GLOBAL THREE-DIMENSIONAL MAGNETOHYDRODYNAMIC SIMULATIONS OF BLACK HOLE ACCRETION DISKS: X-RAY FLARES IN THE PLUNGING REGION

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ABSTRACT

We present the results of three-dimensional global resistive magnetohydrodynamic (MHD) simulations of black hole accretion flows. General relativistic effects are simulated by using the pseudo-Newtonian potential. The initial state is an equilibrium model of a torus threaded by weak toroidal magnetic fields. As the magnetorotational instability (MRI) grows in the torus, mass accretes to the black hole by losing angular momentum. We found that in the innermost plunging region, nonaxisymmetric accretion flow creates bisymmetric spiral magnetic fields and current sheets. Mass accretion along the spiral channel creates one-armed spiral density distribution. Since the accreting matter carries in magnetic fields that are subsequently stretched and amplified as a result of differential rotation, current density increases inside the channel. Magnetic reconnection taking place in the current sheet produces slow-mode shock waves that propagate away from the reconnection site. Magnetic energy release in the innermost plunging region could be the origin of X-ray shots observed in black hole candidates. Numerical simulations reproduced soft X-ray excess preceding the peak of the shots, X-ray hardening at the peak of the shot, and hard X-ray time lags.

Subject headings: accretion, accretion disks — black hole physics — MHD

1. INTRODUCTION

Magnetic fields play essential roles in various activities associated with accretion disks. The most important role is angular momentum transport, which enables the accretion of the disk material. Since Balbus & Hawley (1991) pointed out the importance of the magnetorotational instability (MRI) in accretion disks, the nonlinear growth of the instability has been studied by local three-dimensional magnetohydrodynamic (MHD) simulations (e.g., Hawley, Gammie, & Balbus 1995; Matsumoto & Tajima 1995; Brandenburg et al. 1995) and global MHD simulations (e.g., Armitage 1998; Matsumoto 1999; Hawley 2000; Machida, Hayashi, & Matsumoto 2000).

When general relativistic effects are taken into account, accretion flow in the innermost region of black hole accretion disks changes from a circularly rotating flow to a radially infalling flow. There exists some radius below which almost all energy is advected inward into the black hole. The radial flow is transonic around the last stable orbit. Steady models of black hole accretion disks including this transonic region were constructed under the prescription of -viscosity for optically thick disks (Muchotrzeb & Paczynski 1983; Matsumoto et al. 1984; Abramowicz et al. 1988) and for optically thin disks (e.g., Matsumoto, Kato, & Fukue 1985; Narayan, Kato, & Honma 1997). Global three-dimensional MHD simulations of black hole accretion flows have been carried out by Hawley (2000, 2001), Hawley & Krolik (2001), Haley, Balbus, & Stone (2001), Armitage, Reynolds, & Chiang (2001), Reynolds & Armitage (2001), Hawley & Krolik (2002), Hawley & Balbus (2002), and Krolik & Hawley (2002) by using the pseudo-Newtonian potential (Paczynski & Witta 1980). In black hole accretion flows, they reproduced radial flow profiles similar to those obtained by assuming the phenomenological -viscosity. Hawley & Krolik (2001) showed that the ratio of stress to pressure, which corresponds to the -parameter in conventional disk models, exhibits both systematic gradients and large fluctuations; it rises from in the disk midplane at large radius and 10 near the midplane, well inside the marginally stable radius. Thus the efficiency of angular momentum extraction in the plunging region can be larger than that in transonic disk models (see also Reynolds & Armitage 2001).

Another important mechanism in the innermost region of black hole accretion flow is the magnetic reconnection, which can release magnetic energy even inside the marginally stable radius. In differentially rotating disks, magnetic fields are stretched and amplified as a result of differential rotation. When uniform horizontal field threads an accretion disk, for example, magnetic fields twisted by differential rotation create current sheets, which are subject to magnetic reconnection. Tajima & Gilden (1987) carried out MHD simulations of the formation of current sheets and magnetic reconnection in accretion disks. They applied the mechanism to quasi-periodic oscillations of dwarf novae. Sano & Inutsuka (2001) carried out three-dimensional local MHD simulations of MRI and showed quasi-periodic release of magnetic energy in accretion disks.

