

ARE THE JETS ACCELERATED FROM THE DISK CORONAS IN SOME ACTIVE GALACTIC NUCLEI?

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ABSTRACT

We use a sample of radio-loud active galactic nuclei (AGNs) with estimated central black hole masses to explore AGN jet-formation mechanisms. The jet power of AGNs is estimated from their extended radio luminosity. We find that the jets in several AGNs of this sample are too powerful to be extracted from the standard thin accretion disks or rapidly spinning black holes surrounded by standard thin disks. If advection-dominated accretion flows (ADAFs) are present in these AGNs, their bright optical continuum luminosity cannot be produced by pure ADAFs because of their low accretion rates and low radiation efficiency, unless the ADAFs transit to standard thin disks at some radii R_{tr} . If this is the case, we find that dimensionless accretion rates $\dot{m} = \dot{M}/\dot{M}_{\text{Edd}}$ as high as ≥ 0.05 , and a transition from ADAFs to standard thin disks at rather small radii, $\sim 20GM_{\text{bh}}/c^2$, are required to explain their bright optical continuum emission. We propose that the disk-corona structure is present in at least some AGNs in this sample. The plasmas in the corona are very hot, and the pressure scale height of the corona $H_c \sim R$. Powerful jets with $Q_{\text{jet}} \sim L_{\text{bol}}$ (bolometric luminosity) can form via the large-scale magnetic fields created by dynamo processes in the disk corona of some AGNs. The maximal jet power extractable from the corona, $Q_{\text{jet}}^{\text{max}} \leq 0.6L_c$ (L_c is the corona luminosity), is expected in this jet-formation scenario. The statistical results for the sample of AGNs are consistent with the predictions of this scenario. Finally, the possibility that the jet is driven from a super-Keplerian rotating hot layer located between the corona and the cold disk is discussed. We find that, in principle, this layer could also produce a powerful jet with $Q_{\text{jet}} \sim L_{\text{bol}}$.

Subject headings: accretion, accretion disks — black hole physics — galaxies: active — galaxies: jets

1. INTRODUCTION

Relativistic jets have been observed in many radio-loud active galactic nuclei (AGNs) and are believed to form very close to black holes. In the currently most favored models of the formation of the jet, namely, the Blandford-Payne and Blandford-Znajek mechanisms (Blandford & Payne 1982; Blandford & Znajek 1977), the power is generated through accretion and then extracted from the disk/black hole rotational energy and converted into the kinetic power of the jet. Both of these jet-formation mechanisms predict a link between the accretion disk and the jet. The disk-jet connection has been investigated by many authors using observational data in different ways (e.g., Rawlings & Saunders 1991; Falcke & Biermann 1995; Xu et al. 1999; Cao & Jiang 1999, 2001) and has been found to indicate an intrinsic link between accretion disks and jets. However, this cannot rule out the Blandford-Znajek jet-formation mechanism. Magnetic fields are maintained by the currents in the accretion disk surrounding a rapidly spinning black hole, so the power extracted from a rapidly spinning black hole by the Blandford-Znajek mechanism depends as well on the properties of the disk near the black hole.

It is still unclear which mechanism is responsible for jet formation in AGNs. The relative importance of these two jet-formation mechanisms is explored by different authors (e.g., Ghosh & Abramowicz 1997; Livio et al. 1999, hereafter L99; Cao 2002b). Ghosh & Abramowicz (1997) doubted the importance of the Blandford-Znajek process. For a black hole of given mass and angular momentum, the strength of the Blandford-Znajek process depends crucially on the strength of the poloidal field threading the horizon of the hole. The magnetic field threading a hole should be maintained by the

currents situated in the inner region of the surrounding accretion disk. Ghosh & Abramowicz (1997) argued that the strength of the field threading a black hole had been overestimated. L99 reinvestigated the problem and pointed out that even the calculations of Ghosh & Abramowicz (1997) overestimated the power of the Blandford-Znajek process, since they had overestimated the strength of the large-scale field threading the inner region of an accretion disk, and thus the efficiency of the Blandford-Znajek process. The length scale of the fields created by dynamo processes is of the order of the disk thickness, $\sim H$. The large-scale field can be produced from the small-scale field created by dynamo processes as $B(\lambda) \propto \lambda^{-1}$ for the idealized case, where λ is the length scale of the field (Tout & Pringle 1996; Romanova et al. 1998). Thus, the large-scale field is very weak if the field is created in the thin accretion disks. L99 estimated the maximal jet power extracted from an accretion disk on the assumption that the toroidal field component is of the same order as the poloidal field component at the disk surface. They argued that the maximal jet power extracted from an accretion disk (the Blandford-Payne mechanism) dominates over the maximal power extracted by the Blandford-Znajek process (L99). For the advection-dominated accretion flow (ADAF) cases, the disk thickness $H \sim R$ and the jets can be driven by the large-scale magnetic fields created by dynamo processes (e.g., Armitage & Natarajan 1999). The maximal jet power extracted from rotating black holes and ADAFs was calculated by Armitage & Natarajan (1999) and Meier (2001). Instead of the ADAF model, Merloni & Fabian (2002) proposed that the jets in low-luminosity AGNs may be magnetically (or thermally) driven from the coronas above the geometrically and optically thin disks accreting at low rates. The gas in the corona is almost virialized, and the thickness of the corona is $\sim R$. The

large-scale magnetic fields created by dynamo processes in the corona are significantly stronger than those created in the thin disk, because the corona is much thicker than the cold thin disk, and the disk corona can therefore power a stronger jet than the thin disk. In principle, the maximal jet power extracted for different jet-formation mechanisms can be calculated if the central black hole mass and accretion rate are known.

