OBSERVATIONAL EVIDENCE FOR ACTIVE GALACTIC NUCLEI FEEDBACK AT THE PARSEC SCALE

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ABSTRACT

In a hot accretion flow, the radiation from the innermost region of the flow propagates outward and heats the electrons at large radii via Compton scattering. It has been shown in previous works that if the radiation is strong enough, \( L \gtrsim 2\% L_{\text{Edd}} \), the electrons at the Bondi radius \( (r_B \sim 10^5 r_s) \) will be heated to above the virial temperature; thus, the accretion will be stopped. The accretion will recover after the gas cools down. This results in the oscillation of the black hole activity. In this paper, we show that this mechanism is the origin of the intermittent activity of some compact young radio sources. Such intermittency is required to explain the population of these sources. We calculate the timescales of the black hole oscillation and find that the durations of active and inactive phases are \( 3 \times 10^4 (0.1/\alpha)(M/10^8 M_\odot)(r_B/2 L_{\text{Edd}})^{-1/2} \) yr and \( 10^3 (\alpha/0.1)(M/10^8 M_\odot) \) yr, respectively, consistent with those required to explain observations. Such feedback occurring at the parsec scale should be common in low-luminosity active galactic nuclei and should be considered when we consider their matter and energy output.

Key words: accretion, accretion disks – black hole physics – galaxies: active

1. INTRODUCTION: INTERMITTENT ACTIVITY OF COMPACT RADIO SOURCES

Feedback from active galactic nuclei (AGNs) is now widely believed to play an important role in the formation and evolution of galaxies and to be responsible for many observational results including the \( M_{\text{BH}}–\sigma \) relation and the suppression of star formation in elliptical galaxies (e.g., Silk & Rees 1998; Fabian 1999; King 2003; Sazonov et al. 2005; Murray et al. 2005; Di Matteo et al. 2005; Croton et al. 2006; Ciotti & Ostriker 2001, 2007). While significant progress has been made, many details of the feedback are still uncertain (Ostriker et al. 2010). One approach for improving this situation is to look for more direct observational evidence for feedback and investigate them carefully.

Feedback often causes intermittent activities of AGNs; thus, it should be helpful to look for observational evidence for such intermittency (see a review by Czerny et al. 2009). Among this evidence are the central dominant galaxies in galaxy clusters observed by the Chandra X-ray Observatory and XMM-Newton. X-ray images show cavities and ripples that give evidence for repetitive outbursts. Another example, on which we will focus in this paper, is radio galaxies. Many case studies of radio structure directly show evidence for two or more active periods. Some sources are very compact, only a few kpc in size, which indicates that they are very young. These small radio sources are found to constitute 10–30% of all sources in a flux-limited sample. O’Dea & Baum (1997; see also Snellen et al. 2000) have studied the statistical properties of a combined sample of gigahertz-peaked spectrum sources (GPS) and compact steep spectrum radio sources (CSS). Their results show the existence of far too many compact (young) sources in comparison with the number of galaxies with extended old radio structures. If the total activity period lasts for \( 10^8 \) yr, the number of sources with ages below \( 10^5 \) yr should be roughly \( 10^3 \) times lower than the number of sources older than \( 10^5 \) yr, which is, however, not the case. The most likely explanation for the overabundance is that the activity is intermittent, i.e., the sources undergo an active or outburst phase lasting for \( 10^4 \) yr and that recurs every \( 10^5 \) yr (Reynolds & Begelman 1997; Kaiser et al. 2000; Kunert-Bajraszewska et al. 2005). Detection of several candidates for dying compact sources supports this view (Giroletti et al. 2007; Parma et al. 2007).