In our previous papers (Machida, Hayashi, & Matsumoto 2000; Kawaguchi et al. 2000; Machida et al. 2001), we showed the development of turbulent magnetic fields and current sheets in accretion disks using three-dimensional global MHD simulations assuming ideal MHD. By evaluating the magnetic energy released from the current distribution obtained by ideal MHD simulations, Kawaguchi et al.
(2000) successfully reproduced the 1/\(a^a\)-type time variations observed in black hole candidates.

In addition to the fractal-like 1/\(a^a\) time variations, Negoro et al. (1995) pointed out that the X-ray time variations from Cyg X-1 show large-amplitude X-ray shots. The X-ray shots have a typical time interval of around several seconds and show a symmetric time profile (Negoro et al. 1995; Negoro, Kitamoto, & Mineshige 2001; Negoro & Mineshige 2002). Mannoto et al. (1996) showed by hydrodynamical simulations based on \(a\)-prescription of viscosity that inward propagating waves in advection-dominated black hole accretion flows are reflected around the marginally stable radius and reproduced time-symmetric X-ray profiles. They attributed the origin of rapid (~10 ms) spectral hardening at the peak of X-ray shot to the temperature rise due to acoustic wave reflection.

In this paper we explicitly include electric resistivity in the basic equations and present results of three-dimensional global resistive MHD simulations that show the magnetic energy release following the build up of current sheets in the innermost region of black hole accretion flows.

2. BASIC EQUATIONS AND SIMULATION MODEL

We solved the following resistive MHD equations in cylindrical coordinate system \((\varpi, \phi, z)\) by using a modified Lax-Wendroff scheme with artificial viscosity. We use the normalization \(c = r_a = 1\), where \(c\) is the light speed and \(r_a\) is the Schwarzschild radius:

\[
\frac{\partial\rho}{\partial t} + \mathbf{V}(\rho \mathbf{v}) = 0 ,
\]

\[
\rho \left[ \frac{\partial \mathbf{v}}{\partial t} + \mathbf{v} \cdot \nabla \mathbf{v} \right] = -\nabla P + \mathbf{j} \times \mathbf{B} - \rho \nabla \varphi ,
\]

\[
\frac{\partial \mathbf{B}}{\partial t} = \mathbf{v} \times (\mathbf{v} \times \mathbf{B} - \eta \mathbf{j}) ,
\]

\[
\rho T \frac{dS}{dt} = \eta j^2 ,
\]

where \(\rho\), \(P\), \(\mathbf{v}\), \(\mathbf{B}\), \(\phi\), \(\mathbf{j}\), \(T\), and \(S\) are the density, pressure, velocity, magnetic field, gravitational potential, current density, temperature, and specific entropy, respectively. The specific entropy is expressed as \(S = C_e \ln (P/\rho^C)\), where \(C_e\) is the specific heat at constant volume and \(\gamma\) is the specific heat ratio. A Joule heating term is included in the energy equation. We neglect radiative cooling. We assume the anomalous resistivity \(\eta\) adopted in solar flare simulations (e.g., Yokoyama & Shibata 1994), which mimics the enhancement of resistivity where current density \(j\) is large, as follows:

\[
\eta = \eta_0 \left| \max (\nu_d - \nu_c, 0) \right|^2 ,
\]

where \(\nu_d \equiv j/\rho\) and \(\nu_c\) is the threshold above which anomalous resistivity sets in. We adopt \(\nu_c = 8\), and \(\eta_0 = 5 \times 10^{-4}\).

The initial condition is an equilibrium model of an axisymmetric MHD torus threaded by toroidal magnetic fields surrounding a black hole. We assume pseudo-Newtonian potential \(\phi = -GM/(r - r_a)\) (Paczynski & Witta 1980), where \(G\) is the gravitational constant and \(M\) is the mass of the black hole. We neglect the self-gravity of the gas. At the initial state, the torus is assumed to have a constant specific angular momentum \(L\) and assumed to obey a polytropic equation of state \(P = K \rho^\gamma\), where \(K\) is a constant. We assume that the torus is embedded in a hot, nonrotating, isothermal, spherical halo.