There are different approaches proposed to measure the black hole masses in AGNs. The tight correlation between central black hole mass M_{bh} and stellar dispersion velocity σ of the host galaxy is used to estimate the black hole mass for the sources with measured stellar dispersion velocity (e.g., Ferrarese & Merritt 2000; Gebhardt et al. 2000). The central black hole mass can also be estimated from its host galaxy luminosity by using the correlation between black hole mass M_{bh} and host galaxy luminosity M_R at the R band (e.g., McLure & Dunlop 2002). However, only a small fraction of AGNs have measured stellar dispersion velocity or host galaxy luminosity, which prevents us from using these approaches to estimate the masses of black holes in most AGNs. Kaspi et al. (2000) measured the sizes of broad-line regions (BLRs) in several tens of Seyfert 1 galaxies and quasars from the time delay between their variabilities in optical continuum emission and broad-line emission using the reverberation mapping method (Peterson 1993). The central black hole masses of these AGNs have been estimated from their broad-line widths on the assumption that the motion of the clouds in BLRs is virialized (Kaspi et al. 2000). Kaspi et al. (2000) found a correlation between the measured BLR size and optical continuum luminosity for their radio-quiet AGN sample. As only a small fraction of AGNs have measured BLR sizes, this empirical relation between R_{BLR} and optical continuum luminosity λL_λ is employed to estimate the masses of black holes in AGNs from their broad-line widths (e.g., Laor 2000; McLure & Dunlop 2001; Gu et al. 2001). McLure & Dunlop (2001) compared the virialized $H\beta$ black hole mass estimates with those estimated from the $M_{\text{bh}}-\sigma$ relation for a sample of radio-loud AGNs. They found that the virialized $H\beta$ black hole masses are systematically lower than those derived from the stellar velocity dispersion, and they concluded that the disklike BLR geometry may be in these sources. This is generally consistent with the anti-correlation between broad $H\beta$ line width and R_c (the ratio of the strengths of the radio core to the lobe) for a flat-spectrum radio-loud sample found by Wills & Browne (1986). The inclination angles of the jets in flat-spectrum AGNs are believed to be small with respect to the line of sight, and their optical continuum emission is probably beamed (see Urry & Padovani 1995, and references therein). The black hole mass estimate for flat-spectrum radio-loud AGNs using the broad-line widths and optical continuum luminosities may be lower than their real values, which implies that black hole mass estimated from the broad-line width may be affected by the disklike BLR orientation (McLure & Dunlop 2001).

In this paper, we use a sample of AGNs to explore jet formation. The black hole masses of the sources in this sample are estimated from their optical continuum luminosities and the widths of broad-line $H\beta$. The sample we use in this investigation is radio-selected and includes some sources with very strong jets. In this paper, we focus on how the powerful jets in these sources are formed in the regions near black holes.

2. BLACK HOLE MASSES AND ACCRETION RATES

Kaspi et al. (2000) derived an empirical relation between the BLR size and optical continuum luminosity,

$$R_{\text{BLR}} = (32.9^{+2.0}_{-1.9}) \left[\frac{\lambda L_\lambda(5100 \text{ \AA})}{10^{44} \text{ ergs s}^{-1}} \right]^{0.700 \pm 0.033} \text{ lt-days}, \quad (1)$$

for a sample of Seyfert 1 galaxies and quasars in which the sizes of BLRs are measured with the reverberation mapping method (Peterson 1993). The central black hole masses M_{bh} can be estimated from the velocities v_{BLR} of the clouds in the BLRs,

$$M_{\text{bh}} = \frac{v_{\text{BLR}}^2 R_{\text{BLR}}}{G}, \quad (2)$$

where the motions of the clouds are assumed to be virialized and isotropic. The velocities of the clouds in BLRs, v_{BLR} , are derived from the width of the broad emission lines. For most AGNs, the BLR sizes have not been measured by the reverberation mapping method; the empirical relation (see eq. [1]) is instead used to estimate the BLR sizes. The central black hole masses of AGNs have been estimated from their broad-line widths and optical continuum luminosity (e.g., Laor 2000; McLure & Dunlop 2001; Gu et al. 2001). In this paper, we use the sample of Gu et al. (2001). The sources in their sample are selected from a parent sample consisting of the 1 Jy and Strong Source Surveys S4 and S5 radio catalogs. As this is a radio-selected sample, it includes mostly strong radio sources. We assume that the physics of BLRs in radio-loud quasars is similar to that of radio-quiet quasars in the sample of Kaspi et al. (2000), i.e., the tight correlation between the BLR size and optical ionizing luminosity holds for radio-loud quasars. The central black hole masses of 86 radio-loud AGNs are estimated by using their broad-line width data and optical continuum luminosities. These 86 AGNs consist of 55 flat-spectrum sources and 31 steep-spectrum sources. We adopt the sample of Gu et al. (2001) for our present investigation. All the profile data of broad emission line $H\beta$ and derived black hole masses of these sources are listed in Gu et al. (2001).

For normal bright AGNs, the bolometric luminosity can be estimated from their optical luminosity $L_{\lambda, \text{opt}}$ at 5100 Å by (Kaspi et al. 2000)

$$L_{\text{bol}} \simeq 9\lambda L_{\lambda, \text{opt}}, \quad (3)$$

which is only a rough estimate. The coefficient in equation (3) may vary for AGNs with different spectral energy distributions.

As the central black hole masses are available for the sources in this sample, we can use equation (3) to derive the dimensionless accretion rates $\dot{m}$ of these AGNs. The derived dimensionless accretion rate is given by

$$\dot{m} = \dot{M}/\dot{M}_{\text{Edd}} \simeq L_{\text{bol}}/L_{\text{Edd}}. \quad (4)$$

This is valid for standard thin accretion disks because the radiative efficiency of thin disks is a constant for given black hole spin parameter a . Combining relations (1) and (2), we can obtain

$$M_{\text{bh}} \propto v_{\text{BLR}}^2 L_{\lambda, \text{opt}}^{0.7}. \quad (5)$$

Thus, using relations (3), (4), and (5), we get

$$\dot{m} \propto v_{\text{BLR}}^{-2} L_{\lambda, \text{opt}}^{0.3}, \quad (6)$$

since $L_{\text{Edd}} \propto M_{\text{bh}}$.

The estimated black hole masses should be taken with caution, especially for flat-spectrum AGNs, because we have not considered the orientational effects for either the broad-line width or the beamed optical continuum emission. The fact that the virialized $H\beta$ black hole mass estimates are systematically lower than those estimated from the $M_{\text{bh}}-\sigma$ relation implies that the BLR geometry may be disklike. The black hole masses estimated for flat-spectrum sources using $H\beta$ line widths may be underestimated (McLure & Dunlop 2001). The errors in the present estimate of the dimensionless accretion rate may be caused by the uncertainties in the black hole mass estimate due to orientational effects, the beamed optical continuum emission, and the coefficient in relation (3). How these uncertainties in the estimates of the black hole mass M_{bh} , bolometric luminosity L_{bol} , and the dimensionless accretion rate $\dot{m}$ may affect our conclusions will be discussed in § 6.