To explain the origin of such intermittent activity, Czerny et al. (2009) compiled a sample consisting of 72 GPS with measured age. The age is determined based on their measured hot-spot separation speed (kinematic age) or on the synchrotron cooling time. The majority of the age values are within the range of 200 yr–\( 10^7 \) yr. The monochromatic luminosity \( L_{\text{5GHz}} \) at 5 GHz is available for these sources, with typical values of \( 10^{42} \)–\( 10^{45} \) erg s\(^{-1} \) and an average value of \( 10^{43.6} \) erg s\(^{-1} \). This value can be used to estimate the bolometric luminosities of the accretion flows, \( L_{\text{acc}} \), in the following way (Czerny et al. 2009): For some sources, the 2–10 keV luminosity \( L_x \) is available. It is found that \( L_x \) is comparable to or a factor of \( ~5 \) higher than \( L_{\text{5GHz}} \) (Vink et al. 2006; Siemiginowska et al. 2008). To estimate \( L_{\text{acc}} \) from \( L_x \), we have to assume some template of the spectral energy distribution (SED) of the sources, either luminous AGNs or low-luminosity AGNs (LLAGNs), because their SEDs are characteristically different (Ho 1999). Assuming the SEDs of these GPS sources are similar to those of luminous AGNs, \( L_{\text{acc}} \) is at least 10 times higher than \( L_x \) (Elvis et al. 1994). Assuming a fiducial number of 50, we have \( L_{\text{acc}} \sim 50 L_x \sim 100 L_{\text{5GHz}} \sim 10^{44} \text{–} 10^{47} \) erg s\(^{-1} \) with an averaged value of \( 3 \times 10^{45} \) erg s\(^{-1} \). For several sources whose black hole masses are measured, \( L_{\text{acc}} \sim (0.03–1.5)L_{\text{Edd}} \) with a typical value of \( \sim 0.1 L_{\text{Edd}} \) (Czerny et al. 2009). However, as we argue below, these sources are more likely LLAGNs whose SEDs are distinctly different from luminous AGNs, characterized mainly by the lack of big blue bumps (Ho 1999, 2008). In this case, \( L_{\text{acc}} \sim 10 L_x \) (Ho 1999), so we have \( L_{\text{acc}} \sim (10–50)L_{\text{5GHz}} \sim (10^{43.5} \times 10^{46}) \) erg s\(^{-1} \). The typical Eddington ratio is then \( L_{\text{acc}} \sim 0.02 L_{\text{Edd}} \). Another statistical result from the sample in Czerny et al. (2009) is that the more luminous the sources the younger they are.

The only model proposed so far to explain the intermittent activity of these compact radio sources is the thermal instability

\(^1\text{The caveat is that the observed X-ray emission may not be dominated by the accretion flow (Stawarz et al. 2008).}\)
of a radiation-pressure-dominated standard thin disk (Czerny et al. 2009; see also Janiuk & Czerny 2011). In Section 2, we discuss several problems with this interpretation. We then propose in Section 3 an alternative model, a “global Compton scattering feedback” model based on the work of Yuan et al. (2009b). We assume that these radio sources are powered by hot accretion flows (advection-dominated accretion flow; ADAF) rather than by the standard thin disk. The Compton heating of the accretion flows (advection-dominated accretion flow; ADAF) will cause an oscillation of the black hole activity when \( L_{\text{acc}} \gtrsim 2\%L_{\text{Edd}} \). We calculate the durations of the active (or outburst) and inactive (or quiescent) phases and find that they are consistent with the observations. The last section (Section 4) is devoted to discussions.

2. CRITICISMS OF THE RADIATION PRESSURE INSTABILITY MODEL

Czerny et al. (2009) assume that the accretion flow in these radio sources is described by the standard thin disk. It is generally thought that this type of disk suffers from thermal instability when the accretion rate is above a certain threshold, which corresponds to \( L \sim 0.025L_{\text{Edd}} \) for supermassive black holes (it is \( \sim 0.2L_{\text{Edd}} \) for stellar-mass black holes; e.g., Svensson & Zdziarski 1994; Janiuk et al. 2002). This threshold is by coincidence close to that of the global Compton scattering instability. Assuming the SEDs of these radio sources are similar to those of luminous AGNs (refer to Section 1), they estimated their luminosities and found that all the sources in their sample are above this threshold, so they concluded that they are thermally unstable. As a result of the instability, their accretion rates periodically undergo high and low values. They further assumed this results in the high- and low-energy output of the jets. By adjusting parameters, this model can explain quite well the timescales of both the active and inactive phases.

