GLOBAL COMPTON HEATING AND COOLING IN HOT ACCRETION FLOWS

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ABSTRACT

The hot accretion flow (such as advection-dominated accretion flow) is usually optically thin in the radial direction, therefore the photons produced at one radius can travel for a long distance without being absorbed. These photons can thus heat or cool electrons at other radii via Compton scattering. This effect has been ignored in most previous works on hot accretion flows and is the focus of this paper. If the mass accretion rate is described by \( \dot{M} = M_0(r/r_{\text{out}})^{0.3} \) and \( r_{\text{out}} = 10^3 r_s \), we find that the Compton scattering will play a cooling and heating role at \( r \lesssim 5 \times 10^3 r_s \) and \( r \gtrsim 5 \times 10^3 r_s \), respectively. Specifically, when \( M_0 > 0.1L_{\text{Edd}}/c^2 \), the Compton cooling rate is larger than the local viscous heating rate at certain radius; therefore the cooling effect is important. When \( M_0 > 2L_{\text{Edd}}/c^2 \), the heating effect at \( r_{\text{out}} \) is important. We can obtain the self-consistent steady solution with the global Compton effect included only if \( M_0 \lesssim L_{\text{Edd}}/c^2 \) for \( r_{\text{out}} = 50 r_s \), which corresponds to \( L \lesssim 0.02L_{\text{Edd}} \).

Above this rate the Compton cooling is so strong at the inner region that hot solutions cannot exist. On the other hand, for \( r_{\text{out}} = 10^3 r_s \), we can only get the self-consistent solution when \( M_0 \lesssim L_{\text{Edd}}/c^2 \) and \( L < 0.01L_{\text{Edd}} \). The value of this critical accretion rate is antirelated with the value of \( r_{\text{out}} \). Above this accretion rate, the equilibrium temperature of electrons at \( r_{\text{out}} \) is higher than the virial temperature as a result of strong Compton heating, so the accretion is suppressed. In this case the activity of the black hole will likely “oscillate” between an active and an inactive phase, with the oscillation timescale being the radiative timescale of the gas at \( r_{\text{out}} \).

Key words: accretion, accretion disks – black hole physics – galaxies: active – quasars: general – X-rays: general

1. INTRODUCTION

Compton scattering between photons and electrons is an important process in astrophysics. If the photons are not produced at the same place where the electrons are located, we call it “global Compton scattering.” Momentum and energy of photons and electrons can be exchanged in this process and these two aspects often play an important role in determining the dynamics of the gas flow. On the galactic scale, this so-called radiative feedback mechanism is now believed to be crucial for understanding active galactic nucleus (AGN) feedback on galaxy formation and evolution (e.g., Ciotti & Ostriker 1997, 2001, 2007; Murray et al. 2005; Hopkins et al. 2005). On a smaller scale, the effect of Compton scattering on the dynamics of gas flows surrounding a strongly radiating quasar has been investigated and outflow is produced as a consequence of Compton heating (e.g., Proga et al. 2008). Following the earlier work of Krolik et al. (1981), Mathews & Ferland (1987) considered the Compton heating effect for the broad-line region (BLR) of quasar. This effect is also important for the standard thin disk if the disk is warped or irradiated by a source above the disk plane (e.g., Shakura & Sunyaev 1973; Begelman et al. 1983; Dubus et al. 1999). If the accretion flow is geometrically thick and optically thin, the photons can travel a large distance without being absorbed, therefore the global Compton scattering effect is in principle important. This is the case for spherical accretion and hot accretion flows. The latter includes the advection-dominated accretion flow (ADAF; Narayan & Yi 1994, 1995) and luminous hot accretion flow (LHAF; Yuan 2001, 2003), two types of hot accretion flows corresponding to low and high accretion rates, respectively.

