SUPERCRITICAL ACCRETION FLOWS AROUND BLACK HOLES: TWO-DIMENSIONAL, RADIATION PRESSURE–DOMINATED DISKS WITH PHOTON TRAPPING

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ABSTRACT

The quasi-steady structure of supercritical accretion flows around a black hole is studied based on two-dimensional radiation-hydrodynamic (2D-RHD) simulations. The supercritical flow is composed of two parts: the disk region and the outflow regions above and below the disk. Within the disk region the circular motion and the patchy density structure are observed, which is caused by Kelvin-Helmholtz instability and probably by convection. The mass accretion rate decreases inward, roughly in proportion to the radius, and the remaining part of the disk material leaves the disk to form the outflow because of the strong radiation pressure force. We confirm that photon trapping plays an important role within the disk. Thus, matter can fall onto the black hole at a rate exceeding the Eddington rate. The emission is highly anisotropic and moderately collimated so that the apparent luminosity can exceed the Eddington luminosity by a factor of a few in the face-on view. The mass accretion rate onto the black hole increases with the absorption opacity (metallicity) of the accreting matter. This implies that the black hole tends to grow faster in metal-rich regions, such as in starburst galaxies or star-forming regions.

Subject headings: accretion, accretion disks — black hole physics — hydrodynamics — methods: numerical — radiative transfer

1. INTRODUCTION

It is widely believed that accretion flow onto black holes drives major activities of astrophysical black holes, such as active galactic nuclei (AGNs), Galactic black hole candidates (BHCs), and possibly gamma-ray bursts (GRBs). It is also a common belief that the basic accretion processes and radiation properties can be well described by the standard-disk model of Shakura & Sunyaev (1973). On the other hand, observational facts that cannot fit the standard-disk picture are also being accumulated recently. A good example is intense high-energy emission from black holes, which indicates the presence of very hot plasmas around black holes with temperature $T \sim 10^9$ K. This leads to the ideas of accretion disk corona and/or radiatively inefficient flow (RIAF).

The standard-disk picture breaks down not only in the low-luminosity regimes but also in the high-luminosity regimes, in which the mass accretion rate, $M$, becomes comparable to or exceeds the critical mass accretion rate, $M_{\text{crit}} \equiv L_\text{E}/c^2$, where $L_\text{E}$ is the Eddington luminosity.

Apparently, features of the supercritical flows look similar to those of the standard-type disks, since the disk remains optically thick and thus emits blackbody-like emission, just as the standard-type disks do. Flow dynamics is, however, distinct.

What makes the supercritical accretion flows distinct from the standard-disk-type flow is the presence of photon trapping (Belgelman 1978). Photon trapping occurs when the photon diffusion time, the time for photons to travel from the equatorial plane to the surface, exceeds the accretion timescale. Under such circumstances, photons generated via the viscous process are advected inward with gas flow without being able to go out from the surface immediately. We can define the trapping radius, inside which photon trapping is significant,

$$r_{\text{trap}} \sim \frac{H}{r} r_g,$$

(1)

where $m$ is the mass accretion rate normalized by the critical mass accretion rate, $H$ is the half-thickness of the disk, $r$ is the radius, and $r_g(=2GM/c^2)$ is the Schwarzschild radius, where $M$ is the black hole mass (Ohsuga et al. [2002], hereafter Paper I; see also Begelman [1978] for the spherical case).

It is claimed that photon-trapping effects were incorporated into the slim-disk model by Abramowicz et al. (1988; for a review, see Kato et al. [1998], § 10.1). We have shown previously, however, that the slim-disk model does not fully consider the photon-trapping effects, since it is a radially one-dimensional model, whereas photon trapping is basically a multidimensional effect (see Paper I). We thus need at least a two-dimensional treatment. Indeed, we have demonstrated in Paper I that the photon trapping is grossly underestimated in the slim-disk model.

Let us recall what is lacking in the slim-disk model more explicitly. The equation of energy balance, $Q_{\text{in}} = Q_{\text{rad}} + Q_{\text{adv}}$, is solved in the slim-disk model, where $Q_{\text{vis}}$, $Q_{\text{rad}}$, and $Q_{\text{adv}}$ are the viscous heating, radiative cooling, and advective cooling rates per unit surface, respectively. The problem resides in that the radiative cooling is evaluated under the usual diffusion approximation (in the vertical direction). This approximation may be justified if radial inflow of gas is totally negligible so that photons can mainly diffuse in the vertical direction. However, this is not the case when the diffusion timescale is longer than the accretion timescale. This leads to overestimation of $Q_{\text{rad}}$, and hence, underestimations of $Q_{\text{adv}}$, compared with the correct value. More importantly, photon trapping modifies the spectral energy distribution (SED) as well. We have found that large photon trapping yields spectral softening, because hard photons
that are created deep inside the disk are more effectively trapped than soft photons (Ohsuga et al. 2003). Since our previous definitions are only partly multidimensional, we, as a next step, need to perform a fully two-dimensional analysis of the supercritical accretion flows. This is a major motivation for the present study.

When considering photon-trapping effects, we should also pay attention to the fact that the supercritical accretion flow becomes geometrically thick. Then multidimensional gas motion, such as convective or large-scale circulation, might occur. Furthermore, strong outflow might also be generated at the disk surface via radiation pressure force. Such complex flow motion will influence the radiation energy distribution through the advective energy transport, which, in turn, affects the flow motion via radiation pressure force. We need to carefully solve such strong coupling between radiation and matter.

Two-dimensional radiation-hydrodynamic (2D-RHD) simulations of accretion disks were initiated by Eggum et al. (1987), who assumed equilibrium between gas and radiation. The improved simulations, in which the energy of the gas and radiation are separately treated, were performed by Kley (1989), Okuda et al. (1997), Fujita & Okuda (1998), Kley & Lin (1999), and Okuda & Fujita (2000).

The simulations of supercritical flows around black holes were pioneered by Eggum et al. (1988), who again assumed equilibrium between gas and radiation, and were improved by Okuda (2002). Eggum et al. (1988) showed that mass accretion onto the black hole occurs at a supercritical rate, although the mass accretion and mass outflow rates were still variable in their simulations; that is, a quasi-steady state had not been achieved by their simulations. This also happens in the simulations by Okuda (2002), in which the resulting luminosity slowly decreases with time. He also found that the mass accretion rate is subcritical around the black hole, although the mass is injected from the outer boundary at a supercritical rate. Note that the calculation times of these simulations were only 0.6 and 1.6 s in physical units, respectively, for a black hole mass of 10 $M_\odot$. Notably, these are shorter than the viscous timescale,

$$\tau_{\text{vis}} \sim (5.7 \text{ s}) \left( \frac{M}{10 M_\odot} \right)^{3/2} \left( \frac{r}{100r_g} \right)^{3/2} \left( \frac{\alpha}{0.1} \right)^{-1} \left( \frac{H/r}{0.5} \right)^{-2},$$

where $\alpha$ is the viscosity parameter. More recently, Okuda et al. (2005) performed long-term 2D-RHD calculations of the supercritical accretion. The luminosity and mass accretion rates seem to be quasi-steady in their simulations, although the flow structure is not steady yet. Since the sum of the mass accretion rate at the inner boundary and mass outflow rate at the outer boundary is much smaller than the mass input rate at the disk boundary, the mass within the computational domain continues to increase. In addition, the photon trapping did not appear in their simulations, whereas the mass accretion rate exceeds the critical value.

