THREE-DIMENSIONAL SIMULATIONS OF INFLOWS IRRADIATED BY A PRECESSING ACCRETION DISK IN ACTIVE GALACTIC NUCLEI: FORMATION OF OUTFLOWS

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ABSTRACT

We present three-dimensional hydrodynamical simulations of gas flows in the vicinity of an active galactic nucleus (AGN) powered by a precessing accretion disk. We consider the effects of the radiation force from such a disk on its environment on a relatively large scale (up to ~10 pc). We implicitly include the precessing disk by forcing the disk radiation field to precess around a symmetry axis with a given period (P) and a tilt angle (θ). We study the time evolution of the flows irradiated by the disk and investigate basic dependencies of the flow morphology, mass flux, and angular momentum on different combinations of θ and P. As this is our first attempt to model such three-dimensional gas flows, we consider a simplest form of radiation force, i.e., force due to electron scattering, and neglect the forces due to line and dust scattering/absorption. Furthermore, the gas is assumed to be nearly isothermal. We find that the gas flow settles into a configuration with two components, (1) an equatorial inflow and (2) a bipolar inflow/outflow, with the outflow leaving the system along the poles (the directions of disk normals). However, the flow does not always reach a steady state. We find that the maximum outflow velocity and the kinetic outflow power at the outer boundary can be reduced significantly with increasing θ. We also find that the mass inflow rate across the inner boundary does not change significantly with increasing θ. The amount of the density-weighted mean specific angular momentum deposited in the environment by the precessing disk increases as P approaches the gas free-fall time (t_g) and then decreases as P becomes much larger than t_g. Generally, the characteristics of the flows are closely related to a combination of P and θ, but not to P and θ individually. Our models exhibit helical structures in the weakly collimated outflows. Although on different scales, the model reproduces the Z- or S-shaped density morphology of gas outflows, which are often seen in radio observations of AGNs.

Subject headings: accretion, accretion disks — galaxies: jets — galaxies: kinematics and dynamics — hydrodynamics — methods: numerical

1. INTRODUCTION

Powered by accretion of matter onto a supermassive (10^6–10^10 M_☉) black hole (SMBH), active galactic nuclei (AGNs) release large amounts of energy (e.g., Lynden-Bell 1969) as electromagnetic radiation (10^{10}–10^{14} L_☉) over a wide range of wavelengths, from the X-ray to the radio. The very central location of AGNs in their host galaxies indicates that the radiation from AGNs can play an important role in determining the physical characteristics (e.g., the ionization structure, the gas dynamics, and the density distribution) of their surrounding environment on different scales, i.e., from the scale of the AGN itself to a larger galactic scale, and even to an intergalactic scale (e.g., Quilis et al. 2001; Dalla Vecchia et al. 2004; McNamara et al. 2005; Zanni et al. 2005; Fabian et al. 2006b; Vernaleo & Reynolds 2006). The feedback process of AGNs in the form of mass or energy outflows, in turn, is one of the key elements in galaxy formation/evolutionary models (e.g., Ciotti & Ostriker 1997, 2001, 2007; Silk & Rees 1998; King 2003; Begelman & Nath 2005; Hopkins et al. 2005; Murray et al. 2005; Sazonov et al. 2005; Silk 2005; Springel et al. 2005; Brighenti & Mathews 2006; Fabian et al. 2006a; Fontanot et al. 2006; Thacker et al. 2006; Wang et al. 2006; Tremonti et al. 2007).

Although the AGN outflows can be driven by magnetocentrifugal force (e.g., Blandford & Payne 1982; Emmering et al. 1992; Königl & Kartje 1994; Bottorff et al. 1997) and thermal pressure (e.g., Weymann et al. 1982; Begelman et al. 1991; Everett & Murray 2007), it is the radiation force from the luminous accretion disk that is most likely the dominant force driving winds capable of explaining the blueshifted absorption-line features often seen in the UV and optical spectra of AGNs (e.g., Shlosman et al. 1985; Murray et al. 1995; Proga et al. 2000; Proga & Kallman 2004). In reality, these three forces may interplay and contribute to the dynamics of the outflows in AGNs to somewhat different degrees.

Another complication in the outflow gas dynamics is the presence of dust. The radiation pressure on dust can drive dust outflows, and their dynamics is likely to be coupled with the gas dynamics (e.g., Phinney 1989; Pier & Krolik 1992; Emmering et al. 1992; Laor & Draine 1993; Königl & Kartje 1994; Murray et al. 2005). The AGN environment on relatively large scales (10^2–10^3 pc) is known to be a mixture of gas and dust (e.g., Antonucci 1984; Miller & Goodrich 1990; Awaki et al. 1991; Blanco et al. 1990; Krolik 1999); however, on much smaller scales (<~10 pc), one does not expect much dust to be present, because the temperature of the environment is high (>10^4 K). Concentrating on only the gas component, the dynamics of the outflows on smaller scales was studied by, e.g., Arav et al. (1994) and Proga et al. (2000) in one and two dimensions, respectively.

Radio observations show that a significant fraction of extended extragalactic sources display bending or twisting jets from their host galaxies. For example, Florido et al. (1990) found that ~11% of their sample (368 objects) show antisymmetrically bending jets (S- or Z-shaped morphology), while ~9% show symmetrically bending jets (U-shaped morphology). Similarly, Hutchings et al. (1988) studied the morphology of the radio lobes from 128 quasars (with z < 1) and found that 30% of the sample show signs of bending jets. The bending and misalignment of jets are
also observed on parsec scales in compact radio sources (e.g., Linfield 1981; Appl et al. 1996; Zensus 1997). Examples of the radio maps displaying the $S$- or $Z$-shaped morphology of jets can be found in, e.g., Condon & Mitchell (1984), Hunstead et al. (1984), and Tremblay et al. (2006).

Using data available in the literature, Lu & Zhou (2005) compiled a list (see their Table 1) of 41 known extragalactic radio sources that show evidence of jet precession, along with their jet precession periods ($P$) and the half-opening angle ($\psi$) of jet precession cones. According to this list, a large fraction (67%) of these systems have rather small half-opening angles, i.e., $\psi \leq 15$. A large scatter in the precession periods is found in their sample; however, most of the precession periods are found between $10^4$ and $10^6$ yr (see also Roos 1988). Note that the precession periods are usually too long to be determined directly by variability observations. Typically, the precession periods are found by fitting the radio map with a kinematic jet model (e.g., Gower et al. 1982; Veilleux et al. 1993). Interestingly, Appl et al. (1996) showed that a typical precession period of a tilted massive torus around an SMBH is $\sim 10^6$ yr.