Following Okada, Fukue, & Matsumoto (1989), we assume the square of the Alfvén speed \(v_A^2 = B_z^2/(4\pi \rho)\) \(\equiv \mathcal{H}(\rho)\gamma^2\), where \(B_z\) is the toroidal magnetic field and \(\mathcal{H}\) is a constant. Using this assumption, we can integrate the equation of motion into a potential form,

\[
\Psi(\varpi, z) = -\frac{1}{2} \left( \frac{L^2}{2 \varpi^2} + \frac{1}{\gamma - 1} v_A^2 + \frac{\gamma}{2(\gamma - 1)} v_A^2 \right) = \Psi_h = \text{constant} ,
\]

where \(v_A^2 = \gamma P/\rho\) is the square of the sound speed, \(r = (\varpi^2 + z^2)^{1/2}\), and \(\Psi_h = \Psi(\varpi_h, 0)\). Here the reference radius \(\varpi_h\) is defined as the radius where the rotation speed \(L/\varpi_h\) equals the Keplerian velocity \(v_k = [\varpi_h (\partial \phi/\partial \varpi)_{\varpi=\varpi_h}]^{1/2}\). Using equation (6), we obtain the density distribution

\[
\rho = \rho_h \left( \frac{\max (\Psi_h + 1/[2(r - 1)] - L^2/(2\varpi^2), 0)}{K[\gamma/(\gamma - 1)][1 + \beta_h^{-1} (2\varpi^2/\varpi_h^2)/\gamma^{-1}]} \right)^{1/(\gamma - 1)} ,
\]

where \(\beta_h \equiv (2K/H)/\varpi_h^{2(\gamma - 1)}\) is the ratio of gas pressure to magnetic pressure at \((\varpi, z) = (\varpi_h, 0)\). We set \(\rho_h = 1\). The parameters describing the structure of the MHD torus are \(\gamma, \beta_h, L, \) and \(K\). In this paper we report the results of simulations for parameters \(\varpi_h = 50r_a, \beta_h = 100, \gamma = 5/3, L = (\varpi_h/2)^{1/2} \varpi_h/\varpi_h - 1, \) and \(K = 0.0005\). The density and sound speed of the halo at \(r = \varpi_h\) is taken to be \(\rho_{\text{halo}}/\rho_h = 10^{-3}\) and \(v_\varpi/\varpi_h = 3/5,\) respectively.

The number of mesh points is \((N_\varpi, N_\phi, N_z) = (250, 64, 192)\). The grid size is \(\Delta \varpi = \Delta z = 0.1\) for \(0 < \varpi, z < 10\) and otherwise increases with \(\varpi\) and \(z\). The outer boundaries at \(\varpi = 150\) and at \(z = 70\) are free boundaries where waves can transmit. A periodic boundary condition is imposed for the \(\varpi\)-direction. Since we include the full azimuthal angle \(0 \leq \varphi \leq 2\pi\) in the simulation region, we can follow the evolution of low azimuthal wavenumber nonaxisymmetric structures. We imposed symmetric boundary condition at the equatorial plane. An absorbing boundary condition is applied at the inner boundary by introducing a damping rate,

\[
a_1 = 0.1 \left( 1.0 - \tanh \frac{r - 2 + 5 \Delta \varpi}{2 \Delta \varpi} \right) .
\]

Inside \(r = 2\), the deviation of physical quantity \(q\) from initial value \(q_0\) is artificially reduced with damping rate \(a_1\) as

\[
q^{\text{ew}} = q - a_1 (q - q_0) .
\]

This damping layer serves as a nonreflecting boundary that absorbs accreting mass and waves propagating inside \(r = 2\). Small-amplitude, random perturbations are imposed at \(t = 0\) for the azimuthal velocity.

3. NUMERICAL RESULTS

3.1. Global Structure of Nonradiative Accretion Disk

Figures 1a and 1b show the isosurface of the density \((\rho = 0.4)\) at \(t = 0\) (initial condition) and \(t = 30, 590 \approx 10_0\),
where $t_0 = 2\pi \omega_b / v_{Kb} \approx 3079$ is the one orbital time at $\omega_b$. Owing to the efficient angular momentum redistribution, the torus is deformed into a disk and matter accretes to the central part by losing the angular momentum. In the outer part of the torus, matter gains angular momentum and expands radially. The angular momentum transport is initially driven by the magnetorotational instability. After the nonaxisymmetric MRI grows, the disk region becomes turbulent.