For spherical accretion, the interaction of momentum between photons and electrons sets up the largest possible luminosity the accretion flow could reach, namely the Eddington luminosity \( L_{\text{Edd}} \) (but this limit does not apply when the accretion flow has a nonzero angular momentum; see, e.g., Ohshima & Mineshige 2007). The effect of the energy interaction between photons and electrons in a spherical accretion flow has been investigated by Ostriker et al. (1976). It was found that when the luminosity is larger than a certain value, the outward energetic photons could heat gas flow significantly so that the local sound speed is larger than the escape speed, or, in other words, the temperature is higher than the virial temperature, thus the accretion is suppressed. This effect, due to energy input from the outgoing radiation field, occurs for much lower luminosities than the momentum (Eddington) limit.

In almost all of the previous works on the dynamics of hot accretion flows, only the “local” Compton scattering effect has been considered while the global Compton effect has been neglected. Here “local” means that photons are produced at the same region where the electrons locate. This local Compton scattering serves as the main cooling mechanism of electrons (Compton cooling) and the main mechanism of producing X-ray emission (thermal Comptonization). To our knowledge, the only works considering the global Compton scattering are Esin (1997) and Park & Ostriker (1999, 2001, 2007). Esin (1997) deals with a one-dimensional ADAF and finds that the global Compton heating/cooling is not important and can be neglected. Park & Ostriker (2001, 2007) deal with a two-dimensional flow and focus on the possible production of outflow in the polar region because of strong Compton heating there. Their conclusion is that Compton heating effect is important in many cases, which is quite different from Esin’s result. They do not attempt to obtain the self-consistent solutions.

All their works are based on the self-similar solution of ADAF (Narayan & Yi 1994, 1995). The discrepancy between Esin...
and significant direct electron heating by turbulent dissipation (1997) and Park & Ostriker (1999, 2001, 2007) is likely due to the additional but different assumptions adopted. While the self-similar approximation is quite successful in catching the main spirit of an ADAF, it is not a good approximation when we want to calculate the radiation since order of magnitude error could be produced. This is because the self-similar approximation breaks down at the inner region of the ADAF where most of the radiation comes from. When we consider the effect of global Compton heating/cooling, obviously it is crucial to calculate the exact spectrum from the exact global solution of the accretion flow. This is the main motivation of the present paper. In addition, important theoretical progress on ADAF solutions have been made since its discovery. The two most important ones are the presence of outflow (e.g., Stone & Pringle 2001) and significant direct electron heating by turbulent dissipation (e.g., Quataert & Gruzinov 1999), which is much stronger than the heating by Coulomb collisions between ions and electrons. Both of them are important in determining the dynamics of ADAFs while they have not been properly taken into account in the above works. In the present paper, we will focus on the one-dimensional case, given that the uncertainty of our understanding of the two-dimensional structure of accretion flow is still large. Of course, a full understanding of the two-dimensional case is obviously important and should be a topic of future research.

2. THE IMPORTANCE OF COMPTON HEATING OR COOLING IN HOT ACCRETION FLOWS

2.1. Calculation Method

Consider a canonical hot accretion flow without taking into account the global Compton heating or cooling. Outflow is taken into account by adopting the following radius-dependent mass accretion rate (e.g., Blandford & Begelman 1999):

$$M = -4\pi r H \rho v = \dot{M}_0 \left(\frac{r}{r_{\text{out}}}\right)^s,$$

(1)

where $\dot{M}_0$ is the mass accretion rate at the outer boundary $r_{\text{out}}$. The value of index $s$ describes the strength of the outflow and we use $s = 0.3$ from the detailed modeling of Sgr A* (Yuan et al. 2003). The energy equations for ions and electrons are

$$\rho v \left(\frac{d\epsilon_i}{dr} - \frac{p_i}{\rho^2} \frac{dp_i}{dr}\right) = (1 - \delta)q^+ - q_e,$$

(2)

$$\rho v \left(\frac{d\epsilon_e}{dr} - \frac{p_e}{\rho^2} \frac{dp_e}{dr}\right) = \delta q^+ + q_e - q^-,$$

(3)

where $\epsilon_{i,e}$ is the internal energy of ions and electrons per unit mass of the gas, $q_e$ is the Coulomb energy exchange rate between electrons and ions, $q^+$ is the electron cooling rate, including synchrotron and bremsstrahlung emissions and their local Comptonization, $q^-$ is the net turbulent heating rate, the value of $\delta$ describes the fraction of turbulent heating which directly heats electrons and we use $\delta = 0.5$ again from the modeling of Sgr A* (see also Sharma et al. 2007).