The $S$- and $Z$-shaped morphology seen in the observations mentioned above can be naturally explained by precessing jets. Furthermore, the precessing of jets can occur if the underlying accretion disk is tilted (or warped) with respect to the plane of symmetry. There are at least five known mechanisms that can cause warping and precession in accretion disks: (1) the Bardeen-Petterson effect (Bardeen & Petterson 1975; see also Schreier et al. 1972; Nelson & Papaloizou 2000; Fragile & Anninos 2005; King et al. 2005), (2) tidal interactions in binary black hole (BH) systems (e.g., Roos 1988; Sillanpää et al. 1988; Katz 1997; Romero et al. 2000; Caproni et al. 2004), (3) radiation-driven instability (e.g., Petterson 1977; Pringle 1996; Maloney et al. 1996; Armitage & Pringle 1997), (4) magnetically driven instability (Alv 1980; Lai 2003), and (5) disk-ISM interactions (e.g., Quillen & Bower 1999). Using a small sample of AGNs, Caproni et al. (2006) examined whether mechanisms (1)–(4) are capable of explaining the observed precession periods. Similarly, Tremblay et al. (2006) searched for a possible cause of disk precession and warping of the FR I radio source 3C 449 using mechanisms (2), (3), and (4) above. In general, it is very difficult to determine the exact cause of disk/jet precession for a given AGN system because of large uncertainties in model parameters and observed precession periods (which are also often model dependent).

Kochanek & Hawley (1990) presented a hydrodynamical simulation of jet propagation along the surface of an axisymmetric hollow cone to approximate a jet with fast precession; however, the intrinsically nonaxisymmetric nature of the dynamics of jet precession requires the problem to be solved/simulated in three dimensions. Hydrodynamical simulations of extragalactic radio sources with precessing jets in three full dimensions have been performed by, e.g., Cox et al. (1991), Hardee & Clarke (1992), and Hardee et al. (1994). Typically, in these models, the jets are driven near the origin by a small-amplitude precession to break the symmetry and excite the helical modes of the Kelvin-Helmholtz instability. Careful stability/instability analysis of such simulations has been presented by Hardee et al. (1995). The effect of the magnetic field has been also investigated by, e.g., Hardee & Clarke (1995), while the effect of optically thin radiative cooling on the Kelvin-Helmholtz instability has been investigated by, e.g., Xu et al. (2000). Precession of relativistic jets in three dimensions with or without a magnetic field has been also studied (e.g., Hardee et al. 2001; Hughes et al. 2002; Aloy et al. 2003; Mizuno et al. 2007). On much larger scales, Sternberg & Soker (2007) studied the effect of precessing massive slow jets onto the intergalactic medium (IGM) in a galaxy cluster and found that such jets can inflate a fat bubble in the IGM. In the models mentioned above, the jets themselves are injected on small scales, and the jet propagations are studied. However, it is also possible to model a self-consistent production of a jet and its subsequent propagation. For example, a jet can be produced from infalling matter by radiation pressure due to a luminous accretion disk (e.g., see Proga 2007; Proga et al. 2008 for axisymmetric cases).

Regardless of the exact cause of disk/jet precession, the observations (e.g., Florido et al. 1990; Hutchings et al. 1988) suggest that a significant fraction of AGNs contain warped or precessing disks. One might expect the details of the radiative feedback processes in such systems to be different from the ones predicted by axisymmetry models (e.g., Proga et al. 2000, 2008; Proga 2007). If they differ, then by how much? In this paper, we explore the effects of disk precession on the gas dynamics in the AGN environment by simulating the outflows driven by the radiation force from a luminous precessing accretion disk around a SMBH. Specifically, we examine how the mass accretion rate, the outflow powers (kinetic and thermal), the morphology of the flows, and the specific angular momentum of the gas are affected by the presence of a precessing disk and its radiation field. This is our first step toward a full extension of the axisymmetric radiation-driven wind model of Proga (2007) to a full three-dimensional model.

In the following section, we describe our method and model assumptions, and we give the results of our three-dimensional hydrodynamical simulations in §3. Our conclusions are summarized in §4.

2. Method

Our basic model configuration is shown in Figure 1. The model geometry and the assumptions of the SMBH and the disk are very similar to those in Proga (2007). In Figure 1, a SMBH with mass $M_{\text{BH}}$ and its Schwarzschild radius $r_s = 2GM_{\text{BH}}/c^2$ is placed at the center of the Cartesian coordinate system $(x, y, z)$. The X-ray–emitting corona region is defined as a sphere with radius $r_*$, as shown in the figure. The geometrically thin and optically thick hot accretion disk (e.g., Shkura & Sunyaev 1973) is placed near the $x$-$y$ plane. In the case of an axisymmetric model, the $z$-axis in the figure becomes the axis of symmetry, and the accretion disk is in the $x$-$y$ plane. To simulate the disk precession, we assume that the angular momentum ($J_D$) of the accretion disk is tilted from the $z$-axis by an angle $\Theta$. In other words, the accretion disk is assumed to be tilted by $\Theta$ from the $x$-$y$ plane. Furthermore, the accretion disk, hence its angular momentum $J_D$, is assumed to precess around the $z$-axis with the precession period $P$. The three-dimensional hydrodynamical simulations are performed in the spherical coordinate system $(r, \theta, \phi)$ from the inner boundary $r_i$ to the outer boundary $r_o$. The pole of the spherical coordinate system coincides with the $z$-axis. The radiation forces, from the corona region (the sphere with its radius $r_*$) and the accretion disk, acting on the gas located at a point $(p)$ in the field are assumed to be directed only along the radius. The magnitude of the radiation force due to the corona is assumed to be a function of the radius $r$, but that due to the accretion disk is assumed to be a function of $r$ and the angle ($\theta'$) between the angular momentum of the disk $J_D$ and the position vector $r$, as shown in the figure. The point-source–like approximation for the disk radiation pressure at $p$ is valid when $r_D \ll r_i$. In the following, we describe our radiation hydrodynamics, our implementation of the continuum radiation sources (the corona and disk), and the model parameters and assumptions.
2.1. Hydrodynamics