However, there are several questions about this model. The first is whether jets can be formed from standard thin disks. To answer this question, let us look at the observations of black hole X-ray binaries (BHXBs). It is believed that the physics of BHXBs and AGNs is the same, while the observational data in BHXBs are much better than that in the case of AGNs in terms of formation of jets from different accretion modes. An individual BHXB usually has five states, namely quiescent, low, hard, intermediate, high, and soft, and very high (also called steep power-law states). They are characterized by different spectra and timing features (see a review by McClintock & Remillard 2006). Roughly speaking, the features of the soft state are similar to those of the luminous AGNs, while the hard state is similar to that of LLAGNs. Different states are believed to be caused by different accretion modes. Specifically, the accretion flow in the soft state is believed to be described by the standard thin disk and that in the hard state by the hot accretion flow such as an ADAF (Narayan & Yi 1994, 1995; for reviews on ADAFs see Narayan et al. 1998 and Narayan & McClintock 2008) or luminous hot accretion flow (Yuan 2001; Yuan & Zdziarski 2004). For reviews on the models of BHXBs and LLAGNs, refer to Zdziarski & Gierliński (2004), Done et al. (2007), Narayan (2005), and Yuan (2007). Returning to our question of jet formation, radio observations of BHXBs clearly show that jets exist in the hard state, evidenced by the flat radio spectrum and the elongated radio image, while both features disappear in the soft state (Fender 2006). This strongly indicates that jets cannot be formed in the standard thin disk, but can be in the hot accretion flows. Theoretical studies are consistent with this observational result (e.g., Livio et al. 1999; Meier 2001).

Observations find that the highest luminosity that hard states can reach is \( \sim (10–30)\%L_{\text{Edd}} \). This is roughly consistent with predictions of hot accretion flows (Yuan & Zdziarski 2004; Yuan et al. 2007). Above this luminosity, the cooling in the hot accretion flow becomes so strong that the accretion flow cannot remain hot, and it collapses. On the other hand, strong radio emission has also been detected when BHXBs are more luminous, i.e., in the very high state, GRS 1915+105 and radio-loud quasars may correspond to this “radio-loud very high state.” This, however, does not imply that jets can also be formed in standard thin disks, because the very high state cannot be described by a standard thin disk. Actually, the accretion disk model of the very high state is still an unsolved problem. More importantly, detailed observations indicate that the strong radio emission comes from radio-emitting blobs ejected during the transition from a “hard very high state” to a “soft very high state,” which may be modeled by hot accretion flow-like and thin disk-like models, respectively (McClintock & Remillard 2006). Such ejections are called “episodic jets” to discriminate their many different features as compared to the “continuous jets” (see reviews in Fender & Belloni 2004; Fender et al. 2004). Yuan et al. (2009a) propose a magnetohydrodynamic (MHD) model for the formation of episodic jets, by analogy with the coronal mass ejection (CME) phenomenon in the Sun. Unlike the “continuous jets” that are formed in the presence of large-scale open magnetic fields, episodic jets are formed in the region of closed magnetic fields in the disk corona. According to this model, episodic jets can be formed in hot accretion flows and during their collapse into a cold thin disk. Both direct observations and the theoretical model indicate that episodic jets are not formed in standard thin disks.

A question is then whether the compact radio sources correspond to the hard state or the very high state. The model proposed in Reynolds & Begelman (1997) requires that the duration of the jet launching in the active phase be \( \sim 10^4 \) yr. If the compact radio sources correspond to the very high state, this duration is the timescale of the state transition. Unfortunately, the details of the state transition are still an open question, although it is believed to be associated with the collapse of the hot accretion flow to a cold disk. The observed timescale of the state transition ranges from a few days to \( \sim 100 \) days (Yu & Yan 2009; Zdziarski & Gierliński 2004). Scaling this timescale with black hole mass by a factor of \( 10^7 \) gives \( 10^5–10^6 \) yr, which is one or two orders of magnitude longer than that required by Reynolds & Begelman (1997). Another point against the analogy between the compact radio sources and the very high state is the luminosity. As we state in Section 1, the typical luminosity of these radio sources is \( \sim 0.02L_{\text{Edd}} \). On the other hand, the typical transition luminosity from hard to soft is significantly higher, within the range of \( \sim (2–20)\%L_{\text{Edd}} \). It therefore seems more likely that these compact radio sources correspond to the hard state rather than to the very high state. We would like to point out, however, that it is hard at present to completely rule out the latter possibility.

