THE STRUCTURE AND SPECTRAL FEATURES OF A THIN DISK AND EVAPORATION-FED CORONA IN HIGH-LUMINOSITY ACTIVE GALACTIC NUCLEI

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ABSTRACT

We investigate the accretion process in high-luminosity active galactic nuclei (HLAGNs) in the scenario of the disk evaporation model. Based on this model, the thin disk can extend down to the innermost stable circular orbit (ISCO) at accretion rates higher than 0.02 $M_{\text{Edd}}$ while the corona is weak since part of the coronal gas is cooled by strong inverse Compton scattering of the disk photons. This implies that the corona cannot produce as strong X-ray radiation as observed in HLAGNs with large Eddington ratio. In addition to the viscous heating, other heating to the corona is necessary to interpret HLAGN. In this paper, we assume that a part of accretion energy released in the disk is transported into the corona, heating up the electrons, and is thereby radiated away. For the first time, we compute the corona structure with additional heating, fully taking into account the mass supply to the corona, and find that the corona could indeed survive at higher accretion rates and that its radiation power increases. The spectra composed of bremsstrahlung and Compton radiation are also calculated. Our calculations show that the Compton-dominated spectrum becomes harder with the increase of energy fraction (f) liberating in the corona, and the photon index for hard X-ray (2–10 keV) is 2.2 < $\Gamma$ < 2.7. We discuss possible heating mechanisms for the corona. Combining the energy fraction transported to the corona with the accretion rate by magnetic heating, we find that the hard X-ray spectrum becomes steeper at a larger accretion rate and the bolometric correction factor ($L_{\text{bol}}/L_{2-10\text{keV}}$) increases with increasing accretion rate for $f < 8/35$, which is roughly consistent with the observational results.

Key words: accretion, accretion disks – galaxies: active – X-rays: galaxies

1. INTRODUCTION

Accretion of gas onto the central supermassive black hole is one of the fundamental astrophysical processes responsible for high-efficient energy release in active galactic nuclei (AGNs). According to the difference in the luminosity, AGNs are classified into two types: low-luminosity AGNs (LLAGNs) and high-luminosity AGNs (HLAGNs). The LLAGNs show hard power-law X-rays and/or radio jets, while the HLAGNs are characterized by the big blue bump around the UV waveband, soft X-ray excess, Fe Kα lines at about 6.4 keV, the power-law X-rays with photon index $\Gamma \sim 1.9$, and cutoff at about a few 100 keV. It is speculated that the different characteristics of LLAGNs and HLAGNs are caused by different accretion mechanism at different accretion rates, in a similar way to black hole X-ray binaries. The accretion in LLAGNs is via an advection-dominated accretion flow (ADAF; Ichimaru 1977; Narayan & Yi 1994, 1995a, 1995b) in which most of the accretion energy is stored in the gas as entropy and only a small fraction is radiated in hard X-rays (see Kato et al. 2008 for a review). While in HLAGNs, the commonly accepted accretion model is a hybrid accretion disk; i.e., a cool disk is embedded in a hot corona. The optical to UV emission is from the optically thick and geometrically thin accretion disk (e.g., Shakura & Sunyaev 1973; Lynden-Bell & Pringle 1974), and possibly also the soft X-ray emission in some AGNs (Yuan et al. 2010); while the hard X-ray is from Compton scattering of the soft photons from the disk by the hot electrons in the optically thin corona with $T_e \sim 10^7$ K.

Observations show that the X-ray luminosity contributes a large fraction to the bolometric luminosity in HLAGNs. This can be seen from individual spectra (e.g., Elvis et al. 1994; Ho 2008), as well as the broad band spectral energy distribution of large samples (e.g., Yuan et al. 1998). The hard X-ray bolometric correction, which is defined as the ratio of bolometric luminosity to 2–10 keV luminosity, is typically around 40–70 with an Eddington ratio $L_{\text{bol}}/L_{\text{Edd}}$ ($L_{\text{Edd}} = 1.26 \times 10^{38}(M_{\text{BH}}/10^8 M_\odot)$ erg s$^{-1}$) above 0.1, while below 0.1 it is typically 15–25 (Vasudevan & Fabian 2009), which confirms the importance of X-ray radiation in LLAGNs. Fits to the X-ray spectra of HLAGNs by a Comptonization model yield an electron temperature of $\sim 10^6$ K in the optically thin hot corona (e.g., Liu et al. 2003). A question arises: how can the gas in the corona remain at so high a temperature when it continuously radiates strong X-rays in HLAGNs? Where is the heating energy from?

Theoretically, gas accreted from the interstellar medium or secondary is cold. Some of the gas evaporates to the corona and then accretes inward into the corona on its way to the central black hole. In HLAGNs, gas is dominantly accreted via a thin disk, only a small fraction of gas is accreted in the hot corona. Therefore, the energy released in the corona through accretion is much smaller than that in the disk. This implies that the radiation from the hot corona is much weaker than that from the disk in a steady disk–corona flow. With strong radiation in X-rays as observed in HLAGNs, the corona is overcooled in a thermal timescale and collapses into the disk unless there is another energy supply to the corona besides the viscous heating.

Recent investigations of the spectral energy distribution (SED) in HLAGNs found that the bolometric correction increases with an increasing Eddington ratio, whereas it is independent of the black hole mass (Vasudevan & Fabian 2007, 2009; Lu & Yu 1999; Wang et al. 2004; Liu & Liu 2009). This
correlation provides clues and constraints to possible heating mechanisms to the disk corona in HLAGNs.

In this work, we study the corona flow at high accretion rates in the frame of the disk corona evaporation/condensation model. By including additional heating to the corona, we investigate the detailed interaction, i.e., the mass and energy exchange between disk and corona; we self-consistently determine the corona structure and disk radiation features. The spectrum from the disk and corona is calculated by Monte Carlo simulations. Comparing the spectrum with observational features, we discuss the possible heating mechanism in HLAGNs. Our aim is to interpret the SED and its correlation with Eddington ratio.