We employ three-dimensional hydrodynamical simulations of the outflow from and accretion onto a central part of the AGN, using the ZEUS-MP code (cf. Hayes et al. 2006), a massively parallel message-passing-interface–implemented (MPI-implemented) version of the ZEUS-3D code (cf. Hardee & Clarke 1992; Clarke 1996). The ZEUS-MP is a Eulerian hydrodynamics code that uses the method of finite differencing on a staggered mesh with a second-order–accurate, monotonic advection scheme (Hayes et al. 2006). To compute the structure and evolution of a flow irradiated by a strong continuum radiation of AGNs, we solve the following set of hydrodynamic (HD) equations:

\[
\frac{D\rho}{Dt} + \rho \nabla \cdot \vec{v} = 0, \\
\rho \frac{Dv}{Dt} = -\nabla P + \rho g + \rho g_{\text{rad}}, \\
\rho \frac{Du}{Dt} \left( \frac{u}{\rho} \right) = -P \nabla \cdot \vec{v} + \rho C,
\]

where \(\rho, u, P, \) and \(v\) are the mass density, energy density, pressure, and the velocity of gas, respectively, and \(g\) is the gravitational force per unit mass. The Lagrangian/comoving derivative is defined as \(D/Dt = \partial/\partial t + \vec{v} \cdot \nabla\). We have introduced three new components to the ZEUS-MP in order to treat the gas dynamics more appropriately for the gas flow in and around AGNs. The first is the acceleration due to radiative force per unit mass \(g_{\text{rad}}\) in equation (2), and the second is the effect of radiative cooling (and heating) simply as the net cooling rate \(C\) in equation (3). As these are our first three-dimensional simulations with this code, we consider a simplest case, i.e., \(C = 0\), but \(g_{\text{rad}} \neq 0\). We also use \(\gamma = 1.01\) in the equation of state \(P = (\gamma - 1)u\), where \(\gamma\) is the adiabatic index. In the following, our implementation of \(g_{\text{rad}}\) is described.

2.1.1. Radiation Force

To evaluate the radiative acceleration due to line absorption/scattering, we follow the method of Proga et al. (2000), who applied the modified Castor, Abbott, and Klein (CAK) approximation (Castor et al. 1975). Their model works under the assumption of the Sobolev approximation (e.g., Sobolev 1957; Castor 1970; Lucy 1971); hence, the following conditions are assumed to be valid: (1) the presence of a large velocity gradient in the gas flow and (2) negligible intrinsic line width compared to the Doppler broadening of a line. Following Proga et al. (2000), the radiative acceleration of a unit mass at a point \(\vec{p}\) can be written as

\[
g_{\text{rad}} = \int_{\Omega} (1 + \mathcal{M}) \left[ \frac{\sigma_c J(r, \hat{n})}{c} \right] \hat{n} \, d\Omega,
\]

where \(J\) is the frequency-integrated continuum intensity in the direction \(\hat{n}\) and \(\Omega\) is the solid angle subtended by the source of continuum radiation. Also, \(\sigma_c\) is the electron-scattering cross
section. The force multiplier $M$ is a function of the optical depth parameter $\tau$, which is similar to the Sobolev optical depth (cf. Rybicki & Hummer 1978) and can be written as

$$\tau = \frac{\sigma_t \nu_{th}}{Q},$$

where $\nu_{th}$ is the thermal velocity of the gas and $Q = d\nu/dl$ is the directional derivative of the velocity field in the direction $\hat{n}$, where $dl$ is the line element in the same direction. Furthermore, equation (4) can be simplified greatly when the continuum radiation source is approximated as a point, i.e., when $r \gg r_c$, where $r_c$ is the radius of the radiation source. In our case, we consider the accretion disk, which emits most of its radiation from its innermost part, between $r_i$ and $r_p$ in Figure 1; hence, the condition $r \gg r_c$ is satisfied. Using this approximation, the radiative acceleration $g_{rad}$ will be radial only and be a function of the radial position and polar angle (if the contribution from the disk luminosity is included), i.e., $g_{rad} = g_{rad}(r, \theta)\hat{r}$. This simplification is very useful for our purposes, as it reduces the computational time significantly; hence it enables us to perform large-scale three-dimensional simulations. Unlike Proga et al. (2000), we consider the case in which the radiative acceleration is dominated by the continuum process, i.e., $M = 0$ in equation (4) in this paper, since we initially intend to investigate the basic characteristics of the impact of the disk precession that do not depend on the details of the radiation force model. The models with $M \neq 0$ in equation (4) and $C \neq 0$ will be presented in a forthcoming paper.

2.2. Continuum Radiation Source

As mentioned earlier, we consider two different continuum radiation sources in our models: (1) the accretion disk and (2) the central spherical corona. Since the geometry of the central engine in AGNs is not well understood, we assume that it is composed of a spherically shaped corona with its radius $r_c$ and the innermost part of the accretion disk (cf. Fig. 1). The disk is assumed to be flat, Keplerian, geometrically thin, and optically thick. The disk radiation is assumed to be dominated by the radiation from the disk between the radii $r_i$ and $r_p$, where $r_c = 3r_p$ and $r_i < r_D \ll r_i$ (cf. Fig. 1). Note that the exact size of $r_D$ does not matter as long as it satisfies this condition in order for the point-source approximation mentioned in §2.1 to be valid.

In terms of the disk mass accretion rate ($\dot{M}_D$), the mass of the BH ($M_{BH}$), and the Schwarzschild radius ($r_S$), the total luminosity ($L$) of the system can be written as

$$L = \eta \dot{M}_D c^2,$$

$$= \frac{2\eta G M_{BH} \dot{M}_D}{r_S},$$

where $\eta$ is the rest-mass conversion efficiency (e.g., Shakura & Sunyaev 1973). Following Proga (2007) and Proga et al. (2008), we simply assume the system essentially radiates only in the UV and the X-ray bands. The total luminosity of the system $L$ is then the sum of the UV luminosity $L_{UV}$ and the X-ray luminosity $L_X$, i.e., $L = L_{UV} + L_X$. Furthermore, we assume that the disk radiates only in the UV and that the central corona radiates in the X-ray. The ratio of the disk luminosity ($L_D$) to the total luminosity is parameterized as $f_D = L_D/L$ and that of the corona luminosity ($L_c$) to the total luminosity is parameterized as $f_c = L_c/L$. Consequently, $f_D + f_c = 1$.

In the point-source approximation limit, the radiation flux from the central X-ray corona region can be written as

$$F_c = \frac{L_c}{4\pi r^2},$$

where $r$ is the radial distance from the center (by neglecting the source size). Here we neglect the geometrical obscuration of the corona emission by the accretion disk and vice versa. On the other hand, the disk radiation depends on the polar angle $\theta$ because of the source geometry. Again following Proga (2007) and Proga et al. (2008; see also Proga et al. 1998), the disk intensity $I_D$ is assumed to be radial and $I_D \propto |\cos \theta'|$. It follows that the disk radiation flux at the distance $r$ from the center can be written as

$$F_D = 2|\cos \theta'| \frac{L_D}{4\pi r^2},$$

where $\theta'$ is the angle between the disk normal and the position vector $r$ (cf. Fig. 1). The leading factor, 2, in this expression comes from the normalization of the polar angle dependency. Finally, by using equations (4), (8), and (9), the radiative acceleration term in equation (2) can be written as

$$g_{rad} = \frac{\sigma L}{4\pi c} (f + 2|\cos \theta'| f_D) \hat{r}.$$  

2.3. Precessing Disk

As we noted before, here we do not model the precession of the accretion disk itself, but rather manually force the precession. We do not specify the cause of the precession, either. We simply assume that the disk precession exists and investigate its consequences for the AGN environment. The UV-emitting portion of the disk spans from $r_i$ to $r_D$ (cf. Fig. 1). We assume that $r_D \ll r_i$, where $r_i$ is the inner radius of the computational domain of the hydrodynamic simulations. This means that the disk itself is not in the computational domain. The effect of the precessing disk is included as the precessing radiation field in the hydrodynamics of the gas (through eq. [2]).