To summarize, despite the interesting simulations being made so far by two groups, the quasi-steady structure of the supercritical accretion flows still remains an open issue. Furthermore, interesting issues related to supercritical flow, including the photon-trapping effects, the dependence of the observed luminosity on various viewing angles, and the effects of metallicity on the flow structure, have not been investigated previously. This is a motivation for the present study.

Here we report for the first time the quasi-steady structure of the supercritical disk accretion flows around black holes, which were revealed by 2D-RHD simulations. Through the present simulations we mainly aim at understanding the dynamics of the viscous flow in the vicinity of a black hole of $r \lesssim 100r_g$. Basic equations and our model are explained in § 2, and the numerical methods are described in § 3. We then display the quasi-steady flow structure and study the photon-trapping effects in the simulated flow in § 4. Finally, §§ 5 and 6 are devoted to discussion and conclusions.

### 2. Basic Equations and Assumptions

We solve the full set of RHD equations including the viscosity term. We use spherical polar coordinates $(r, \theta, \varphi)$, where $r$ is the radial distance, $\theta$ is the polar angle, and $\varphi$ is the azimuthal angle. In the present study, we assume that the flow is non–self-gravitating, reflection symmetric relative to the equatorial plane (with $\theta = \pi/2$), and axisymmetric with respect to the rotation axis (i.e., $\partial/\partial\varphi = 0$). We describe the gravitational field of the black hole in terms of pseudo-Newtonian hydrodynamics, in which the gravitational potential is given by $\Psi(r) = -GM/(r-r_s)$, as was introduced by Paczynsky & Wiita (1980). The basic equations are the continuity equation,

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0,$$

the equations of motion,

$$\frac{\partial(p\mathbf{v})}{\partial t} + \nabla \cdot (p\mathbf{v} \mathbf{v}) = -\frac{\partial p}{\partial r} + \rho \left[ \frac{v_r^2 + v_\theta^2}{r} + \frac{GM}{(r-r_s)^2} \right] + f_r + q_r,$$

$$\frac{\partial(p\mathbf{r} \mathbf{v})}{\partial t} + \nabla \cdot (p\mathbf{r} \mathbf{v} \mathbf{v}) = -\frac{\partial p}{\partial \theta} + \rho v_\theta^2 \cot \theta + r f_\theta + r q_\theta,$$

and the energy equation of the gas,

$$\frac{\partial \rho e}{\partial t} + \nabla \cdot (\rho e \mathbf{v}) = -p \nabla \cdot \mathbf{v} - 4\pi \kappa B + c \kappa E_0 + \Phi_{\text{vis}},$$

and the energy equation of the radiation (Mihalas & Mihalas 1984),

$$\frac{\partial E_0}{\partial t} + \nabla \cdot (E_0 \mathbf{v}) = -\nabla \cdot \mathbf{F}_0 - \nabla \cdot \mathbf{P}_0 + 4\pi \kappa B - c \kappa E_0.$$

Here $\rho$ is the gas mass density, $\mathbf{v} = (v_r, v_\theta, v_\varphi)$ is the velocity, $p$ is the gas pressure, $e$ is the internal energy density of the gas, $B$ is the blackbody intensity, $E_0$ is the radiation energy density, $\mathbf{F}_0$ is the radiation flux, $\mathbf{P}_0$ is the radiation pressure tensor, $\kappa$ is the absorption opacity, $q = (q_r, q_\theta, q_\varphi)$ is the viscous force, and $\Phi_{\text{vis}}$ is the viscous dissipative function. The radiation force, $f_{\text{rad}} = (f_r, f_\theta)$, is given by

$$f_{\text{rad}} = \frac{\chi}{c} \mathbf{F}_0,$$

where $\chi = (\kappa + \rho \sigma \tau/m_p)$ is the total opacity, where $\sigma \tau$ is the Thomson scattering cross section and $m_p$ is the proton mass.

As the equation of state, we use

$$p = (\gamma - 1) e,$$
where $\gamma$ is the specific heat ratio. The temperature of the gas, $T$, can then be calculated from

$$p = \frac{\rho k_B T}{\mu m_p},$$

(11)

where $k_B$ is the Boltzmann constant and $\mu$ is the mean molecular weight.

To complete the set of equations, we apply flux-limited diffusion (FLD) approximation developed by Levermore and Pomraning (1981). In this framework, the radiation flux is written as

$$F_0 = -\frac{c}{\chi} \nabla E_0,$$

(12)

with the flux limiter, $\lambda$, and the radiation pressure tensor is expressed in terms of the radiation energy density via

$$P_0 = f E_0,$$

(13)

where $f$ is the Eddington tensor. Here the flux limiter is given by

$$\lambda = \frac{2 + R}{6 + 3R + R^2},$$

(14)

using the dimensionless quantity, $R = |\nabla E_0|/(\chi E_0)$. The components of the Eddington tensor are

$$f = \frac{1}{2} (1 - f) I + \frac{1}{2} (3f - 1) n n,$$

(15)

where $f$ is the Eddington factor,

$$f = \lambda + \lambda^2 R^2,$$

(16)

and $n$ is the unit vector in the direction of the radiation energy density gradient,

$$n = \frac{\nabla E_0}{|\nabla E_0|}.$$

(17)

This approximation holds in both the optically thick and thin regimes. In the optically thick limit, we find $\lambda \rightarrow 1/3$ and $f \rightarrow 1/3$ because of $R \rightarrow 0$. In the optically thin limit of $R \rightarrow \infty$, on the other hand, we have $|F_0| = c E_0$. These give correct relations in the optically thick diffusion limit and optically thin streaming limit, respectively.

For the absorption opacity, we consider the free-absorption, $\kappa_{\text{ff}}$, and bound-free absorption, $\kappa_{\text{bf}}$, as

$$\kappa = \kappa_{\text{ff}} + \kappa_{\text{bf}},$$

(18)

where $\kappa_{\text{ff}}$ and $\kappa_{\text{bf}}$ are given by

$$\kappa_{\text{ff}} = 1.7 \times 10^{-25} T^{-7/2} \left( \frac{\rho}{m_p} \right)^2 \text{ cm}^{-1},$$

(19)

and

$$\kappa_{\text{bf}} = 4.8 \times 10^{-24} T^{-7/2} \left( \frac{\rho}{m_p} \right)^2 \left( \frac{Z}{Z_\odot} \right) \text{ cm}^{-1},$$

(20)

where $Z$ is the metallicity (Hayashi et al. 1962; Rybicki & Lightman 1979).