Figure 2a shows the azimuthally averaged density contours and the momentum vectors at $t = 30, 590$. Circulating motions similar to those observed in ideal MHD simulations (e.g., Hawley & Krolik 2001; Machida et al. 2001) appear near the equatorial plane ($\omega \sim 12$). In the innermost region ($\omega < 5$) radial infall becomes significant. Near the surface layer around $\omega \sim 7$, a shock wave front is formed and higher angular momentum gas blows out into the disk corona. Figure 2b shows the azimuthally averaged density distribution at $t = 0, t = 21,502$, and $t = 30,590$.

Figures 3a and 3c show the equatorial density distribution. Figures 3b and 3d show magnetic field lines projected onto the equatorial plane. In the outer region (Figs. 3a and 3b), the magnetic field lines are tightly wound and show turbulent structure. In the inner region (Figs. 3c and 3d), magnetic field lines are less turbulent and globally show a bisymmetric spiral (BSS) shape. A one-armed spiral mode

![Fig. 1.](image1) Fig. 1.—(a) Initial model of the torus. (b) Isosurface of density ($\rho = 0.4$) at $t = 30,590$. Numbers show the $(x, y, z)$ coordinate.

![Fig. 2.](image2) Fig. 2.—(a) Azimuthally averaged density distribution log $\rho$ (color scale) and the poloidal momentum vectors (arrows) at $t = 30,590$. Long arrows are plotted in dark colors. (b) Azimuthally averaged density distribution at $t = 0, t = 21,502$, and $t = 30,590$. 

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dominates in the density distribution. Figure 4 shows a three-dimensional structure of magnetic field lines at $t = 30,590$ around the surface of the innermost region ($\varpi < 10$) of the disk. White curves show the magnetic field lines. The yellow plane shows the equatorial plane, and the blue surface is the isosurface of the density ($\rho = 0.4$). Magnetic field lines have significant $z$-components and show helical structure near the surface of the disk. Matter swirls into the black hole along these field lines. Bisymmetric spiral magnetic fields are created owing to the infall of the disk material.

3.2. Time Evolution and Angular Momentum Redistribution

The amplification of magnetic fields and its saturation occurs similarly to those reported by Hawley & Krolik (2001, 2002). Figure 5a shows the time evolution of the ratio of the gas pressure to magnetic pressure $\beta = \langle P \rangle / \langle B^2 / 8\pi \rangle$, where the angle brackets mean the volume average. The solid curve and the dashed curve show the average in the inner region ($4 \leq \varpi \leq 10, 0 \leq z \leq 1$) and in the outer region ($20 \leq \varpi \leq 40, 0 \leq z \leq 3$), respectively. The disk stays in a quasi–steady state with $\beta \sim 10$ for timescales much longer than the dynamical timescale. Figure 5b shows the time evolution of the ratio of the Maxwell stress to pressure $\alpha_B \equiv -\langle B_\varpi B_z / 4\pi \rangle / \langle P \rangle$ averaged in the inner region and in the outer region. We found $\alpha_B \simeq 0.1$ in the inner region and $\alpha_B \simeq 0.02$ in the outer region. Figure 5c shows the time variation of the accretion rate at $\varpi = 2.5$. The accretion rate increases with time even at the end of simulation because much mass is still stored in the original torus ($\varpi_b \sim 50$). It is beyond the scope of this paper to continue simulations on timescales long enough for the black hole to accrete most of the mass of the initial torus.
The Maxwell stress in Figure 5b shows spikes similar to those observed in local simulations. Sano & Inutsuka (2001) used local three-dimensional resistive MHD simulations to show that these spikes are created by current dissipation in channel flows that appear in the nonlinear stage of the MRI. Although the BSS magnetic fields that appear in our simulation are not exactly the same as channel flows in a disk initially threaded by vertical magnetic fields, the BSS fields have current layers where the magnetic field lines change their direction. Magnetic reconnection can take place in such current layers. We call this region the BSS channel.