HLAGNs include radio-loud and radio-quiet Seyfert 1s and QSOs. Since the radio-loud AGNs are assumed to be powered by the jet and the beaming effect will dilute the X-ray radiation from the central accretion engine, here we focus on an accretion by the jet and the beaming effect will dilute the X-ray radiation.

Comparing the spectrum with observational features, we discuss the detailed interaction, i.e., the mass and energy exchange in HLAGNs. Our aim is to compare both radiation and mass coupling between disk and corona and study the detailed vertical structure of the coronal flow above a thin disk, supposing some fraction \( f \) of the gravitational energy is added into the corona.

2.2. Physical Processes and Differential Equations

The corona can be described by the following equations. The equation of state is

\[
P = \frac{9\eta_\rho}{2\mu}(T_e + T_i),
\]

where \( \mu = 0.62 \) is the molecular weight assuming a standard chemical composition \((X = 0.75, Y = 0.25)\) for the corona. Here, we assume \( n_i = n_e \), which is strictly true only for a pure hydrogen plasma.

The equation of continuity is

\[
\frac{d}{dz}(\rho v_z) = \frac{\eta_M}{R} \frac{\rho v_R}{R^2 + z^2} \rho v_z,
\]

where \( v_R = -a(V_s^2/\Omega R) \) is the radial component of velocity, \( V_s = (P/\rho)^{1/2} \) is the isothermal sound speed, and \( v_e = \sqrt{(GM/R)(1 + (z^2/R^2))^{-3/2}} \) is the angular component of velocity. Here, the partial derivative of mass flux with respect to radius is approximated as, \((1/R)(\partial \rho v_R/\partial R) \approx -\eta_M(2/R)\rho v_R\), where \( \eta_M \) depends on the net mass gain/loss rate through the radial boundaries (Meyer-Hofmeister & Meyer 2003; Liu et al. 2004). In this work, \( \eta_M = 1 \) is taken for the case that no mass enters through the outer boundary and all the mass flowing in the corona is contributed by the evaporation (Meyer et al. 2000a).

The equation of the \( z \)-component of momentum is

\[
\rho v_z \frac{dv_z}{dz} = \frac{-dP}{dz} - \rho \frac{GMz}{(R^2 + z^2)^{3/2}}.
\]

In the hot corona, because of a heavier mass of ions than that of electrons, the viscous heating raises the ion temperature first and ions are cooled by coulomb collision with electrons and both radial and vertical advection. Here, we do not include direct heat to electrons. The ion thermal conduction is not taken into account because of the long mean free path compared to the vertical scale height of the corona. So the energy equation for ions is

\[
\frac{d}{dz} \left\{ \frac{\rho v_z}{2} \left[ \frac{v^2}{2} + \frac{\gamma}{\gamma - 1} \frac{P_i}{\rho i} - \frac{GM}{R^2 + z^2} \right] \right\} = \frac{3}{2} \alpha P \Omega - q_{he}
\]

\[
+ \frac{2}{R} \rho \frac{\rho v_R}{R^2 + z^2} \left[ \frac{v^2}{2} + \frac{\gamma}{\gamma - 1} \frac{P_i}{\rho i} - \frac{GM}{R^2 + z^2} \right] - 2 \frac{\rho v_z}{R^2 + z^2} \left[ \frac{v^2}{2} + \frac{\gamma}{\gamma - 1} \frac{P_i}{\rho i} - \frac{GM}{R^2 + z^2} \right],
\]

where \( n_e \) is the energy modification parameter, \( \eta_e = \eta_M + 0.5 \). The difference between \( \eta_i \) and \( \eta_e \) comes from the derivative of energy flux with respect to radius, \((1/R)(\partial \rho v_R \eta_e/\partial R) \approx -\eta_M(2/R)\rho v_R \epsilon + \rho v_R (\partial \epsilon/\partial R) = -(\eta_M + 0.5)(2/R)\rho v_R \epsilon,
\]
where $\varepsilon \equiv v^2/2 + (\gamma/\gamma - 1)\rho_0 - GM/(R^2 + z^2)^{1/2}$ and its derivative $(\partial \varepsilon/\partial R) \approx -\varepsilon/R$ since the potential, kinetic, and thermal specific energies all scale approximately as $1/R$ (Meyer-Hofmeister & Meyer 2003; Liu et al. 2004). In Equation (4), $q_\text{ue}$ is the exchange rate of energy between electrons and ions through coulomb collision and is described as

$$q_\text{ue} = \left(\frac{2}{\pi}\right)\frac{3}{2} \frac{m_e}{m_i} \ln \Lambda \sigma_T \epsilon n_2 \left(\frac{\kappa T_e}{\kappa T_i} - 1\right) + \frac{T_e}{T_i^2},$$

where $m_i$ and $m_e$ are the proton and electron masses, $\kappa$ is the Boltzmann constant, $c$ is the light speed, $\sigma_T$ is the Thomson scattering cross section, and $\ln \Lambda = 20$ is the Coulomb logarithm.

Different from the previous works, besides the viscous heating, a fraction of local gravitational energy is assumed to directly heat the electrons in corona, i.e.,

$$Q_\text{add} = \frac{3GM \dot{M}}{8\pi R^3} \left[1 - (3R_e/R)^{1/2}\right] \times f,$$

where $\dot{M}$ is the total accretion rate in the disk corona. The physical meaning of this additional heating term will be discussed in Section 4. To simplify the calculations, this additional heating flux is supposed to distribute along the vertical direction in a form similar to that of Compton cooling, i.e.,

$$q_\text{add} = \frac{4\kappa T_e}{m_e c^2} n_e \sigma_T c a T_e^4 \frac{\sigma T_e}{2},$$

where

$$T_e = \frac{\kappa T_e}{m_e c^2} \left(1 + \frac{m_e T_i}{m_i T_e}\right),$$

in which $T_s = \kappa T_e / m_e c^2 \left(1 + m_e T_i / m_i T_e\right)$.

