THE LARGE-SCALE MAGNETIC FIELDS OF ADVECTION-DOMINATED ACCRETION FLOWS

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ABSTRACT

We calculate the advection/diffusion of the large-scale magnetic field threading an advection-dominated accretion flow (ADAF) and find that the magnetic field can be dragged inward by the accretion flow efficiently if the magnetic Prandtl number $\mathcal{P}_m = \eta/\nu \sim 1$. This is due to the large radial velocity of the ADAF. It is found that the magnetic pressure can be as high as $\sim 50\%$ of the gas pressure in the inner region of the ADAF close to the black hole horizon, even if the external imposed homogeneous vertical field strength is $\lesssim 5\%$ of the gas pressure at the outer radius of the ADAF, which is caused by the gas in the ADAF plunging rapidly to the black hole within the marginal stable circular orbit. In the inner region of the ADAF, the accretion flow is significantly pressured in the vertical direction by the magnetic fields, and therefore its gas pressure can be two orders of magnitude higher than that in the ADAF without magnetic fields. This means that the magnetic field strength near the black hole is underestimated by assuming equipartition between magnetic and gas pressure with the conventional ADAF model. Our results show that the magnetic field strength of the flow near the black hole horizon can be more than one order of magnitude higher than that in the ADAF at $\sim 3R_g$ ($R_g = 2GM/c^2$), which implies that the Blandford–Znajek mechanism could be more important than the Blandford–Payne mechanism for ADAFs. We find that the accretion flow is decelerated near the black hole by the magnetic field when the external imposed field is strong enough or the gas pressure of the flow is low at the outer radius, or both. This corresponds to a critical accretion rate, below which the accretion flow will be arrested by the magnetic field near the black hole for a given external imposed field. In this case, the gas may accrete as magnetically confined blobs diffusing through field lines in the region very close to the black hole horizon, similar to those in compact stars. Our calculations are also valid for the case that the inner ADAF connects to the outer cold thin disk at a certain radius. In this case, the advection of the external fields is quite inefficient in the outer thin disk due to its low radial velocity, and the field lines thread the disk almost vertically, while these field lines can be efficiently dragged inward by the radial motion of the inner ADAF.

Key words: accretion, accretion disks – galaxies: active – galaxies: jets – galaxies: magnetic fields

Online-only material: color figures

1. INTRODUCTION

The winds driven from the accretion disk through the magnetic field lines threading the disk have been considered as promising explanations for jets/outflows observed in different types of sources, such as active galactic nuclei (AGNs), X-ray binaries, and young stellar objects (see reviews in Spruit 1996, 2010; Königl & Pudritz 2000; Pudritz et al. 2007). In this model, the ordered magnetic field corotates with the gases in the disk, and the jets are powered by the gravitation energy released by accretion of the gases through the ordered field threading the disk (Blandford & Payne 1982). The ordered large-scale magnetic field threading the disk is the crucial ingredient in this model. In most previous works, the strength of the magnetic field is simply assumed to scale with the gas/radiation pressure of the accretion disk (e.g., Moderski & Sikora 1996; Ghosh & Abramowicz 1997; Livio et al. 1999; Armitage & Natarajan 1999; Nemmen et al. 2007; Wu & Cao 2008; McNamara et al. 2009). It was suggested that the external generated large-scale poloidal field (e.g., the field originating from the interstellar medium) would be dragged inward by the accretion plasma, while the field will diffuse outward at the same time (Bisnovatyi-Kogan & Ruzmaikin 1974, 1976; van Ballegooijen 1989; Lubow et al. 1994; Ogilvie & Livio 2001). In this case, the final steady magnetic field configuration can be derived, in which the inward advection of the field lines is balanced by the outward movement of field lines due to magnetic diffusion (Lubow et al. 1994). This means that the magnetic field configuration is predominantly determined by the radial velocity of the disk and magnetic diffusivity. As the radial velocity of the accretion disk is roughly proportional to the kinematic viscosity $\nu$, the magnetic field configuration is sensitive to the magnetic Prandtl number $\mathcal{P}_m = \eta/\nu \sim 1$, where $\eta$ is magnetic diffusivity.