Figure 6 shows the distribution of density, pressure, radial velocity, and specific angular momentum near the equatorial plane averaged azimuthally, vertically (0 < z < 3), and in time. The solid and dashed curves show average in 29,000 < t < 31,000 and 28,000 < t < 29,000, respectively. Around ω ≃ 10, the radial structure changes from a dense torus to an accreting disk (Fig. 6a). The equatorial density increases with time because the torus flattens. Figure 6b shows the averaged gas pressure and the magnetic pressure. In the innermost region (2 < ω < 3), magnetic pressure is about 30% of the gas pressure. Figure 6c shows the distribution of the radial velocity. In the inflowing region (ω ≤ 10), the radial velocity is roughly proportional to ω⁻³ and exceeds 0.1c around the radius of marginally stable orbit (ω ≃ 3). The radial velocity in the outer region (ω > 10) is much smaller than the sound speed. Figure 6d shows the specific angular momentum averaged in the vertical direction, in the azimuthal direction, and in time. The dotted curve shows the Keplerian angular momentum distribution. Although the distribution of initial angular momentum is uniform, the angular momentum distribution becomes nearly Keplerian in the radial range 3 < ω < 100.

Figure 7 shows the time development of the density averaged azimuthally and vertically (0 ≤ z ≤ 0.3). The color scale shows the logarithmic density, the range of which is −0.5 (blue) ≤ log ρ ≤ 0 (pink). The interface between the circularly rotating dense disk and radially infalling flow (blue region) moves inward from ω ≃ 30 at t ≃ 24,000 to
Fig. 6.—Radial profiles of density, gas pressure, magnetic pressure, radial velocity, and specific angular momentum averaged in the azimuthal direction, in the vertical direction (0 < z < 0.3) and in time. The solid curves show quantities averaged in 29,000 < t < 31,000, and the dashed curves show quantities averaged in 28,000 < t < 29,000. The dotted curve in (d) shows the Keplerian angular momentum distribution.

Fig. 7.—Time development of the azimuthally averaged density at the equatorial plane. The range of the color scale is −0.5 (blue) ≤ log ρ ≤ 0 (pink).

Fig. 8.—Time development of the azimuthally averaged radial velocity at the equatorial plane Δv_r = v_r − (v_r), where (v_r) = −max(3π⁻¹, 0.003). The range of the gray scale is −0.07 (dark) ≤ Δv_r ≤ 0.07 (light).
Fig. 9.—Positions of Lagrange test particles on the equatorial plane. Large symbols (1–6) trace the same particles.
$\omega \simeq 3$ at $t \sim 26,000$. The radius of the interface oscillates around $\omega \simeq 10$. The infall of dense matter from the inner region of the torus ($\omega \simeq 20$) takes place intermittently at $t \sim 26,000$ and $t \sim 29,000$. The interval of the infall is nearly equal to the rotation period of the dense torus ($t_0 \sim 3079$). The intermittency arises from the oscillation of the torus excited by the angular momentum exchange between the infalling matter and the torus material; the dense torus gains angular momentum from the infalling blob as a result of the angular momentum conservation. Note that in Figure 7 the location of the density maximum of the torus is pushed outward following the infall of dense blobs. Short-timescale oscillations in the inner region are overlaid on this global oscillation.

Figure 8 shows the time development of the azimuthally averaged radial velocity relative to the mean radial velocity approximated as $\langle v_r \rangle = -\max(3\omega^{-3}, 0.003)$ at the equatorial plane. The light and dark areas show the regions where $\Delta v_r = v_r - \langle v_r \rangle > 0$ and $\Delta v_r < 0$, respectively. The innermost region ($2 \leq \omega \leq 10$) shows quasi-periodic oscillations with frequency $f \sim \kappa_{\text{max}}/2\pi \sim 1140(M/M_\odot)^{-1}$ (Hz), where $\kappa_{\text{max}}$ is the maximum of epicyclic frequency, as already indicated by one-dimensional simulations (e.g., Matsumoto, Kato, & Honma 1988; Kato, Fukue, & Mineshige 1998). Significant inflow occurs around $t = 25,000, 27,000, 28,000, 29,500,$ and $31,000$.