3. JET POWER

The jet power can be estimated from low-frequency radio luminosity by

$$Q_{\text{jet}} \simeq 3 \times 10^{38} f^{3/2} L_{\text{ext}, 151}^{6/7} \text{ W}, \quad (7)$$

where $L_{\text{ext}, 151}$ is the extended radio luminosity at 151 MHz in units of $10^{28} \text{ W Hz}^{-1} \text{ sr}^{-1}$ (Willott et al. 1999). Willott et al. (1999) argue that the normalization is very uncertain and introduce the factor f to account for these uncertainties. They use a wide variety of arguments to suggest that $1 \leq f \leq 20$. In this paper, we conservatively adopt the lower limit $f = 1$. For some flat-spectrum sources, the radio/optical continuum emission is strongly beamed toward us because of the relativistic motions of the sources and small viewing angles of the jets with respect to the line of sight. We find that the spectra of some flat-spectrum quasars are flat even at a wavelength around 151 MHz, which implies that the radio emission from these sources at 151 MHz is still dominated by the core emission rather than the extended emission. Thus, the observed low-frequency radio emission at 151 MHz may still be Doppler beamed. We therefore use the extended radio emission measured by the Very Large Array (VLA)¹ to estimate the jet power (see Cao 2003). The VLA observations are usually performed at a frequency much higher than 151 MHz. The extended radio emission measured by the VLA has to be K -corrected to 151 MHz in the rest frame of the source assuming $\alpha_e = 0.8$ ($f_\nu \propto \nu^{-\alpha_e}$) (Cassaro et al. 1999). Apart from steep-spectrum AGNs, we find 41 flat-spectrum AGNs with extended emission data in this sample.

4. JET-FORMATION MECHANISMS

L99 estimated the maximal strength of the large-scale fields driving the jets and the maximal jet power extracted from the rapidly spinning black hole or accretion disk on the assumption that the fields are amplified by dynamo processes. Following L99, we find that the maximal jet power extracted from

a rotating black hole or an accretion disk can be calculated for a standard thin disk if the black hole mass M_{bh} and accretion rate $\dot{m}$ are specified.

4.1. Jet Power Extracted from Standard Thin Disks

The maximal power of the jet accelerated by a magnetized accretion disk is

$$L_{\text{BP}}^{\text{max}} = 4\pi \int \frac{B_{\text{pd}}^2}{4\pi} R^2 \Omega(R) dR, \quad (8)$$

where $B_{\text{pd}} \sim B_\varphi$ is assumed, and B_{pd} is the strength of the large-scale ordered poloidal field at the disk surface.

The strength of the field at the disk surface is usually assumed to scale with the pressure of the disk, as in Ghosh & Abramowicz (1997). However, L99 pointed out that the large-scale field can be produced from the small-scale field created by dynamo processes as $B(\lambda) \propto \lambda^{-1}$ for the idealized case, where λ is the length scale of the field (Tout & Pringle 1996; Romanova et al. 1998). Thus, the large-scale field threading the disk that is related to the field produced by dynamo processes can be approximated by (L99)

$$B_{\text{pd}} \sim \frac{H}{R} B_{\text{dynamo}}. \quad (9)$$

The dimensionless pressure scale height of the disk H/R is given by (Laor & Netzer 1989)

$$\frac{H}{R} = 15.0 \dot{m} r^{-1} c_2, \quad (10)$$

where the coefficient c_2 is defined in Novikov & Thorne (1973, hereafter NT73), and the dimensionless quantities are defined by

$$r = \frac{R}{R_G}, \quad R_G = \frac{GM_{\text{bh}}}{c^2}, \quad \dot{m} = \frac{\dot{M}}{\dot{M}_{\text{Edd}}},$$

and

$$\dot{M}_{\text{Edd}} = \frac{L_{\text{Edd}}}{\eta_{\text{eff}} c^2} = 1.39 \times 10^{15} m \text{ kg s}^{-1}, \quad m = \frac{M_{\text{bh}}}{M_\odot}, \quad (11)$$

where $\eta_{\text{eff}} = 0.1$ is adopted.

The dimensionless scale height of the disk H/R is in principle a function of R , and it reaches a maximal value in the inner region of the disk (Laor & Netzer 1989). We adopt the maximal value of H/R in the estimate of large-scale field strength B_{pd} at the disk surface.

In L99, the strength of the magnetic field produced by dynamo processes in the disk is given by

$$\frac{B_{\text{dynamo}}^2}{4\pi} \sim \frac{W}{2H}, \quad (12)$$

where W is the integrated shear stress of the disk and H is the scale height of the disk. For a relativistic accretion disk, the integrated shear stress is given by equation (5.6.14a) in NT73. Equation (12) can be rewritten as

$$B_{\text{dynamo}} = 3.56 \times 10^8 r^{-3/4} m^{-1/2} A^{-1} B E^{1/2} \text{ G}, \quad (13)$$

where A , B , and E are general relativistic correction factors defined in NT73.

¹ The VLA is operated by the National Radio Astronomy Observatory, which is a facility of the National Science Foundation, operated under cooperative agreement by Associated Universities, Inc.

In standard accretion disk models, the angular velocity of the matter in the disk is usually very close to Keplerian velocity. For a relativistic accretion disk surrounding a rotating black hole, the Keplerian angular velocity is given by

$$\Omega(r) = 2.034 \times 10^5 \frac{1}{m(r^{3/2} + a)} \text{ s}^{-1}, \quad (14)$$

where a is dimensionless specific angular momentum of a rotating black hole.

Using equations (9)–(14), the maximal power of the jet accelerated from a magnetized disk can be found by integrating equation (8), if some parameters m , $\dot{m}$, and a are specified.

As discussed in L99, the power extracted from a rotating black hole by the Blandford-Znajek process is determined by the hole mass M_{bh} , the spin of the hole a , and the strength of the poloidal field threading the horizon of a rotating hole B_{ph} :

$$L_{\text{BZ}}^{\text{max}} = \frac{1}{32} \omega_F^2 B_{\perp}^2 R_h^2 c a^2, \quad (15)$$

for a black hole of mass M_{bh} and dimensionless angular momentum a , with a magnetic field $B_{\perp}$ normal to the horizon at R_h . Here the factor $\omega_F^2 \equiv \Omega_F(\Omega_h - \Omega_F)/\Omega_h^2$ depends on the angular velocity of field lines Ω_F relative to that of the hole, Ω_h . In order to estimate the maximal power extractable from a spinning black hole, we adopt $\omega_F = 1/2$. As the field $B_{\perp}$ is maintained by the currents in the accretion disk surrounding the hole, the strength of $B_{\perp}$ should be of the same order as that in the inner edge of the disk, and $B_{\perp} \simeq B_{\text{pd}}(r_{\text{in}})$ is therefore adopted.