We assume that the disk is tilted from the $x$-$y$ plane (in the Cartesian coordinate system) by $\Theta$ as in Figure 1. Equivalently, the disk angular momentum $J_D$ (assuming a flat uniform Keplerian disk) deviates from the $z$-axis by $\Theta$. Furthermore, we assume that $J_D$ precesses around the $z$-axis with precession period $P$. With these assumptions, the components of $J_D$ in the Cartesian coordinate system can be written as

$$J_{Dx} = J_D \sin \Theta \cos \frac{2\pi t}{P},$$

$$J_{Dy} = J_D \sin \Theta \sin \frac{2\pi t}{P},$$

$$J_{Dz} = J_D \cos \Theta,$$

where $t$ is the time measured from the beginning of the hydrodynamic simulations. Here we set $J_D$ to be in the $x$-$z$ plane (as shown in Fig. 1) at $t = 0$. By setting $\Theta = 0$, the model reduces to an asymmetric case as in Proga (2007). To compute the radiative acceleration as expressed in equation (10), one requires the angle between $J_D$ and the position vector $r$ at which the set of the HD equations (eqs. [1], [2], and [3]) are solved. This can be obtained simply by finding the inner product of $J_D$ and $r$. 
2.4. Model Setup

In all models presented here, the following ranges of the coordinates are adopted: $r_{i} \leq r \leq r_{o}$, $0 \leq \theta \leq \pi$, and $0 \leq \phi < 2\pi$, where $r_{i} = 500\ r_{s}$ and $r_{o} = 2.5 \times 10^{5}\ r_{s}$. The radius of the central and spherical X-ray corona region $r_{c}$ coincides with the inner radius of the accretion disk (Fig. 1). In our simulations, the polar and azimuthal angle ranges are divided into 128 and 64 zones and are equally spaced. In the $r$-direction, the grid is divided into 128 zones in which the zone size ratio is fixed at $\Delta r_{k+1} / \Delta r_{k} = 1.04$.

For the initial conditions, the density and the temperature of gas are set uniformly, i.e., $\rho = \rho_{0}$ and $T = T_{0}$ everywhere in the computational domain, where $\rho_{0} = 1.0 \times 10^{-21}\ \text{g cm}^{-3}$ and $T_{0} = 2 \times 10^{7}\ \text{K}$ throughout this paper. The initial velocity of the gas is simply set to zero everywhere.

At the inner and outer boundaries, we apply the outflow (free-to-outflow) boundary conditions, in which the field values are extrapolated beyond the boundaries using the values of the ghost zones residing outside the normal computational zones (see Stone & Norman [1992] for more details). At the outer boundary, all HD quantities (except the radial velocity) are fixed to their initial values (e.g., $T = T_{0}$ and $\rho = \rho_{0}$) during the evolution of each model. The radial velocity components are allowed to vary. Proga (2007) applied these conditions to represent a steady flow condition at the outer boundary. They found that this technique leads to a solution that relaxes to a steady state in both spherical and non-spherical accretion with an outflow (see also Proga & Begelman 2003). This imitates the condition in which a continuous supply of gas is available at the outer boundary.

3. RESULTS

3.1. Reference Values

We consider four different cases that have different combinations of the disk tilt angle ($\theta$) and the disk precession period ($P$), as summarized in Table 1. The following parameters are common to all the models presented here and are exactly the same as in Proga (2007). We assume that the central BH is nonrotating and has mass $M_{\text{BH}} = 10^{8}\ M_{\odot}$. The size of the disk inner radius is assumed to be $r_{s} = 3r_{S} = 8.8 \times 10^{11}\ \text{cm}$ (cf. § 2.4). The mass accretion rate ($M_{a}$) onto the central SMBH and the rest-mass conversion efficiency ($\eta$) are assumed to be $1 \times 10^{26}\ \text{g s}^{-1}$ and 0.0833, respectively. With these parameters, the corresponding accretion luminosity of the system is $L = 7.5 \times 10^{45}\ \text{ergs s}^{-1} = 2 \times 10^{12}\ L_{\odot}$. Equivalently, the system has the Eddington number $\Gamma = 0.6$, where $\Gamma \equiv L/L_{\text{Edd}}$ and $L_{\text{Edd}}$ is the Eddington luminosity of the Schwarzschild BH, i.e., $4\pi GM_{\text{BH}}/c^{2}$. The fractions of the luminosity in the UV ($f_{\text{UV}}$) and that in the X-ray ($f_{\text{X}}$) are fixed at 0.95 and 0.05, respectively, as in Proga (2007; see their run C).

Important reference physical quantities relevant to our systems are as follows. The Compton radius, $R_{c} \equiv GM_{\text{BH}}\mu m_{p}/kT_{C}$, is $8 \times 10^{18}\ \text{cm}$, or equivalently $9 \times 10^{4}\ r_{s}$, where $T_{C}$, $\mu$, and $m_{p}$ are the Compton temperature, the mean molecular weight of gas, and the proton mass, respectively. Here we assume that the gas temperature at infinity is $T_{\infty} = T_{C} = 2 \times 10^{7}\ \text{K}$; hence, the corresponding speed of sound at infinity is $c_{\infty} = (\gamma kT_{C}/\mu m_{p})^{1/2} = 4 \times 10^{7}\ \text{cm s}^{-1}$, where $\gamma$ is the adiabatic index. In this paper, $\gamma = 1.01$ (almost isothermal) is adopted to imitate a gas in Compton equilibrium with the radiation field. The corresponding Bondi radius (Bondi 1952) is $r_{B} = GM_{\text{BH}}/c^{2} = 4.8 \times 10^{18}\ \text{cm}$. The Bondi accretion rate (for the isothermal flow) is $M_{B} = 3.3 \times 10^{23}\ \text{g s}^{-1} = 0.52 M_{\odot}\ \text{yr}^{-1}$. The corresponding free-fall time ($t_{f}$) of gas from the Bondi radius to the inner boundary is $2.1 \times 10^{11}\ \text{s} = 7.0 \times 10^{3}\ \text{yr}$, which is about 2.3 times smaller than the precession period used for models II and III and about 23 times smaller than that of model IV (cf. Table 1).