Figure 9 shows the position of the Lagrange test particles uniformly distributed in the equatorial plane inside $\omega < 18$ at $t = 30,319$. The spiral arms that appear in Figure 9 correspond to the dense arms in the density distribution (see Fig. 3). The diamond-shaped symbols trace the same test particles. The motion of the particles show the development of nonaxisymmetric MRI. The particle labeled 4, for example, has lower angular momentum than other particles and

Fig. 10.—Density distribution in the region $2 \leq \omega \leq 10$. The color scale in each panel shows the scale of log $\rho$. Numbers at the upper right corner of each panel show time. Arrows indicate the reconnection jet.
spirally infalling particles around the black hole is intermittent rather than time steady. Figure 8 indicates that accretion to the black hole is intermittent rather than time steady.

Figure 8 shows that evacuated region is created around $x = 5, y = -5$ at $t = 30, 590$. The sharp ridge of the density distribution around $x = 0, y = -7$ in Figure 3b indicates the interaction between the outgoing wave and the infalling matter. In order to show the propagation of the outgoing wave, we show in Figure 10 the distribution of density. At $t = 30, 570$, a dense region (orange) appears and propagates outward. As indicated by the arrows, density enhancement develops in the evacuated region and propagates in the direction opposite to the rotation of the disk. In the next subsection, we show that the motions of these density enhancements are due to magnetic reconnection.

### 3.3. Magnetic Reconnection in the Plunging Region

Figure 11 shows the time evolution of the current density distribution and magnetic field lines in the innermost region ($-10 \leq x \leq 10, -5 \leq y \leq 5$). In the outer region, where magnetic field lines show turbulent structure (Fig. 3b), the isocontours of current density have fractal-like distributions (Kawaguchi et al. 2000). On the other hand, long current sheets are formed in the inner infalling region ($\varpi \leq 8$), where magnetic field lines show a BSS shape (Fig. 3d). Magnetic energy is accumulated in these current sheets. The red areas in Figure 11 are active regions where the current density is high. Active region A is a narrow peak in the current distribution (current sheet). The white arrows labeled A trace the motion of this current sheet. Active region B has a cusp in the current distribution. The white arrows labeled B indicate the tip of the reconnecting magnetic field lines behind the reconnection point. At $t = 30, 590$, the current layer in active region B coincides with the density enhancement shown in Figure 10b. The lower panels in Figure 11 show that the cusp-shaped current layer B propagates in the direction relatively opposite to the direction of disk rotation and that current dissipation takes place. These are the magnetic flaring sites where magnetic energy is released by magnetic reconnection. The double peak in the current distribution that splits from the cusp is a characteristic feature of slow shocks associated with magnetic reconnection (see Priest 1982).

Figure 12 shows the time evolution of magnetic field lines. Closed magnetic loops (magnetic islands) formed around the cusp of magnetic field lines indicate that magnetic reconnection is taking place in the current sheet. Figure 12 indicates that magnetic reconnection in an accretion disk has two types: (1) magnetic reconnection taking place inside an elongating magnetic loop and (2) magnetic reconnection driven by the interaction of two magnetic loops.

In Figure 13 we schematically show the magnetic reconnections in solar corona and in accretion disks. The arcade-type solar flares (e.g., Tsuneta et al. 1992) that accompany loop eruption belong to the loop-elongation type. Solar microflares driven by loop interaction produce X-ray brightening and X-ray jets (see, e.g., Shibata et al. 1992). The loop-elongation–type reconnection can be seen in active region A at $t = 30, 610$ in Figures 11 and 12c. A backward-propagating reconnection jet appears, as indicated by the arrows in Figure 10c. Active region B in Figure 11 is subject to both types of reconnection. At $t = 30, 590$, active region B collides with the preceding magnetic loop (Fig. 12b) and produces intense current layers. Magnetic reconnection taking place in this current layer released more energy than active region A because current layer B has larger volume. Although loop interaction releases less energy than arcade eruptions in the solar corona, loop interaction can release more energy in accretion disks because the rotational energy of the disk is converted to magnetic energy by such interactions.