Parker (1979) argued that $\nu \sim \eta \sim l v_t$ ($l$ is the largest eddy size and $v_t$ is turnover velocity), and $\mathcal{P}_m = \eta/\nu \sim 1$, is expected in isotropic turbulence. This issue was explored by different authors using numerical simulations (e.g., Yousef et al. 2003; Lesur & Longaretti 2009; Fromang & Stone 2009; Guan & Gammie 2009), which all suggest that the magnetic Prandtl number should be around unity. A suitable magnetic field configuration is crucial for launching a jet from the accretion disk. More specifically, the angle of field lines inclined to the mid-plane of the disk is required to be less than $\sim 60^\circ$ for launching jets from a Keplerian cold disk (Blandford & Payne 1982; Cao & Spruit 1994). This critical angle could be larger than $60^\circ$ for the accretion disk surrounding a rapidly spinning black hole (Cao 1997), which indicates that the spin of the black hole may help to launch jets centrifugally by cold magnetized disks (Cao 1997; Sadowski & Sikora 2010). Lubow et al. (1994) explored the final steady magnetic field configuration with the balance between advection and diffusion of the large-scale magnetic field and found that significant inward dragging of fields occurs only if $\mathcal{P}_m \lesssim H/R$ is satisfied ($H$ is the scale height of the disk at radius $R$). This means that the dragging of the external fields is always unimportant for thin disks, as $\mathcal{P}_m \sim 1$ is suggested; however, the advection of magnetic fields may be efficient for
a geometrically thick accretion disk with $H \sim R$ (e.g., Ferreira & Petrucci 2011). An alternative model was suggested by Spruit & Uzdensky (2005) for advection of the external field in the disk, in which turbulent diffusion is reduced by bundles of the large-scale magnetic field (Stehle & Spruit 2001). Lovelace et al. (2009) suggested that the field can be efficiently advected inward based on the assumption of the surface layer of the accretion disk to be nonturbulent. The general relativistic magnetohydrodynamic simulation of an accretion torus embedded in a large-scale magnetic field showed that the mass is accreted mainly within the accretion disk, and the magnetic field flux is carried by the motions in the low-density corona (Beckwith et al. 2009).

A low mass accretion rate $\dot{m}$ may lead the accretion flows to be advection-dominated (Narayan & Yi 1994; 1995). Advection-dominated accretion flows (ADAFs) are suggested to be present in low-luminosity AGNs (see Narayan 2002 for a review, and references therein), FR I radio galaxies (e.g., Ghisellini & Celotti 2001; Cao & Rawlings 2004), or BL Lac objects (e.g., Cao 2003). The ADAF model can successfully explain most observational features of low-luminosity AGNs and black hole X-ray binaries in the low/hard state (e.g., Lasota et al. 1996; Gammie et al. 1999; Quataert et al. 1999; Xu & Cao 2009). ADAFs are hot and geometrically thick, and have relatively higher radial velocities than thin accretion disks (Narayan & Yi 1994, 1995). This implies that the advection of the external fields in ADAFs may be more efficient than that in thin disks.

In this work, we explore the advection/diffusion of the large-scale magnetic fields threading an ADAF, in which the compression of the accretion flow in the vertical direction by the magnetic field is properly considered.

2. MODEL

2.1. The Structure of the ADAF

The dynamics of a steady ADAF with large-scale magnetic fields are described by a set of differential equations, namely, the continuity equation, the radial and azimuthal momentum equations, and the energy equation. In this work, we use cylindrical coordinates $(R, \phi, z)$.

The continuity equation is

$$\frac{d}{dR}(\rho RH_v) = 0,$$

where mass loss rate in the winds from the ADAF is neglected.

The radial momentum equation is

$$\frac{dR}{v_R} \frac{dv_R}{dR} = -\left(\Omega_k^2 - \Omega^2\right) R - \frac{1}{\rho} \frac{d}{dR} \left(\rho c_s^2\right) + \frac{B_R^2 B_z}{2\pi \Sigma} - \frac{B_z H \partial B_z}{2\pi \Sigma} \partial R,$$

where $B_R^2$ is the radial component of the large-scale magnetic fields at the disk surface, and the pseudo-Newtonian potential is adopted to simulate the gravity of a non-rotating black hole (Paczyński & Wiita 1980). The term $B_R^2 B_z / 2\pi \Sigma$ is the radial component of the force caused by curvature of the field lines, which is derived from $\int H(B_z/4\pi \Sigma) \partial B_R/\partial z dz$ by approximating $\partial B_R/\partial z \simeq B_R^2 / H$. The Keplerian angular velocity is given by

$$\Omega_k^2(R) = \frac{GM}{(R - R_g)^3 R}.$$

The radial magnetic fields will be sheared into azimuthal fields by the differential rotation of the flow, which leads to magnetorotational instability (MRI) and the MRI-driven turbulence (Balbus & Hawley 1991, 1998). For simplicity, we assume that the conventional $\alpha$-viscosity can describe the angular momentum transportation in the accretion flow caused by MRI-driven turbulence. The angular momentum equation is