The second question is of the thermal instability of the radiation-dominated thin disk. As we mentioned at the beginning of this section, when the accretion rate is higher than a threshold at which the pressure in the accretion flow is dominated by the radiation pressure, the standard thin disk was thought to be thermally unstable (Shakura & Sunyaev 1976;
Piran 1978). Time-dependent global hydrodynamical calculations (e.g., Honma et al. 1992; Janiuk et al. 2002; Li et al. 2007) found that the local thermal instability will result in “limit-cycle” behavior. However, Gierliński & Done (2004) show that observations of the soft state of BHXBs, which has widely been believed to be described by the standard thin disk model, challenge the above prediction. The luminosity of the high state sources compiled in their sample spans the range 0.01 \( \lesssim L/L_{\text{Edd}} \lesssim 0.5 \), with the highest well exceeding the theoretical stable limit. Thus, we should expect to see limit-cycle behavior. In contrast to this expectation, however, observations show little variability, which convincingly indicates that they are thermally stable. Recently, the non-existence of the thermal instability of the radiation-dominated thin disk has also been directly shown by three-dimensional MHD numerical simulation of the shearing box (Hirose et al. 2009).

Some theoretical efforts have been made to solve this puzzle. One explanation has been proposed by Hirose et al. (2009). They argue that although the stress is proportional to the total pressure,\(^2\) the stress fluctuation precedes pressure fluctuations, contrary to the usual supposition that pressure controls the magnetic stress. This explains the thermal stability. An alternative explanation has been proposed recently by Zheng et al. (2011). They argue that the previous analytical analysis of thermal stability neglects the role of magnetic pressure, which usually accounts for \( \sim 10\% \) of the total pressure, and thus is dynamically important. They show that if the magnetic pressure decreases with response to the increase in temperature of the accretion flow, the threshold of the accretion rate above which the disk becomes unstable increases significantly compared to the case in which the magnetic pressure is not considered. The physical reason is that in this case the dependence of turbulent dissipation heating on temperature becomes weaker.

3. COMPTON SCATTERING FEEDBACK MODEL

We propose that the intermittent activity of compact radio sources is caused by the “global Compton scattering” feedback effect in hot accretion flows. The idea of global Compton scattering was first proposed by Ostriker et al. (1976) and Cowie et al. (1978) in the context of spherical accretion and was later investigated by Esin (1997), Park & Ostriker (2001, 2007), and Yuan et al. (2009b, hereafter YXO09) for rotating accretion flows. Here we briefly review the main results of YXO09. Most of the radiation comes from the innermost region of the accretion flow, \( \sim 10r_s \). The radiation will propagate outward and heat the electrons at large radii via Compton scattering. If the luminosity is high enough, the electrons will be heated above their virial temperature, thus we cannot obtain a steady solution. In this case, YXO09 argue that the black hole activity will oscillate between an active and an inactive phase. The quantitative results are as follows: Given the existence of outflow in hot accretion flows, the mass accretion rate is a function of radius, \( M(r) = \dot{M}_{\text{out}} r / r_{\text{out}} \), with \( \dot{M}_{\text{out}} \) being the accretion rate at the outer boundary \( r_{\text{out}} \). In YXO09 we adopted \( s = 0.3 \), following the modeling result of the supermassive black hole in Sgr A* (Yuan et al. 2003). We set \( r_{\text{out}} = 10^3 r_s \), since this is the Bondi radius for typical parameters of black hole mass \( M = 10^8 M_\odot \) and interstellar medium (ISM) temperature \( T = 10^7 \) K, as in the present work. We find that when \( \dot{M}_{\text{out}} \gtrsim \dot{M}_{\text{Edd}} \), or equivalently \( L \gtrsim 2\% L_{\text{Edd}} \), the electron temperature at \( 10^3 r_s \) will be heated to higher than the virial value, thus no steady accretion solution exists. This radius is called the virial radius. Its value is found to be roughly anti-correlated with \( r_{\text{vir}} \), \( \sim 10^3 r_s \), which is higher than the virial value. Thus, this phase finishes when the gas inside \( 10^3 r_s \) is depleted; the source then enters into the inactive phase. At the inactive phase, the gas is very hot, with \( T \sim 10^9 \) K; thus, the mass accretion rate and the luminosity are very low. This phase finishes when the hot gas cools down.

\(^3\) In this regard, we note that Czerny et al. (2009) assume that the stress is proportional to the geometrical mean between the gas and the total pressure.