Here we assume that only the $r^2$-component of the viscous stress tensor, which plays important roles for the transport of the angular momentum and heating of the disk plasma, is nonzero,

$$\tau_{rr} = \eta \frac{\partial}{\partial r} \left( \frac{v_r}{r} \right),$$

(21)

in the present study, where $\eta$ is the dynamical viscosity coefficient. Then the radial and polar components of the viscous force are null ($q_r = q_\theta = 0$), and the right-hand side of equation (6) is described as

$$r (\sin \theta) q_\phi = \frac{1}{r^2} \frac{\partial}{\partial r} \left[ r^2 (\sin \theta) \tau_{rr} \right].$$

(22)

The viscous dissipative function is given by

$$\Phi_{\text{vis}} = \eta \left[ \frac{r}{r^2} \left( \frac{v_r}{r} \right) \right]^2.$$  

(23)

Finally, we prescribe the dynamical viscosity coefficient as a function of the pressure

$$\eta = \frac{\alpha}{\Omega_K} \frac{p + \lambda E_0}{\Omega_K},$$

(24)

where $\Omega_K$ is the Keplerian angular speed. It is a modified $\alpha$ prescription of the viscosity, which is proposed by Shakura & Sunyaev (1973). In this form, the dynamical viscosity coefficient is proportional to the total pressure because of $\lambda \rightarrow 1/3$ in the optically thick regime. In the optically thin regime, by contrast, we find $\eta = \alpha p/3 \Omega_K$ [or the kinematic viscosity is $\nu = \eta/\rho = \alpha c_s^2/\Omega_K$], where $c_s \equiv (p/\rho)^{1/2}$ is the isothermal sound velocity], since $\lambda$ vanishes in this limit.

Here we need to remark that the realistic formalism about the viscosity should be investigated from the magnetohydrodynamic point of view, since the dominant sources of viscosity would be chaotic magnetic fields and turbulence in the gas flow (e.g., Machida et al. 2000; Stone & Pringle 2001).

3. NUMERICAL METHODS

3.1. The Code

We numerically solve the set of RHD equations shown in § 2 using an explicit-implicit finite difference scheme on the Eulerian grids. Our methods and boundary conditions are similar to those of Okuda (2002), but we adopt a different initial condition and have carried out long-term calculations in order to examine a quasi-steady structure of the supercritical accretion flows. Since we assume axisymmetry, as well as reflection symmetry, the computational domain can be restricted to one quadrant of the meridional plane. The domain consists of spherical shells of $r_{\text{in}} \leq r \leq r_{\text{out}}$ and $0 \leq \theta \leq \pi/2$, where $r_{\text{in}}$ and $r_{\text{out}}$ are the radial coordinates of the inner and outer boundaries of the computational domain, respectively, and the domain is divided into $N_r \times N_\theta$ grid cells, amounting to $96 \times 96$ meshes. The $N_r$ grid points in the radial direction are equally spaced logarithmically, while the $N_\theta$ grids are equally distributed in such a way to achieve $\Delta \cos \Theta = 1/N_\theta$. All the physical quantities are defined at each cell center.

We divide the numerical procedure for the finite-difference equations into the following steps: (1) hydrodynamic terms for ideal fluid, (2) advection term in the energy equation of the radiation, (3) radiation flux term in the radiation energy equation, (4) gas-radiation interaction terms in the energy equations.
of the radiation and gas, and (5) viscous terms in the momentum equation and energy equations of the gas. Steps 1 and 2 are solved with the explicit method, while steps 3–5 are treated based on the implicit method. In the first step, we use the computational hydrodynamic code Virginia Hydrodynamics One. It is based on the piecewise parabolic method (Colella & Woodward 1984). Equations (3)–(7), except the viscous terms and the gas-radiation interaction terms of equation (7), are solved in this step. In the second step, an integral formulation is used to generate a conservative differencing scheme for the advection term of equation (8). The energy transport by the radiation flux term is solved in the third step. The radiation energy density is updated again using the BiCGSTAB method for a matrix inversion. In the fourth step, we solve the gas-radiation interaction terms in equation (7). All the terms on the right-hand side of equation (8), except for the radiative flux term, are also treated in this step. The radiation energy and gas energy are advanced simultaneously. The method used in this step is basically the same as that described by Turner & Stone (2001). In the final step, we update the azimuthal component of the velocity by solving the viscosity term in equation (6) with the Gauss-Jordan elimination for a matrix inversion. The viscous dissipative function is also calculated in this step.

Throughout the present study, we assume $M = 10 M_\odot$, $\alpha = 0.1$, $\gamma = 5/3$, and $\mu = 0.5$. The size of the computational domain is set at $(r_{\text{in}}, r_{\text{out}}) = (3r_g, 500r_g)$. (Here it is noted that our conclusion does not change so much, even if we employ $r_{\text{in}} = 2r_g$.)

3.2. Boundary and Initial Conditions

We employ the absorbing inner boundary condition so that the density, gas pressure, and velocity are smoothly dumped (e.g., Kato et al. 2004). Here it is noted that our simulation results are not sensitive to the inner boundary condition or to the location of the inner boundary. The results do not change even if we use the free boundary conditions, for which all the matter and waves can transmit freely. The radiation flux is set at $F_\phi = cE_0$ at the inner boundary; that is, we apply the condition of the optically thin limit at the inner boundary.

The outer boundary at $r = r_{\text{out}}$ is divided into two parts: the disk part ($\theta \geq 0.45\pi$) and the part above the disk ($\theta < 0.45\pi$). Throughout the outer-disk boundary we assume that matter is continuously injected into the computational domain. We set the injected mass accretion rate (mass input rate) so as to be constant at the boundary. In the present study, we explore the cases of $\dot{m}_{\text{input}} = 300, 1000,$ and 3000, where $\dot{m}_{\text{input}}$ is the mass input rate normalized by the critical mass accretion rate. The injected matter is supposed to rotate with sub-Keplerian speed, and it has a specific angular momentum corresponding to the Keplerian angular momentum at $r = 100r_g$, since our main purpose of the present study is to investigate the viscous accretion flows within $r \leq 100r_g$. By setting the Keplerian radius ($100r_g$) much smaller than the radius of the outer boundary, we can prevent the outer boundary conditions from directly affecting the accretion flow within $100r_g$. With such a large outer boundary, we can reproduce complex motions like circulation around $100r_g$ in the disk; that is, the matter can transiently go out across the radius of $100r_g$ and return. At the outer boundary region above the accretion disk, we use free boundary conditions and allow for matter to go out but not to come in. If the radial velocity is negative at the outermost grid, it is automatically set at zero. We also employ the radiation flux in the optically thin limit, $F_\phi = cE_0$, at the outer boundary.

With respect to the rotation axis we assume that $\rho, p, v_r$, and $E_0$ are symmetric, while $v_\theta$ and $v_\phi$ are antisymmetric. On the equatorial plane, on the other hand, $\rho, p, v_r, v_\phi$, and $E_0$ are symmetric, and $v_\theta$ is antisymmetric.