4.2. Jet Power Extracted from the Disk Coronas

In this paper, we only consider the case of the jets being magnetically driven by the fields created in the coronas of the disks, as recently suggested by Merloni & Fabian (2002) for the black holes accreting at low rates.

In the disk-corona scenario, the cold disk and the hot corona above the disk are in pressure equilibrium. Most gravitational energy of the accretion matter is released in the hot corona (Haardt & Maraschi 1991; Kusunose & Mineshige 1994; Svensson & Zdziarski 1994). A small fraction of the soft photons from the cold disk is Compton upscattered to X-ray photons by the hot electrons in the corona. Even if the magnetic pressure is in equipartition with the gas pressure, the radiation of the Compton scattering dominates over the synchrotron radiation in the corona (e.g., Liu et al. 2002). Roughly half of the scattered X-ray photons illuminate the cold disk (Haardt & Maraschi 1991; Nakamura & Osaki 1993; Cao et al. 1998; Kawaguchi et al. 2001). Most soft photons from the disk leave the system without being scattered in the corona, and they are observed as optical/UV continuum. The cold disk in the disk-corona scenario has roughly half the brightness of a standard disk (SD) without a corona, if both disks are accreting at the same rate (Haardt & Maraschi 1991). The gases in the coronas are nearly virialized, their thermal velocity $v_{\text{th}} \sim (GM/R)^{1/2}$. Thus, the thickness of the corona $H_c \sim c_s/\Omega_K \sim R$, as the sound speed $v_s \sim v_{\text{th}}$ (Nakamura & Osaki 1993). The maximal magnetic stresses created by the dynamo processes are

$$\frac{B_{\text{dyn}}^2}{4\pi} \sim \frac{W_c}{2H_c}, \quad (16)$$

where $W_c = 2H_c w_c$ is the integrated shear stress of the corona (Shakura & Sunyaev 1973; L99). As the scale height of the corona $H_c \sim R$, the fields created by the dynamo processes have length scales of R . The maximal jet power that can be extracted from the corona in unit surface area is

$$q_j^{\text{max}} \sim \frac{B_{\text{dyn}}^2}{4\pi} R \Omega_K(R). \quad (17)$$

The viscous dissipation in the corona is

$$f_{\text{vis}}^+ = \frac{1}{2} W_c R \left| \frac{d\Omega}{dr} \right| = \frac{3}{4} W_c \Omega_K(R) \quad (18)$$

in unit surface area (Shakura & Sunyaev 1973). Combining equations (16)–(18) and noting $H_c \sim R$ for the corona, we have

$$q_j^{\text{max}} \sim \frac{2}{3} f_{\text{vis}}^+. \quad (19)$$

Integrating equation (19) over the surface of the corona on the assumption of local equilibrium between the radiation and energy dissipation in the corona, we obtain

$$Q_{\text{jet}}^{\text{max}} \sim \frac{2}{3} L_c, \quad (20)$$

where L_c is corona luminosity. In the disk-corona scenario, almost all gravitational energy of the accretion matter is released in the corona, i.e., $L_c \simeq L_{\text{bol}}$ (Haardt & Maraschi 1991). The emission of the corona is mainly in X-ray bands, half of which leave the disk-corona system and are observed as X-ray emission and the remaining half of which illuminate the cold disk to reradiate in optical/UV bands (Haardt & Maraschi 1991). Thus, we have

$$L_X \simeq \frac{L_{\text{bol}}}{2} \quad (21)$$

and

$$Q_{\text{jet}}^{\text{max}} \sim \frac{2}{3} L_{\text{bol}} \simeq \frac{4}{3} L_X, \quad (22)$$

which indicate that the jets can be more efficiently accelerated from the coronas above the disks than directly from the thin disks (see the discussion of the jet accelerated by the fields of the thin disks in § 1).

5. SPECTRA OF THE DISKS

The ADAF will transit to a standard optically thick, geometrically thin accretion disk at a radius outside R_{tr} (Esin et al. 1997). Here we consider a general case, i.e., an ADAF is present near the black hole and transits to a cold SD beyond the transition radius R_{tr} (Esin et al. 1997).

The flux that results from viscous dissipation in the outer region of the disk is

$$F_{\text{vis}}(R) \simeq \frac{3GM_{\text{bh}}\dot{M}}{8\pi R^3}, \quad (23)$$

which is a good approximation for $R_{\text{tr}} \gg R_{\text{in}}$. The local disk temperature of the thin cold disk is

$$T_{\text{disk}}(R) = \frac{F_{\text{vis}}^{1/4}(R)}{\sigma_B^{1/4}}, \quad (24)$$

by assuming local blackbody emission. In order to calculate the disk spectrum, we include an empirical color correction for the disk thermal emission as a function of radius. The correction has the form (Chiang 2002)

$$f_{\text{col}}(T_{\text{disk}}) = f_{\infty} - \frac{(f_{\infty} - 1)[1 + \exp(-\nu_b/\Delta\nu)]}{1 + \exp[(\nu_p - \nu_b)/\Delta\nu]}, \quad (25)$$

where $\nu_p \equiv 2.82k_B T_{\text{disk}}/h$ is the peak frequency of blackbody emission with temperature T_{disk} . This expression for f_{col} goes from unity at low temperatures to f_{∞} at high temperatures with a transition at $\nu_b \approx \nu_p$. Chiang (2002) found that $f_{\infty} = 2.3$ and $\nu_b = \Delta\nu = 5 \times 10^{15}$ Hz can well reproduce the model disk spectra of Hubeny et al. (2001). The disk spectra can therefore be calculated by

$$L_{\nu} = 8\pi^2 \left(\frac{GM}{c^2} \right)^2 \frac{h\nu^3}{c^2} \int_{r_{\text{tr}}}^{\infty} \frac{rdr}{f_{\text{col}}^4 [\exp(h\nu/f_{\text{col}}k_B T_{\text{disk}}) - 1]}. \quad (26)$$

In this ADAF+SD scenario, the ionizing luminosity from the disk is a combination of the emission from the inner ADAF and outer SD regions. The calculations of Cao (2002a) indicate that the optical continuum emission from the inner ADAF region is negligible compared with that from the outer SD region if R_{tr} is around tens to several hundreds of Schwarzschild radii, because of the low radiation efficiency of ADAFs.

