Long-Time Simulations of Astrophysical Jets: Energy Structure and Quasi-Periodic Ejection

Ahmed IBRAHIM\textsuperscript{1,2} and Kazunari SHIBATA\textsuperscript{1}

\textsuperscript{1}Kwasan Observatory, Kyoto University, 17 Ohmine-cho Kita-Kazan, Yamashina-ku, Kyoto, 607-8471
\texttt{aaeakf1@kwasan.kyoto-u.ac.jp}, \texttt{shibata@kwasan.kyoto-u.ac.jp}

\textsuperscript{2}Al-Azhar University, Faculty of Science, Department of Astronomy and Meteorology, PO Box 11884, Naser City, Cairo, Egypt

(Received 2007 January 4; accepted 2007 February 10)

Abstract

We have performed self-consistent 2.5-dimensional nonsteady MHD numerical simulations of jet formation as long as possible, including the dynamics of accretion disks. Previous simulations showed that, in the case where the calculation time of the simulations is very short as compared with the time scale of observed jets, there is no significant difference between the characteristics of the nonsteady and steady MHD simulations. Thus, we have investigated long-time evolutions of the mass-accretion rate, mass-outflow rate, jet velocity, and various energy fluxes. We found that the ejection of a jet is quasi-periodic. The period of the ejection, $T_{\text{ejection}}$, is related to the time needed for the initial magnetic field, $B_0$, to be twisted to generate a toroidal field, $T_{\text{ejection}} \propto V_A^{-1} \propto B_0^{-1} \propto E_{\text{mg}}^{-1/2}$, where $V_A$ is the Alfvén velocity and $E_{\text{mg}}$ the initial magnetic energy. We compared our results with both the steady-state theory and the previous 2.5-dimensional nonsteady MHD simulations. We found that the time-averaged velocity of the jet, $V_{\text{jet,avg}}$, is $\sim 0.1 V_K$ and $\sim 0.1 V_{\text{jet,max}}$, where $V_K$ is the Keplerian velocity at $(r, z) = (1, 0)$ and $V_{\text{jet,max}}$ the maximum velocity of the jet. Nevertheless, the characteristics of our simulations are consistent with those of the steady solution and previous short-time simulations. We found that the dependence of the time-averaged velocity and the mass-outflow rate, $\dot{M}_{\text{w,avg}}$, on the initial magnetic field are approximately $V_{\text{jet,avg}} \propto B_0^{0.3}$ and $\dot{M}_{\text{w,avg}} \propto B_0^{0.32}$, respectively.

Key words: galaxies: jets — ISM: jets and outflows — magnetohydrodynamics: MHD — methods: numerical

1. Introduction

Astrophysical jets are fast and well-collimated flows, observed in a wide variety of astronomical systems. These episodic outflows are associated with young stellar objects (YSOs) (e.g., Burrows 1996; Bachiller 1996; Reipurth et al. 2002; Curiel et al. 2006), massive black holes in active galactic nuclei (AGNs) (e.g., Biretta et al. 1995; Junor et al. 1999; Jiang 2002; Curiel et al. 2006), massive black holes in active galaxies: jets — ISM: jets and outflows — magnetohydrodynamics: MHD — methods: numerical

energy and angular momentum are removed magnetically from the accretion disk by field lines anchored to the disk surface and extending to large distances. Blandford and Payne (1982), using self-similar solutions of steady and axisymmetric MHD equations and assuming a cold Keplerian accretion disk, showed that a centrifugally driven outflow of matter from the disk is possible, if the angle between the field line and the disk surface is less than 60°.

The literature is rich with papers concerning MHD models of jet formation from accretion disks. It seems worth sketching some of the most important work: (e.g., Pudritz & Norman 1986; Sakurai 1987; Lovelace et al. 1991; Contopoulos & Lovelace 1994; Najita & Shu 1994; Fendt & Camenzind 1996). All of these studies were based on the theory of steady and axisymmetric MHD winds developed by Weber and Davis (1967) for the solar wind. Cao and Spruit (1994) examined the mass-outflow rate of magnetically driven jets, and studied the solution that passes through the slow magnetosonic point (see also Li 1995). They approved that the angle between the accretion disk and the field line is very important for achieving a high mass-outflow rate.

The acceleration and collimation of the jet have been studied in the steady state (e.g., Sauty et al. 2004; Bogovalov & Tsinganos 2005). In recent years, many simulations taking other physical processes into consideration, e.g., magnetic diffusion (Kuwabara et al. 2000, 2005; Fendt & Čemeljić 2002; Casse & Keppens 2002, 2004), the dynamo process in the accretion disk (von Rekowski et al. 2003), and the radiation force (Proga 2003), have been performed. Koide, Shibata, and Kudoh (1999) showed that the characteristics of a magnetically driven jet in a general relativistic MHD simulation are similar to those of a nonrelativistic MHD jet (Shibata & Uchida 1986). In addition to these studies that considered the initial magnetic field as being large scale, there are several papers that considered the evolution of a stellar magnetic dipole in interaction with a diffusive accretion disk. Hayashi, Shibata, and Matsumoto (1996) observed the magnetic reconnection and evolution of X-ray flares during the first rotational periods in their numerical simulations. In a one-dimensional analytically steady state solution, Kudoh and Shibata (1995, 1997a)
introduced new classifications as:

1. Magneto-centrifugally driven jets, when the magnetic field is strong; in that case they found that the jet velocity, $V_{\text{jet}}$, and the mass-outflow rate, $M_w = \frac{dM_w}{dt}$, depend on the magnetic energy, $E_{\text{magn}}$, as $V_{\text{jet}} \propto B_0^{1/3}$ and $M_w \sim \text{constant}$.

(2) Magnetic pressure-driven jets, when the magnetic field is weak; in that case $V_{\text{jet}} \propto B_0^{1/6}$ and $M_w \propto B_0^{0.5}$; this was confirmed by Ustyugova et al. (1999).

In most of theoretical models of jets from accretion disks, however, accretion disks are treated as boundary conditions. Accretion disks only play a key role in supplying energy and mass to the jets, and neither the accretion flow nor the internal structure of disks are considered in these models. Since the disk itself is not treated, such simulations may last over hundreds of Keplerian periods. This idea was first applied by Ustyugova et al. (1995). Extending this work, Romanova et al. (1997) found a stationary final state of a slowly collimating disk wind in the case of a split-monopole initial field structure after 100 Keplerian periods. Ouyed and Pudritz (1997a,b) presented time-dependent simulations of the jet formation from a Keplerian disk. For a certain (already collimating) initial magnetic field distribution, a stationary state of the jet flow was obtained after about 400 Keplerian periods of the inner disk with an increased degree of collimation. Ouyed, Clarke, and Pudritz (2003) investigated the problem of jet stability and magnetic collimation by extending the axisymmetric simulations to 3D.