We first get the exact global solution of the hot accretion flow so that we know all the quantities such as density and temperature as a function of radius. This requires us to solve the set of equations describing the conservations of mass (Equation (1)), energy (Equations (2) and (3)), and momentum (Yuan et al. 2003). The global solution should satisfy the outer boundary condition at the outer boundary $r_{\text{out}}$, the inner boundary condition at the horizon, and a sonic point condition at the sonic point. To calculate the rate of “global” Compton heating/cooling of electrons at a given radius $r$, we need to know the spectrum received at $r$ emitted by the whole flow. This requires us to solve the radiative transfer equations along the radial direction, which is complicated when scattering is important. For simplicity, here we deal with the scattering in a simple way and write the received spectrum at $r$ emitted by the flow inside of $r$ as

$$F_v^\text{in}(r) = \int_{r_s}^{r} e^{-\tau} \frac{1}{4\pi r^2} \frac{dL_v(r')}{dr'} dr'.$$

(4)

Here, $\tau$ is the scattering optical depth from $r'$ to $r$, $\tau = \int_{r'}^{r} \sigma T n_e dr'$, and $dL_v(r')$ is the emitted monochromatic luminosity from a shell at $r'$ with thickness $dr'$ and height $H(r')$. It includes synchrotron and bremsstrahlung emissions and their local Comptonization. We approximate the calculation of the unscattered part of $dL_v(r')$ by solving the radiative transfer along the vertical direction of ADAFs adopting a two-stream approximation (Mannomo et al. 1997),

$$dL_v^\text{out}(r') = \frac{4\pi^2}{\sqrt{3}} B_v [1 - \exp(-2\sqrt{3}\tau_v^*)] r' dr'.$$

(5)

Here, $B_v$ denotes the Planck spectrum, $\tau_v^* \equiv (\pi/2)^{1/2} k_\nu (0) H(r')$ is the optical depth for absorption in the vertical direction with $k_\nu (0)$ being the absorption coefficient on the equatorial plane. The free–free absorption and synchrotron self-absorption are included in this way. The strength of the magnetic field in the accretion flow is determined by a parameter $\beta$ defined as the ratio of the gas pressure to the magnetic pressure and we set $\beta = 9$. For the calculation of the Compton scattered part of $dL_v(r')$, we use the approach of Coppi & Blandford (1990; Equation (2.2)). The integration in Equation (4) begins from the black hole horizon $r_s \equiv 2GM/c^2$.

The spectrum received at $r$ emitted by the flow outside of $r$ is (Park & Ostriker 2007)

$$F_v^\text{out}(r) = \int_{r}^{r_{\text{out}}} \frac{e^{-\tau}}{4\pi r'H(r')} r' \ln \frac{r' + r}{r' - r} \frac{dL_v(r')}{dr'} dr'.$$

(6)

The total spectrum received at $r$ is the sum of $F_v^\text{in}(r)$ and $F_v^\text{out}(r)$. We assume that the electrons have a Maxwell distribution with temperature $T_e (\theta_e \equiv kT_e/m_e c^2)$ and the energy of the photon before scattering is $\epsilon \equiv h\nu/m_e c^2$. Since the electrons in the hot accretion flow are relativistic at the innermost region and the peak photon energy $\epsilon \approx 1$, to calculate the average energy of a scattered photon, we use the following exact form which is valid for any photon energy and electron temperature ( Guilbert 1986):

$$(\epsilon_1) = \epsilon + \frac{\sigma_T}{2K_2(1/\theta_e)\sigma} \int_{-\infty}^{+\infty} (\theta_e + \sin \phi - \epsilon) G(\epsilon_1 e^{\phi} e^{2\phi}) \times \exp \left(\frac{-\cos \phi}{\theta_e}\right) d\phi,$$