Therefore, the energy equation for both the ions and electrons is

$$\frac{d}{dz} \left\{\rho v_z \left[\frac{v^2}{2} + \frac{\gamma P}{\gamma - 1} - \frac{GM}{(R^2 + z^2)^{1/2}}\right] + F_c\right\}$$

$$= \frac{3}{2} \rho \dot{M} \Omega + q_\text{add} - n_e n_i L(T_e) - q_{\text{Cmp}}$$

$$+ \frac{\eta_T}{R^2 + z^2} \rho v_R \left[\frac{v^2}{2} + \frac{\gamma P}{\gamma - 1} - \frac{GM}{(R^2 + z^2)^{1/2}}\right]$$

$$- \frac{2 z}{R^2 + z^2} \rho v_z \left[\frac{v^2}{2} + \frac{\gamma P}{\gamma - 1} - \frac{GM}{(R^2 + z^2)^{1/2}}\right] + F_c.$$

(11)

In this equation, $n_e n_i L(T_e)$ is the bremsstrahlung cooling rate and $F_c$ is the thermal conduction (Spitzer 1962),

$$F_c = -\kappa_0 T_e \frac{d T_e}{dz},$$

with $\kappa_0 = 10^{-6}$ erg s$^{-1}$ cm$^{-1}$ K$^{-2}$ for fully ionized plasma.

For the Compton cooling rate, $q_{\text{Cmp}}$, we mainly consider inverse Compton scattering of the soft photons from the disk, while the ones contributed from the bremsstrahlung cooling are not included:

$$q_{\text{Cmp}} = \frac{4\kappa T_e}{m_e c^2} n_e \sigma_T c a T_e^4 \frac{\sigma T_e}{2},$$

where $a$ is the radiation constant and $T_{\text{eff}}$ is the effective temperature of the underlying thin disk, fulfilling

$$\sigma T_{\text{eff}}^4 = \frac{3GM M_d - f M}{8\pi R^3} \left[1 - (3R_e/R)^{1/2}\right],$$

where $M_d$ is the mass accretion rate in the thin disk, which depends on the distance because of evaporation,

$$M_d(R) = M - M_{\text{evap}}(R),$$

with $M_{\text{evap}}(R)$ as the integrated evaporation rate from the outer edge of the disk to the distance $R$, which can be approximated by the one-zone evaporation rate, $M_{\text{evap}} = 2\pi R^2 (\rho v_z) h$ (see Liu et al. 2002a). The negative term associated with $-fM$ in Equation (14) corresponds to the transportation of accretion energy from the disk to the corona, which results in a lower temperature in the disk. Here, we have not considered back reaction (heating of a thin disk by coronal illumination) and its justification will be discussed later.

These five differential equations, Equations (2), (3), (4), (11), and (12), which contain five variables $P(z)$, $T(z)$, $T_e(z)$, $F_c(z)$, and $\rho v_z$, can be solved with five boundary conditions.

2.3. Boundary Conditions and Numerical Method of Computation

At the lower boundary, the temperature of the gas should be the effective temperature of the accretion disk. Liu et al. (1995) showed that the coronal temperature increases from effective temperature to $10^{6.5}$ K in a very thin layer of nearly constant pressure. Its physics is described by the balance between thermal conduction and radiation loss. So a relation can be established between temperature and heat flux which can be scaled according to the pressure. Combining with Shmeleva–Syrovatskii relation as $F_c = -((\kappa_0 L(T) T^{3/2})^{1/2}/\kappa) P$ (Shmeleva & Syrovatskii 1973), the lower boundary conditions can be approximated (Meyer et al. 2000a) as

$$T_i = T_e = 10^{6.5} \text{ K},$$

$$F_c = -2.73 \times 10^6 P \text{ at } z = z_0.$$

There is no pressure and no heat flux at infinity, which requires sound transition at some height $z = z_1$. So we constrain the upper boundary as

$$F_c = 0 \text{ and } v_z = V_z \text{ at } z = z_1.$$

With such boundary conditions, assuming an initial value $\lambda$ and a pair of lower boundary values for $P_0$ and $(\rho v_z)_{h}$, we start the integration along $z$. If the trial values for $P_0$ and $(\rho v_z)_{h}$ fulfill the upper boundary conditions, then the initial value of $P_0$ and $(\rho v_z)_{h}$ can be taken as the solutions of the differential equations. The value of $f$ can then be calculated out from Equation (10). The results for a series of $f = 0.0, 0.1, 0.2, 0.3, 0.4$ are given in the following section. When $f = 0.0$, which means that there is no other energy heating the corona except the viscous heating, this result is similar to the former works and will be compared with the results for $f > 0.0$. 

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2.4. Spectra from the Disk and Corona

With the detailed calculations of the corona structure, the temperature and density in the corona and the effective temperature are determined. We then calculate the spectrum from such a disk and corona for a given structure.

In the corona, the emissions are from bremsstrahlung radiation and inverse Compton scattering of disk photons. For the bremsstrahlung cooling spectrum, we adopt the Eddington approximation and two-stream approximation (Rybicki & Lightman 1979; Mannmoto et al. 1997) to calculate the radiation and inverse Compton scattering of disk photons. For a given soft photon from the cold disk, then we calculate the penetrate spectrum. The remaining weight ω0 of passing through the corona slab.

First, we set the initial weight $\omega_0 = 1$ for a given soft photon from the cold disk. Then, we calculated the escape probability $P_0$ of passing through the corona slab. The value of $\omega_0 P_0$ is the transmitted portion and is recorded to calculate the penetrate spectrum. The remaining weight $\omega_1 = \omega_0 P_0 (1 - P_0)$ is the portion that undergoes more than one scattering. Suppose that $P_n$ is the escape probability after the $n$th scattering, then $\omega_0 P_n$ is the transmitted portion of photons after the $n$th scattering. The remaining portion $\omega_n (1 - P_n)$ undergoes the $(n+1)$th scattering. This calculation is continued until the weight $\omega_n$ becomes sufficiently small. In our calculation, the number of processes calculated is 100,000 and $\omega_n < 10^{-4}$. Finally, we obtain the transmitted spectrum and the soft spectrum.