On the other hand, the first disk and wind numerical calculations were published by Uchida and Shibata (1985). These pioneering simulations were carried out for magnetized disks with a sub-Keplerian rotation, designed to mimic the collapse of an initially more extended and slowly rotating object. Their models developed a rapid radial collapse in which the initially poloidal field threading the disk winds up due to differential rotation of the collapsing disk. The vertical Alfvén speed being smaller than the free-fall collapse speed implies that a strong, vertical toroidal field pressure gradient, $\partial (B^2_z/8\pi)/\partial z$, must rapidly build up. The resulting vertical pressure gradient results in the transient ejection of coronal material above and below the disk as the spring uncoils. These outflows, while very suggestive, represent more transient responses to the initial state. Shibata and Uchida (1986) investigated the detailed properties of these jets. They found that the velocity of the jet was typically on the order of the disk’s Keplerian velocity, and increased with increasing the magnetic field strength in the same manner as the scaling law of Michel’s (1969) solution. Matsumoto et al. (1996) carried out 2D MHD simulations of a torus threaded by poloidal magnetic fields. They confirmed the results obtained by Shibata and Uchida (1986). Kudoh, Matsumoto, and Shibata (1998) studied the formation mechanism of jets from geometrically thin disks and the dependance of the initial magnetic field strength, $B_0$, in detail by performing self-consistent 2.5-dimensional, nonsteady, ideal MHD simulations, including the dynamics of the disk. They found that the ejection point is determined by the effective potential resulting from the gravitational and centrifugal forces along the field lines (Blandford & Payne 1982), and also that the velocity and mass-outflow rate are consistent with those predicted by the steady theory (Kudoh & Shibata 1995, 1997a).

Kato, Kudoh, and Shibata (2002) performed 2.5-dimensional, axisymmetric, ideal MHD simulations of jets from geometrically thin disks for Keplerian and sub-Keplerian cases over a wide range of initial magnetic field strengths. Kigure and Shibata (2005) investigated the problem of jet formation and stability by using 3-dimensional MHD simulations. To investigate the stability of an MHD jet, they introduced a perturbation to the accretion disk with a nonaxisymmetric sinusoidal or random fluctuation of rotational velocity. In both perturbation cases, a nonaxisymmetric structure with $m = 2$ appears in the jet, where $m$ is the azimuthal wavenumber. They concluded that this structure seems to originate in the accretion disk.

We note that all calculations including the treatment of the disk structure and the ideal MHD (e.g., Shibata & Uchida 1986, 1987, 1989, 1990; Matsumoto et al. 1996; Kudoh et al. 1998, 2002; Kato et al. 2002; Kigure & Shibata 2005) could be performed only for 1–2 Keplerian periods of an inner disk. At this point, we emphasize that the observed kinematic time scale of protostellar jets can be as large as $10^3–10^4$ yr, corresponding to $5 \times 10^5–5 \times 10^6$ stellar rotational periods (and inner disk rotations)! For example, proper-motion measurements for HH 30 jet (Burrows et al. 1996) give a knot velocity of about $100–300$ km s$^{-1}$ and a knot production rate of about 0.4 knots per year. Assuming a similar jet velocity along the whole jet extending to 0.25 pc (López et al. 1995), the kinematic age is about 1000 yr.

In order to have access to the observed time scale of jets, and to know whether the jet formation becomes a quasi-steady state, we performed long-time 2.5-dimensional MHD simulations of jets.

We solved the dynamics of the disk as in Shibata and Uchida (1986), Kudoh, Matsumoto, and Shibata (1998), and their following studies. We also wanted to know whether the time-averaged physical quantities have the same characteristics as those in the steady model and previous simulations (Kudoh et al. 1998; Kato et al. 2002). Our calculation time was about 20 times longer than those of previous simulations. Also, the simulation box was sufficiently large to minimize the effect of the top and side boundary conditions.

2. Numerical Method

2.1. Assumptions and Basic Equations

We have used the same numerical techniques as used by Kudoh, Matsumoto, and Shibata (1998). We solved the following ideal MHD equations numerically in a 2.5-dimensional approximation. The ideal MHD equations in cylindrical coordinates $(r, \varphi, z)$ are:

$$\frac{\partial \rho}{\partial t} + v_r \frac{\partial \rho}{\partial r} + v_z \frac{\partial \rho}{\partial z} = -\rho \left[ \frac{\partial}{\partial r} \left( rv_r \right) + \frac{\partial v_z}{\partial z} \right],$$

$$\frac{\partial v_r}{\partial t} + v_r \frac{\partial v_r}{\partial r} + v_z \frac{\partial v_z}{\partial z} = -\frac{1}{\rho} \frac{\partial p}{\partial r} + \frac{v_z^2}{r} - \frac{\partial \Psi}{\partial r} - \frac{1}{4\pi \rho} \times \left[ B_x \frac{\partial B_z}{\partial r} + B_y \frac{\partial B_z}{\partial r} \right] + \frac{B_z}{4\pi \rho} \frac{\partial \Psi}{\partial z},$$

(2)
\[
\frac{\partial (rv_r)}{\partial t} + v_r \frac{\partial (rv_r)}{\partial r} + v_z \frac{\partial (rv_r)}{\partial z} + v_r \frac{\partial (rv_r)}{\partial r}
= \frac{1}{4\pi \rho} \left[ B_r \frac{\partial}{\partial r} (r B_r) + B_z \frac{\partial}{\partial z} (r B_r) \right].
\]

(3)

where \( r_0 \) is the unit length (see table 1). At the location where \((r, z) = (r_0, 0)\), the density of the disk is maximum. The parameter \( \alpha \) is defined as \((\gamma V_{K0}^2 / V_{sc}^2)\), where \( V_{sc} \) is the sound velocity in the corona, \( V_{K0} = (GM / r_0)^{1/2} \) the Keplerian velocity at radius \( r_0 \), and \( \rho_c \) the coronal density at radius \( r_0 \). We used \( \alpha = 1 \) and \( \rho_c / \rho_0 = 10^{-3} \) throughout this work, where \( \rho_0 \) is the initial density at \((r, z) = (r_0, 0)\). The initial magnetic field was assumed to be uniform and parallel to the rotation axis of the disk; \((B_r, B_\theta, B_z) = (0, 0, B_0)\).

Indeed, the presence of a large-scale magnetic field in accretion-disk-jet systems is observed in AGN jets and protostellar jets (Ray et al. 1997; Gabuzda et al. 2004; Pushkarev et al. 2005; Vlemmings et al. 2006), and is assumed theoretically in jet-launching models (e.g., Casse & Keppens 2004, and references therein).

2.3. Boundary Conditions

The symmetry is imposed for the variables \( \rho, p, v_r, v_z \), and \( B_z \) around the axis \((r = 0)\), while antisymmetry is imposed for the variables \( v_r, v_\theta, B_r, B_\theta \). The side \((r = R_{\text{max}})\), top \((z = Z_{\text{max}})\) and bottom \((z = Z_{\text{min}})\) surfaces are assumed to be free boundaries. In order to avoid a singularity at the origin, the region around \( r = z = 0 \) is treated by softening the gravitational potential as

\[
\Psi = \begin{cases} 
-GM / (r^2 + z^2)^{1/2} & \text{for } \epsilon < (r^2 + z^2)^{1/2}, \\
-GM \left(1/\epsilon - [(r^2 + z^2)^{1/2} - \epsilon] / \epsilon^2 \right) & \text{for } 0.5 \epsilon < (r^2 + z^2)^{1/2} \leq \epsilon, \\
-1.5GM / \epsilon & \text{for } (r^2 + z^2)^{1/2} \leq 0.5 \epsilon,
\end{cases}
\]

(14)

where \( \epsilon \) was equal to 0.2 \( r_0 \) throughout this work.