We start the calculations with a hot, rarefied, and optically thin atmosphere. Initially there is no cold, dense disk in the computational domain. The initial atmosphere is constructed to approximately achieve hydrostatic equilibrium in the radial direction; namely, its density profile is given by

$$\rho = \rho_{\text{out}} \exp \left[ \frac{\mu m_T G M}{k \theta T_{\text{out}}} \left( \frac{r_{\text{out}}}{r} - 1 \right) \right],$$

where $\rho_{\text{out}}$ is the density at the outer boundary. We employ $\rho_{\text{out}} = 10^{-17}$ g cm$^{-3}$ and $T = 10^{11}$ K in the present study. Since this atmospheric gas is finally ejected out of the computational domain, it does not affect the resulting quasi-steady structure.

3.3. Time Step

The time step is restricted by the Courant-Friedrichs-Levi condition. We set the time step as

$$\Delta t = \xi \min \left( \frac{\Delta r}{|v_r| + c_s}, \frac{r \Delta \theta}{|v_\theta| + c_s} \right),$$

where $\xi$ is a parameter and $\Delta r$ and $\Delta \theta$ are the cell sizes in the radial and polar directions, respectively. Although we only show the results for the cases with $\xi = 0.4$, we also simulated the cases with $\xi = 0.1$ or 0.05, confirming that our conclusions are not altered by this change.

4. RESULTS

4.1. Quasi-steady Structure

We first represent time evolution of 2D-RHD simulations. Overall evolution is divided into two distinct phases: the accretion phase and the quasi-steady phase.

The mass is continuously injected through the outer-disk boundary and creates continuous gas inflow because of the gravity force by the central black hole. Since the angular momentum of the injected mass is set equal to the Keplerian angular momentum at $r = 100r_g$, it is natural that the gas tends to accumulate around the regions with a radius of $100r_g$ by degrees (see Fig. 1). This is the accumulation phase. Since Eggum et al. (1988) and Okuda (2002) started calculations with a cold, dense disk, this phase did not appear in their simulations. Eventually, the viscosity starts to work so that the angular momentum of the gas can be transported outward, which drives inflow gas motion in a quasi-steady fashion. This is the quasi-steady phase.

Such a two-step evolution is clear in Figure 2. In this figure, we show the time evolution of the normalized mass accretion rate onto the central black hole, $\dot{m} = M/(L_e/c^2)$ (thick solid curve), the luminosity, $L/L_e$ (thin solid curve), the viscous heating rate, $L_{\text{vis}}/L_e$ (dotted curve), and the total mass contained within the computational domain, $m_{\text{total}} = M_{\text{total}}/10^{19}$ g (dashed curve). Here we set the mass input rate to $\dot{m}_{\text{input}} = 1000$ and the metallicity to $Z = 1 Z_\odot$. Both the accretion rate and luminosity steadily increase until $t \lesssim 7$ s. We also notice that the mass accretion rate is much smaller than the mass input rate ($\dot{m}_{\text{input}} = 1000$), and that the total mass within the computational domain rapidly increases with time, indicating that mass accumulation really occurs in this phase. Then the quasi-steady accretion phase starts (at $t \gtrsim 7$ s) when the mass accretion rate exceeds the critical value, $\dot{m} > 1$, and all the physical quantities
stay nearly constant. (The constant $m_{\text{total}}$ implies that the sum of the mass accretion rate at the inner boundary and mass outflow rate at the outer boundary is equal to the mass input rate.) We can thus conclude that we could for the first time succeed in reproducing the quasi-steady state of the supercritical accretion flows with 2D-RHD simulations. The critical time ($t_\text{sep}$) separating these two phases roughly coincides with the viscous timescale ($t_\text{viscous}$). It is natural that the quasi-steady structure is obtained within 10 s, which is the viscous timescale at the Keplerian radius ($r \sim 100r_g$), although the size of the computational domain ($500r_g$) is much larger. This is because the injected matter accretes from the outer boundary to the Keplerian radius with a free-fall velocity so that the evolutionary timescale of the outer regions is given by the free-fall timescale at $r = 500r_g$, $(1.6 s) (r/500r_g)^{3/2}[M/(10 M_\odot)]$. This is certainly shorter than the viscous timescale at $r = 100r_g$.

Let us next examine the quasi-steady structure in some detail. Figure 3 displays the cross-sectional view of the density distributions (colors), overlaid with the velocity vectors (arrows) at $t = 10$ s. Here the $x$- and $y$-axes are $R = r \sin \theta$ and $z = r \cos \theta$, respectively. We understand from this figure that the flow structure is roughly divided into two regions: the disk region around the equatorial plane (orange) and the outflow region above the inflow region (blue). Roughly, the boundary is at $z/R \sim 0.8$; that is, the disk is geometrically and optically thick, as was predicted by the slim-disk model. However, the density distribution definitely deviates from that of the slim-disk model, since it is neither smooth nor plane parallel in the vertical direction. We can even see a number of cavities in Figure 3. The flow pattern is also complex, although the slim-disk model predicts the simple convergence flow. We found the prominent circular motion within the disk and the strong outflow that is generated at the disk surface. The patchy structure around the boundary between the dense inflow region and the rarefied outflow region seems to be caused by the Kelvin-Helmholtz (K-H) instability (discussed below). The less dense gas in the outflow region penetrates into the disk body because of the K-H instability, thus forming the cavities.

To understand the dynamics of our entire calculation domain, we show the two-dimensional density distribution in the whole region in Figure 4. We see from this figure that the density is
larger around the equatorial plane (\(|z| \leq 10r_g\)) than at large altitudes (\(|z| \gtrsim 10r_g\)) in the outer region, \(R \gtrsim 100r_g\). This reflects that the injected mass accretes along the equatorial plane. On the other hand, the viscous accretion disk forms inside \(r \lesssim 100r_g\). We focus on the viscous accretion flows in the present study. The situation is apparently similar to that studied by Chakrabarti (1996), although they were concerned with subcritical flow, and their study did not use multidimensional or radiation-hydrodynamic simulations.

Figure 5 shows the two-dimensional distributions of radiation energy density (top left), the ratio of the radiation energy to the internal energy of gas (top right), the gas temperature (bottom left), and the radial velocity normalized by the escape velocity (bottom right) on the \(R-z\) plane. As shown in the top left panel, the radiation energy distribution roughly coincides with the gas density distribution. That is, the radiation energy tends to be larger around the equatorial plane than that around the rotation axis. Since the radiation energy distribution is smoothed due to the radiative diffusion within the disk, there is no cavity found in the radiation energy distribution, which makes a marked difference from the density distribution (see Fig. 3).

Radiation energy greatly exceeds the gas energy in the entire region, including the outflow region in our simulations (Fig. 5, top right). This confirms that most of the region is radiation pressure dominated and that strong radiation pressure supports the geometrically thick disk and drives the outflow.

The gas temperature distribution shown in the bottom left panel shows relatively low temperatures in the disk region, compared with those in the outflow region, although the viscous heating rate is much larger in the former than in the latter. This can be understood, because radiative cooling is more effective in a dense region as a consequence of its strong density dependence. The gas in the disk region is heated up by the viscous heating, and the processed energy is effectively converted into the radiation energy. Therefore, gas temperature does not rise so much within the disk. Conversely, the gas cannot emit effectively in the outflow region, since both free-free emissivity and bound-free emissivity are more sensitive to the density than the gas temperature, \(\propto \rho^2T^{1/2}\). It can be understood by the comparison between the bottom left and the top right panels. The radiation temperature (\(\propto E_r^{1/4}\)) in the outflow region is lower than that in the disk region, contrary to the gas temperature profile.