3. NUMERIC RESULTS

To study AGNs, we fix the black hole mass to be $M_{BH} = 10^8 M_\odot$ throughout our calculations and the Eddington accretion rate is $M_{\text{Edd}} = 1.39 \times 10^{26} \text{gs}^{-1}$. The viscous parameter $\alpha$ in the corona is fixed to be 0.3.

3.1. The Properties of Corona under Viscous Heating

First, we set $\lambda = 0.0$, meaning $f = 0$, which is the case when the corona is heated by only viscous friction. We calculate the corona structures at $R = 1000 R_*$ for $\dot{m} = 0.02$ and $\dot{m} = 0.2$ as shown in the Figure 1; here, $\dot{m} = \dot{M}/M_{\text{Edd}}$, which is the total accretion rate. Comparing the results in two panels, we find that the vertical profiles of the coronal quantities are similar for different accretion rates. From the lower boundary upward, the pressure ($P$) and the vertical mass flow ($\rho v_z$) decrease dramatically in a thin layer. The downward heat flux ($-F_c$) increases in this layer until a maximum is reached and then decreases along $z$ and can be neglected on the upper boundary.

Figure 1. Vertical structure of the corona around a black hole at a distance $R = 1000 R_*$ for accretion rates $\dot{m} = 0.02$ (left panel) and $\dot{m} = 0.2$ (right panel). Here, $z$ is the height above the equatorial plane of the accretion flow; $T_e$ and $T_i$ are the electron temperature and ion temperature, respectively, while $T_{\text{vir}}$ is the virial temperature used for scaling. $F_c$, $\dot{m} v_z$, $P$, and $\rho v_z$ are the heat flux, vertical mass flow rate per unit area, total pressure, and vertical velocity scaled by the sound speed, respectively. The quantities at the lower boundary are marked with subscript 0.

that described by Pozdniakov et al. (1977). In the code, we introduce the weight $\omega$ to efficiently calculate the effects of multiple scattering. First, we set the initial weight $\omega_0 = 1$ for a given soft photon from the cold disk. Then, we calculated the escape probability $P_0$ of passing through the corona slab. The value of $\omega_0 P_0$ is the transmitted portion and is recorded to calculate the penetrate spectrum. The remaining weight $\omega_1 = \omega_0 P_0 (1 - P_0)$ is the portion that undergoes more than one scattering. Suppose that $P_n$ is the escape probability after the $n$th scattering, then $\omega_0 P_n$ is the transmitted portion of photons after the $n$th scattering. The remaining portion $\omega_n (1 - P_n)$ undergoes the $(n+1)$th scattering. This calculation is continued until the weight $\omega_n$ becomes sufficiently small. In our calculation, the number of processes calculated is 100,000 and $\omega_n < 10^{-4}$. Finally, we obtain the transmitted spectrum and the soft spectrum.
The temperatures of electrons ($T_e$) and ions ($T_i$), starting at a coupled value of $10^6$ K at the coronal lower layer, increase with height and decouple at about $z/R \sim 0.1$. Upon this height, $T_i$ still increases up to $T_i \sim 0.3 T_{vir}$ and $T_e$ keeps almost the same value ($\sim 0.05 T_{vir}$) throughout the corona, where $T_{vir} = GM_{bh}\mu/\kappa R = 3.37 \times 10^{12} (R/R_s)^{-1}$ K is the virial temperature. These typical features in disk corona, which are also shown in earlier works (e.g., Meyer et al. 2000b), indicate that the corona above a thin disk undergoes very steep changes in temperature and density and cannot be simply averaged in the vertical direction.

We also calculate the evaporation rate $m_{evap} = M_{evap}/M_{fold}$ along distances, which represents the mass flowing rate in the corona. The results are shown in Figure 2. Similar to the previous works (e.g., Liu at al. 2002a), the coronal accretion rate increases toward the central black hole, reaches a maximum at a few hundred Schwarzschild radii, and then drops very quickly in the inner region because of strong inverse Compton scattering, i.e., a large fraction of the corona gas condensing onto the disk near the black hole. Comparing the three evaporation curves for accretion rate, $m = 0.02$, 0.2, and 0.4, we find that the maximum evaporation rate decreases with increasing accretion rate and the corona gas begins condensing at outer regions for higher accretion rate. The temperature in electrons, as shown in Figure 1, is lower at higher accretion rates. This indicates that the corona becomes very weak at high accretion rates, like in HLAGNs.

We can roughly estimate the radiation from the disk and corona under viscous heating. As shown in Figure 2, the maximum corona accretion rate is $m_{evap,\max} = 4 \times 10^{-3}$ at $\sim 1250 R_s$ for $m = 0.2$. Supposing that the corona keeps accreting at this accretion rate, the radiation strength of the corona relative to the disk is

$$\frac{L_e}{L_d} = \frac{\eta M_c c^2}{\eta M_d c^2} \leq \frac{m_{evap,\max}}{m - m_{evap,\max}} = 0.02. \quad (19)$$

Here, we have supposed that radiative efficiency $\eta$ in the corona is equal to that in the disk. From Equation (19), we can see that the corona is very weak and the radiation is dominated by the disk. In fact, the luminosity ratio between the corona and disk is much smaller than the above value since more and more coronal gas condenses into the disk at smaller distance ($R < 1250 R_s$).

In order to produce strong X-ray in the corona, we consider that some fraction of gravitational energy is transferred to the corona, providing an additional heating mechanism against the strong Compton cooling.

3.2. The Properties of Corona with Additional Heating

3.2.1. Vertical Structure

Assuming a fraction $f (=0.1, 0.2, 0.3, 0.4)$ of accretion energy is added to the corona, we perform numerical calculations for accretion rates of $m = 0.1, 0.2$, and 0.5.