These boundary conditions are different from those of MHD simulations of jets from accretion disks performed by Ouyed and Pudritz (1997a), Romanove et al. (1997), Meier et al. (1997), Ustyugova et al. (1999), and Pudritz, Rogers, and Ouyed (2006). They assumed an inflowing fixed boundary condition (i.e., the angular momentum is continually injected from the boundary). Therefore, their numerical simulations allowed no disk accretion due to magnetic braking, nor the growth of a magneto-rotational instability.

2.4. Numerical Schemes

Our numerical simulations were carried out using the cubic-interpolated pseudoparticle (CIP) method (Yabe & Aoki 1991; Yabe et al. 1991) and the method of characteristics-constrained transport (MOCCT) (Evans & Hawley 1988; Stone & Norman 1992). We used MOCCT to solve the magnetic-induction equation, while we used CIP to solve the others. A full description of the CIP-MOCCT scheme is given in Kudoh, Matsumoto, and Shibata (1999). We used the two-dimensional version of a scheme used by Kudoh and Shibata (1997b).

2.5. Parameters

The physical quantities in our simulations were normalized with their typical initial values at \((r, z) = (r_0, 0)\), where the
initial density of the disk had the maximum value of \( \rho_0 \) at 
\((r, z) = (r_0, 0)\). The description of the normalized units for each 
variable is summarized in table 1. To set the initial conditions 
we need to describe two nondimensional parameters:

\[
E_{\text{th}} = \frac{V_{s0}^2}{\gamma V_{K0}^2},
\]

\[
E_{\text{mg}} = \frac{V_{A0}^2}{V_{K0}^2}.
\]

Here, \( E_{\text{th}} \) is the initial thermal energy, which was constant 
throughout the present work, at a value of 0.018, \( V_{s0} = (\gamma \rho_0 / \rho_0)^{1/2} \) the sound velocity, \( V_{A0} = B_0/(4\pi \rho_0)^{1/2} \) the 
Alfvén velocity, and \( \rho_0 \) the initial pressure at \((r, z) = (r_0, 0)\).

Table 2 includes the initial parameters that we used in this work. 
In order to let our simulations continue as long as possible, we adopted \( L = L_0 \rho^{0.45} \) (\( L_0 = 1.00 \)) and \( \alpha = 1 \), which 
made the disk thicker and wider than that of Kato, Kudoh, and 
Shibata (2002). From a numerical point of view, thick disks 
are better resolved than thin ones, because they include much 
matter, which makes the time of simulations longer, and their 
internal structures are easier to study.

The size of our computational domain, 
\((R_{\text{min}}/r_0 - R_{\text{max}}/r_0, Z_{\text{min}}/r_0 - Z_{\text{max}}/r_0) = (0.0-24.8, -65.75-65.75)\), was 
widthier than in previous simulations in order to minimize the 
boundary condition effects. In this domain, the grid spacing is 
divided uniformly, \((\Delta r, \Delta z) = (0.05, 0.01)\), for the region \( r/r_0 < 1 \) and 
\( z/r_0 < 1 \), and stretched in \( r \) and \( z \) for \( r/r_0 > 1 \) or \( z/r_0 > 1 \). 
The minimum grid size is 0.05\( r_0 \) in the \( r \)-direction and 0.01\( r_0 \) 
in the \( z \)-direction. The maximum grid size in the \( r \)-direction 
is 0.1\( r_0 \) and 0.4\( r_0 \) in the \( z \)-direction; the number of grid points 
used in this work was 269 \times 639.

3. Numerical Results

The time evolutions of the density and temperature \((T = \gamma p/\rho)\) 
distributions of the models 1 \((E_{\text{mg}} = 2 \times 10^{-4})\), 
4 \((E_{\text{mg}} = 2 \times 10^{-5})\), and 6 \((E_{\text{mg}} = 2 \times 10^{-6})\) are shown in 
figures 1, 2, and 3, respectively. The arrows describe the 
poloidal velocity and the white lines the poloidal magnetic 
field. The time unit, \( t = 2\pi \simeq 6.28 \), corresponds to one 
Keplerian orbit at \((r, z) = (1, 0)\). We notice that the evolutions 
for all models that look similar to each other.

We now show the case of model 4 (see figure 2). In the 
first Keplerian orbit of the disk, the magnetic field lines are 
twisted by disk rotation, and the toroidal field is continuously 
generated at the disk surface. That field propagates up into 
the corona along the large-scale magnetic field as a torsional 
Alfvén wave \((t = 13)\). The magnetic field lines inside the disk 
are deformed by the disk rotation. The deformed magnetic 
field brakes the disk rotation, which leads to extract angular 
momentum from the disk matter. The disk matter falls to the 
central object, while the Alfvén wave transports the angular 
momentum to the corona. The surface layer is more affected 
by the loss of angular momentum than the other layers inside 
the disk, so that the disk matter falls faster than the equatorial 
part. This is avalanche-like accretion that was studied inten-
sively by Matsumoto et al. (1996). Because only a small frac-
tion of the accreted matter is ejected into the bipolar direc-
tion due to the Lorentz force in the relaxing twist, both the 
density and the pressure of the inner region increase. Within 
this region, the angular momentum is transferred from the high 
to the low-density parts. The disk matter on the disk surface 
is ejected as a jet \((t = 35)\). At \( t = 82 \), due to the magnetorota-
tional instability the channel flow becomes clear. Initially, both 
the gas pressure and the magneto-centrifugal force drive and 
accelerate the outflow (below the Alfvén surface). After that, 
when the toroidal field generated as the result of differential 
rotation of the accretion disk is accumulated, the acceleration 
is due to the magnetic pressure gradient, \( \partial(B_z^2/8\pi)/\partial z \). To 
illustrate the twist level of the magnetic field lines, we show in 
figure 4 the time variation of the ratio of the toroidal magnetic 
field, \( B_z \), to the poloidal magnetic field, \( B_p \), along \( r = 0.225 \). 
Figure 4 illustrates that in the case of an initial weak magnetic 
field the twisting field, \( B_z \), becomes more dominant, and both 
the Alfvén point and the slow point become near to the accre-
tion disk (Pelletier & Pudritz 1992; Kudoh & Shibata 1997a). 
In both weak and strong initial magnetic fields a strong twist 
appears at \( t = 15 \), and propagates outward when the jet is 
ejected.

The outflow consists of both the material that is initially in 
the corona and the material from the disk. The channel flow 
continues to grow in the disk, and jets are ejected continu-
ously and intermittently \((t = 82)\). Jet ejection and accretion still 
continue at the last stage of evolution in our simulations \((t = 115)\). We can also see that the magnetic field lines entwine with
Fig. 1. Temporal evolution of an ideal accretion disk threaded by a poloidal magnetic field. The color levels represent the density level (upper panels) and temperature level (middle panels), while the solid lines stand for poloidal magnetic field lines. The time unit labeling each snapshot, \(t = 2\pi \sim 6.28\), corresponds to one Keplerian orbit at \((r, z) = (1, 0)\). After a few rotations, outflows are escaping from the disk, and the outflow and accretion remain until the last moment of our simulation time. The square refers to the region analyzed by Kudoh, Matsumoto, and Shibata (1998).

Each other, and that magnetic turbulent flow develops. We can see that magnetic islands are created by the magnetic reconnection in the disk. Since we assumed ideal MHD, the magnetic reconnection is caused by numerical diffusion.