The bottom right panel indicates the radial velocity normalized by the escape velocity. The gas moves toward the black hole or flows outward slowly in the disk regions (Fig. 5, blue area). The white color indicates that the velocity exceeds the escape velocity in this area. It is found that the gas is accelerated through the radiation pressure and is blown away to a large distance (see also Fig. 5, top right). Such a flow component will be identified as a strong disk wind. Note that the outflow presented here is distinct in nature from that known as a bounce jet (Chen et al. 1997), which arises because of a bounce of free-falling low angular momentum material when it goes through the centrifugal barrier at small \(r\).

The outflow will also produce large absorption in the emergent spectra. We also notice strong velocity shear at the boundary between the disk region and the outflow region. This complex density profile around the disk surface as shown in Figure 3 is explained as a consequence of the K-H instability.

The growth timescale of the K-H instability is roughly given by

\[
T_{KH} \approx \frac{1}{k v_r} \left( \frac{\rho_{disk}}{\rho_{out}} \right)^{1/2},
\]

where \(k\) is the wavenumber and \(\rho_{disk}/\rho_{out}\) is the density ratio of the disk region to the outflow region at the disk boundary. Here we assume an incompressible fluid as well as \(\rho_{disk} \gg \rho_{out}\) and neglect the gravity. Also, the viscosity is not considered, since the \(r\theta\)-component of the viscous stress tensor is set at zero in the present simulations. By setting \(k = 2 \pi/10r_g\), \(v_r = 0.1c\), and \(\rho_{disk}/\rho_{out} = 10\), we find \(T_{KH} \sim 5 \times 10^{-3}\) s. Note that this is shorter than the escape time, \(r/v_r = 0.1\) s, for \(r = 100r_g\) and \(v_r = 0.1c\). That is, there is ample time for the K-H instability to grow before the material flows outward.

The mass accretion rates as a function of the radius are displayed in Figure 6 for the case of \(m_{input} = 1000\) and \(Z = 1 Z_\odot\). The solid, dotted, and dashed curves indicate the \(\dot{m}\) profiles at elapsed times of \(t = 10, 30, \) and \(50\) s, respectively. Here the mass accretion rate at each radius is evaluated by

\[
\dot{m}_r = - \frac{c^2}{2E_r} \int 2\pi r^2 \rho(r, \theta) \min[0, v_r(r, \theta)] \sin \theta d\theta.
\]

It is found that the mass accretion rates are not constant in the radial direction but decrease inward, as the flow approaches the black hole. Roughly, we find \(\dot{m}_r \propto r\). In addition, we find that the radial \(\dot{m}_r\) profile remains nearly the same after 7 s, from which we can conclude that the flow is in a quasi-steady state. This \(\dot{m}_r\) change is caused by a cooperation between wind mass loss around the disk surface and the circular motion deep inside the disk. Note that accretion rates are assumed to be constant in space in the slim-disk model formulation.

The multidimensional numerical simulations of RIAFs have shown that the accretion disks are convectively unstable, and thereby the circular motion is driven within the flow (e.g., Igumenshchik & Abramowicz 1999; Stone et al. 1999; McKinney & Gammie...
The convection might cooperate with the K-H instability in generation of the circular motion, as well as the cavities found in our simulations. The simulations of RIAFs have also revealed that $\dot{m}_r$ increases with radius as $\dot{m}_r \propto r^{3.4-1}$, which is similar to our value, and was attributed to the circular motion, as well as to the bipolar outflows (Stone et al. [1999]; see also Narayan et al. [2000] for a self-similar solution). Here we need to emphasize that although the flow structures look similar at first glance, the basic physical processes are distinct from those of RIAF simulations. That is, entropy is carried by fluid in the unstable, rotating magnetized gas.
RIAF, whereas it is mostly by photons in the present case. In addition, since the RIAF simulations basically do not take into account the radiative cooling, they would overestimate the driving force of the outflows.

4.2. Photon Trapping

The photon trapping characterizes the supercritical accretion flows. It works to reduce the energy conversion efficiency, \( \dot{L}/\dot{M}c^2 \) (see, e.g., Paper I; Ohsuga et al. 2003). As a result, the flow luminosity becomes insensitive to the mass accretion rate when \( \dot{m} \gg 1 \). However, simple stream lines were assumed in the previous studies on supercritical flows, including the slim-disk model. Even if we do not consider possible multidimensional gas motions, photon trapping works to some degree, leading to a reduction in the energy conversion efficiency. Such multidimensional gas motion and associated reaction of the radiation can be calculated in the present RHD simulations, since we have calculated the radiation energy transport, being coupled with the gas dynamics.

Figure 7 plots the luminosity as a function of the mass accretion rate onto the black hole. The filled squares, circles, and triangles indicate the results for \( Z = 10, 1, \) and \( 0 Z_{\odot} \), respectively, with different mass input rates \( \dot{m}_{\text{input}} = 300, 1000, \) and 3000 from left to right. We also indicated in the same figure the luminosity calculated based on the slim-disk model (Watarai et al. 2000) and the one based on model A of Paper I with dashed and dotted curves, respectively. In Paper I, a simple model for the accretion flow is employed, and the luminosity, carefully taking into account the photon trapping, is calculated by solving energy transport inside the accretion flows. (More precisely, the luminosity plotted in Fig. 7 is the corrected one, which fixes initial small errors in their model A [see Fig. 1 in Paper I].) Here it should be stressed that the mass accretion rate onto the black hole is not an input parameter but is calculated dynamically in the present simulations, although it was a parameter in both model A of Paper I and the slim-disk model. The resulting \( \dot{m} \) profile was shown in Figure 6.

It is evident in Figure 7 that the calculated luminosity agrees more with model A of Paper I than with the slim-disk model, in all the cases. This proves that the two-dimensional effects of photon trapping are really significant in the supercritical accretion flows. This result is, in a sense, surprising. In model A in Paper I, we assumed that the viscous heating occurs only in the vicinity of the equatorial plane, although the gas might be heated up at a high altitude. Since the photons emitted at deep inside the disk tend to be more effectively trapped in the flow, we anticipated that the photon-trapping effects would be reduced in the realistic situation, compared with model A.

We also argued in Paper I that photon-trapping effects may be attenuated by the presence of large-scale circulation motion, which could help photon diffusion motion and thus considerably reduce photon traveling time to the surface of the flow. However, our current results show significant photon-trapping effects even when we explicitly include complex flow motions.

Figure 7 also shows that the mass accretion rate onto the black hole increases with an increase of the metallicity (for fixed mass input rates). In our simulations the gas with higher metallicity has larger absorption opacities so that the gas energy can be more effectively converted into the radiation energy, yielding smaller gas pressure than in the metal-poor case. The gas could be effectively blown away by the strong gas pressure. However, there is a countereffect. The gas with large absorption opacity enhances the radiation energy. Enhanced radiation flux and large opacity can drive strong radiation pressure–driven outflows. The physical cause will be investigated in the future.