(7)

with $G(\epsilon) \equiv g_0(\epsilon) - g_1(\epsilon)$ and the cross section for scattering

$$\sigma(\epsilon, \theta_e) = \frac{\sigma_T}{2K_2(1/\theta_e)} \int_{-\infty}^{+\infty} g_0(\epsilon_1 e^{\phi} e^{2\phi}) \exp \left(\frac{-\cos \phi}{\theta_e}\right) d\phi.$$
Here, $K_2(x)$ is a modified Bessel function of second order and
\[
g_{n}(y) \equiv \frac{3}{8} \int_{0}^{\infty} \left( t(t - 2) + 1 + ty + \frac{1}{1 + ty} \right) \frac{dt}{(1 + ty)^{n+2}}. \tag{9}
\]
In the Thompson limit, Equations (7) and (8) are transformed into the familiar form of
\[
\langle \epsilon_1 \rangle = \epsilon + \epsilon \frac{4kT_e - \epsilon m_e c^2}{m_e c^2}, \tag{10}
\]
and
\[
\sigma(\epsilon, \theta_e) = \sigma_T. \tag{11}
\]
The number of scattering in a region of the accretion flow with unit width in the radial direction and scattering optical depth $\tau_{es} \equiv \sigma(\epsilon, \theta_e) n_e$ is
\[
N = \tau_{es}, \tag{12}
\]
with $\theta_e$ and $n_e$ being the temperature and number density of electrons in that region. The Compton heating (cooling) rate in that region (with unit radial length) is then
\[
q_{\text{comp}} = \int N \left[ F_v^{\text{in}}(r) + F_v^{\text{out}}(r) \right] \frac{\epsilon}{\epsilon} \langle \epsilon_1 \rangle d\nu. \tag{13}
\]
Note that in the above equation we actually use the moment of intensity “$\nu$” not “$F$”. Following Park & Ostriker (2007), we formally define a “radiation temperature” (or “Compton temperature”; see also Levich & Sunyaev 1970; Krolik et al. 1981) as
\[
\theta_c \equiv \frac{\int \left[ F_v^{\text{in}}(r) + F_v^{\text{out}}(r) \right] h\nu d\nu}{4m_e c^2 \int \left[ F_v^{\text{in}}(r) + F_v^{\text{out}}(r) \right] d\nu}. \tag{14}
\]
Under this definition, whether the Compton scattering plays a heating or cooling effect roughly depends on whether the electron temperature $\theta_e$ is larger or smaller than $\theta_c$. In the Thompson limit, the Compton heating/cooling rate is exactly proportional to $(\theta_e - \theta_c)$.

It is important to note here that the radiation temperature is obtained from the flux distribution by weighting with the factor $h\nu$. It physically represents the equilibrium between Compton heating and cooling. For a typical quasar spectrum, it corresponds to $2 \times 10^7 \text{K}$ or several keV where the bulk of the radiation is emitted in the UV or, in some cases, IR portions of the spectrum (Mathews & Ferland 1987; Sazonov et al. 2004). As we will see below (e.g., Figure 3(c)), the radiation temperature of an ADAF spectrum is much higher because of its different spectrum.

### 2.2. Results

The dominant heating term of electrons in Equation (3) is $q_{\text{shc.e}} \equiv dq^s$. We compare the rate of global Compton heating/cooling with $q_{\text{shc.e}}$ and the results are shown in Figure 1 for $M = M_0 / r_0^{0.3}$ with $M_0 = 0.1, 1$, and $2 M_{\odot}$ ($M_{\odot} \equiv L_{\odot} c^2$). At large radii, $r \gtrsim 5 \times 10^3 r_s$, Compton scattering heats electrons; while at small radii, $r \lesssim 5 \times 10^3 r_s$, it cools electrons. This is of course because the radiation temperature $\theta_c$ is lower (higher) than the electron temperature $\theta_e$ at the small (large) radii, as shown by the right panel of Figure 1.

We can see from the left panel of Figure 1 that the Compton effect is important when $M_0 \gtrsim 0.1 M_{\odot}$. In this case, its cooling effect cannot be neglected. The corresponding accretion rate at the black hole horizon is $\sim 10^{-2} M_{\odot}$ and the corresponding luminosity is $\sim 5 \times 10^4 L_{\odot}$. The “lowest” value of $M_0$ above which Compton heating effect is important is a function of $r_{\text{out}}$.