3.2. Comparison of Axisymmetric Models in Two and Three Dimensions

Before we proceed to the main precession disk models, we briefly compare our axisymmetric model (model I) with the axisymmetric models presented earlier by Proga (2007), who used model parameters very similar to those in our model I. The main differences here are in the treatment of the radiation force and that in the radiative heating/cooling. As mentioned earlier, we set the force multiplier $M = 0$ (in eq. [4]) and the net cooling rate $C = 0$ (in eq. [3]), while Proga (2007) used nonzero values of those two terms. In our model I, the adiabatic index is set to $\gamma = 1.01$ (essentially isothermal), but their models use $\gamma = 5/3$. However, Proga (2007) found that their run A is nearly isothermal although $\gamma = 5/3$ was used (see their Fig. 1). Another important difference is the numerical codes used. Proga (2007) used the ZEUS-2D code (Stone & Norman 1992).

The overall geometry of the flow in model I (Figs. 2 and 3) is similar to those in Proga (2007). The matter accretes onto the central BH near the equatorial plane, and strong outflows occur in the polar direction. The collimation of our model is relatively weak compared to their run C, which uses exactly the same disk and corona luminosities as our model I. The wider bipolar outflow pattern seen here resembles that of their run A, which has the highest X-ray heating. The difference and the resemblance seen here are caused by the following two key factors: (1) a nearly isothermal equation of state and (2) no radiative cooling ($C = 0$) in our model. These conditions keep the gas warm everywhere in the computational domain, and the temperature is essentially that set at the outer boundary ($T_{\infty} = T_{C} = 2 \times 10^{7}\ \text{K}$). This results in a situation very similar to that in run A of Proga (2007), in which the gas temperature is also relatively high because of the high X-ray heating and cooling. The high temperature, hence the ionization state of the gas, makes the line force in their model very inefficient, resulting in the situation in which the gas is almost entirely driven by the continuum process (electron scattering) and thermal effects, just as in our model I.
Although not shown here, we have also checked the internal consistency of the ZEUS-MP (three-dimensional) code by running the axisymmetric models (model I) in both two- and three-dimensional modes. We find that the results from the two runs agree with each other in all aspects, e.g., inflow and outflow geometry, density distribution, velocity, and mass accretion and outflow rates.

3.3. Dependence on the Disk Tilt Angle $\Theta$

We now examine the model dependency on disk tilt angle ($\Theta$), while keeping all other parameters fixed. The results from models I, II, and III (cf., Table 1), which use $\Theta = 0^\circ$, $5^\circ$, and $15^\circ$, respectively, are compared for this purpose. Note that the observations suggest that a large fraction of AGNs have rather small $\Theta$, i.e., $\Theta \lesssim 15^\circ$ (e.g., Lu & Zhou 2005).

Figure 2 shows the slices of the density and velocity fields (in the $y = 0$ plane) from snapshots of our four simulations. The snapshots are chosen at the time when the models reached a (semi)steady state for models I and II. As we see later, the flow never reaches a steady state in model III; therefore, we chose the snapshot of the model at the time when the flow pattern is typical of a whole simulation time sequence. While accretion occurs mainly in the equatorial plane ($z = 0$) for model I, it occurs in an inclined plane with a pitch angle (the angle between the equatorial plane and the accretion plane) similar to the disk inclination angle, for models II and III. In the precessing disk models (II
and III), the deviation from the axisymmetric is clearly seen in both the density distribution of gas and the shapes of the Mach number contours. The corresponding three-dimensional density and Mach number contour surfaces of these models are also shown in Figures 3 and 4. The morphology of the density distribution seen in Figure 2 resembles that of the Z-shaped (for model II) and the S-shaped (for model III) radio jets (e.g., Condon & Mitchell 1984; Hunstead et al. 1984; Tremblay et al. 2006), although on different scales. Obviously, the difference between the Z- and S-shapes are simply due to the difference in the viewing angles. Unlike the MHD precessing jet models, the bending structures of the density distributions seen here are shaped by the geometry of the sonic surfaces. When accreting material from the outer boundary encounters the relatively low density but high-speed outflowing gas launched by the radiation force from the inner part, the gas becomes compressed, forming higher density regions. The flows in the bending density structure itself are rather complex (especially in model III), but the direction of the flow becomes outward (in radial direction) as they approach the sonic surface (excluding the one shaped like a disk formed by the accreting gas in the inner region). Relatively large curvatures of the flows seen in both the density and the Mach number contours of models II and III can be also understood from the fact that the precession period used in these models ($P = 16,000$ yr)
is comparable to the gas free-fall time \( t_{ff} = 7000 \text{ yr} \), cf. § 3.1. The curvatures or the twists of the weakly collimated bipolar flows can be clearly seen in the three-dimensional representation of these models in Figures 3 and 4.

We compute the mass fluxes as a function of radius for a quantitative comparison of the characteristics of the flows in the models. Following Proga (2007), the net mass flux \( \dot{M}_{\text{net}} \), the inflow mass flux \( \dot{M}_{\text{in}} \), and the outflow mass flux \( \dot{M}_{\text{out}} \) can be computed from

\[
\dot{M}(r) = \int_{s} \rho v \cdot da
\]

\[
= r^2 \int_{4\pi} \rho v_r \, d\Omega,
\]

where \( v_r \) is the radial component of the velocity \( \mathbf{v} \). In the equation above, \( \dot{M} = \dot{M}_{\text{net}} \) if all \( v_r \) are included. Similarly, \( \dot{M} = \dot{M}_{\text{in}} \) for \( v_r < 0 \) and \( \dot{M} = \dot{M}_{\text{out}} \) for \( v_r > 0 \). Also, \( da = r^2 \sin \theta \, d\theta \, d\phi \) and \( d\Omega = \sin \theta \, d\theta \, d\phi \). Similarly, we further define the outflow power in the form of kinetic energy \( (P_k) \) and that in the form of thermal energy \( (P_{\text{th}}) \) as functions of the radius, i.e.,

\[
P_k(r) = r^2 \int_{4\pi} \rho v_r^3 \, d\Omega,
\]

\[
P_{\text{th}}(r) = r^2 \int_{4\pi} \rho u v_r \, d\Omega,
\]

where \( v_r > 0 \).