Figure 1. Schematic figure of the beginnings of the active (a) and inactive (b) phases. At the active phase, the accreting gas inside \( 10^3 r_s \) is virial; the gas outside this radius is heated by the global Compton scattering toward \( T \sim 10^9 \) K, which is higher than the virial value. Thus, this phase finishes when the gas inside \( 10^3 r_s \) is depleted; the source then enters into the inactive phase. At the inactive phase, the gas is very hot, with \( T \sim 10^9 \) K; thus, the mass accretion rate and the luminosity are very low. This phase finishes when the hot gas cools down.

We assume the mass of the black hole \( M = 10^8 M_\odot \). The density and temperature of the ISM at Bondi radius are denoted as \( n_{\text{ISM}} \) and \( T_{\text{ISM}} \sim 10^7 \) K. This high temperature of \( T_{\text{ISM}} \) is supported by the direct observations of some galactic centers, such as Sgr A* (Baganoff et al. 2003), some nearby galaxies (Pellegrini 2005), and the radio galaxy M 87 (e.g., Di Matteo
et al. 2003). Theoretically, this could be because of the shock heating due to the collision between stellar winds (Quataert 2004). The outer boundary of the accretion flow is set at the Bondi radius, \( r_{OB} \sim r_B \sim GM/c^2 \sim 10^5 r_s \).

As we state above, when \( L \sim 2\% L_{Edd} \) (or equivalently \( M_{out} \sim M_{Edd} \)), the gas at \( r_{vir} \sim r_B \) will be heated to above the virial temperature and will thus diffuse. The gas within \( r_{vir} \) can still be accreted, which powers the high-luminosity active phase. The duration of this phase is determined by the accretion timescale at \( r_{vir} \).

\[
t_{\text{act}} \sim t_{\text{acc}}(r_{vir}) \sim \frac{r_{vir}}{v_J} \sim 3 \times 10^4 M_8 \left( \frac{\alpha}{0.1} \right)^{-1} \left( \frac{M_{\text{out}}}{M_{Edd}} \right)^{-1} \text{yr}
\]

(Narayan et al. 1998). For typical values of \( M_{\text{out}} \sim M_{Edd} \), as indicated by the observed \( L \sim 2\% L_{Edd} \), this timescale is in good agreement with the required duration of the active phase in Reynolds & Begelman (1997). Moreover, the sample compiled in Czerny et al. (2009) indicates that sources with higher luminosity have younger ages. This is easy to understand from Equation (1), because for an ADAF we roughly have \( L \sim M_{\text{out}}^2 \), and thus \( t_{\text{act}} \sim 3 \times 10^4 M_8 (0.1/\alpha) (L/0.02 L_{Edd})^{-1/2} \text{yr} \).

The temperature profile within \( r_{vir} \sim 10^5 r_s \) has been calculated in YX009. Now we estimate the temperature of the gas beyond \( 10^5 r_s \), which is required to calculate the duration of the inactive phase. The temperature is mainly determined by the Compton heating, viscous dissipation, and bremsstrahlung cooling (for simplicity, we do not consider other possible heating mechanisms). Cooling is usually negligible in an ADAF. In the non-relativistic limit, the Compton heating reads (YX009):

\[
q_{\text{comp}} = n_e (\theta_e - \theta_T) F \sigma_T \sim n_e \theta_e \sigma_T \propto r^{-3.2}. \tag{2}
\]

Here, \( \theta_e \equiv k T_e / m_e c^2 \), \( T_e \) is the radiation temperature characterizing the spectrum of radiation from the ADAF, \( F \) is the radiation flux, and \( \sigma_T \) is the Thomson cross section. The value of \( T_e \) can be well approximated to be (YX009):

\[
T_e \sim 10^9 \text{ K}. \tag{3}
\]

Note that this value is much larger than the typical value for the spectrum of luminous quasars, where \( T_e \sim 10^7 \text{ K} \) (Sazonov et al. 2005). This is of course because of the significant difference between the SEDs of luminous AGNs and LLAGNs (Ho 1999, 2008; Yuan et al. 2009b). The viscous heating rate is given by

\[
q_{\text{vis}} \propto n_e \theta_e^{1/2} H r^2 (d\Omega/dv)^2 \propto r^{-3.7}. \tag{4}
\]