For $r_{\text{out}} = 10^4 r_s$, this value is $\sim 2 M_{\odot}$ and the corresponding luminosity is $\sim 2 \times 10^4 L_{\odot}$. When $r_{\text{out}}$ is larger, the value of $M_0$ is lower. In reality $r_{\text{out}}$ usually has a largest feasible value. If the ADAF starts out from a transition from a standard thin disk, $r_{\text{out}}$ equals the transition radius. On the other hand, if the accretion flow starts out as an ADAF such as in our Galactic center, $r_{\text{out}}$ should be determined by the Bondi radius. Outside the Bondi radius, the effect of Compton heating is not so clear, because matching an ADAF solution to one with proper boundary condition at infinity is an unsolved problem.

Our result that Compton scattering heats electrons at large radii while cools electrons at small radii is qualitatively consistent with both Esin (1997) and Park & Ostriker (2001, 2007). However, Esin (1997) found that the Compton heating rate is always smaller even than the Coulomb collision heating rate and therefore is negligible. This is different from our results and Park & Ostriker (2001, 2007). The reason may come from some oversimplifications and the different (old) ADAF model adopted in Esin (1997). Overall, we see that for extended solutions ($r_{\text{out}} \gtrsim 10^4 r_s$), both Compton cooling in the inner parts and Compton heating in the outer parts dramatically alter the solutions when $L \gtrsim 10^{-2} L_{\odot}$.

### 3. THE SELF-CONSISTENT SOLUTIONS

The above result indicates that we should take into account the effect of global Compton heating/cooling when we calculate the global solution of the hot accretion flow when $M$ is relatively large. This has not been studied in Esin (1997) and Park & Ostriker (1999, 2001, 2007). We use an iteration method to achieve this. We first solve the global solution without considering the global Compton effect, calculating the rate of Compton heating/cooling at each radius as described above, $q_{\text{comp}}$. We then include this term in the energy equations of electrons,
\[
\rho \nu \left( \frac{d\epsilon_e}{dr} - \frac{p_e}{\rho^2} \frac{d\rho}{dr} \right) = dq^s + q_{\text{shc.e}} - q - q_{\text{comp}}, \tag{15}
\]
and calculate the “new” global solution of the accretion flow based on this “new” equation. Then we get a new Compton heating/cooling rate. If the new rate is not equal to the guessed value we replace the guessed value with the new one and repeat this procedure until they are equal. However, we must emphasize that the solution obtained by the above approach is actually not exactly “self-consistent.” We use Equation (13) to calculate $q_{\text{comp}}$. But in Equation (13), only the local Compton scattering is considered when calculating $F_v^{\text{in}} + F_v^{\text{out}}$, and the global scattering is difficult to include because we do not know the spectrum emitted at other radii which again requires us to consider global scattering. This is difficult to deal with even using the iteration approach. The best way to solve this problem is by Monte Carlo simulation combined with iteration method. This is beyond the scope of this paper and will be our next work. On the other hand, we believe our result should be a good zeroth-order approximation to the real solution.

Bearing this in mind, Figures 2 and 3 show the calculation results. Figures 2(a)–2(c) are for a stellar mass black hole with black hole mass $M = 10 M_\odot$ and $M = (r/50r_s)^{0.3} M_{\odot}$. Figure 2(a) shows the electron temperature of the global solution with (dashed line) and without (solid line) the global Compton
scattering effect included. Because Compton scattering plays a cooling role at small radii, we see that the electron temperature decreases after the Compton effect is taken into account as we expect.