6. RESULTS

In Figure 1, we plot the relation between $L_{\text{bol}}/L_{\text{Edd}}$ and $Q_{\text{jet}}/L_{\text{bol}}$. The jet power of AGNs is estimated from the low-frequency extended radio luminosity by using relation (7). The maximal jet power can be extracted by the magnetic fields created in the standard accretion disks (*solid line*), and the rapidly spinning black holes of $a = 0.95$ (*dotted line*) are calculated by using the method described in § 4.1. The flat-spectrum and steep-spectrum sources are plotted as squares and circles, respectively. The sources shown above the solid line (*filled squares and circles*) in Figure 1 are referred to as high jet power sources. This indicates that the Blandford-Payne mechanism is unable to produce the level of jet power observed in the high jet power sources if only standard thin accretion disks are present in these sources. The jets in almost all sources in this sample cannot be powered only by the Blandford-Znajek mechanism, even if the black holes in these sources are rapidly spinning at $a = 0.95$.

As mentioned in § 2, there are uncertainties in the estimates of the black hole mass M_{bh} , bolometric luminosity L_{bol} , and the dimensionless accretion rate $\dot{m}$. The beamed optical continuum emission from the jets may lead to overestimates of bolometric luminosity L_{bol} , especially for flat-spectrum sources, which would shift the real locations of the sources in Figure 1 away from the solid line representing the maximal jet power extracted by the Blandford-Payne mechanism. If the disklike BLR geometry is present in these sources, their black hole masses may be underestimated (McLure & Dunlop 2001), and the real locations of the sources would also be shifted away from the solid maximal jet power line in Figure 1. The uncertainty of the coefficient in relation (3) can also lead to errors in the estimate of bolometric luminosity L_{bol} . If the present coefficient in relation (3) is overestimated, the locations of high jet power sources will be shifted away from the solid maximal jet power line. The jets in all of these high jet power sources can be extracted by the Blandford-Payne

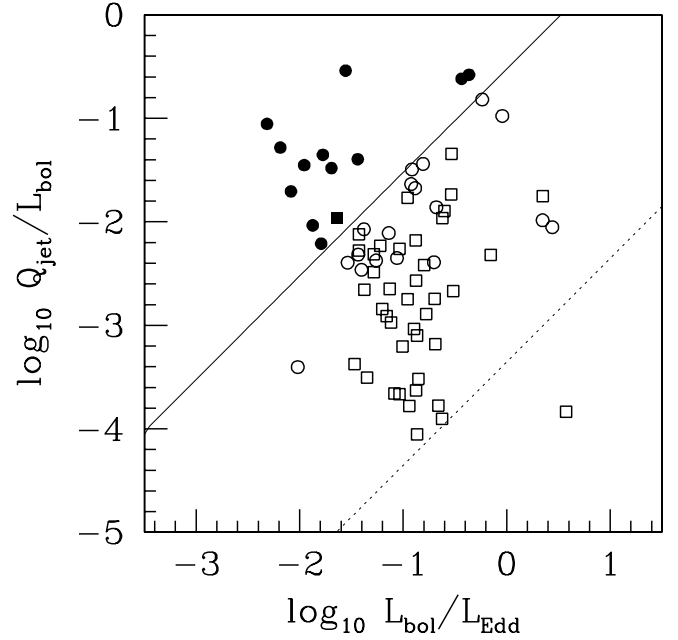


FIG. 1.—Plot of $L_{\text{bol}}/L_{\text{Edd}}$ vs. $Q_{\text{jet}}/L_{\text{bol}}$. The squares represent flat-spectrum sources, and the circles represent steep-spectrum sources. The solid line represents the maximal jet power extracted from a standard thin accretion disk (the Blandford-Payne mechanism), while the dotted line represents the maximal jet power extracted from a rapidly spinning black hole $a = 0.95$ surrounded by a standard thin accretion disk (the Blandford-Znajek mechanism). The sources shown above the solid line (*filled squares and circles*) are referred to as high jet power sources.

mechanism for thin disk cases only if the coefficient in relation (3) is significantly underestimated at more than 1 order of magnitude, i.e., >100 is required instead of 9. Thus, the conclusion that the jets in these high jet power sources cannot be accelerated by the Blandford-Payne mechanism for standard thin disk cases will not be altered by the uncertainties of the estimates of the black hole mass M_{bh} , bolometric luminosity L_{bol} , or dimensionless accretion rate $\dot{m}$, unless the present coefficient in relation (3) is underestimated at an order of magnitude, which seems impossible.

The relation between black hole mass M_{bh} and optical continuum luminosity L_{λ} at $5100 \text{ \AA}$ is plotted in Figure 2. The optical continuum emission from a pure ADAF can be calculated using the approach proposed by Mahadevan (1997). We use the same approach proposed by Cao (2002a) to calculate the maximal optical continuum emission from an ADAF as a function of black hole mass M_{bh} . The maximal optical continuum emission requires the parameter $\beta = 0.5$, which describes the magnetic field strength with respect to gas pressure, and viscosity $\alpha = 1$. Changing the value of accretion rate $\dot{m}$, we can find the maximal optical continuum luminosity for a given black hole mass (see Cao 2002a for details). We find that all sources have optical continuum luminosity higher than the maximal optical luminosity expected from pure ADAFs, which implies that the emission from pure ADAFs is unable to explain the optical ionizing luminosity of these sources. We then consider another possibility, i.e., that the ADAF transits to a standard thin disk outside the transition radius R_{tr} . The theoretical calculations of the optical continuum luminosity for such ADAF+SD systems with different accretion rates and transition radii are also plotted in Figure 2 (the calculations are carried out as described in § 5). We find that ADAFs may be present in the inner region of the disk in