In both Equations (2) and (4), the scaling of an ADAF (with outflow) \( n_e \propto r^{-3/2} \), \( v_J = r^{-1/2} \), and \( \theta_e \equiv k T_e / m_e c^2 \propto r^{-1} \) are adopted. At \( r \sim 10^5 r_s \), we have \( q_{\text{comp}} \gtrsim q_{\text{vis}} \), since this causes the electrons to be heated above the virial temperature. Equations (2) and (4) then imply that the heating is dominated by Compton scattering at \( r \gtrsim 10^5 r_s \). The evolution of electron temperature is then described by

\[
n_e k \frac{dT_e}{dt} = n_e (\theta_e - \theta_T) F \sigma_T. \tag{5}
\]

The solution is

\[
T_e = T_{e0} e^{-\alpha t} + T_s (1 - e^{-\alpha t}). \tag{6}
\]

Here, \( T_{e0} \) is the electron temperature at the beginning of the active phase and \( A = \frac{L_{e0}}{\dot{M} \sigma_T m_e c^2} \). The heating timescale is the duration of the active phase \( t_{\text{act}} \), so the spatial range in which electrons can be heated to \( T_s \) is determined by

\[
A t_{\text{act}} \approx 1. \tag{7}
\]

This gives

\[
r \sim \left( \frac{L \sigma_T t_{\text{act}}}{4 \pi m_e c^2} \right)^{1/2} \sim 10^6 r_s. \tag{8}
\]

In the last calculation the typical values of \( L \sim 2\% L_{Edd} \) and \( t_{\text{act}} \) from Equation (1) are adopted. Therefore, we obtain that the gas up to \( \sim 10^6 r_s \) can be heated to Compton temperature \( T_s \) after the active phase. Note that we neglect the bremsstrahlung cooling in our calculation, so this value should be regarded as an upper limit.

In correspondence with the increase of the temperature, the density of the gas should become much lower than that of the unheated ISM. The accretion of this low-density gas results in the significant decrease of the mass accretion rate and luminosity. The source then enters the inactive phase. We denote the density as \( n_{\text{inact}} \) and estimate its value by the pressure balance of this gas with the surrounding ISM, whose density and temperature are denoted as \( n_{\text{ISM}} \) and \( T_{\text{ISM}}(\sim 10^7 \text{ K}) \):

\[
n_{\text{inact}} = n_{\text{ISM}} \frac{T_{\text{ISM}}}{T_s} = 10^{-2} n_{\text{ISM}}. \tag{9}
\]

We estimate the mass accretion rate with the Bondi accretion theory, so

\[
r_{\text{Bondi}} \sim \frac{GM}{c^2}, \tag{10}
\]

and

\[
M_{\text{inact}} \sim \alpha M_{\text{Bondi}} \approx 4\pi r_{\text{Bondi}}^2 n_{\text{inact}} m_p c_s (r_{\text{Bondi}}) \propto T^{-3/2} n_{\text{inact}}. \tag{11}
\]

Here, \( \alpha \) is the viscous parameter and \( c_s (r_{\text{Bondi}}) \) is the sound speed at \( r_{\text{Bondi}} \). The accretion rate is up to \( \sim 10^5 \) times lower in the inactive phase compared to the active phase, which is why the source becomes inactive.

The Compton-heated gas from \( r_{\text{Bondi}} \) to \( 10^6 r_s \) will cool with time. The duration of the inactive phase should be determined by the depletion of the heated gas due to accretion, or the cooling of this heated gas, whichever is shorter. The former is hard to estimate since \( 10^6 r_s \) is well out of the influence sphere of the black hole. It must, however, be much longer than the accretion timescale determined by the ADAF theory, i.e.,

\[
t_{\text{acc}} \gg \frac{10^5 r_s}{v_J} \sim 10^5 M_8 \left( \frac{\alpha}{0.1} \right)^{-1} \text{yr}. \tag{12}
\]

The cooling timescale reads

\[
t_{\text{cool}} = \frac{n_{\text{inact}} T_s}{\dot{j}_{\text{brem}}} \approx 6 \times 10^3 T_s^{1/2} n_{\text{inact}} \text{yr}. \tag{13}
\]