Figure 2(b) shows the ratio of Compton heating and local viscous heating of the electrons, \( q_{\text{comp}}/q_{\text{vis,e}} \), before (solid line) and after (dashed line) the global Compton effect is included. It is interesting to note that at \( r \gtrsim 4r_s \), the absolute value of this ratio is between 1 and 4. Since we typically have \( q_{\text{e}} \ll q_{\text{vis,e}} \), this implies that the right-hand side of Equation (15) is negative in that region, i.e., the viscous heating of electrons is smaller than its radiative cooling \( (q^- - q_{\text{comp}}) \). In other words, the energy advection of electrons plays a heating role, just like the ions in the LHAF solution (Yuan 2001). In the inner region of \( r \lesssim 4r_s \), where most of the radiation comes from, the absolute value of \( q_{\text{comp}}/q_{\text{vis,e}} \sim 0.5 \). We find in this case \( (q^- - q_{\text{comp}}) \sim q^- \). So in the innermost region the viscous heating of electrons is equal to its radiative cooling \( (q^- - q_{\text{comp}}) \).

Obviously, after taking into account the global Compton cooling, the radiative efficiency will increase for a given accretion rate. But this does not mean that the highest luminosity \( L_{\text{max}} \) a hot accretion flow can produce will increase. The main heating mechanism of electrons are viscous heating and compression work (the second term in the right-hand side of Equation (15)) while the main cooling comes from \( (q^- - q_{\text{comp}}) \). The highest accretion rate beyond which a hot solution no longer exists is determined by the balance between heating and cooling. The heating term is roughly proportional to \( M \) while the cooling term roughly to \( M^2 \) since Compton scattering is a two-body collision process. This is why a hot accretion solution has a highest \( M \). Obviously, when the global Compton cooling \( q_{\text{comp}} \) is included, the cooling becomes stronger compared to the case of only local cooling \( q^- \), thus the balance between cooling and heating will occur at a lower \( M \). Actually \( M_0 = L_{\text{Edd}}/c^2 \) as shown in Figure 2 is almost the highest accretion rate at which we can get the self-consistent hot solution, which is a factor of 2–3 lower than the highest rate when the global Compton effect is not taken into account. When \( M_0 \) is higher, we find that we are not able to get the self-consistent solution since the flow will collapse due to the strong radiative cooling. The decrease of the highest \( M \) results in the decrease of \( L_{\text{max}} \). Our calculation shows that \( L_{\text{max}} \) decreases by a factor of \( \sim 2 \), i.e., we now have \( L_{\text{max}} \sim 3\% L_{\text{Edd}} \).

We now check how different the spectrum produced by the self-consistent solution is compared to the spectrum produced by the “old” solution. Figure 2(c) shows the spectra from the hot accretion flow before (solid line) and after (dashed line) the Compton effect is taken into account. We see that both the luminosity and the cutoff energy of the spectrum (i.e., the corresponding frequency of the peak of the spectrum) decrease because of the global Compton cooling. This is of course because the electron temperature of the self-consistent solution decreases compared to the “old” solution. We would like to emphasize again that only local seed photons are considered when we calculate the spectrum although we do consider the global scattering in calculating the dynamics. Our calculation shows that \( -q_{\text{comp}} \sim q^- \), so we expect that when the global Compton scattering is considered, the luminosity of the “exact” self-consistent solution will be \( \sim 2 \) times higher than that shown by the dashed line in Figure 2(c). But the slope and the cutoff...
energy will not change because they are irrelevant to the amount of seed photons.

Figures 3(a)–3(c) are similar to Figures 2(a)–2(c), but are for a supermassive black hole with $M = 10^8 M_\odot$ and accretion rate $\dot{M} = 0.5(r/10^4 r_s)^{0.3}\dot{M}_{\text{Edd}}$. We see from the figures that the electron temperature decreases after the global Compton scattering effect is taken into account as we expect, because in most regions the Compton scattering will cool the electrons. Correspondingly, the luminosity of the accretion flow also decreases by roughly a factor of 2 and the cutoff energy of the spectrum also becomes smaller.