4. Time Evolutions of Some Physical Quantities of Jets

One of the most important purposes of this study was to clarify whether the jet ejection becomes a steady state. In figures 5, 6, and 9 we show the time evolutions of the mass-outflow rate, \(M_w\), the mass-accretion rate, \(\dot{M}_a\), and the toroidal magnetic-field energy, \(E_{mg}\), for model 1 (\(E_{mg} = 2 \times 10^{-4}\)), model 4 (\(E_{mg} = 2 \times 10^{-5}\)), and model 6 (\(E_{mg} = 2 \times 10^{-6}\)).

\[
M_w = \pi \left[ \int_0^1 \rho v_x(r, z = z_p) r dr - \int_0^1 \rho v_x(r, z = -z_p) r dr \right]
\]

at \(z = 4\),

\[
\dot{M}_a = -\pi \int_{-1}^1 \rho v_r(r = 1.0, z) r dz.
\]

where \(z_p\) is the \(z\)-coordinate for jets to pass through. The mass-accretion rate for various models, \(\dot{M}_a\), is defined by

It is not easy to define \(M_w\) and \(\dot{M}_a\), for the following reasons: (1) ejection and accretion are very nonsteady, and sometimes the ejected matter falls back onto the disk; (2) when the initial magnetic field is weak, the jet-ejection point goes away from the equatorial plane after a long-time simulation. Kuwabara et al. (2000) and Kudoh, Matsumoto, and Shibata (2002) found that mass accretion and mass ejection take place intermittently if there is no diffusion in the disk. The simulation times of both of them ranged from one orbit in the case of Kuwabara et al. (2000) to three orbits in the case of Kudoh, Matsumoto, and Shibata (2002). Our simulations included more than 20 orbits.
Fig. 2. Temporal evolution of an ideal accretion disk threaded by a poloidal magnetic field. The color levels represent the density level (upper panels) and temperature level (lower panels) while the solid lines stand for poloidal magnetic field lines. The time unit labeling each snapshot, \( t \sim 2\pi \sim 6.28 \), corresponds to one Keplerian orbit at \( (r, z) = (1.0) \). After a few rotations, outflows are escaping from the disk, and the outflow and accretion remain until the last moment of our simulation time.

Fig. 3. Temporal evolution of an ideal accretion disk threaded by a poloidal magnetic field. The color levels represent the density level (upper panels) and temperature level (lower panels), while the solid lines stand for poloidal magnetic filed lines. The time unit labeling each snapshot, \( t \sim 2\pi \sim 6.28 \), corresponds to one Keplerian orbit at \( (r, z) = (1.0) \). After a few rotations, outflows are escaping from the disk, and the outflow and accretion remain until the last moment of our simulation time.
Fig. 4. Ratio of $B_\phi$ to $B_p$ along $r = 0.225$ at three different time shots, for three magnetic field energies. The twisting magnetic field becomes more prominent in the case of a weak magnetic field.

for different models. From figure 5, it is clear that the mass-ejection flux is still intermittent until the last stage of evolution. The intermittent ejection is clear in all models in figure 5, but there are some differences among them. The first one, in the case of a strong magnetic field, the intermittency at the beginning of the simulation is almost similar to that at the end. However, in the case of a weak field, as the simulation goes on, the ejection-mass flux becomes very nonsteady. Not only does the intermittency increase in the case of a weak magnetic field, but also the absolute amount of the mass flux increases. However, in the case of a strong magnetic field we notice the opposite situation. Figure 6 shows the accretion rates for models 1, 4, and 6. The mass accretion, like mass ejection, was intermittent until the last stage of our simulations. In both weak and strong initial magnetic fields the mass accretion was intermittent. The general trend of the intermittency of the mass accretion is similar to that of the mass ejection. That similarity is clearer, for example, in model 4 than in the other models. In model 4 the absolute value of mass accretion at the beginning of the simulation is smaller than that at the end. We notice such a trend in the mass ejection; the same trend also occurred in the other models.

Because the mass-accretion rate is highly variable, the plot at one radius in time shows only a little of overall characteristics of the accretion. We therefore show values that were averaged over space or time, in order to obtain a better understanding of the overall accretion within the disk.

Figure 7 shows the accretion rate against time, averaged between $r = 0.52$ and $r = 10$. This region was chosen since the accretion has a small value at larger radii. This shows that even though the accretion rate is highly intermittent, it is predominantly positive. The behavior of the accretion is more complex in the case of a weak initial magnetic field after 10 orbits, because the evolution of the disk takes a long time until it becomes more turbulent. The accretion, in the case of a strong initial magnetic field, is positive during the first 5 orbits. During that time, part of the subtracted angular momentum is ejected as a jet. The other part is transported to large distances by the magnetic stresses. This is exactly what we see in figure 8, which shows the accretion rate as a relation of the radius of the disk, averaged over the first 10 orbits. Accordingly, we can see that the accretion is positive in the inner part of the disk and the accretion is negative beyond a radius of 2–4. The negative accretion continues for about 5 orbits. The accretion becomes positive again because of continuous ejection of the jet, which means continuous subtraction of the angular momentum.

Figure 9 shows the toroidal magnetic energy for models 1, 4, and 6. The toroidal magnetic field energy, $E_{mg}$, for various models is defined by

$$E_{mg} = \pi \left[ \int_0^1 \frac{B^2_\psi (r, z = z_0)}{8\pi} r dr \right] + \frac{1}{\pi} \int_0^1 \frac{B^2_\psi (r, z = -z_0)}{8\pi} r dr. \quad (19)$$

The intermittency is clear and the trend is also similar to that of the mass ejection. Hence, we tried to find the relation between the mass accretion, ejection, and magnetic field. The ejections are intermittent and seem to have periods for all models, and the periodicity seems to be different in those models with

---

Fig. 5. Time variation of the mass-outflow rate for three models; the outflow is still intermittent after a long-time simulation for different values of the initial magnetic field energy.
Fig. 6. Time variation of the mass-accretion rate for three models; the accretion is still intermittent after a long-time simulation for different values of the initial magnetic field energy. One Keplerian orbit at \((r, z) = (1.0) \sim 2\pi \sim 6.28\) time unit.

Fig. 7. Accretion rate, averaged between \(r = 0.52\) and \(r = 10\), against time for different values of the initial magnetic field energy. One Keplerian orbit at \((r, z) = (1.0) \sim 2\pi \sim 6.28\) time unit.
different initial magnetic fields.