We start our calculations from a distance of 100 $R_s$ to the ISCO region ($R = 3.1 R_s$ for convenience). It has been shown in Section 3.1 that the maximum corona accretion rate can only reach $\sim 4 \times 10^{-3} M_{\text{fold}}$ at $\sim 1000 R_s$ for $m = 0.2$. In the inner region, strong Compton cooling leads to condensation of the corona. Whether the corona can exist at distances inside 100$R_s$ depends on how much accretion energy is added to the corona. From our test calculations, we find that for $f = 0.2$, a corona can still survive in the region from 100 $R_s$ to 3.1 $R_s$. The coronal vertical structure is shown in Figure 3. The profile is similar to the case for $f = 0$ at large distances, though the absolute values are different.

In order to study the corona properties with different energy input, we calculate the coronal accretion rate for $f = 0.1, 0.2, 0.3, 0.4$, where the total mass supply rates are assumed to be $m = 0.2, 0.5$, respectively. The results are plotted in Figures 4 and 5. One can see from Figure 4 that the accretion rate in the corona decreases with decreasing distance as a consequence of condensation, while the electron temperature and scattering optical depth are nearly the same at distances up to 100 $R_s$, as shown in Figure 5. Besides, Figure 4 also shows that the larger the fraction $f$ is, the larger the gas flow becomes in the corona. This is because more energy goes directly to heat the electrons in the corona to a higher temperature, leading to more conductive flux and hence more evaporation or less condensation. Thus, more gas is kept in and accreted through the corona. This is confirmed by the higher electron temperature and larger scattering optical depth for larger $f$, as shown in Figure 5.

Comparing the two panels in Figure 4, we understand that the coronal accretion rates are insensitive to the mass supply rate as long as the fraction of accretion energy released in the corona is the same. Calculations show that the electron temperature and scattering optical depth of the corona do not vary with the accretion rate either. This can be understood as a consequence of energy balance in the corona. For a given $f$, when the mass supply rate ($\dot{m}$) increases, the heating rate added to the corona also increases ($\propto f \dot{m}$). On the other hand, the cooling rate by inverse Compton scattering also increases with the increasing density of soft photon ($\propto (1 - f) \dot{m}$), if $f \dot{m} \gg m_{evap}$. If the additional heat and Compton cooling are the dominant energy terms for the coronal electrons, then there should be no variation in both the evaporation rate and Compton y-parameter. Therefore, the electron temperature does not vary either.

3.2.2. Spectra of the Disk and Corona

Deriving the structure of our model at each radius, we get the radius-dependent electron temperature and electron scattering optical depth, which are the input parameters for bremsstrahlung and Compton spectrum. The bremsstrahlung radiation is proportional to $n_e^2$ but the inverse Compton scattering is proportional to $n_e$. So from the vertical structure, we can see that the bremsstrahlung dominates in the lower layer of the corona and the upper layer is dominated by the Compton
Figure 3. Vertical structure of corona for $\dot{m} = 0.2$, $f = 0.2$ at $R = 100 R_s$ (left) and $R = 20 R_s$ (right).

Figure 4. Radial distribution of corona accretion rate for $\dot{m} = 0.2$ (left) and $\dot{m} = 0.5$ (right) with different $f$.

Figure 5. Radial distribution of scattering optical depth ($\tau$) and electron temperature ($T_e$) for $\dot{m} = 0.2$ with different $f$. 
radiation. Since the temperature and number density in the upper corona still depend on the vertical height, we need to choose an appropriate temperature and density at every given distance for calculating the Compton spectrum with a Monte Carlo simulation. Here, we constrain them by the luminosity, i.e., when the luminosity derived from the Monte Carlo simulation equals the one derived from the coronal structure calculation, the chosen temperature and density are thought to be correct and we take this spectrum as the true spectrum from the corona.

Figure 6 shows the spectra for different $f$, $f = 0.1$, 0.2, 0.3, and different accretion rates, $\dot{m} = 0.2$, 0.5. The spectrum is Compton dominated in the high accretion rate. The larger the $f$ is, the stronger both the Compton and the bremsstrahlung radiation become, since the coronal density and temperature increase with increased heating to the corona; while the optical-UV radiation, which is from the disk, slightly decreases with increasing $f$. The spectrum of hard X-ray in the range of $2$–$10$ keV is marked by the two vertical dashed lines. The hard X-ray spectrum is harder for higher $f$ and $L_{\text{bol}}/L_{2\text{-}10\text{keV}}$ decreases with $f$. For $f = 0.1$, the spectrum seems so steep that the corona radiation is still very weak. The photon index for the hard X-ray is in the range of $2.2$–$2.7$ for $0.2 \leq f \leq 0.4$, which roughly fits the spectrum of HLAGNs. Table 1 lists the photon index for $2$–$10$ keV and the ratio of $L_{\text{bol}}/L_{2\text{-}10\text{keV}}$. For the different accretion rates, the spectrum index is nearly the same with the same $f$, which can also be seen from the total spectrum shown in Figure 7. Comparing the two panels in Figure 7, we understand that the spectrum varies with the fraction of accretion energy liberating in the corona, while it does not sensitively depend on $\dot{m}$, as long as $f$ is independent on the accretion rate. This can be understood from our structure calculating results; namely, for a given $f$, the electron temperature in the corona, $T_e$, and the scattering optical depth, $\tau$, do not change as the accretion rate, though the disk radiation increases and hence Compton radiation also increases with increasing accretion rate.

The photon index and the luminosity ratio, as listed in Table 1, indicate that a fraction of accretion energy is necessary in order to explain observations in HLAGNs. For $0.2 \leq f \leq 0.4$, the predicted photon index is around $2.2$–$2.7$, and the bolometric correction from hard X-rays $L_{\text{bol}}/L_{2\text{-}10\text{keV}}$ is in the range of $28$–$228$. These are roughly consistent with observations (Vasudevan & Fabian 2007, 2009; Shemmer et al. 2006 and references therein; Zhou & Zhao 2010). In Figure 8, we show how the photon index in X-rays varies with bolometric luminosity. The figure reveals that the bolometric correction is larger for a steeper X-ray spectrum, which seems to be a common feature for objects with different Eddington ratios. This is a consequence of the decrease of energy fraction released in the corona. We now raise the questions, what are the underlying physics driving the accretion energy to release in the corona? Does the energy fraction $f$ depend on the accretion rate or the black hole mass? We discuss this in the following section.