Figure 14 shows the time evolution of temperature. Magnetic energy is deposited into the thermal energy and heats the plasma around the reconnection site. Filamentary-
Fig. 12.—Magnetic field lines projected onto the equatorial plane. Numbers show time. The size of the box is $-5 < x < 5$ and $-5 < y < 5$.

Fig. 13.—Schematic picture of magnetic reconnection in solar flares and in accretion disks.
shaped hot regions that coincide with current sheets trace the slow shocks associated with magnetic reconnection.

Figure 15a shows the time evolution of the volume-integrated Joule heating rate $E_j = \int \eta j^2 dV$ in $2 \leq \omega \leq 6$ and $0 \leq z \leq 10$. Figure 15b shows the time evolution of $\int j^2 dV$. Figure 15c shows the magnetic energy $E_m = \int (B^2/8\pi) dV$. Figure 15d shows the accretion rate at $\omega = 2.5$. Arrows indicate $t = 30, 590$ and $t = 30, 630$. The largest flare occurs around this time interval. Since we turn on anomalous resistivity only when $j/\rho > (j/\rho)_c$, magnetic energy is released after the accretion rate decreases and the innermost region is rarefied. The Joule dissipation rate $\eta j^2$ tends to be zero when the matter density of the BSS channel is high. This is the reason that $E_j$ has maximum after the peak of $\int j^2 dV$. Magnetic energy decreases between $t = 30, 590$ and $t = 30, 630$, as expected from the magnetic reconnection model. Only a small fraction ($<1\%$) of the released magnetic energy goes into the Joule heating because $\eta \leq 10^{-4}$ in this simulation.

The released magnetic energy is converted into thermal energy by slow shock (see Fig. 14). The volume-integrated current density correlates well with the mass accretion rate. This is because magnetic fields are stretched and amplified by mass accretion. In other words, parts of the gravitational energy of the accreting matter are stored in the current sheet and released by magnetic reconnection.

Next let us compare the X-ray light curve obtained by numerical simulations with observations of X-ray shots. Figure 16 shows the time variation of the X-ray luminosity computed by the optically thin bremsstrahlung luminosity $F_X = \int \rho^2 T^{7/2} f(T) dV$ (dashed curve) and Joule heating rate $E_j = \int \eta j^2 dV$ (solid curve), where the integration is over the innermost region ($2 \leq \omega \leq 6, 0 \leq z \leq 10$). We introduce the cutoff function $f(T)$, where $f(T) = 1$ when $0.01 < T < 0.1$ and $f(T) = 0$ otherwise. The vertical scale of Figure 16 is arbitrary because $F_X$ and $E_j$ depend on the unit density, $\rho_0$, and the electric resistivity, $\eta_0$ (see eq. [5]).
As the matter in dense spiral arms accretes to the innermost region, X-ray luminosity gradually increases and peaks at around $t/C_{24}^{30}; 100$. The timescale of the increase of X-ray luminosity is the accretion timescale of dense blobs, and typically $t/C_{24}^{103} = c/C_{24}^{0.01} M = M/C_{12}^{-s}$ in our simulation (see Fig. 7). As the dense blob is swallowed into the black hole, X-ray luminosity decreases. Subsequently, magnetic reconnection takes place in the rarefied region and releases the magnetic energy. The soft X-ray excess precedes the magnetic reconnection by about $500 r_g/c \sim 5 M/\dot{M}_o$ ms.

### 4. SUMMARY AND DISCUSSION

We have shown that in the innermost plunging region of black hole accretion disks, the growth of nonaxisymmetric MRI creates BSS magnetic fields. Mass accretion proceeds along these spiral channels. As a result, the density distribution tends to show one-armed spiral structure.

A current sheet is created inside the BSS channel because the magnetic field changes its direction inside the channel. The current sheet is built up by converting the gravitational energy of the accreting gas into magnetic energy. In other words, magnetic energy is accumulated by mass accretion. When the BSS channel is rarefied after dense blobs infall, magnetic reconnection takes place and converts the magnetic energy into heat. This mechanism is analogous to the accumulation of magnetic energy by mass motion preceding solar flares and protostellar flares (e.g., Hayashi, Shibata, & Matsumoto 1996). Magnetic reconnection can also take place through the interaction of two long current sheets. These X-ray flares in the plunging region could be the origin of the large-amplitude X-ray time variabilities characteristic of black hole candidates.