Next, we checked the times of the ejection peaks for different models, and studied the relation between the peak time and the time needed for the toroidal field to be accumulated and untwisted in the vertical direction by carrying the mass flux. From the initial conditions, \( B_z = B_0 = \text{constant} \) and \( B_r = B_\phi = 0 \). With the rotation of the disk by the angular velocity, \( \Omega \), the toroidal field, \( B_\phi \), is generated from \( B_z \).

\[
\frac{\partial B_\phi}{\partial T} \approx \Omega B_z. \tag{20}
\]

By integrating equation (20),

\[
B_\phi \approx \Omega B_z T. \tag{21}
\]

The ejection occurs when the magnetic energy equals the rotational energy, then

\[
\frac{B_\phi^2}{8\pi} \approx \frac{1}{2} \rho V_\phi^2 \approx \frac{1}{2} \rho V_K^2. \tag{22}
\]

A combination of equations (21) and (22) yields:

\[
\frac{\Omega^2 B_z^2 T^2}{8\pi} \approx \frac{1}{2} \rho V_\phi^2 \approx \frac{1}{2} \rho V_K^2, \tag{23}
\]

\[
T_{\text{ejection}} \approx \sqrt{\frac{4\pi\rho V_K}{\Omega B_z}} \approx \frac{1}{V_A} \frac{V_K}{\Omega}. \tag{24}
\]

\[
T_{\text{ejection}} \propto \frac{1}{V_A} \propto \frac{1}{B_0} \propto E_{\text{mg}}^{-\frac{1}{2}}. \tag{25}
\]

where \( V_A \) is the Alfvén speed, \( B_0 \) the strength of the initial magnetic field, and \( T_{\text{ejection}} \) the expected theoretical ejection time corresponding to the peak of the mass-ejection rate, as shown in figure 10. Figure 11 shows the relation of the initial magnetic field strength and the average time interval between the peaks of the mass-ejection rate of models 1, 3, 4, and 6. In figure 11 the squares correspond to our numerical values and the solid line is the best fit to them with the following dependence:

\[
T_{\text{aver,ejec}} \propto E_{\text{mg}}^{-0.3}. \tag{26}
\]

The dotted line corresponds to the analytical relation, equation (25). We notice that the dependence of the ejection time on the initial magnetic field is in agreement with our analytical expectation, equation (25). The small deviation from the analytical relation comes from our assumption that both the density and the magnetic field will remain constant during the simulation. However, the situation is that both the density and the magnetic field increase with time during the simulation. By taking the time evolution of both the density and the magnetic field into account (inside the disk), we obtained from equation (24)

\[
T_{\text{ejection}} \propto \sqrt{\frac{\rho}{B_z}} \propto \frac{1}{V_A}. \tag{27}
\]

Both the density and the \( B_z \) component of the magnetic field increase steeply in the case of a strong initial magnetic field. On the other hand, the density almost remains constant in the case of a weak initial magnetic field, while \( B_z \) increases. Then, as \( B_z \) increases with time, the ejection time becomes shorter than that in the case of a constant initial magnetic field. However, the time evolution of the density will make an opposite effect in the case of a strong initial magnetic field, i.e., the effect of the increase in \( B_z \) will be weakened by the effect of the increase in \( \rho \). On the other hand, in the case of a weak initial magnetic field, the effect of an increase of the density is too
small to weaken the effect of an increase of the magnetic field. Hence, the evolution effect of $B_z$ is more prominent in the case of a weak initial magnetic field. We thus think that the evolution of both the density and the magnetic field may explain the discrepancy between the analytical dependence of the ejection time and the initial magnetic field. Consequently, when we consider the average magnetic field instead of the initial magnetic field in figure 11a, the agreement is better, as we can see from figure 11b. In this case the best fit is

$$T_{\text{avg,ejec}} \propto E_{m_\text{g}}^{-0.36}.$$  

\section{5. Dependences of the Averaged Velocities and Mass Outflow on the Initial Magnetic Field Strength}

The previous simulations (e.g., Kudoh et al. 1998; Kigure & Shibata 2005) pointed out that both of the mass-outflow rates and the velocities of the magnetically driven jets depend on the strength of the magnetic field. As we mentioned in the introduction of this paper, the previous simulations could be performed for 1–2 Keplerian periods of the inner disk. What is the relation between the jet velocity and the magnetic field strength if we performed long simulations? Long simulations give us the ability to study the average values for the jet velocity, mass-accretion rate, mass-ejection rate, and the toroidal magnetic field with the initial magnetic field strength, instead of the maximum values. Figure 12a illustrates the time-averaged jet velocity, $V_{\text{jet,avg}}$, as a function of the initial magnetic field, $E_{m_\text{g}}$. We define $V_{\text{jet,avg}}$ by

$$V_{\text{jet,avg}} = \frac{1}{T_{e}} \int_{0}^{T_{e}} \int_{0}^{4} v_z(r,z) dr dt \left/ 10T_{\text{orbit}} \right.,$$  

where $T_e$ is the end of the simulation time, $T_{\text{orbit}}$ the time of one disk rotation, and $10T_{\text{orbit}}$ the sum of the time over which

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure9.png}
\caption{Time variation of the toroidal energy for three models; the outflow is still intermittent, similar to mass outflow, after a long-time simulation for different values of the initial magnetic field energy. One Keplerian orbit at $(r,z) = (1,0) \sim 2\pi \sim 6.28$ time unit.}
\end{figure}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure10.png}
\caption{Mass flux for four models as a function of the time. The number corresponds to the maximum flux at which we measure the corresponding time (peak time) and plot this time with the initial magnetic field in figure 11.}
\end{figure}
Fig. 11. We show the relations of the average time interval between the mass-ejection peaks to the initial magnetic field strength (a), and to the average magnetic field (b). The solid line corresponds to the best fit to the numerical result and the dotted line to the analytical relation [equation (25)].

Fig. 12. (a) Time-averaged velocities, $V_{\text{jet, avg}}$, as a function of $E_{\text{mg}}$. (b) $V_{\text{jet, max}}$ as a function of $E_{\text{mg}}$. (c) Time-averaged velocities as a function of $E_{\text{mg}}/(\dot{M}_w)$. (d) $V_{\text{jet, max}}$ as a function of $E_{\text{mg}}/(\dot{M}_w)$.
the speed of the jet is averaged. The figure shows that the averaged-jet velocity is on the order of a tenth of the Keplerian speed for a wide range of $E_{mg}$ in the disk, and that its dependence on $E_{mg}$, as shown in figure 12a, is approximately

$$V_{jet,\text{avg}} \propto E_{mg}^{0.15}. \quad (30)$$

Figure 12b illustrates the maximum-jet velocities, $V_{jet,\text{max}}$, as a function of $E_{mg}$. The solid line shows

$$V_{jet,\text{max}} \propto E_{mg}^{0.17}. \quad (31)$$

Figures 12c and 12d illustrate the time-averaged jet velocities and the maximum jet velocities, respectively, as a function of $E_{mg}/M_w$. The solid line shows

$$V_{jet,\text{avg}} \propto (E_{mg}/M_w)^{0.16}, \quad (32)$$

$$V_{jet,\text{max}} \propto (E_{mg}/M_w)^{0.17}. \quad (33)$$

We notice that (in the case of a weak magnetic field) the two relations are consistent with Michel’s steady solution, $v_z \propto (B_0^2/M_w)^{1/3}$, although the jet and accretion do not reach the steady state.