### 4. DISCUSSION

#### 4.1. Heating Mechanism in the Corona—Magnetic Heating?

Studying the magnetic field in the accretion flow has a long history, either with magnetohydrodynamics (MHD) simulations (e.g., Balbus & Hawley 1998; Matsumoto 1999; Stone & Pringle 2001; Machida et al. 2001; Machida & Matsumoto 2003; Kato et al. 2004; Kawanaka et al. 2008 and references therein; Ohuga et al. 2009; Ohuga & Mineshige 2011; Penna et al. 2010) or numerical simulations to interpret observations (Liu et al. 2002b, 2003; Cao 2009).

There are good reasons to believe that the magnetic field plays an important role in transporting the accretion energy from the disk to the corona (Haardt & Maraschi 1991). Liu et al. (2002b, 2003) investigated the Parker instability in the accretion flow, following the model for the solar corona (e.g., Yokoyama...
& Shibata 2001), and found that the accretion energy can be transported to the corona and heats the electrons in the corona. In the frame of disk corona with mass evaporation, Qian et al. (2007) analyzed the coronal structure and evaporation features with the magnetic field and found that the maximum evaporation rate stays more or less constant (with the magnetic field and found that the maximum evaporation rate decreases with increasing radius corresponding to the maximum evaporation rate decreases with increasing magnetic field. Qiao & Liu (2009) investigated the influence of a viscosity parameter $\alpha$ in the corona and suggested that the coronal radiation is largely enhanced when viscosity increases.

However, it is not very clear how the magnetic field connects with viscosity in the evaporation-fed corona, though some investigations of accretion flows show that the magnetic field and viscosity are associated (e.g., Balbus & Hawley 1991; Matsumoto & Tajima 1995; Abramowicz et al. 1996; Balbus 2003). Here, we discuss whether the predicted spectrum can be consistent with observations if it is the magnetic field that brings a fraction of gravitational energy to the corona.

Let us assume that the magnetic field is generated in the disk by dynamo action. Then, the magnetic loops can emerge out of the disk by Parker buoyant instability. In this process, the accretion energy stored in the magnetic field is transported to the corona and heats the electrons in the corona. According to the equipartition theorem between the magnetic pressure and gas pressure, we set $P_B = B^2/8\pi = (1/\beta)P_{\text{g,d}}$, where $P_B$ and $P_{\text{g,d}}$ are magnetic pressure and gas pressure in the disk. Then, the magnetic energy flux is

$$F_B = \frac{B^2}{4\pi} \sqrt{\frac{B^2}{4\pi \rho_d}} = \left(2P_B\right)^{1/2} \rho_d^{-1} = \left(\frac{2}{\beta} P_{\text{g,d}}\right)^{1/2} \rho_d^{-1}. \ \ \ (20)$$

If this part of the magnetic energy flux is the additional heating source for the corona, then the ratio of this part of energy to total gravitational energy is

$$f = \frac{F_B}{F_{\text{tot}}} = \left(\frac{2}{\beta} P_{\text{g,d}}\right)^{1/2} \rho_d^{-1} \sqrt{F_{\text{tot}}}, \ \ \ \ (21)$$

where $F_{\text{tot}}$ is the total accretion flux through the accretion process, $F_{\text{tot}} = (3GM/M/8\pi R^3)[1 - (3R_s/R)^{1/2}] = 8.56 \times 10^{26} \mu \text{m}^{-1} r^{-3} \Phi$, where $\Phi = 1 - (3/R)^{1/2}$.

In the thin disk, the pressure of the disk is

$$P = P_{\text{g,d}} + P_B + P_{\text{g,d}} = \left(1 + \frac{1}{\beta}\right) \frac{\rho_d kT_d}{\mu m_p} + \frac{aT_d^4}{3}, \ \ \ \ (22)$$

where $\mu = 0.62$ is the averaged molecular weight. Since some fraction ($f$) of gravitational energy is carried into the corona by magnetic activity, the energy equation in the disk is

$$\frac{4acT_d^4}{3\tau} = (1 - f) \times F_{\text{tot}} = (1 - f) \frac{3GM}{8\pi R^3} \left[1 - (3R_s/R)^{1/2}\right]. \ \ \ (23)$$

where $\tau = (\kappa_{\text{es}} + \kappa_0 \rho_d T_d^{3.5})\Sigma$, $\kappa_{\text{es}} = 0.4 \text{ cm}^2 \text{ g}^{-1}$, $\kappa_0 = 6.4 \times 10^{22} \text{ cm}^2 \text{ g}^{-1}$, and $\Sigma = 2\rho_d H_d$ is the surface density of the disk. The angular momentum equation in the disk is

$$\frac{M}{3\pi} \Phi = \nu \Sigma, \ \ \ \ \ (24)$$

where $\nu = (2/3)ac_h H_d$, the viscous velocity $c_h = \Omega H_d$. 

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**Figure 7.** Total spectrum contributed from a multi-blackbody from the disk, bremsstrahlung, and inverse Compton scattering for $f = 0.2$ (left panel), and $f = 0.4$ (right panel).

**Figure 8.** Relation between the photon index ($\Gamma$) and bolometric correction ($L_{\text{bol}}/L_{\text{2-10 keV}}$). It shows that $\Gamma$ increases with $L_{\text{bol}}/L_{\text{2-10 keV}}$, which is a consequence of decreasing energy fraction released in the corona.
If the gas pressure in the disk is dominant over radiation pressure, $P_{\text{g,d}} \gg P_{r,d}$, and Thomson scattering is much higher than free–free absorption, $\sigma_T \gg \sigma_{\text{ff}}$, then we have $\tau \approx \kappa_{\text{es}} \Sigma$, which is similar to case (b) in Shakura & Sunyaev (1973), except that $r = R/R_c$ is taken here instead of $r = R/3 R_c$. We can derive the disk quantities. The density in the disk is

$$\rho_d = 21.43(\alpha m)^{-\frac{3}{2}}(\dot{m} \Phi)^{\frac{1}{2}} r^{-\frac{31}{2}} \left(1 + \frac{1}{\beta^2}\right)^{-\frac{3}{4}} (1 - f)^{-\frac{3}{2}}, \quad (25)$$

and the gas pressure is

$$P_{\text{g,d}} = 1.91 \times 10^{18}(\alpha m)^{-\frac{5}{2}}(\dot{m} \Phi)^{\frac{1}{2}} r^{-\frac{31}{2}} \left(1 + \frac{1}{\beta^2}\right)^{-\frac{5}{4}} (1 - f)^{-\frac{5}{2}}. \quad (26)$$