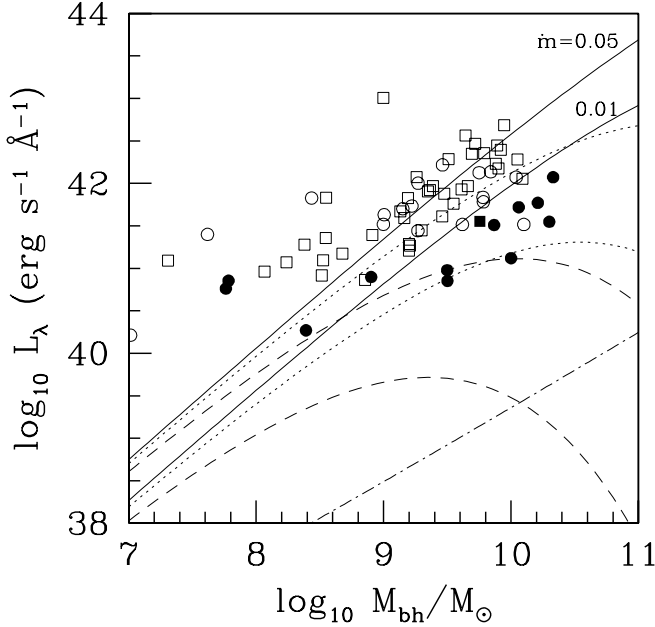


FIG. 2.—Black hole mass M_{bh} vs. optical luminosity L_{λ} at 5100 Å. The dash-dotted line represents the maximal optical continuum luminosity from pure ADAFs. The solid lines represent SDs accreting at $\dot{m} = 0.01$ and 0.05, respectively. The dotted lines and the dashed lines represent the ADAF+SD models with different transition radii $R_{\text{tr}} = 20GM/c^2$ and $50GM/c^2$, respectively. The upper lines are plotted for $\dot{m} = 0.05$, while the lower lines are plotted for $\dot{m} = 0.01$. The symbols are the same as in Fig. 1.

these high jet power sources only if the critical accretion rates $\dot{m}_{\text{crit}}$ are as high as ≥ 0.05 . For flat-spectrum sources, the optical continuum emission from the accretion disks may be over-estimated, and the accretion rates inferred from Figure 2 may be larger than their real values. The beaming effect is believed to be less important with respect to the line of sight for steep-spectrum sources (because of their relatively larger angles) than for flat-spectrum sources, and it is therefore suggested that the optical continuum emission of steep-spectrum sources is mainly from their accretion disks (Serjeant et al. 1998). The high jet power sources in this sample are steep-spectrum sources, with the exception of one flat-spectrum source. Thus, the estimated dimensionless accretion rates for these high jet power sources should be less affected by the beamed jet emission than those of flat-spectrum sources, and the main results for the accretion rates and transition radii for these high jet power sources should not be changed significantly.

The relation between bolometric luminosity L_{bol} and jet power Q_{jet} is plotted in Figure 3. The jet power is lower than the bolometric luminosity for all sources in this sample, which would still hold even if the bolometric luminosity is over-estimated by less than an order of magnitude for flat-spectrum sources.

In order to test relation (22) between $Q_{\text{jet}}^{\text{max}}$ and total X-ray luminosity L_X from the disk coronas predicted by the scenario in which the jets are magnetically accelerated by the fields created in the disk coronas, we roughly convert X-ray luminosity in 0.1–2.4 keV, $L_{X,0.1-2.4 \text{ keV}}$, detected by *ROSAT* (Voges et al. 1999) to the total X-ray luminosity, assuming the X-ray continuum emission to be in the range 0.01–100 keV with a mean energy spectral index of $\alpha_X = 1$ ($f_X \propto E^{-\alpha_X}$). This is consistent with the general features of the theoretical calculations for the X-ray spectra from the disk coronas (e.g., Nakamura & Osaki 1993). In this case, the total X-ray

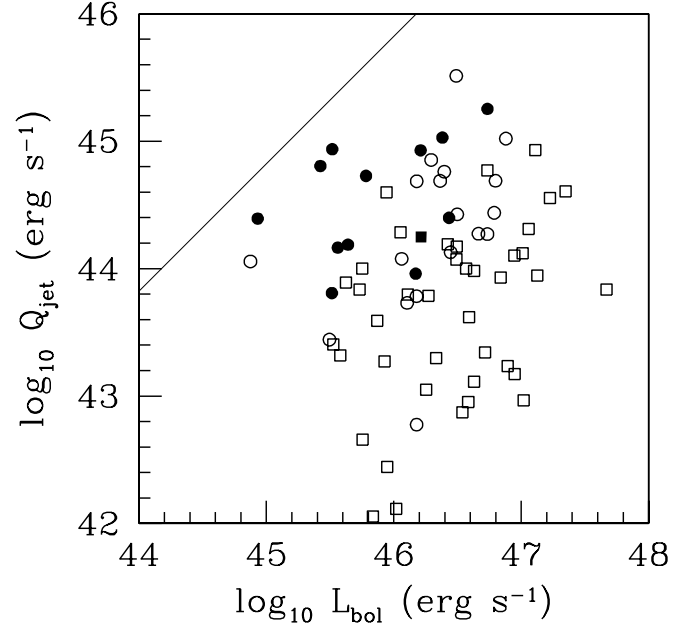


FIG. 3.—Relation between bolometric luminosity L_{bol} and jet power Q_{jet} . The line represents $Q_{\text{jet}} = 2/3 L_{\text{bol}}$. The symbols are the same as in Fig. 1.

luminosity $L_X \simeq 2.9L_{X,0.1-2.4 \text{ keV}}$. Figure 4 shows the relation between K -corrected core radio luminosity $L_{c,5G}$ and X-ray luminosity L_X . The linear regression considering errors in both coordinates gives (Press et al. 1992)

$$\log_{10} L_X = 0.866 \log_{10} L_{c,5G} + 7.517 \quad (27)$$

for 40 flat-spectrum sources (the solid line in Fig. 4) and

$$\log_{10} L_X = 0.371 \log_{10} L_{c,5G} + 29.272 \quad (28)$$

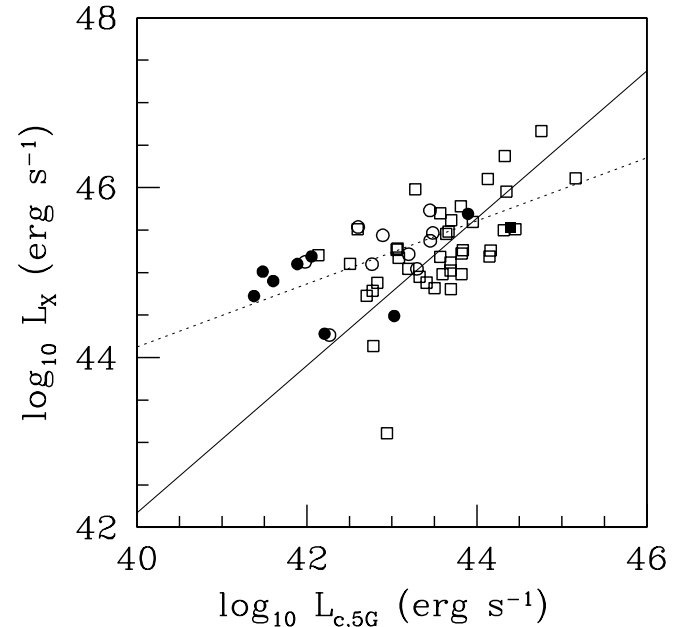


FIG. 4.—Relation between X-ray luminosity L_X and core radio luminosity L_c at 5 GHz. The solid and dotted lines represent the linear regressions for flat-spectrum and steep-spectrum sources, respectively. The symbols are the same as in Fig. 1.