Let us see more explicitly how significant photon trapping is. Figure 8 compares the distributions of the radial component of the radiative flux in the comoving frame (top) and that in the inertial frame (bottom) at \( t = 10 \) s. Other parameters are the same; the mass accretion rate is \( \dot{m}_{\text{input}} = 1000 \), and the metallicity is \( Z = 1 Z_{\odot} \). The former radiative flux \( (F_0^r) \) is roughly...
proportional to the radial gradient of the radiation energy distribution. On the other hand, the latter flux \( F^r \) includes the advective transport of radiation energy \( (v_i E_0) \) in addition to the former flux; namely, we approximately have \( F^r \sim F^r_0 + v_i E_0 \). In other words, differences between the two panels represent how significant photon trapping is. In fact, we see a significant difference between the two panels; whereas the blue area (indicating large radiation flux) is restricted to the vicinity of the black hole in the top panel, it is more widely spread over the entire disk region in the bottom panel. This is a direct manifestation of the photon-trapping effects.

### 4.3. Collimated Emission

Since the supercritical accretion flows are geometrically and optically thick, the observed images and luminosity should strongly depend on the viewing (inclination) angle. We calculate the intensity map using the monochromatic radiation transfer equation,

\[
I \cdot \nabla I = \rho \Gamma^3 \left( 1 - \frac{v \cdot I_0}{c} \right) \left( \frac{c B}{4 \pi m_p} E_0 - \chi I_0 \right),
\]

where \( I \) is the specific intensity, \( \Gamma \equiv (1 - v^2/c^2)^{-1/2} \) is the Lorentz factor, and \( I = (\sin \Theta \cos \Phi, \sin \Theta \sin \Phi, \cos \Theta) \) and \( I_0 = \Gamma^{-1}(1 - v \cdot I/c)^{-1}[I - (\Gamma/c)v + (\Gamma - 1/v^2)(I \cdot v)v] \) are the directional cosine in the inertial and comoving frame, respectively, where \( \Theta \) and \( \Phi \) are the azimuthal and polar angles, respectively, of radiation propagation in the inertial frame. Here we assumed isotropic scattering.

The calculated effective temperature \( T_{\text{eff}} \) maps at \( t = 50 \text{ s} \) are shown in Figure 9 for various viewing angles: \( \cos i = 1/8, 3/8, 5/8, \) and \( 7/8 \). The other parameters are \( m_{\text{input}} = 1000 \) and \( Z = 1 \ Z_\odot \). On the observer’s screen, the black hole lies at the center of the Cartesian coordinate \( (X, Y) \). Since we use the spherical computational domain with a size of \( 500 r_g \) in the present study, the four blue corners in each map lie outside the domain. In the face-on view (Fig. 9, top left; \( \cos i = 7/8 \)), the effective temperature exceeds \( 3 \times 10^7 \text{ K} \) within \( 100 r_g \) and amounts to \( \sim 10^7 \text{ K} \) in the central region. Such a high-temperature region disappears in the top right panel (the case with \( \cos i = 3/8 \)). In this panel, the most luminous region is found on the top side (not at the center) and is elongated in the vertical direction. This is caused by an occultation of the innermost part of the flow by the outer parts. (See Fukue [2000] and Watarai et al. [2005] for self-occultation effects based on the slim-disk model.) Since the accretion flow is both geometrically and optically thick, the innermost region cannot be seen for large inclination angles, \( i \).

The emission from the supercritical accretion flows is mildly collimated for the same reason. Figure 10 represents the viewing angle dependence of the isotropic luminosity (normalized by bolometric luminosity) for \( Z = 1 \ Z_\odot \) (circles) and \( 10 \ Z_\odot \) (squares). The other parameters are \( t = 50 \text{ s} \) and \( m_{\text{input}} = 1000 \). Here the isotropic luminosity is calculated by assuming an isotropic radiation field, \( L(i) = 4 \pi D^2 I(i) \zeta \), where \( D \) is the distance from an observer and \( \zeta \) is a revised factor. The luminosity calculated by solving the radiation transfer equation does not always coincide with that evaluated by the FLD method. Here we assume \( \zeta = 1.7 \) to fit these luminosities. As expected, the isotropic luminosity is quite sensitive to the viewing angle, indicating that the emission from the flow is mildly collimated. If the flow shape would be flat and (geometrically) thin and if no collimation would occur, the observed luminosity should vary along the cosine curve, as is indicated by the dotted curve.

We also find that the angle dependence of the isotropic luminosity is more enhanced in the large-metallicity case of \( Z = 10 \ Z_\odot \) than in that of \( Z = 1 \ Z_\odot \). This is because of different density distributions. As we have already mentioned, the mass outflow rate is smaller in the high-metallicity case than in the
low-metallicity case. As a result, the density contrast between the disk region and the outflow region above the disk gets larger with an increase in the metallicity. Accordingly, the hot innermost region of the high-metallicity flows are easier to observe because of smaller optical depth in the outflow region for small viewing angles.

The supercritical accretion flows would be identified as very high $L/L_E$ objects in the face-on case, since the emission is mildly collimated in the polar direction and the bolometric luminosity itself exceeds the Eddington luminosity. The former effect is enhanced in the case of high metallicity (large absorption opacity). It is also true that supercritical flow may be

Fig. 9.—Surface temperature images for the model for different viewing angles: $\cos i = 1/8, 3/8, 5/8, \text{and } 7/8$ in the top left, top right, bottom left, and the bottom right panels, respectively. The origin of the Cartesian coordinate on the observer’s screen, $(X, Y) = (0, 0)$, is set at the location of the black hole. The black hole is at the center. The adopted parameters are $\dot{m}_{\text{input}} = 1000$ and $Z = 1 Z_\odot$.
identified as a fairly subcritical source, if the viewing angle is large (Watarai et al. 2005).

5. DISCUSSION

5.1. Comparison with the Slim-Disk Model

5.1.1. Flow Structure

Here we directly compare the structure of the supercritical accretion flows based on our two-dimensional simulations with that of the slim-disk model. Figure 11 shows the time-average density profile (top), radial and rotation velocity profiles (middle), and radiative temperature profile (bottom) on the equatorial plane in the quasi-steady accretion phase of $t = 40–50$ s. The adopted parameters are $\dot{m}_{\text{input}} = 1000$ and $Z = 1 Z_\odot$. The resulting profiles of the one-dimensional numerical study for the slim-disk model are represented by the dashed curves (Watarai et al. 2000). The thin solid curve in the top and bottom panels indicates the slope of the self-similar solution of the slim-disk model. The dotted curve in the middle panel is the Keplerian velocity.