With this disk pressure and density, the energy fraction can be obtained from Equation (21),

$$f = 1.88 a^{-1}(1 + \beta)^{-\frac{1}{2}}. \quad (27)$$

One can see that $f$ is independent of the accretion rate if the disk is gas pressure dominated. This implies that a constant fraction of disk accretion energy is transported to the corona at different accretion rates if equipartition coefficient $\beta$ does not vary with accretion rates. If this is the case for the additional heating to the corona, then the model spectrum will not vary with the accretion rate, as shown in Figure 7.

In fact, the accretion disk around a supermassive black hole is often dominated by radiation pressure. In the case of $P_{\text{g,d}} \ll P_{r,d}$ and $\sigma_T \gg \sigma_{\text{ff}}$, the density and pressure in the disk are expressed as

$$\rho_d = 1.14 \times 10^{-6}(\alpha m)^{-\frac{3}{2}}(\dot{m} \Phi)^{\frac{1}{2}} r^{-\frac{31}{2}} (1 - f)^{-\frac{3}{2}}, \quad (28)$$

$\dot{P}_{\text{g,d}} = 6.21 \times 10^9(\alpha m)^{-\frac{5}{2}}(\dot{m} \Phi)^{\frac{1}{2}} r^{-\frac{31}{2}} (1 - f)^{-\frac{5}{2}}, \quad (29)$

which are similar to the solution of case (a) in Shakura & Sunyaev (1973). Then, we can determine $f$ from Equation (21) as

$$f(1 - f)^{\frac{3}{2}} = 1.52 \times 10^{-9} a^{-\frac{11}{4}} m^{-\frac{3}{4}} (\dot{m} \Phi)^{\frac{1}{4}} r^{-\frac{35}{4}} \beta^{-\frac{3}{2}}. \quad (30)$$

Apparently, in a radiation-pressure-dominated disk, $f$ depends on the accretion rate. The function $f(1 - f)^{2/3}$ in Equation (30) reaches a maximum of 0.095 at $f = 8/35$, which requires $1.52 \times 10^{-9} a^{-11/8} m^{-3/8} (\dot{m} \Phi)^{-3} r^{33/16} \beta^{-3/2} \leq 0.095$. This indicates that the accretion rate must be larger than a certain value in order to get a solution for $f$ from Equation (30). If this condition $1.52 \times 10^{-9} a^{-11/8} m^{-3/8} (\dot{m} \Phi)^{-3} r^{33/16} \beta^{-3/2} \leq 0.095$ is fulfilled, then there are two sets of solutions for $f$ (we saw this bimodal nature in Liu et al. 2002b). As shown in Figure 9, one set of solutions is in the range of $0 < f < (8/35)$, where $f$ decreases with increasing accretion rate $\dot{m}$, while the other one is in the range of $(8/35) < f < 1$, where $f$ increases with increasing $\dot{m}$. The relation between the energy fraction and the accretion rate, i.e., $f(\dot{m})$, depends on the viscosity and magnetic field. In the left panel of Figure 9, we show the effect of viscosity by assuming $\beta = 1$ under the theory of the equipartition of gas pressure and magnetic pressure. Actually, the local and global MHD simulation in the radiation-dominated accretion flow reveals that magnetic energy density may exceed gas energy density but be kept below radiation energy density, which means that $\beta$ can be less than 1.0 (e.g., Turner et al. 2003; Ohsuga et al. 2009; Ohsuga & Mineshige 2011). The effect of magnetic fields is plotted in the right panel of Figure 9. It shows that the influence of the magnetic field is significant. The stronger the magnetic field, the larger $\dot{m}$ corresponding to $f = 8/35$. For a given accretion rate, for example, $\dot{m} = 0.4$, $f$ is larger with stronger magnetic field (i.e., smaller $\beta$) in the lower part of the curve ($f < 8/35$).

Observed spectral features provide a clue to the possible range of $f$. If $f < (8/35)$, then the fraction of accretion energy transferred to the corona decreases with accretion rate. As a consequence, the hard X-ray spectrum becomes softer at high accretion rates, as shown in Figure 10. However, if a large fraction of accretion energy ($f > (8/35)$) is transferred to the corona, $f$ increases with accretion rate, consequently, the spectrum is harder at higher accretion rates. Comparing the model prediction and observations, a small fraction of accretion energy liberation in the corona is preferred, that is, $<8/35$. In Figure 10, we plot the spectrum for different accretion rates. The dashed lines represent a spectrum for a low accretion rate, and the dotted line is for a relatively higher accretion rate, and the
In our corona, we assume that all of the viscously dissipated energy heats ions only. However, direct heating to the electrons can also increase the corona emission. According to the detailed modeling to Sgr A* (Yuan et al. 2003), a significant fraction of the viscously dissipated energy should heat electrons directly. Such a result was later supported by the numerical simulation done by Sharma et al. (2007). Nevertheless, this effect is negligible if much larger additional heating to the corona is assumed. When a strong disk exists under a corona as in luminous AGN, a small fraction of disk accretion energy transported to coronal electrons can dominate the coronal heating. For instance, assuming $m = 0.2$ and $f = 0.1$, the additional heating to the corona is 0.02 (in units of Eddington luminosity); while the accretion rate in the corona, as shown in Figure 4, is $10^{-4}$ to $10^{-3}$ in the major radiation region, which produces viscous heat of $10^{-4}$ to $10^{-3}$ (in units of Eddington luminosity). Even if all of this viscous heating directly goes to electrons, it is much smaller than the additional heating 0.02. The ratio is in the range of 0.5% to 5%. This ratio is nearly the same for a higher $f$ since the accretion rate in the corona increases with increasing $f$. For an accretion rate higher than 0.2, the ratio is smaller, which can be derived from the accretion rate of the corona shown in the right panel of Figure 4. Therefore, the neglect of viscous heating directly to electrons is a reasonable approximation.