Figure 13a illustrates the time-averaged mass-accretion rates of the disk as a function of $E_{mg}$. The figure shows that its dependence on $E_{mg}$ is about

$$<\dot{M}_a> \propto E_{mg}^{0.6}. \quad (34)$$

Figure 13b illustrates the time-averaged mass-ejection rates of our jets as a function of $E_{mg}$. Its dependence on $E_{mg}$ is approximately

$$<\dot{M}_w> \propto E_{mg}^{0.16}. \quad (35)$$

Figure 13c illustrates the time-averaged toroidal magnetic energy as a function of $E_{mg}$. Its dependence on $E_{mg}$ is approximately

$$<E_{mgt}> \propto E_{mg}^{0.5}. \quad (36)$$

Figure 13d illustrates the ratio of the time-averaged mass-outflow rate of the jet to the time-averaged mass-accretion rate of the disk, $<M_w>/<M_a>$, as a function of $E_{mg}$.

$$<M_w>/<M_a> \propto E_{mg}^{-0.4}. \quad (37)$$

5.1. Long-Time Evolution of the Jet Velocity

Figure 14 shows the long-time evolution of the jet velocity, which is calculated through the jet region $r=0-1$. Figure 14a shows the averaged jet velocity and figure 14b the maximum jet velocity. Both of them never reach a steady state, and have the same characteristics of the variation. The jet velocity in the case of a strong initial magnetic field reaches the maximum value earlier than in the weak filed case, but after the first five orbits the jet velocities for all initial magnetic fields are similar. After the first ten orbits the jet velocity for the weak initial magnetic field becomes higher than that for the strong one.

We notice that in all models the jet velocity severely decreases after the first three orbits. We think that the high jet velocity in the first stage of evolution results from both effects of initial and boundary conditions. Also, at the first stage of the simulation the ejection point of the jet is near to
the gravitational center. With the simulation going the density and pressure of the central region of the disk increase, which leads to a decrease or stopping of accretion within that region. As a result, the ejection point of the jet becomes farther from the gravitational center, so that the jet velocity decreases. Also, the effect of a softening gravitational potential near to the gravitational center leads to a decrease of the jet velocity. While the simulation is running, the magnetic field lines are trapped inside the softening gravitational potential. Consequently, they lose their angular momentum, and their rotational velocities become much smaller, which leads to a decrease of the jet velocity.

5.2. Jet Driving Forces

Figure 15 shows the time evolution of powers of the jet, i.e., the Poynting flux, $F_{pj}$; the kinetic flux, $F_{kj}$; and the enthalpy flux, $F_{en,j}$, which are described as:

$$F_{pj} = \int_0^1 2\pi r \frac{c}{4\pi} (E \times B)_z \, dr \quad \text{at} \quad z = 4,$$

$$F_{kj} = \int_0^1 2\pi rv^2 v_z \, dr \quad \text{at} \quad z = 4,$$

$$F_{en,j} = \int_0^1 2\pi rv^2 \left( \frac{\gamma - 1}{\gamma} \right) \, dr \quad \text{at} \quad z = 4,$$

where $z$ is the $z$-coordinate for the jet to pass through.

The Poynting flux, $F_{pj}$, plays the dominant role in driving and accelerating the jet during the initial evolution of our simulation in the case of a strong magnetic field, model 1, $E_{mg} = 2 \times 10^{-4}$, until $t = 50$; after that it suddenly becomes very weak.

In the case of a weak initial magnetic field, model 5, $E_{mg} = 5 \times 10^{-6}$, during the initial stage of the evolution, the enthalpy flux is the dominant one until $t \sim 50$, and both the kinetic flux and the Poynting flux have nearly the same values. After $t \sim 50$, the dominant energy flux is the Poynting flux, and the other two fluxes have nearly the same values. At the last stage of evolution, the Poynting flux decreases again and the enthalpy flux becomes more dominant. In the case of a very weak initial magnetic field, model 9, $E_{mg} = 2 \times 10^{-7}$, the enthalpy flux is the most dominant until the last stage of the simulation, except after $t \sim 225$ when it decreases sharply.

Figure 16 shows the relation between the average enthalpy flux, $F_{en,j,avg}$, the kinetic flux, $F_{kj,avg}$, the Poynting flux, $F_{pj,avg}$, and the initial magnetic field strength. The average fluxes are defined as:

$$F_{pj,avg} = -\int_0^T \int_0^1 2\pi r \frac{c}{4\pi} (E \times B)_z \, dr \, dt \int 10T_{crit} \quad \text{at} \quad z = 4,$$

$$F_{kj,avg} = \int_0^T \int_0^1 2\pi rv^2 v_z \, dr \, dt \int 10T_{crit} \quad \text{at} \quad z = 4,$$

$$F_{en,j,avg} = \int_0^T \int_0^1 2\pi rv^2 \left( \frac{\gamma - 1}{\gamma} \right) \, dr \, dt \int 10T_{crit} \quad \text{at} \quad z = 4,$$

where $T_e$ is the end of the simulation time and $10T_{crit}$ the sum of integration.
of the time over which the flux of the jet is averaged. Figure 16 shows that the averaged enthalpy flux has the same value whatever is the initial magnetic field strength, whereas both the averaged kinetic flux and the averaged Poynting flux increase with the strength of the initial magnetic field.

Kudoh and Shibata (1997a) showed that the dominant energy of a jet depends on the strength of the magnetic field. When the poloidal component of the magnetic field is $B_p \propto r^{-2}$, the fast magnetosonic point appears far from the Alfvén point and the dominant energy of the jet is the Poynting flux. In our simulations, the initial magnetic field is uniform. In such models, the fast magnetosonic point is located far from the Alfvén point (Kuwabara et al. 2000).

6. Radial Jet Structure

Figures 17, 18, and 19 show the radial dependence of the mass and energy flux of the jet. The mass flux definition is described by equation (18), the toroidal magnetic energy by equation (19) and the Poynting flux, kinetic flux, and enthalpy flux by equations (41)–(43), respectively. Figure 17 shows the radial profiles of the density, poloidal velocity, and toroidal field. The radial dependence of the density shows that the peak density (which defines the jet) in models 1 and 4 (stronger initial magnetic field cases) is greater than that of model 6 (a weaker initial magnetic field case). The radial dependence of the poloidal velocity shows that the collimated flows show higher velocities closer to the disk axis. The radial dependence of the toroidal magnetic field shows that the maximum value is close to the disk axis, as expected for a collimated jet. Figure 18 shows the radial profiles of the Poynting, enthalpy, and kinetic fluxes. We plot the time-averaged flux in annular rings around the symmetry axis ($z$-axis) to show the radial dependence of each energy flux. We show the radial dependence also by taking the different strengths of the initial magnetic field to show their effect. In the case of model 4, the collimations of both the mass flux and the kinetic flux are clear. In figure 18 we show the calculated flux at a fixed $r$.

Fig. 16. Time-averaged values of the enthalpy, kinetic, and Poynting fluxes as a function of the initial magnetic field strength.

Fig. 17. Radial profiles of averaged physical quantities. The cuts were taken at $z = 4$. The radial dependence of the density for each initial magnetic field shows that the peak density (which defines the jet) in the first two initial magnetic fields is greater than that in the last weak initial magnetic field. The radial dependence of the poloidal velocity shows that the collimated flows show higher velocities closer to the disk axis. The radial dependence of the toroidal magnetic field shows that the maximum value is close to the disk axis, as expected for a collimated jet.
height \((z = 4)\). Next we discuss calculations of the flux at different heights. At each height we calculated the position of the maximum of each energy flux. Figure 19 shows the \(r\)-coordinate versus the \(z\)-coordinate of the maximum of the mass flux, toroidal energy, and kinetic flux for models 1 and 3. This figure shows how the maximum of each flux at different heights progresses in the \(z\)-direction. In this figure we can notice that the collimation degree is different for different kinds of the flux and the initial magnetic field. In the case of a strong initial magnetic field, collimation is achieved earlier than in the weak field case. Also, collimation is achieved after some height over the accretion disk.