Fig. 4.—As in Fig. 3, but for models III (top) and IV (bottom).
The resulting mass fluxes and the outflow powers of the models are summarized in Figure 5. In all cases, the mass inflow flux exceeds the mass outflow rate at all radii. For models I and II, the net mass fluxes ($\dot{M}_{\text{net}}$) are almost constant at all radii, indicating that the flows in these models are steady. Despite the presence of the disk precession in model II, the flow becomes steady. The density distribution and the velocity field become almost constant in the coordinate system corotating with the disk precession period. On the other hand, $\dot{M}_{\text{net}}$ for model III does not remain constant as $r$ becomes larger ($r > 10^{18}$ cm) because of the unsteady nature of the flow (cf. Figs. 2 and 4). As the disk tilt angle $\Theta$ increases, the direction of the outflows, which are normally along the polar directions (the $\pm z$-directions) unless the disk is tilted, moves toward the equatorial plane (the $x$-$z$ plane), where the flow is predominantly inward. These opposite flows makes it harder for the outflowing gas to reach the outer boundary. Furthermore, since the disk is precessing, the direction of the outflow is constantly changing. This results in continuous collisions between the inflowing and outflowing gas, especially for a larger $\Theta$ model. The net mass fluxes at the inner boundary $\dot{M}_{\text{net}}(r_i)$ are

Fig. 5.—Comparison of mass flux and energy flux as a function of radius for models I (top left), II (top right), III (bottom left), and IV (bottom right). The plot for each model is subdivided into two panels: top, mass flux and bottom, energy flux. In the mass flux plots, the inflow (dashed line; $\dot{M}_{\text{in}}$), outflow (solid line; $\dot{M}_{\text{out}}$), and net mass fluxes (dotted line; $\dot{M}_{\text{net}}$), as defined in equation (15), are separately plotted as functions of radial distance from the center. The absolute values of $\dot{M}_{\text{in}}$ and $\dot{M}_{\text{net}}$ are plotted here, since they are negative at all radii. In the energy flux plots, the kinetic energy (solid line) and the thermal energy (dotted line), defined in eqs. (16) and (17), respectively, are shown. Note that the time slices of the model simulations used here to compute the fluxes are same as those in Fig. 2.
the speed is very significant (disk tilt angle indicate inflow) decreases slightly, but not significantly as the pressure gradient force in these models; however, the displacement of the center of mass remains fairly small ($\Delta v_r = -20 \text{ km s}^{-1}$) as $\theta$ changes from $0^\circ$ to $5^\circ$.

Figure 5 also shows the outflow powers ($P_k$ and $P_{in}$) of the models as a function of radius, as defined in equations (16) and (17). As for the mass flow curves in the same figure, the dependency of the energy flux curves on radius for models I, II, and III are very similar to each other. Small but noticeable deviations of the curves for model III from those for models I and II are seen at large radii ($r > 3 \times 10^{18} \text{ cm}$). The figure shows that in all three models, the outflow power is dominated by the density (in the computational domain) caused by the (almost) isothermal equation of state and the temperature ($2.0 \times 10^7 \text{ K}$) fixed at the outer boundary. The kinetic powers or the radiation forces are not as significant as the pressure gradient force in these models, however, their importance cannot be ignored, since they shape the geometry of the outflow because they strongly depend on the polar angle position of a point in the computational field. We also note that as $\theta$ increases, the kinetic power at the outer boundary $P_k(r_o)$ decreases significantly; e.g., $P_k(r_o)$ of model III is 3 times smaller than that of the symmetric model, model I (see Table 1).

Next, we examine the degree of nonaxisymmetry in models I, II, and III, and seek for any obvious dependency on $\theta$. For this purpose, we compute the center of mass (CM) of the gas in the planes perpendicular to the $z$-axis as a function of $z$, i.e., $x_c(z)$ and $y_c(z)$, which are defined as

$$x_c(z) = \int_{r_o}^{r_i} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \delta(z - z') \rho(x, y, z') dx' dy' dz',$$

and

$$y_c(z) = \int_{r_o}^{r_i} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \delta(z - z') \rho(x, y, z') dx' dy' dz',$$

where

$$m(z) = \int_{r_o}^{r_i} \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \delta(z - z') \rho(x, y, z') dx' dy' dz,$$

and $\alpha(y, z) = (r_o^2 - z^2 - y^2)^{1/2}$. The results are shown in Figure 6. As expected, the CM position remains constant and on the $z$-axis [$x_c(z) = 0$ and $y_c(z) = 0$] for model I, as this is an axisymmetric model. For both models II and III, the maximum amount of deviation for each component of the CM ($\{x_c, y_c\}$) is about 0.3 pc, which is relatively small compared to the outer boundary radius ($r_o = 7.1$ pc). The $x_c$ and $y_c$ curves are anti-symmetric about the $z = 0$ position, since our model accretion disk, hence the radiation force, is symmetric about the origin of the coordinate system. The plot also shows that the positions of the maxima and minima in the $x_c$ and $y_c$ curves do not coincide, but rather they are shifted in both the $+z$ and $-z$ directions. This clearly demonstrates a helical or twisting nature of the flows, as one can also simply see it in the three-dimensional density and Mach number contour plots in Figures 3 and 4.

To summarize, as the tilt angle of the disk precession $\theta$ increases, reductions of the maximum outflow velocity ($v_r$) and the kinetic outflow power ($P_k$) at the outer boundary $r_o$ occur, as a consequence of the stronger interactions between the outflowing and inflowing gas as $\theta$ increases. The net mass inflow flux ($M_{net}$) at the inner boundary does not strongly depend on $\theta$. The thermal outflow energy power dominates the kinetic outflow power in our models here because of the high temperature fixed at the outer boundary and because the gas is (almost) isothermal. The flows of models II and III show helical structures; however, the radius of the helices (based on the CM positions along the $z$-axis) does not change greatly as $\theta$ increases from $5^\circ$ to $15^\circ$.

### 3.4. Dependency on the Disk Precession Period $P$

We now examine the dependency of the model on the disk precession period ($P$). We vary the value of $P$ while fixing the disk tilt angle to $\theta = 5^\circ$. For this purpose, we compare models I, II, and IV, as summarized in Table 1. The precession periods $P$ are $\infty$, $1.6 \times 10^4$, and $1.6 \times 10^5$ yr, respectively, for models I, II,
and IV. In the units of the free-fall time \(t_{ff} = 7.0 \times 10^3 \) yr from the Bondi radius (§ 3.1), they are 2, 3, and 23, respectively. Note that the observations suggest that typical values of the jet precession period are \(P = 10^4-10^6 \) yr (cf. Table 1 in Lu & Zhou 2005).