We have also calculated the average energy of the photons emitted by the self-consistent solution shown in Figure 3(c), 

$$h\langle \nu \rangle \equiv \int L_\nu d\nu / \int (L_\nu / h\nu) d\nu,$$

and the corresponding energy of the radiation temperature $\theta_\nu$ at $r_{\text{out}}$, $h\nu_\nu \equiv m_e c^2 \theta_\nu$. The results are $\sim 1 \text{ eV}$ and $100 \text{ keV}$, respectively, and they are shown by two arrows in Figure 3(c). These values are much higher than that of a typical quasar spectrum where, e.g., $h\nu_\nu$ is only several keV (Mathews & Ferland 1987; Sazonov et al. 2004).

The spectrum shown in Figure 3(c) extends to very high energy, $\gtrsim \text{MeV}$. Observationally, the $e$-folding energy of the average power-law X-ray spectrum observed by Ginga, OSSE, and EXSOSAT of radio-quiet Seyfert 1s is $E_c = 0.7^{+0.6}_{-0.3}$ MeV (Zdziarski et al. 1995; Gondek et al. 1996), which is consistent with the model given the (large) error bar. Better data are required to constrain the theoretical model.

As we state in Section 2.2, Compton heating effect at large radii is another obstacle for us to obtain the self-consistent solution. For $M_\odot = \dot{M}_{\text{Edd}}$, if $r_{\text{out}} \gtrsim 10^5 r_s$, we find that the Compton heating effect around $r_{\text{out}}$ is so strong that the equilibrium temperature of electrons would be higher than the virial value defined as $5/2k T_{\text{vir}} = G M m_p / r$, which will in turn make the ion temperature also higher than the virial one due
to the Coulomb coupling between them. In this case, the gas is unbound and thus cannot be accreted. The corresponding highest luminosity in this case is $\sim 2\% L_{\text{Edd}}$. This value is similar to that obtained by Ostriker et al. (1976) and Park & Ostriker (2001). From Figures 1 and 3(b), we expect that $M_0$ and the critical $r_{\text{out}}$ (signed as $r_{\text{virial}}$) beyond which the equilibrium temperature is higher than the virial temperature are roughly anticorrelated, i.e., a lower $M_0$ corresponds to a larger $r_{\text{virial}}$. Exact estimation of the relation between $M_0$ and $r_{\text{virial}}$ is not straightforward. This is because we need to know the radiation temperature $\theta_x$ as a function of $M_0$ which requires numerical calculations. Note that from Figures 1 and 3(b) the minimum value of $r_{\text{virial}}$ should be larger than $\sim 5 \times 10^3 r_s$.

Although no steady self-consistent solution exists due to the strong Compton heating at and beyond $r_{\text{virial}}$, an “oscillation” of the activity of the black hole is expected (e.g., Cowie et al. 1978; Ciotti & Ostriker 2007). When the accretion rate is high, only the gas inside of $r_{\text{virial}}$ can be accreted. This active phase will last for a timescale of the accretion timescale at $r_{\text{virial}}$. Then all the gas will be used up and the active phase stops. In this case, Compton heating also stops so the gas outside of $r_{\text{virial}}$ will be cooled by radiation and be accreted again, and the cycle repeats. The time the nonactive phase will last is determined by the radiative timescale of the gas at $r_{\text{virial}}$, since it is longer than the accretion timescale there. An alternative consequence of the strong Compton heating at large radii is that the accretion can be self-regulated by irradiating the outer flow (Shakura & Sunyaev 1973). That is, the strong Compton heating will not completely stop the accretion, but only decrease the accretion rate. This then reduces the energy release in the inner part, which in turn reduces the irradiation. A multidimensional numerical simulation is required to solve this issue and accurate time-dependence is needed as well since steady solutions may not be stable.

Figure 3. Compton effect for a supermassive black hole, with black hole mass $M = 10^8 M_\odot$ and mass accretion rate $\dot{M} = 0.5(\frac{r_{\text{out}}}{10^{14}})^{0.3} M_{\text{Edd}}$. The solid and dashed lines are for the solutions before and after the global Compton effect is taken into account. (a) Electron temperature profile. (b) Ratio of the rate of Compton heating/cooling and the turbulent heating of electrons. (c) Spectrum of the accretion flows. The two arrows show the values of the average energy and energy-weighted energy of photons emitted by the self-consistent solution.
For the massive black holes seen in the nuclei of most galaxies, the Compton heated interruption of the high-luminosity states should be typical if a hot accretion flow exists there. We now know that the accretion flow in low-luminosity AGNs is of this type (see Yuan 2007 and Ho 2008 for reviews). For luminous AGNs such as quasars, although people incline to think that it is a standard thin disk which is optically thick, many problems remain for this model (e.g., Shlosman et al. 1990; Koratkar & Blaes 1999). If the actual accretion flow is radially optically thin to Compton scattering, similar to the hot accretion flow, our analysis applies.