7. Summary and Discussion

In this paper we have shown that a long-time 2.5-dimensional MHD numerical simulation of magnetized accretion flows leads to an intermittent jet-like outflow and that the flows never reach a steady state. Our simulation is an extension of the work of Matsumoto et al. (1996) and Kudoh, Matsumoto, and Shibata (1998). Both of them performed time-dependent 2.5-dimensional MHD numerical simulations of jets from accretion disks, including the dynamical processes within the disk. They showed that the ejection mechanism of the jets is the same as that in the steady theory; the centrifugal force along the poloidal field line accelerates the jets within an Alfvén radius and, above the Alfvén radius, the jet is accelerated by the magnetic pressure. The simulations for both of them last only for one inner disk rotation. What does happen to the characteristics of the jets if the simulations become longer? The answer to this question is given in this paper.

7.1. Physical Meaning of Time Evolution for \(\dot{M}_w\), \(\dot{M}_a\), and \(E_{\text{mgt}}\)

Cosmic jets are ubiquitous, being quite often associated with newborn stars, X-ray binaries, active galactic nuclei, and gamma-ray bursts. In all such cases, jets and disks seem to be interrelated. Not only jets need disks in order to provide them with the ejected plasma and magnetic fields, but also disks need jets in order that the accreted plasma can get rid of its excess angular momentum to accrete. Observationally, sufficient evidence has already been accumulated for such a correlation. For example, in star-forming regions an apparent
correlation is found between accretion diagnostics and outflow signatures (Hartigan et al. 1995). Hence, our current understanding is that jets are fed by the material of an accretion disk surrounding the central object. The main controller between the mass accretion rate, $\dot{M}_w$, and the mass ejection rate, $\dot{M}_a$, is the magnetic-field strength. Matsumoto et al. (1996) studied the dependence on the strength of the initial magnetic field. They showed that the ratio of the mass ejected as a jet to the total mass of the flux tube was about 10% of the accreting mass. Figures 5 and 11 show that the time evolutions of the mass and toroidal magnetic field outflows are quasi-periodic, and that the periodicity of the jet can be related to the time needed for the initial magnetic field to be twisted to generate a toroidal field. Sato (2003) found that these periodicities are about $2\pi$. The toroidal magnetic field energy increases as the magnetic field line is twisted and accumulated, because the magnetic-field line is dragged with infalling and rotating gases. Then, the piled-up energy is released by a magnetically driven outflow, which is triggered by magnetic reconnection, so that the mass outflow is intermittent and has some periodicities.

Figure 13d shows that the average $\dot{M}_w$ and $\dot{M}_a$ (averaged over 10 rotations) are closely related. The magnetic-field strength is the controller of the relation between $\dot{M}_w$ and $\dot{M}_a$. When the initial magnetic energy is weak, the ratio between $\dot{M}_w$ and $\dot{M}_a$ is high. The ratio decreases with the strength of the initial magnetic field. Pelletier and Pudritz (1992) stated that if the magnetic field of the disk is reduced, there is a very large increase in the mass-loss rate in the wind. The point is that since a slower wind is driven in weaker disk fields, one must provide much more mass in order to carry off the same amount of disk angular momentum. In the case of a strong initial magnetic field the ratio reaches a constant value. These dependences are consistent with the results of the steady solution of Kudoh and Shibata (1997a).

7.2. Dependences on the Strength of the Initial Magnetic Field in the Weak-Field Case

Kudoh and Shibata (1997a) derived the dependence of the mass-outflow rate, $\dot{M}_w$, and the mass-accretion rate, $\dot{M}_a$, on the strength of the initial magnetic field using a semianalytical method,

$$\dot{M}_w \propto E_{mg}^{1/2},$$
$$\dot{M}_a \propto E_{mg}. \quad (44)$$

Furthermore, Michel’s scaling law (Michel 1969) is written as

$$V_{\text{jet}} \sim \left( \frac{\Phi^2 \Psi^2}{M_w} \right)^{1/3} \sim \left( \frac{E_{\text{mg}}}{M_w} \right)^{1/3}. \quad (46)$$

On the other hand, we obtain

$$\frac{M_{w,\text{avg}}}{M_{a,\text{avg}}} \propto E_{mg}^{-0.4}, \quad (47)$$
$$V_{\text{jet,avg}} \propto E_{mg}^{0.15}. \quad (48)$$

Our results shown in figures 12 and 13 are consistent with not only the steady solutions, but also the results of Kudoh, Matsumoto, and Shibata (1998), Kato, Kudoh, and Shibata (2002), and Kigure and Shibata (2005). However, there is a remarkable difference between ours and theirs. They showed only maximum values for $\dot{M}_w$, $\dot{M}_a$, and $V_{\text{jet}}$, while we examined time-averaged values. We find that $V_{\text{jet,avg}}$ and $M_{w,\text{avg}}$ are 0.1 times smaller than $V_{\text{jet, max}}$ and $M_{w, \text{ max}}$ in the case of Kudoh, Matsumoto, and Shibata (1998).

8. Conclusions

We performed time-dependent, 2.5-dimensional, axially symmetric, MHD numerical simulations of jets for many orbital periods. Our simulations solved the dynamical process of the accretion disk for a longer time than that of the previous simulations. Because we wanted to know whether the jet continues to be ejected and becomes steady (or quasi-steady) state, we investigated the time variations of $\dot{M}_w$, $\dot{M}_a$, and $E_{\text{mgt}}$. We also wanted to know the dependences of $M_{w,\text{avg}}$, $M_{a,\text{avg}}$, and $V_{\text{jet,avg}}$ on the initial magnetic field strength, and compared our results with the steady theory and the cases of Kudoh, Matsumoto, and Shibata (1998) and Kato, Kudoh, and Shibata (2002).

We summarize our results in the following:

1. In all models, the ejection of jets is intermittent.
2. The ejection of jets has a period that is related to the toroidal-field formation. There are also relationships between $\dot{M}_w$ and $\dot{M}_a$, and between $M_{w,\text{avg}}$ and $E_{\text{mgt}}$; specifically, their time variations are similar.
3. The dependences of $M_{w,\text{avg}}$, $M_{a,\text{avg}}$, and $E_{\text{mgt}}$ on the strength of the initial magnetic field are consistent with those in the steady theory and the cases of Kudoh, Matsumoto, and Shibata (1998) and Kato, Kudoh, and Shibata (2002). In all cases, however, $V_{\text{jet,avg}}$ and $M_{w,\text{avg}}$ are 0.1 times smaller than $V_{\text{jet, max}}$ and $M_{w, \text{ max}}$ in the cases of Kudoh, Matsumoto, and Shibata (1998), Kato, Kudoh, and Shibata (2002), and Kigure and Shibata (2005).
This work was supported by a Grant-in-Aid for the global COE program “The Next Generation of Physics, Spun from Universality and Emergence” from the Ministry of Education, Culture, Sports, Science and Technology (MEXT). Numerical computations were carried out on VPP5000 at the Astronomy Data Analysis Center of the National Astronomical Observatory, Japan (project ID: iaa31c), an inter-university research institute of astronomy operated by the MEXT.