The density profile is shown by the solid curve in the top panel. The density peak around $70R_g$ is made by accumulation of the injected matter, which occurs because of the centrifugal barrier. We compare the density profile in the viscous accretion regime with that of numerical and self-similar solutions of the slim-disk model. In this panel, it is found that the slope of the density profile of our simulations is steeper than that of the numerical solution of the slim-disk model. We also find that the slope is close to that of the self-similar solution of the slim disk, $\rho \propto R^{-3/2}$ (Watarai & Fukue 1999; see also Spruit et al. 1987), whereas the profile becomes flatter in the vicinity of the black hole ($R < 5r_g$). Here we note that the slope in our simulations is steeper than that of the RIAFs. The RIAF simulations showed the density profile to be $\rho \propto R^{-1/2}$ (McKinney & Gammie 2002).

In the middle panel, the radial and rotation velocities are plotted by the circles and the solid curve, respectively. The dotted curve indicates the Keplerian velocity, $v_K = (GM/R)^{1/2}/(R - r_g)$. It is found that the rotation velocity nicely agrees with the Keplerian velocity, as well as with the result of the numerical solution of the slim disk. However, the radial velocity profile is not smooth and largely deviates from that of the slim-disk value, whose slope agrees with that of the Keplerian velocity (dotted curve) in the free-fall regime. The radial velocity remarkably varies at $R = 5r_g - 70r_g$, and the radial velocity sometimes becomes negative, around $R = 6r_g - 20r_g$, and $60r_g$. Such a complex $v_r$ profile is formed by the prominent circular motions around the equatorial plane (see Fig. 3).

The bottom panel shows the radiation temperature profile, where the radiation temperature is defined as $T_r = (E/R)^{1/4}$, where $a$ is the radiation constant. As shown in this panel, the profile nicely agrees with that of the one-dimensional numerical solution and is flatter than that of the self-similar solution of the slim-disk model, $T_r \propto R^{-5/8}$.

To sum up, we can thus conclude that the quasi-steady structure of the supercritical accretion flows is apparently similar to but, precisely speaking, deviates from the prediction of the slim-disk model.

5.1.2. Effective Temperature

Next, we represent again the effective temperature profile in order to directly compare with the prediction of the slim-disk model, although the two-dimensional images of the effective temperature have already been shown in Figure 9. Figure 12 represents the effective temperatures at $Y = 0$ as a function of $X$ for $i = 0$ and $\pi/6$, where $X$ and $Y$ are the horizontal and vertical coordinates on the observer’s screen (see Fig. 9). The other parameters are $\dot{m}_{\text{input}} = 1000, Z = 1 Z_\odot$, and $t = 50$ s. It is found
that the slope of the effective temperature profile for $i = 0$ becomes flatter than that of the slim-disk model, $T_{\text{eff}} \propto R^{-1.2}$, within $30r_g$. In contrast, we found that the profile is consistent with the slim-disk model in the outer region. As we have already mentioned, the effective temperature, as well as the luminosity, is quite sensitive to the viewing angle, since the accretion flow is both geometrically and optically thick. We found that the effective temperature profile for $i = \pi/6$ is almost flat.

The slim-disk model succeeds in reproducing the observed behavior of ultraluminous X-ray sources (ULXs) and narrow-line Seyfert 1 galaxies (NLS1s), which are thought to be candidates of near-critical or supercritical flows (Watarai et al. 2001; Mineshige et al. 2000; Kawaguchi 2003). The SEDs are produced based on the blackbody emission with the effective temperature coupled with some modification. We here note that the temperature profile obtained by our current simulations shows some deviations from that of the slim-disk model.

5.2. Blueshifted Absorption by Outflow Matter

One characteristic feature of the supercritical accretion flow is its generation of the radiation pressure–driven wind. Our simulations show that a large quantity of gas is blown away by the strong radiation pressure above and below the accretion disk. Such outflow material, which is thought to be highly ionized by the radiation from the accretion flow, would absorb the continuum emission. Therefore, the supercritical accretion objects have blueshifted absorption lines by the highly ionized ions. Such blueshifted absorption in high-ionization lines was observed in the UV spectra of broad absorption line quasars (Weymann et al. 1981; Becker et al. 2000). Moreover, Pounds et al. (2003a, 2003b) and Reeves et al. (2003) reported by the blueshifted X-ray absorption lines that the highly ionized matter has outflow velocity on the order of $\sim 0.1c$ in quasars, and the flow is likely to be optically thick to the electron scattering within $r \sim 100r_g$. In our simulations, the typical outflow velocity is $0.1c$, and the optical depth of the outflow regions is around $1.2(\rho/10^{-3} \text{ g cm}^{-3})(r/100r_g)$. (This estimated value does not depend on the black hole mass because of $\rho \propto M$ and $r \propto M^{-1}$.) These are remarkably close to the observed results, and therefore our simulations basically account for the origins of high-velocity outflows. Detailed comparison with the observations is left for future work.

5.3. Growth of Supermassive Black Holes

We reveal by 2D-RHD simulations that the black hole can swallow the gas at the supercritical accretion rate, although the luminosity exceeds the Eddington luminosity. This result implies that the black hole can rapidly grow, and the growth timescale is given by $M/M = 4.5 \times 10^8 (\dot{m}/100)^{-1}$ yr. Thus, if the supermassive black holes in the galactic nuclei are built up by the supercritical accretion, such a growing phase would be quite short because of a short growth timescale.

Umemura (2001) suggested the radiation-hydrodynamic model for the formation of the supermassive black holes, in which the supermassive black holes grow in the ultraluminous infrared galaxies (ULIRGs) due to the mass accretion caused by the radiation drag. Based on this model, Kawakatu et al. (2003) also claim that protogasars, which have relatively small black holes and show only narrow emission lines, appear just after the ULIRG phase, $t \sim 10^8-10^9$ yr. However, such a protogasar phase might not be observed if the black hole grows at the supercritical accretion rate in the ULIRG phase. A seed black hole in the ULIRG can evolve to a supermassive black hole via the supercritical accretion by the end of the ULIRG phase, since the mass accretion rate from the circumnuclear regions can exceed the critical rate due to the effective radiation drag. Thus, we presumably observe the grown-up quasars after the ULIRG phase, and protogasars are obscured by huge amounts of the dust in the ULIRG.

As shown in Figure 6, our simulations show that the black hole can effectively grow in the case of metal-rich accreting gas. This implies that supercritical accretion flows tend to emerge in metal-rich objects like starburst galaxies or star-forming regions. This tendency is also consistent with the observed results that NLS1s, which are thought to be near-critical or supercritical accretion objects, are metal-rich objects (Nagao et al. 2002; Shenmer & Netzer 2002; Shenmer et al. 2004). It is proposed that NLS1s accrete at supercritical rates based on study of the optical band. (Kawaguchi 2003; Collin & Kawaguchi 2004).

The growth of the black hole may gradually slow down, since the density and the absorption opacity of the supercritical accretion flows are small around the massive black hole, if the normalized mass accretion rate does not change so much. Observational constraints are proposed by Yu & Tremaine (2002), by which most luminous quasars must be in the subcritical accretion phase. They showed that the luminous quasars have an energy conversion efficiency of $L/Mc^2 \gtrsim 0.1$, based on the study of the local black hole density and the luminosity function of the quasars. Since this efficiency is much smaller than 0.1 in the supercritical accretion flows, most luminous quasars must be in the subcritical accretion phase (see also Soltan 1982).