Figure 2 shows similarities between models IV and I in their morphology of the density distribution and Mach number surfaces. At a given time, the flow in model IV is almost axisymmetric, and the symmetry axis is tilted also by \(\Theta = 5^\circ\) from the z-axis. This is caused by the relatively long precession period for model IV compared to the dynamical timescale or the gas free-fall timescale \(t_{ff}\). The curvature or helical motion of the gas is not significant, and it does not greatly affect the overall morphology of the flow, except for the outermost part of the flow, where the flow is slightly turbulent due to the shear of the slowly precessing flow and the outer boundary. This can be clearly seen in the three-dimensional plots Figure 4. As the precession period \(P\) becomes shorter and comparable to \(t_{ff}\) (as in model II), the flow shows more curvature and the helical structures.

The mass flux curves (cf. eq. [15]) for model IV in Figure 5 also show that the flows in models I and IV are very similar. The overall characteristics of the curves are also similar to that of model II. In fact, the net mass flux \(\dot{M}_{\text{net}}\), the inflow mass flux \(\dot{M}_{\text{in}}\), and the outflow mass flux \(\dot{M}_{\text{out}}\) at the outer boundary of model IV are identical to those of model I (see Table 1). Also note that models I, II, and IV all have the same \(\dot{M}_{\text{in}}\) value at the inner boundary; i.e., the mass inflow rate across the inner boundary is insensitive to the change in the precession period for \(\Theta = 5^\circ\).

The similarity between models I and IV can be also seen in the outflow powers, \(P_k\) and \(P_{\text{th}}\). Figure 5 shows that \(P_k\) and \(P_{\text{th}}\) as a function of radius for model IV are almost identical to those of model I. The \(P_k\) and \(P_{\text{th}}\) values at the outer boundary are indeed identical (Table 1). A slight increase in the maximum outflow velocity at the outer boundary \(v_{r}\) for model IV, compared to model I. The kinetic outflow power \(P_k\) and the maximum outflow velocity at the outer boundary \(v_{r}\) decrease as the precession period becomes comparable to \(t_{ff}\). The CM positions \(x_\text{c}\) and \(y_\text{c}\) as a function of \(z\) (see § 3.3) for model IV are shown in Figure 6. Compared to model II, the maximum displacement of the CM is about 40 times smaller in model IV, i.e., \(|x_\text{c}| < 0.075\) pc and \(|y_\text{c}| < 0.075\) pc. The \(x_\text{c}\) curve for model IV shows a rather complex pattern compared to that in model II. This and the visual inspection of the density and the Mach number contour surfaces in Figure 4 indicate that the bipolar outflows are slightly twisted, but do not have a clear helical structure.

### 3.5. Time Evolution of Mass Accretion/Outflow Rates and Angular Momentum

Next, we examine the variability or steadiness of the flows in each model by monitoring the mass fluxes at the outer boundary, as in equation (15), and the angular momentum of the system as a function of time. For the latter, we compute the density-weighted mean specific angular momentum \(j_r\) of the system, defined as

\[
j_r = \frac{\int_V (\rho (r \times v)) dV}{\int_V \rho dV},
\]

where the denominator is simply the total mass of the gas in the computational domain. Note that the radiation force (eq. [10]) and the gravitational force are in the radial direction only. Consequently, they do not exert a torque on the system; hence, they do not directly contribute to the change in the angular momentum of the system. The system can gain the angular momentum in the following way. In our models, the strength of the disk radiation field depends on the angle measured from the disk normal (cf. eq. [9]). This causes gas pressure gradients in the azimuthal direction and contributes to the angular momentum of gas locally, forming vorticity. Because the radiation field precesses, the precessing outflow will cause the gas with the preferred sign of vorticity to escape from the outer boundary, resulting in a change in the net angular momentum of the gas in the computational domain.

Figure 7 shows \(\dot{M}_{\text{in}}(r_\text{o})\), \(\dot{M}_{\text{out}}(r_\text{o})\), and \(j_r\) for models I, II, III, and IV as a function of time. For models I, II, and IV, both the mass fluxes and \(j_r\) approach asymptotic values by \(t = 7 \times 10^{12}\) s. Small oscillations of \(j_r\) around the asymptotic values are seen for models II and IV. On the other hand, the mass fluxes of model III have much larger amplitudes for the oscillations around an asymptotic value. By visual inspection, their oscillations do not seem to have a clear periodicity. We performed the Lomb-Scargle periodogram analysis (e.g., Horne & Baliunas 1986; Press et al. 1992) on the \(\dot{M}_{\text{in}}(r_\text{o})\), \(\dot{M}_{\text{out}}(r_\text{o})\), and \(j_r\) curves for model III. Only \(\dot{M}_{\text{out}}(r_\text{o})\) shows a relatively strong signal at \(P_{\text{L-S}} = 1.36 \times 10^{12}\) s, which is about 2.7 times longer than the precession period of model III. On the other hand, the \(\dot{M}_{\text{in}}(r_\text{o})\) and \(j_r\) curves do not have any obvious period, but they are rather stochastic.

As mentioned in § 3.3, as the disk tilt angle \(\Theta\) increases, the direction of the outflows, which normally are along the polar directions, unless the disk is tilted, move toward the equatorial plane (the x-z plane), where the flow is predominantly inward. In addition, the precession of the disk causes the direction of the outflow to change constantly, causing constant creation of the shock between the inflowing and the outflowing gas. This leads to a very unstable flow of the gas at all times for models with larger \(\Theta\), e.g., model III with \(\Theta = 15^\circ\). The flow, of course, can be stabilized if the precession period is increased to a value much larger than the free-fall time \(t_{ff}\).

The amount of the (density weighted) mean specific angular momentum deposited in the gas by the precessing disk (measured by \(j_{d}\)) is largest in models II and III (see Table 1), and that in model IV is about 4 times less than those in models II and III. For all models, a time-averaged value (by using the last \(2 \times 10^{12}\) s of the simulation) of \(j_{d}\) is used in Table 1. It seems that the faster the disk precesses, the larger the amount of angular momentum transferred to the environment; however, this trend does not continue as we increase the disk precession speed even faster. Although not shown here, a model with exactly the same set of parameters as in model II, but with \(P = 1600\) yr (10 times faster rotation), showed that the value of \(j_{d}\) decreases to \(\approx 0.01\), which is even smaller than that of model IV (with \(P = 160,000\) yr). This indicates that the amount of angular momentum deposited in the gas depends on how close the precession period is to the dynamical timescale of the flow.

In principle, although it is possible to model the change in the angular momentum of the accretion disk itself through the transfer of angular momentum from the environment, we ignored this effect for simplicity (this is also a limitation of our model). To model the interaction of the disk angular momentum and the angular momentum of the surrounding gas properly, we need to model the dynamics of the gas in the accretion disk itself, as well as the dynamics of the gas, which is on a much larger scale than in our models here. This is computationally challenging with our current code, since we have to resolve the length scale of the
innermost part of the accretion disk ($\sim 10^{-2}$ pc) to the large-scale outflow/inflow gas ($\sim 10$ pc).