This kind of oscillation does not apply to stellar mass black holes in our Galaxy. This is because the prerequisite for such oscillation is that the accretion rate is large and the hot accretion flow extends to large radii. For a stellar mass black hole, the accretion material comes from the companion star and it starts out as a standard thin disk. In the hard state the standard disk does not extend to the innermost stable circular orbit but is replaced by a hot accretion flow within a transition radius \( r_T \). However, when the accretion rate is high, \( r_T \) is small (Yuan & Narayan 2004). So Compton scattering cools rather than heats the hot accretion flow. But on the other hand, the photons emitted by the hot accretion flow will heat the cool electrons in the standard disk at \( r_T \) and thus will change the dynamics of the transition. This effect has never been noted and could be a topic of future work.

4. SUMMARY AND DISCUSSION

For a geometrically thick and optically thin hot accretion flow, the photons can travel for a long distance without being absorbed, and thus be able to heat or cool electrons via Compton scattering. We investigate this global Compton scattering effect and find that for an accretion rate described by \( \dot{M} = M_0 (r/r_{out})^{0.5} \) the Compton cooling effect will be important when \( M_0 \gtrsim 0.1L_{Edd}/c^2 \); while the Compton heating effect will be important when \( M_0 \gtrsim 2L_{Edd}/c^2 \) and \( r_{out} = 10^7 r_s \). Specifically, the scattering heats electrons at \( r > 5 \times 10^3 r_s \) while cools electrons at \( r < 5 \times 10^3 r_s \). If \( r_{out} \) is larger, the critical \( M_0 \) above which the Compton heating effect is important will become lower.

We have successfully obtained the self-consistent steady solution with this effect included for \( M_0 \lesssim L_{Edd}/c^2 \) and \( r_{out} = 50 r_s \). But when \( M_0 \gtrsim L_{Edd}/c^2 \) and \( L \gtrsim 2L_{Edd} \), we fail because of the strong radiative cooling (local plus global Compton scattering). It is also difficult to get the self-consistent solution when \( M_0 \gtrsim L_{Edd}/c^2 \) (\( L > 1\%L_{Edd} \)) and \( r_{out} \gtrsim 10^5 r_s \). This is because, in this case, the Compton heating is so strong at and beyond \( r_{out} \) that the equilibrium electron temperature there will be higher than the virial temperature. More generally, we expect that the radius where the equilibrium temperature due to the Compton heating is equal to the virial temperature, \( r_{viral} \), is anticorrelated with \( M_{out} \). We argue that the black hole will manifest an oscillation of the activity in the case that we fail to get the steady solution. The period will be the radiative timescale of the gas at \( r_{viral} \).

All our discussions so far are for a one-dimensional (but not spherical) accretion flow. Although big uncertainties exist for the vertical structure of accretion flow, we are certain that when \( M \) is high, the scattering will be important, and consequently much of the luminosity will “leak out” perpendicular to the accretion flow as in the standard thin disk. This will have two effects. One is that the highest luminosity up to which we can get the self-consistent solution with the global Compton effect included will be higher. In addition, the Compton heating will be stronger in the vertical direction than in the equatorial plane of the flow. As a result, strong wind will be launched as pointed out by Park & Ostriker (2001, 2007) and found by Proga et al. (2008). All of the described effects are likely to become significant for AGN accretion flows having \( L > 10^{-2}L_{Edd} \) and optically thin (Section 3 for discussion of this possibility in luminous AGNs), which we know from recent applications of the Soltan argument (Yu & Tremain 2002), are the phases during which most massive black hole growth occurs.

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