Let us compare our numerical results with X-ray observations of Cyg X-1 (Negoro et al. 1995). In X-ray–hard states, X-ray flux from Cyg X-1 shows the largest shots, the interval of which is typically several seconds (Negoro & Mineshige 2002). In our simulation, the largest shot occurred at $0.3 M/\dot{M}_o$ s after the initial torus is set up at $\varpi \sim 50$. The magnetic reconnection produces outgoing

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**Fig. 15.** Time evolution of (a) the volume-integrated Joule heating rate, (b) the volume-integrated current density, (c) the magnetic energy, and (d) the accretion rate at the $\varpi = 2.5$. Arrows indicate $t = 30, 590$ and $t = 30, 630$. 

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waves and cusp-shaped current sheets. We have shown that the mass accretion from the circularly rotating dense torus is not time steady but intermittent. The interval of mass feeding from the torus is the rotation period of the torus (Fig. 7a). In our simulation, this timescale corresponds to $t_0 \sim 3000 \rho_g/c \sim 0.03 M/\dot{M}_\odot$ s. If the torus locates at $\rho_g = 500 \rho_g$, this interval $t_0 \sim 0.9 M/\dot{M}_\odot$ s approaches the interval of observed shots if the mass of the black hole is 10 $M_\odot$. Once dense blobs accrete to the innermost region, the profile of the shot is essentially determined by the local processes in the innermost region. Thus the profile of one shot may not depend on the location of the center of the torus. However, the interval of the shot depends on $\omega_0$. Smaller X-ray shots and fluctuations can be created by magnetic reconnections ubiquitous in accretion disks. Observations of Cyg X-1 indicate that the averaged shot profile is time symmetric (Negoro, Kitatmoto, & Mineshige 2001) and X-rays become hard at the peak of the shot.

Figure 17 schematically shows the mechanism of X-ray shots. The upper left-hand panel shows the accretion stage when dense blobs fall into the innermost region. As the density in the innermost region increases, the soft X-ray luminosity increases. Current sheets are formed inside the BSS channel, and magnetic energy is accumulated. After the blobs are swallowed into the black hole, the current sheet in the BSS channel is rarefied. The soft X-ray luminosity begins to decrease because X-ray-emitting gas is depleted. The largest magnetic reconnection takes place in the rarefied current sheet by loop elongation or by loop interaction. Since the released magnetic energy is converted to heat, X-ray luminosity will suddenly increase and then decay exponentially. In order to compute the light curve after magnetic reconnection, we will have to include the effects of heat conduction along the magnetic loop, which may evaporate the disk material (see the numerical simulation of solar flares by Yokoyama & Shibata 2001). We would like to report the results of such simulations in the future.

The BSS magnetic fields in the innermost region of black hole accretion disks contribute to the angular momentum transport of the disk material. In this region, angular momentum transport due to this ordered spiral magnetic fields is more efficient than the angular momentum transport by turbulent magnetic fields.

The dynamical effects of magnetic fields in the plunging region of black hole accretion flows have been investigated analytically by Krolik (1999), Gammie (1999), and Agol & Krolik (1998, 2000) and have been simulated by Hawley & Krolik (2001, 2002) and Krolik & Hawley (2002). Our numerical results support their conclusion that the ratio of stress to pressure ($\alpha$) has a systematic gradient with radius and has larger values well inside the plunging region.

In the innermost region of the disk, the accretion flow is dominated by radial advection. The radial structure of the innermost region ($\rho \leq 10 \rho_g$) of accretion flow obtained by three-dimensional MHD simulations approaches the global transonic solution of optically thin disks with the viscous parameter $\alpha \sim 0.1$ (e.g., Narayan, Kato, & Honma 1997). Direct numerical simulations such as the one we present in this paper have the potential to do much more than seek agreement with $\alpha$-models. They can predict the time variation of black hole accretion flows without assuming the phenomenological $\alpha$-parameter. In order to confirm the applicability of numerically obtained accretion flows to black hole candidates, it is essential to compute the X-ray spectrum from numerical results and compare them with observations. We would like to report the results of such analysis in subsequent papers.

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