4. CONCLUSIONS

We have studied the dynamics of the gas under the influences of the gravity of SMBHs and the radiation force from the luminous accretion disk around the SMBH. The rotational axis of the disk was assumed to be tilted with respect to the symmetry axis, with a given angle $\Theta$ and a precession period $P$ (cf. Fig. 1). We have investigated the dependency of the flow morphology, mass accretion/outflow rates, and angular momentum of the flows for different combinations of $\Theta$ and $P$. This is a natural extension of similar but more comprehensive two-dimensional radiation hydrodynamics models of AGN outflow models by Proga et al. (2000) and Proga (2007). As this is our first attempt to model such gas dynamics in three full dimensions, we have used the reduced set of physical models described in Proga (2007); i.e., the radiation force due to line and dust scattering/absorption and the radiative cooling/heating are omitted. In the following, we summarize our main findings from this investigation.
1. Our assumption of the adiabatic index ($\gamma = 1.01$) keeps the mean temperature of the gas in the computational domain relatively high ($\sim 2 \times 10^7$ K), which is essentially determined by the outer boundary condition. For our axisymmetric model (model I; Figs. 2 and 3), this results in the flow morphology very similar to the model, with a relatively high X-ray heating (see run A in Proga 2007) in which the line force is inefficient because of the high gas temperature and hence the high ionization state of the gas.

2. Although on different scales, we were able to reproduce the Z- or S-shaped density morphology of the gas outflows (Fig. 2), which are often seen in the radio observations of AGNs (e.g., Florido et al. 1990; Hutchings et al. 1988). The bending structure seen here is shaped by the sonic surfaces. When accreting material from the outer boundary encounters the relatively low-density but high-speed outflowing gas launched by the radiation force from the inner part, the gas becomes compressed and forms higher density regions.

3. As the tilt angle of the disk precession $\Theta$ increases, the reduction of the maximum outflow velocity ($v_o$) and the kinetic outflow power $P_k$ at the outer boundary $r_o$ decrease as a consequence of the stronger interactions between the outflowing and inflowing gas (Table 1). The net mass inflow rate ($\dot{M}_{\text{net}}$) at the inner boundary does not change significantly with increasing $\Theta$.

4. A relatively high efficiency of the outflow ($\eta = \dot{M}_{\text{out}}/\dot{M}_{\text{in}}$) by the radiation pressure was observed in our models (70%-80%; see also Table 1) for an Eddington number (1) of 0.6 here. The conversion efficiency $\eta$ (from the outflow to inflow) is about the same for models I and II, but it increased slightly (~12%) for model III, which has the highest disk tilt angle.

5. The thermal outflow energy power dominates the kinetic outflow power (Fig. 5) in the models presented here because of the high temperature of the flow (as mentioned above).

6. The flows of models II and III show helical structures (cf. Figs. 3 and 4); however, the radius of the helices does not change as $\Theta$ increases from $5^\circ$ to $15^\circ$, based on the locations of the center of mass (Fig. 6) of the planes perpendicular to the axis of symmetry (the z-axis in Fig. 1). We leave it for a future investigation to test whether these trends continue as $\Theta$ becomes larger than $15^\circ$.

7. The characteristics of the flows are closely related to a combination of $P$ and $Q$, but not to $P$ and $Q$ individually. Even with a relatively large disk tilt angle $\Theta$, if the precession period is much larger than the dynamical timescale of a system, the flow geometry obviously becomes almost axisymmetric (cf. model IV in Figs. 2 and 4).

8. The gas dynamics of a model with a relative large disk tilt angle ($\Theta = 15^\circ$) with a precession period comparable to the gas free-fall time ($t_f$) of the system (e.g., model III) does not reach a steady state because the outflows driven by the luminous accretion disk constantly collide with the inflowing/accreting gas as the disk precesses, hence as the outflow direction changes.

9. The amount of the density-weighted mean specific angular momentum ($j_{\text{hel}}$) deposited by the precessing disk is largest for models II and III (Table 1; Fig. 7), for which the precession period is comparable to $t_f$.

The models represented here are mainly for exploratory purposes, to examine the basic model dependencies on $\Theta$ and $P$, and have a relatively simple set of physics, but in three full dimensions. In the follow-up paper, we will improve on our model by including the physics omitted here (the line scattering/absorption, dust scattering/absorption, and the radiative cooling/heating) as in the two-dimensional models of, e.g., Proga (2007).

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REFERENCES

Aloy, M.-Á., Martí, J.-M., Gómez, J.-L., Agudo, L., Müller, E., & Ibáñez, J.-M. 2003, ApJ, L109
Aly, J. J. 1980, A&A, 86, 192
Antonucci, R. R. J. 1984, ApJ, 278, 499
Appi, S., Sol, H., & Vicente, L. 1996, A&A, 310, 419
Arav, N., Li, Z.-Y., & Begelman, M. C. 1994, ApJ, 432, 62
Armitage, P. J., & Pringle, J. E. 1997, ApJ, 488, L47
Awaki, H., Koyama, K., Inoue, H., & Halpern, J. P. 1991, PASJ, 43, 195
Bardeen, J. M., & Petterson, J. A. 1975, ApJ, 195, L65
Begelman, M., & Nath, B. B. 2005, MNRAS, 361, 1387
Begelman, M. C., & Sikora, M. 1991, ApJ, 382, 416
Bardeen, J. M., & Petterson, J. A. 1975, ApJ, 195, L65
Begelman, M., & Nath, B. B. 2005, MNRAS, 361, 1387
Blanco, P. R., Ward, M. J., & Wright, G. S. 1990, MNRAS, 242, P4
Blandford, R. D., & Payne, D. G. 1982, MNRAS, 199, 883
Bondi, H. 1952, MNRAS, 112, 195
Brighenti, F., & Mathews, W. G. 2006, ApJ, 643, 120
Caproni, A., Livio, M., Abraham, Z., & Mosquera Cuesta, H. J. 2006, ApJ, 651, 112
Caproni, A., Mosquera Cuesta, H. J., & Abraham, Z. 2004, ApJ, 616, L99
Castor, J. I. 1970, MNRAS, 149, 111
Castor, J. I., Abbott, D. C., & Klein, R. I. 1975, ApJ, 195, 157
Ciotti, L., & Ostriker, J. P. 1997, ApJ, 487, L105
———. 2001, ApJ, 551, 131
———. 2002, ApJ, 565, 1038
Clarke, D. A. 1996, ApJ, 457, 291
Condon, J. J., & Mitchell, K. J. 1984, ApJ, 276, 472
Cox, C. I., Gull, S. F., & Scheuer, P. A. G. 1991, MNRAS, 252, 558

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