Here, \( \dot{j}_{\text{brem}} \) is the bremsstrahlung emissivity. We have \( T_s \sim 10^8 \text{ K} \) and \( n_{\text{inact}} \sim 10^{-2} n_{\text{ISM}} \), but the value of \( \dot{j}_{\text{brem}} \) remains unknown. From our statement above, we know that the condition for the oscillation of the black hole activity is that the luminosity of the active phase is \( L \gtrsim 2\% L_{Edd} \), or equivalently \( M(10^5 r_s) \gtrsim M_{Edd} \). Given that \( M_{\text{out}}(10^5 r_s) \sim \alpha M_{\text{Bondi}} \sim 4\pi r_{\text{Bondi}}^2 n_{\text{inact}} m_p \), we obtain \( n_{\text{ISM}} \gtrsim 10^9 (\alpha / 0.1)^{-1} M_8^{-1} \text{ cm}^{-3} \). Therefore, we have

\[
t_{\text{cool}} = 10^5 M_8 \left( \frac{\alpha}{0.1} \right) \left( \frac{T_s}{10^9 \text{ K}} \right)^{1/2} \text{yr}. \tag{13}
\]
This determines the duration of the inactive phase because it is much shorter than the accretion timescale at $10^5 r_s$ (Equation (12)). This value is in good agreement with the duration of the inactive phase required in Reynolds & Begelman (1997).

4. DISCUSSION: COMPARISONS WITH OTHER WORKS

In this paper, we propose an AGN feedback mechanism occurring at $r_{ya} \sim 10^5 r_s$ from the central black hole, i.e., around the parsec scale. The conditions for this effect to be present are that the accretion flow is described by hot accretion flow (such as an ADAF) and the accretion luminosity $L \gtrsim 2\times10^8L_{Edd}$. The influence of this effect is that the black hole will oscillate between an active and an inactive phase, which last for $3 \times 10^2 (0.1/\alpha)(M/10^8 M_\odot)(L/2\times10^8L_{Edd})^{1/2}$ yr and $10^3 (0.1/0.1)(M/10^8 M_\odot)$ yr, respectively. Such short-timescale oscillation should be common in radio sources with $L \gtrsim 2\times10^8L_{Edd}$. This result implies that even though the fueling rate of AGNs is constant at large scales, the radiation (and also matter) output from the AGNs oscillates. This behavior should be taken into account when we consider the effects of AGN feedback at larger scales, especially those treating the small-scale accretion as the “sub-grid” input of their simulations.

Now let us compare our study with other works. One is the series of works by Ciotti & Ostriker (2001, 2007). In their works, they focus on the radiative heating by the central AGN to the ISM, but on scales much larger than ours, from several pc to kpc. More physics, such as supernova heating, is included there. Similar to our results, they also find that when the luminosity of the central AGNs is high, in their case close to $L_{Edd}$, oscillation occurs. However, the timescale of oscillation is much longer, $\sim 10^3-10^4$ yr. One important reason is that the radiation temperature adopted in their model is much smaller, $\sim 10^6$ K. Physically, such a low $T_x$ corresponds to the typical spectrum of a quasar (Sazonov et al. 2005). When the accretion rate is lower than a certain value, the accretion flow should make a transition from the standard thin disk to an ADAF, as argued in Narayan & Yi (1995) and evidenced by the transition from a soft to a hard state in BHXBs. The spectrum emitted by ADAFs is significantly different from that of quasars, most significantly characterized by the lack of big blue bumps, as we point out in Section 1 (Ho 1999, 2008). The corresponding radiation temperature of such spectra is typically $T_x \sim 10^7$ K (YXO09).

Another possibly important point is the inner boundary of the simulation domain. From YXO09, the global Compton scattering effect is already important at $\sim 10^2 r_s (M/M_{Edd})^{-1}$, i.e., at sub-parsec or parsec scales. In Ciotti & Ostriker’s works, they typically set the inner boundary at several parsec. Thus, some important feedback effect may be missed. Therefore, it will be interesting to repeat the simulations of Ciotti & Ostriker (2001, 2007), this time taking into account the change of $T_x$ at low $M$ and the location of the inner boundary. This work is in progress.

Another series of works are by Proga and his collaborators (e.g., Proga et al. 2000; Proga 2007; Kurosawa & Proga 2009). Similar to our work, they focus on the sub-parsec and parsec scales. Unlike our work, but similar to Ciotti & Ostriker (2001, 2007), they assume that the central radiation spectrum is quasar-like. Another difference from our work is that they focus on cold accretion flows and higher mass accretion rates. However, since oscillation behavior is also found in their works, if the accretion rate decreases to such a low value that the accretion flow makes a transition from a thin disk to a hot accretion flow, $T_x \sim 10^6$ K should again be adopted.

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