References

Abramowicz, M., Jaroszynski, M., & Sikora, M. 1978, A&A, 63, 221
Bachiller, R. 1996, ARA&A, 34, 111
Biretta, J. A., Zhou, F., & Owen, F. N. 1995, ApJ, 447, 582
Blandford, R. D., & Payne, D. G. 1982, MNRAS, 199, 883
Bogovalov, S., & Tsinganos, K. 2005, MNRAS, 357, 918
Burrows, C. J., et al. 1996, ApJ, 473, 437
Cao, X., & Spruit, H. C. 1994, A&A, 287, 80
Casse, F., & Keppens, R. 2002, ApJ, 581, 988
Casse, F., & Keppens, R. 2004, ApJ, 601, 90
Contopoulos, J., & Lovelace, R. V. E. 1994, ApJ, 429, 139
Curiel, S., et al. 2006, ApJ, 638, 878
Evans, C. R., & Hawley, J. F. 1988, ApJ, 332, 659
Fendt, C., & Camenzind, M. 1996, A&A, 313, 591
Fendt, C., & Čemeljić, M. 2002, A&A, 395, 1045
Gabuzda, D. C., Murray, E., & Cronin, P. 2004, MNRAS, 351, 89
Hartigan, P., Edwards, S., & Ghandour, L. 1995, ApJ, 452, 736
Hayashi, M. R., Shibata, K., & Matsumoto, R. 1996, ApJ, 468, L37
Jiang, D. R., & Hong, X. Y. 2003, Acta Astron. Sin., 44, 282
Junor, W., Biretta, J. A., & Livio, M. 1999, Nature, 401, 891
Kato, S. X., Kudoh, T., & Shibata, K. 2002, ApJ, 565, 1035
Kudoh, T., & Shibata, K. 1995, ApJ, 452, L41
Kudoh, T., & Shibata, K. 1997a, ApJ, 474, 362
Kudoh, T., & Shibata, K. 1997b, ApJ, 476, 632
Kuwabara, T., Shibata, K., & Matsumoto, R. 2000, PASJ, 52, 1109
Kuwabara, T., Shibata, K., & Uchida, Y. 1995, PASJ, 48, 619
Kudoh, T., Matsumoto, R., & Shibata, K. 1998, ApJ, 508, 186
Kudoh, T., Matsumoto, R., & Shibata, K. 1999, Comput. Fluid Dyn., J., 8, 56
Kudoh, T., Matsumoto, R., & Shibata, K. 2002, PASJ, 54, 121
Kudoh, T., & Shibata, K. 1995, ApJ, 452, L41
Kudoh, T., & Shibata, K. 1997a, ApJ, 474, 362
Kudoh, T., & Shibata, K. 1997b, ApJ, 476, 632
Kudoh, T., & Shibata, K. 1999, Comput. Fluid Dyn., J., 8, 56
Kudoh, T., Matsumoto, R., & Shibata, K. 2002, PASJ, 54, 121
Kudoh, T., & Shibata, K. 1995, ApJ, 452, L41
Kudoh, T., & Shibata, K. 1997a, ApJ, 474, 362
Kudoh, T., & Shibata, K. 1997b, ApJ, 476, 632
Kuwabara, T., Shibata, K., Kudoh, T., & Matsumoto, R. 2000, PASJ, 52, 1109
Kuwabara, T., Shibata, K., Kudoh, T., & Matsumoto, R. 2005, ApJ, 621, 921
Li, Z.-Y. 1995, ApJ, 444, 848
Lister, M. L. 2005, ASP Conf. Ser., 340, 20
López, J. A., Vázquez, R., & Rodríguez, L. F. 1995, ApJ, 455L, 63L
Lovelace, R. V. E., Berk, H. L., & Contopoulos, J. 1991, ApJ, 379, 696
Matsumoto, R., Uchida, Y., Hirose, S., Shibata, K., Hayashi, M. R., Ferrari, A., Bodo, G., & Norman, C. 1996, ApJ, 461, 115
Meier, D. L., Edgington, S., Godon, P., Payne, D. G., & Lind, K. R. 1997, Nature, 388, 350
Michel, F. C. 1969, ApJ, 158, 727
Migliari, S., Tomsick, J. A., Maccarone, T. J., Gallo, E., Fender, R. P., Nelemans, G., & Russell, D. M. 2006, ApJ, 643, L41
Mirabel, I. F., & Rodríguez, L. F. 1994, Nature, 371, 46
Najita, J. R., & Shu, F. H. 1994, ApJ, 429, 808
Ouyed, R., Clarke, D. A., & Padur, R. C. 2003, ApJ, 582, 292
Ouyed, R., & Padur, R. C. 1997a, ApJ, 482, 712
Ouyed, R., & Padur, R. C. 1997b, ApJ, 484, 794
Pellitteri, G., & Padur, R. C. 1992, ApJ, 394, 117
Proga, D. 2003, ApJ, 585, 406
Pudritz, R. E., & Norman, C. A. 1986, ApJ, 301, 571
Pushkarev, A. B., Gabuzda, D. C., Vetukhovskaya, Yu. N., & Yakimov, V. E. 2005, MNRAS, 356, 859
Ray, T. P., Muxlow, T. B. W., Axon, D. J., Brown, A., Corcoran, D., Dyson, J., & Mundt, R. 1997, Nature, 385, 415
Reipurth, B., Heathcote, S., Morse, J., Hartigan, P., & Bally, J. 2002, AJ, 123, 362
Romanova, M. M., Ustyugova, G. V., Koldoba, A. V., Chechetkin, V. M., & Lovelace, R. V. E. 1997, ApJ, 482, 708
Sakurai, T. 1987, PASJ, 39, 821
Sato, K. 2003, Master’s Thesis, Kyoto University
Sauty, C., Trussoni, E., & Tsinganos, K. 2004, A&A, 421, 797
Shibata, K., & Uchida, Y. 1985, PASJ, 37, 515
Shibata, K., & Uchida, Y. 1986, PASJ, 38, 631
Shibata, K., & Uchida, Y. 1987, PASJ, 39, 559
Shibata, K., & Uchida, Y. 1989, in Theory of Accretion Disks, ed. F. Meyer, W. J. Duschl, J. Frank, & E. Meyer-Hofmeister (Dordrecht: Kluwer Academic Publishers), 65
Shibata, K., & Uchida, Y. 1990, PASJ, 42, 39
Stone, J. M., & Norman, M. L. 1992, ApJS, 80, 791
Tingay, S. J., et al. 1995, Nature, 374, 141
Uchida, Y., & Shibata, K. 1985, PASJ, 37, 515
Ustyugova, G. V., Koldoba, A. V., Romanova, M. M., Chechetkin, V. M., & Lovelace, R. V. E. 1995, ApJ, 439, L39
Ustyugova, G. V., Koldoba, A. V., Romanova, M. M., Chechetkin, V. M., & Lovelace, R. V. E. 1999, ApJ, 516, 221
Vlemmings, W. H. T., Diamond, P. J., & Imai, H. 2006, Nature, 440, 58
von Rekowski, B., Brandenburg, A., Dobler, W., & Shukurov, A. 2003, A&A, 398, 825
Weber, E. J., & Davis, L., Jr. 1967, ApJ, 148, 217
Yabe, T., & Aoki, T. 1991, Comput. Phys. Commun., 66, 219
Yabe, T., Ishikawa, T., Wang, P. Y., Aoki, T., Kadota, Y., & Ikeda, F. 1991, Comp. Phys. Commun., 66, 233