4.3. Disk Wind at High Accretion Rate

When the accretion is greater than 0.1 $M_{\text{Edd}}$, the line-driven wind and radiation-pressure-driven wind will be important, which results in the accretion rate decreasing with the decrease in radius as $\dot{M} \propto R^a$ ($0.0 < a < 1.0$) (see Proga et al. 2000; Ohsuga et al. 2005). It is clear that with wind loss, the soft-photon field becomes weaker than the case without winds. Therefore, the Compton cooling is weaker, leading to less condensation. When the system reaches equilibrium, the spectrum can be harder. This is the case without additional energy input. In this work, additional energy is added into the corona and we find that the larger the $f$ is, the harder the spectrum is. Since $0 < a < 1$, the energy generating rate $\propto (G M M/R) \propto (M_{\text{out}}/R^{1-a})$ is still dominated by the innermost region. Thus, the wind loss causes a decrease of disk luminosity and hence a harder spectrum. However, if wind mass loss is very significant, then there could no longer be an optically thick disk and the effects are large.

4.4. Irradiation

A corona lying above a cool disk can illuminate the disk, which increases the energy density of soft photons, leading to stronger coronal radiation. This effect can be significant in the case of a strong corona above a truncated disk. However, for a relatively weak corona above a strong disk, as in the case of HLAGNs, the effect of irradiation is unimportant. We estimate this effect as follows.

Photons emitted from corona are roughly isotropic, about half of which go down to a slab-like disk, irradiating the disk. The irradiating flux is

$$ F_{\text{irr}} = \frac{1}{2} F_c (1 - a) = \frac{1}{2} F_{\text{tot}} (1 - a), $$

where $a$ is the albedo, which is usually taken as 0.15 (e.g., Zdziarski et al. 1999). This flux is scattered/absorbed in the disk and eventually re-emitted, contributing to the soft photon field. The ratio of this illuminating flux to the disk radiation, $F_d = F_{\text{tot}} (1 - f)$, is

$$ \frac{F_{\text{irr}}}{F_d} = \frac{\frac{1}{2} f (1 - a)}{(1 - f)}. $$

This ratio reaches 1 when $f > 0.7$, indicating that the irradiation becomes dominant only when a large fraction of accretion energy is released in the corona. Therefore, we do not include the irradiation in this work in order to make the numerical calculations simple.

4.5. Comparison with the MHD Results

Hirose et al. (2006) performed the shearing box numerical simulation with radiative transport, focusing on the vertical structure of the standard thin disk around $300 R_g$ ($R_g = \text{...
They found that there is actually very little dissipation in the coronal region as the corona is magnetically dominated. Consequently, the corona is too weak to explain the hard-X-ray emission observed in AGNs. The disk corona evaporation/condensation model without additional heating gives similar results to an MHD simulation (for details see Meyer-Hofmeister et al. 2012). When a corona is heated only by viscous dissipation within the corona, efficient inverse Compton scattering of the disk photons leads to overheating of the corona. As a consequence, part of the coronal gas condenses into the disk, leaving a very weak corona.

To prevent the corona from collapsing, we assume in this paper that a fraction of accretion energy is transported from the disk and released into the corona, which conflicts with the MHD simulation results. We note that in the MHD simulation of Hirose et al. (2006), it was presumed that the disk is dominated by gas pressure, which is reasonable for studying the vertical structure of black hole X-ray binaries in about 300R_g. However, we are investigating the inner region of an AGN disk, where the disk is radiation-pressure dominant. Recently, Blaes et al. (2011) simulated such a kind of accretion flow and found that the thermodynamics of a radiation-pressure dominant disk differ significantly from those envisaged in standard static models of accretion disks. In radiation-dominant plasma, the buoyant motions become significant for the overall energetics of the plasma. The photons’ outward advection becomes comparable to radiative diffusion, and the associated vertical expansion work balances vertical heat transport and dissipation. At the present stage, it is not clear how much discrepancy exists between the MHD simulation and our results.

4.6. Luminous Hot Accretion Flow as an Alternative Model

Based on the thermal stability analysis presented in Yuan (2003), luminous hot accretion flows (LHAFs; Yuan 2001) could also have a two-phase structure. Yuan & Zdziarski (2004) calculated the emitted luminosity from such a kind of two-phase accretion flow and found that it can produce X-ray luminosity as high as 10% Eddington luminosity. This is an alternative model for AGNs with intermediate spectrum and Eddington ratio. In our work, we aim at a more general case for HLAGNs with luminosity up to Eddington value, and calculate the broad waveband spectra from both the disk and the corona.

5. CONCLUSION

We applied the disk evaporation model to the HLAGNs. To explain the typical spectra of HLAGNs, we conclude that there should be additional heating to the corona to prevent it from overheating by strong inverse Compton scattering.

We assume that a fraction of gravitational energy (f) is liberated in the corona and calculate the corona structure. Then, the spectrum from the disk and corona is calculated by a Monte Carlo simulation. We found that the hard X-ray is dominated by Compton cooling and the photon index for hard X-ray in 2–10keV is 2.2 < \Gamma < 2.7. For a given accretion rate, \Gamma and L_{\text{bol}}/L_{2–10\text{keV}} decrease with the increase of f. We discuss a possible mechanism for the presumed heating to the corona, that is, the magnetic heating. With equipartition of magnetic energy to the gas energy in the disk, we derive the fraction of accretion energy stored in magnetic field. Assuming this energy is released in the corona, the model predicts that, for f < (8/35), the hard X-ray becomes softer at a higher accretion rate and L_{\text{bol}}/L_{2–10\text{keV}} increases, which is roughly consistent with the observational results.