5.4. Future Work

5.4.1. Spectral Energy Distribution

Throughout the present study, we use the frequency-integrated energy equation of the radiation (see eq. [8]). By solving the
monochromatic radiation transfer equation, we can calculate the effective temperature profiles and show them in Figure 9. However, the emergent SEDs of the supercritical accretion flows are not a simple superposition of blackbody spectra with various effective temperatures (Ohsuga et al. 2003). The photons generated deep inside the disk have difficulty reaching the disk surface and thus move downward with the gas. Furthermore, most of the photons generated in the vicinity of the black hole will be immediately swallowed by the black hole and cannot contribute to the emergent SED. Frequency-dependent radiation-hydrodynamic simulations are necessary to resolve this issue.

According to the current simulations, the hot outflow with temperature of $10^9$–$10^{10}$ K appears above the disk. The Compton $\gamma$-parameter of this outflow region is comparable to or larger than unity, meaning that the inverse Compton scattering cannot be ignored. Seed soft photons emitted from the disk region should be Compton upscattered, producing high-energy, nonthermal emission that contributes to the emergent SED. The corona above the supercritical accretion disk has been claimed to fit the observed SEDs of NLS1s and the very high state BHs (Wang & Netzer 2003; Kubota & Done 2004), although the formation mechanism of the corona is poorly known. The hot outflow shown in our simulations might resolve this issue.

Here we should note that the Comptonization process is not included in our simulations. If the Compton cooling were ineffective, the hot outflow might be cooled to some extent. The bulk Comptonization may also contribute to the production of the high-energy photons because of large radial velocities, $v_r > 0.1c$, near the black hole. The detailed study of the SED with Comptonization is, however, beyond the scope of this paper and should be explored in future work.

### 5.4.2. Viscosity

We assumed that only the $r\phi$-component of the viscous stress tensor is nonzero while all the other components vanish, because the $r\phi$-component plays an essential role for the angular momentum transport within the disk. If the $r\theta$- and $\theta\phi$-components were nonzero, the structure near the disk surface might change. Figure 3 clearly shows abrupt velocity changes across the boundary between the disk and the outflow regions, which are thought to promote the K-H instability. The growth of this instability might be suppressed by the $r\theta$- and $\theta\phi$-components of viscosity. These components might also partially suppress the circular motion in the disk. Stone et al. (1999) discovered by the simulations of RIAFs that the flow structure differs from that given by Igumenshchev & Abramowicz (1999). It implies that the inclusion of the components of the viscous stress tensor suppresses the convection, since the $r\theta$-component is excluded in Stone et al. (1999). This will be examined in a future study.

More importantly, we need coupled MHD and RHD simulations, since the source of disk viscosity is likely to be of magnetic origins (Hawley et al. 2001; Machida et al. 2001; see Balbus 2003 for a review). Recently, local radiation MHD simulations of the accretion flow have been performed by Turner et al. (2003) and Turner (2004). We definitely need global simulations to include global field effects.

### 5.4.3. The FLD Approximation

Finally, we need to comment on the FLD approximation. It is known that the FLD approximation does not always give good results for the regions with moderate optical thickness. The FLD approximation is thought to be valid in the disk region, which has a large optical depth, but the spherical shape of the radiation energy density distribution above the disk might be artificial.

In addition, the radiation drag force cannot be treated in the FLD approximation, whereas it might play an important role for the transport of the angular momentum within the outflow region. Since the timescale of the angular momentum transport via the radiation drag, which is given by

$$\frac{M_0}{L/c^2} \approx (1.6 \, \text{s}) \left( \frac{L}{4L_E} \right)^{-1} \left( \frac{\rho}{10^{-8} \, \text{g cm}^{-3}} \right) \times \left( \frac{r}{200r_g} \right)^3 \left( \frac{M}{10 M_\odot} \right),$$

(30)

(where $M_\odot$ is the mass within the outflow region), is about 10 times longer than the escape time, $r/r_\odot = (0.2 \, \text{s})(r/200r_g)/(v_r/0.1c)^{-1}[M/(10 M_\odot)]$, roughly 10% of the angular momentum could be extracted in the outflow region. (The exact expressions for the radiation drag are found in the literature [Umemura et al. 1997; Fukue et al. 1997; Ohsuga et al. 1999].)

Radiation drag arises where there exists a large velocity difference between radiation sources and the irradiated matter. In the FLD approximation, however, the radiation flux is determined solely by the local gradient of the radiation energy, and photon trajectories cannot be properly considered. Thus, the FLD approximation cannot treat the radiation drag force, in principle. It would be better to solve the radiation transfer equations without using this approximation.

Begelman (2002) suggested that strong density inhomogeneities on scales much smaller than the disk scale height are formed due to the photon bubble instability in the radiation pressure–dominated accretion disks, and the disk can remain geometrically thin even as the maximum luminosity exceeds the Eddington luminosity by a factor of 100. However, the FLD approximation is not a suitable method for investigating the radiation fields in such inhomogeneous structure. The detailed study of the photon bubble instability should be explored in future work.

### 6. Conclusions

By performing the 2D-RHD simulations, we, for the first time, investigate the quasi-steady structure of the supercritical accretion flows around a black hole with particular attention being paid to the photon-trapping effects. We have obtained several new findings.

1. The quasi-steady structure of the supercritical flow is divided into two parts: the disk region (with mostly inflow) and the outflow region above the disk. The two regions are separated by a sharp density jump. The gas outflow driven by the strong radiation pressure is produced around the rotation axis. Furthermore, there exist velocity shears at the boundary, which cause K-H instability around the disk surface, producing the patchy structure as well as the circular motion within the disk region. Convection may also be responsible for such inhomogeneous structure.

2. The photon trapping plays an important role in the supercritical accretion regime. The advective energy transport is substantial, and the large number of photons generated inside the disk is swallowed by the black hole without being radiated away.

3. Our 2D-RHD model shows some differences from the slim-disk model. The slim-disk model assumes a simple convergence flow, while our simulations revealed rather complex gas motion and structure. We also found that the mass accretion rate is not
constant in space but decreases as matter accretes, roughly as $\dot{m} \propto r$, as a result of wind mass loss and circular motion. The calculated luminosity of the flows agrees more with the prediction of Paper I rather than with that of the slim-disk model.

4. The emission of the supercritical accretion flows is moderately collimated. The apparent luminosity could become more than 10 times larger than the Eddington luminosity. The supercritical accretion flows would be identified as high-$L/L_E$ objects in the face-on view but not, if the viewing angle is large, for which self-occultation tends to reduce the total luminosity and the maximum flow temperature.

5. The mass accretion rate increases with the absorption opacity (metallicity) of the accreting matter. It implies that the black hole tends to grow up faster in metal-rich regions, such as in starburst galaxies or star-forming regions. In addition, the growth of the black hole may gradually slow down, since the density and the absorption opacity of the supercritical accretion flows are small around the massive black hole, if the normalized mass accretion rate is kept constant.

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