectra are roughly in agreement with the observed spectra of Seyfert galaxies. Moreover, we find significant variability of the power-law X-ray emission. The power-law X-ray emission flux predicted from our model changes by a few tens of percent on timescales of the orbital period near the last stable orbit, which is about \( 10^3(M/10^8M_\odot) \) s, while the power-law index nor the cutoff energy do not change considerably. This variability comes purely from the fluctuation of the coronal flow around \( r \sim 20r_S \), where the scattering optical depth of the coronal flow attains its largest value in its structure. This fluctuation is driven by turbulence as the result of MRI, and its amplitude is large enough to explain the observed X-ray variability in Seyfert galaxies (e.g., Miniutti et al. 2007). In the
UV/soft X-ray band, however, we cannot see the bump-like structure that is often found in the spectra of AGNs. The same problem occurred in the other disk–corona models (Shimura et al. 1995; Liu et al. 2003). This bump is supposed to be thermal radiation coming directly from the optically thick disk, while in those disk–corona models the disk is wholly covered with corona. In order to reproduce the UV/soft X-ray bump with a model, the coronal structure may need to be patchy or locally concentrated. In the present study we did not consider the viewing angle dependence of the spectra. If we observe this accretion-flow system face-on, then the scattering optical depth along the line of sight would be smaller than in the case observed with a non-zero viewing angle, and the optically thick component should be observed more clearly.

Iron Kα fluorescent line emissions would occur when the disk is irradiated by X-rays from the corona. The profile of this line emission would be broadened due to relativistic effects (see Reynolds & Nowak 2003 for a review). The detailed structure of this line profile is determined by the line emissivity distribution on the disk, which depends on the spectra of local X-ray irradiation from the corona.

Recently, Suzaku observed the broadened iron line profile from MCG –6–30–15 in detail; they implied that the emissivity profile should be as steep as $F_{\text{line}}(r) \propto r^{-4.4}$ (Miniutti et al. 2007). This is not easy to understand, since standard disks are known to produce a continuum flux of $F_{\text{cont}}(r) \propto r^{-3}$. Some authors have proposed that this is a manifestation of Kerr black-hole effects (e.g., Wilms et al. 2001). Such a steep emissivity profile can be reproduced by thermal Comptonization in the corona, if the irradiating corona has temperature and density profiles that arise inward, as was pointed out by Kawanaka et al. (2005). This is because the fraction of high-energy photons increases inward, so that the number of irradiating photons that are capable of producing iron fluorescent line emissions should increase inward more rapidly than $r^{-3}$. As for the MHD corona model, which we adopted in this study, however, the temperature and density have flatter profiles (see KMS04), so such a steep line emissivity profile is not expected as long as we consider only iron fluorescent-line photons are produced in the transition zone between the hot corona and the cold disk, where density and temperature abruptly change, the heat conduction and the mass evaporation would be important as the mass and energy exchanging processes (Meyer & Meyer-Hofmeister 1994; Liu et al. 2002a). The evaporation of photospheric material was actually shown to be essential in the context of solar flares (see e.g., Yokoyama & Shibata 1994). By including these effects in the simulation, we will be able to obtain a more realistic model of the coronal structure and dynamics. To what extent the disk–corona structure can extend to the inner region is important when we consider relativistic acceleration) around magnetic reconnection flares and hard X-ray emission processes in the MHD coronal flow simulation, and then we can know the local X-ray irradiation onto the disk, which may lead us to understanding the observed broad iron line profiles and their time variabilities from the microphysics of the X-ray emission processes.

We can check the consistency of the assumption that the MHD coronal flow that we adopted in the calculation is cooled dominantly by advection, by comparing the heating rate [$\sim GMm_p\Omega/(2r)$; $\Omega$ is the angular velocity of the flow] and the cooling rate ($\sim$Comptonization luminosity per electron). According to the result of our calculation, the innermost region ($0 < r \lesssim 20r_H$) is not cooled efficiently by inverse Compton scattering. However, in the outer region ($r \gtrsim 20r_H$) the local heating rate and the cooling rate are comparable, which means that the assumption of an advection-dominated flow is only marginally justified. Generally, when Compton cooling becomes efficient the coronal temperature will become lower, and then Compton upscattering will be inefficient. This means that the power-law X-ray emission would not be generated so much from this region. As long as we consider the spectrum and the energy flux in X-ray and higher energy band, to which the photons from the innermost region mainly contribute, this inconsistency would hardly affect the results.

Finally, we should mention the interaction between the corona and the underlying disk. In the transition zone between the hot corona and the cold disk, where density and temperature abruptly change, the heat conduction and the mass evaporation would be important as the mass and energy exchanging processes (Meyer & Meyer-Hofmeister 1994; Liu et al. 2002a). The evaporation of photospheric material was actually shown to be essential in the context of solar flares (see e.g., Yokoyama & Shibata 1994). By including these effects in the simulation, we will be able to obtain a more realistic model of the coronal structure and dynamics. To what extent the disk–corona structure can extend to the inner region is important when we consider any relativistic skewing of the iron-line emission profile, because most of the observations of iron line profiles imply that the iron fluorescence-line photons are emitted from around/inside the last stable orbit. Detailed analyses of such disk–corona interactions and their observational implications are left as future work.

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