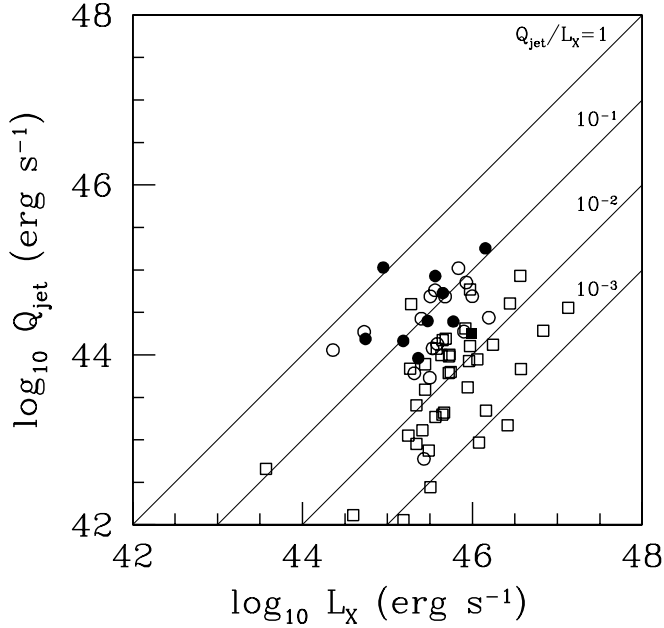


FIG. 5.—Relation between X-ray luminosity L_X and jet power Q_{jet} . The lines represent $Q_{\text{jet}}/L_X = 10^{-3}, 10^{-2}, 10^{-1}$, and 1, respectively. The symbols are the same as in Fig. 1.

for 18 steep-spectrum sources (the dotted line in Fig. 4). The Spearman correlation analyses (Press et al. 1992) show that the correlation for flat-spectrum sources is at the significant level of 99.98%, while it becomes 98.5% for steep-spectrum sources. The different slopes of the correlations between L_X and $L_{c,5G}$ for steep-spectrum and flat-spectrum sources are obvious, which may imply that the observed X-ray emission from these two kinds of AGNs has different origins.

We plot the relation between X-ray luminosity L_X and jet power Q_{jet} in Figure 5. We find that almost all sources have $Q_{\text{jet}} < L_X$, except for one steep-spectrum source with $Q_{\text{jet}} \sim L_X$.

7. DISCUSSION

There are 13 sources with jet power above the maximal jet power expected to be extracted from standard thin disks (see Fig. 1); i.e., the fields of thin disks are unable to drive such strong jets in these high jet power sources. The jet can be accelerated from an ADAF more efficiently than from a standard thin disk, although an accurate calculation of the maximal jet power extracted from an ADAF is still unavailable (Meier 2001). Meier (2001) pointed out that general relativistic effects may be important, and a rough estimate of jet power shows that the maximal jet power extracted from an ADAF as high as 10^{45} – 10^{46} ergs s $^{-1}$ may be possible for a rapidly spinning massive black hole ($\sim 10^9 M_\odot$). If ADAFs are present in these sources, they will have fainter optical continuum emission than SDs as a result of the lower accretion rates and lower radiation efficiency of ADAFs. We find that pure ADAFs are unable to produce observed bright optical continuum luminosity for all sources in this sample (see Fig. 2). However, this cannot rule out the presence of ADAFs in these sources, because the ADAF may transit to a standard thin disk in the region outside the transition radius R_{tr} . In this case, the standard thin disk in the outer region may be responsible for the observed optical continuum emission. The ADAF+SD scenario is tested against observational optical

continuum luminosity in Figure 2. As 12 of 13 high jet power sources are steep-spectrum AGNs, the optical continuum emission is believed to be mainly from the accretion disks (Serjeant et al. 1998). We find that relatively high accretion rates $\dot{m} \geq 0.05$ are required to explain the optical continuum luminosity of these high jet power sources if the ADAFs are truncated at rather small radii (e.g., ≤ 20 Schwarzschild radii). If the transition radii are larger than this value, even higher accretion rates are required. The exact value of the critical accretion rate $\dot{m}_{\text{crit}}$ is still unclear for an ADAF, which depends on the value of the disk viscosity parameter α , i.e., $\dot{m}_{\text{crit}} \simeq 0.28\alpha^2$ (Mahadevan 1997). Recent three-dimensional MHD simulations suggest that the viscosity α in the discs is ~ 0.1 (Armitage 1998) or ~ 0.05 – 0.2 (Hawley & Balbus 2002). If a value of $\alpha = 0.2$ is conservatively adopted, the accretion rates $\dot{m} \leq \dot{m}_{\text{crit}} \sim 0.01$ are required for ADAFs. This probably implies that ADAFs are not present in these high jet power sources.

The bolometric luminosity of an ADAF can be significantly lower than the jet power extracted from or the rapidly spinning black hole surrounded by the ADAF, because the radiation efficiency of ADAFs is much lower than the efficiency of power channeled into the jet (Armitage & Natarajan 1999; Meier 2001). We have not found any source in our sample with $Q_{\text{jet}} > L_{\text{bol}}$ (see Fig. 3). This implies that standard thin disks, at least in the outer regions of the disks, are required to produce the observed bright bolometric luminosity, even if ADAFs are present in these sources, as in the case of $\alpha > 0.2$, although it seems inconsistent with MHD simulations (Hawley & Balbus 2002).

The ADAF solutions can be modified to adiabatic inflow-outflow solutions (ADIOSs) if a powerful wind is present to carry away mass, angular momentum, and energy from the accreting gas (Blandford & Begelman 1999). In this case, the accretion rate of the disk is a function of the radius r instead of a constant accretion rate along r for a pure ADAF. For an ADIOS flow, the gas swallowed by the black hole is only a small fraction of that supplied, as most of the gas is carried away in the wind before it reaches the black hole. The ADIOS flow has similar local structure to that of the ADAF accreting at the same rate, so the accretion rate $\dot{m}$ for an ADIOS flow at any radius should be smaller than the critical rate $\dot{m}_{\text{crit}}$ of ADAFs. The ADIOS flow is fainter than that of an ADAF, if they are accreting at the same rate at the outer radius, because the accretion rate of an ADIOS flow decreases while the gas is flowing onto the hole (e.g., Quataert & Narayan 1999; Chang et al. 2002). Thus, even if an ADIOS flow, rather than an ADAF, is present in the inner region of the disk, the optical continuum emission is still dominated by the emission from the standard thin disk in the outer region, as it is in the ADAF case (Cao 2002a). Our present spectral analyses for ADAF+SD systems are valid even for ADIOS+SD systems. For convection-dominated accretion flows (CDAFs), most of the gas circulates in convection eddies rather than accreting onto the black hole (Narayan et al. 2000; Quataert & Gruzinov 2000). The accretion rates of CDAFs are much smaller than those of non-convecting ADAFs, and CDAFs are very faint (Ball et al. 2001). This implies that the accretion rate in a standard thin disk surrounding a CDAF should be much smaller than the critical accretion rate $\dot{m}_{\text{crit}}$ in ADAFs, if a steady accretion flow is assumed. The observed bright optical continuum emission in these high jet power sources requires the black holes to accrete at rather high rates, which seems to rule out the possibility of CDAFs surrounded by standard thin disks.

The powerful jets in these high jet power sources may be accelerated from ADAFs (or ADIOS flows) surrounded by standard thin disks outside ~ 20 Schwarzschild radii accreting at $\dot{m} \geq 0.05$, which is required by observed bright optical continuum emission in these sources. However, such high accretion rates require a high viscosity α , which seems inconsistent with MHD simulations (Hawley & Balbus 2002). Here we propose an alternative explanation, i.e., the jets, at least in these high jet power sources, are accelerated from the coronas of the disks. In this case, the maximal power of the jet accelerated from the corona of the disk can be as high as its bolometric luminosity (see eq. [22] and Fig. 3), and the strong jets in these high jet power sources can be naturally explained by this scenario. The cold disk irradiated by the corona above can produce bright optical continuum emission as observed in these high jet power sources.

It should be noted that the calculations carried out in § 4.2 are based on the assumption of a perfect parallel corona structure. The real corona above the cold disk may not always have perfect parallel structure. The time variations and spatial inhomogeneity can be caused by the magnetic fields (e.g., Kawaguchi et al. 2000). Thus, the maximal jet power extractable from the corona may be lower than that estimated by equation (22) because of the inhomogeneous corona structure. Most high jet power sources in our sample have $Q_{\text{jet}}/L_{\text{bol}} \leq 0.1$, while three sources have $Q_{\text{jet}}/L_{\text{bol}} \simeq 0.3$ (see Fig. 1). The disk-corona scenario may still be a reasonable explanation for jet formation in these high jet power sources.

The emission from the coronas is mainly in X-ray bands, and this scenario predicts a relation between the maximal power of the jet accelerated from the corona and the X-ray luminosity of the corona (eq. [22]). This prediction can be tested against observations. However, a difficulty arises from the fact that the observed X-ray emission from these radio-loud AGNs may be a mixture of the emission from the jets and accretion disk coronas. If the X-ray emission from the jets in radio-loud AGNs dominates over that from the disk coronas, a correlation between radio core emission and X-ray emission is expected, because both of them are from the jets and are Doppler beamed. Figure 4 gives different correlations between core radio luminosity $L_{c,5G}$ and X-ray luminosity L_X for flat-spectrum and steep-spectrum AGNs. Falcke et al. (2004) used LINERs, FR I galaxies, and BL Lac objects in their investigations of the radio–X-ray correlation, in which ADAFs are believed to be present and the radio and X-ray emission are believed to be dominated by the emission from the jets. The correlation found in this work between radio and X-ray luminosities for flat-spectrum sources is roughly consistent with the correlation given by Falcke et al. (2004). The core dominance parameter R_c varies over several orders of magnitude

for the sources in this sample (Cao & Jiang 2001), which reflects the Doppler factors spread over a large range. In this paper, we do not intend to explore this correlation in detail, as done by Falcke et al. (2004), because of unknown Doppler beaming factors of the sources in this sample. Nevertheless, the different slopes of the correlations for flat-spectrum and steep-spectrum sources do indicate that the different origins of the X-ray emission are these two different kinds of sources. Compared with flat-spectrum sources, most steep-spectrum sources have brighter X-ray emission than their flat-spectrum counterparts with similar radio core emission (see Fig. 4). This may imply that the X-ray emission from the disk coronas, rather than from the jets, is dominant for steep-spectrum AGNs. If this is the case, the observed X-ray emission from steep-spectrum sources is probably from the disk coronas. We can tentatively use Figure 5 to test our estimate of the relation between $Q_{\text{jet}}^{\text{max}}$ and L_X (eq. [22]) expected for the jets magnetically accelerated from the coronas of the disks. Figure 5 shows that $Q_{\text{jet}} < L_X$ is satisfied for almost all sources, with the exception of one with $Q_{\text{jet}} \sim L_X$, which is consistent with the scenario of the jets being magnetically accelerated from the disk coronas. If this scenario is correct, it implies that the factor f in equation (7) describing the uncertainty of the jet power estimated from the extended radio luminosity should be close to unity, at least for these high jet power sources (see Figs. 3 and 5).

Recently, Hujeirat et al. (2002) proposed that the jet is launched from a layer governed by a highly diffuse, super-Keplerian rotating plasma that is thermally dominated by virial hot and magnetized ions. This layer is located between the accretion disk and the corona surrounding the nucleus. In this model, most of the accretion energy can be converted into magnetic and kinetic energies that go into powering the jet (Hujeirat et al. 2003; Hujeirat 2004), and the powerful jet with $Q_{\text{jet}} \sim L_{\text{bol}}$ can form naturally in such a layer. This model, in principle, could explain the powerful jets of the high jet power sources in our present sample, although the detailed numerical model calculations for the jets in these sources are beyond the scope of this paper.

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