$$\frac{d\Omega}{dR} = \frac{v_R(j - j_m)}{\nu R^2},$$

where $j$ is the specific angular momentum of the flow, $j_m$ is the specific angular momentum of the gas swallowed by the black hole, and the $\alpha$-viscosity,

$$\nu = \alpha c_s H,$$

is adopted. The angular momentum equation (4) reduces to an algebraic equation,

$$j = \frac{v_R R j_m}{\alpha c_s H + v_R R},$$

by assuming $d\Omega / dR \simeq -\Omega / R$.

The energy equation is

$$\Sigma v_R T \frac{d s}{dR} = Q^+ - Q^-,$$

where $T$ is the temperature and $s$ is the entropy of the gas in the ADAF. Equation (7) can be re-written as

$$\frac{2}{(\gamma - 1)c_s dR} \frac{dc_s}{dR} = \frac{1}{\rho dR} + \frac{f v_R}{\alpha c_s^2 R^2 H}(j - j_m)^2,$$

where $a$ and $b$ are constants.
where $\gamma$ is the ratio of specific heats and the parameter $f$ describes the fraction of the dissipated energy advected in the flow. Substituting the continuity equation (1) into Equation (8), we have

$$
\frac{2}{(\gamma - 1)c_s} \frac{dc_s}{dR} = -\frac{1}{R} \frac{1}{H} \frac{dH}{dR} - \frac{1}{v_R} \frac{dv_R}{dR} + \frac{f v_R}{\alpha c_s R^2 H} (j - j_m)^2.
$$

(9)

The vertical structure of the accretion flow is significantly altered in the presence of large-scale magnetic fields. In this work, we calculate the vertical structure of the ADAF following the approach given in Cao & Spruit (2002). Assuming the accretion flow to be isothermal in the vertical direction, we can calculate the vertical structure of the ADAF with

$$
\frac{c_s^2 d\rho(z)}{dz} = -\rho(z) \Omega_K^2 z - \frac{B_R \partial B_R}{4\pi} \frac{\partial z}{\partial z} + \frac{B_R \partial B_z}{4\pi} \frac{\partial \Omega}{\partial R},
$$

(10)

provided the shape of the magnetic field lines in the accretion flow is known. An additional term containing $\partial B_z/\partial R$, which is the vertical component of the force caused by curvature of the field lines, is included, because $H \sim R$ holds at least in the outer region of ADAFs. For geometrically thin accretion disks ($H \ll R$), $\partial B_z/\partial z \gg \partial B_z/\partial R$, and the term containing $\partial B_z/\partial R$ is therefore neglected (Cao & Spruit 2002).

In principle, the field line shape is computable by solving the radial and vertical momentum equations with suitable boundary conditions (see Cao & Spruit 2002 for a detailed discussion). For the isothermal case, an approximate analytical expression is proposed for the shape of the field lines in the flow:

$$
R - R_i = \frac{H}{\kappa_0 \eta_i} \left(1 + \eta_i^2 + \eta_i^2 z^2 H^{-1}(2)\right)^{1/2} - \frac{H}{\kappa_0 \eta_i^2} (1 - \eta_i^2 z^2 H^{-1/2}),
$$

(11)

where $H$ is the scale height of the disk, $\eta_i = \tanh(1)$, and the inclination of the field line $\kappa_i = B_z/B_R$ at the disk surface $z = H$ (Cao & Spruit 2002). This expression can reproduce the basic features of the Kippenhahn–Schlüter model (Kippenhahn & Schlüter 1957) well either for weak or strong field cases. The vertical distribution of the accretion disk with magnetic fields is not exactly a Gaussian distribution of an isothermal disk. In this case, we define the disk scale height $H$ as $\rho(H) = \rho(0) \exp(-1/2)$, the disk scale height $H$ can then be evaluated numerically with the given magnetic field line shape (11). As done by Cao & Spruit (2002), we use a fitting formula to calculate the scale height of the disk in the rest of this work,

$$
\frac{H}{R} = \frac{1}{2} \left( \frac{4c_s^2}{R^2 \Omega_K^2} + f_1 z \right)^{1/2} - \frac{1}{2} f_1,
$$

(12)

where

$$
f_1 = \frac{1}{2(1 - e^{-1/2}) \kappa_0} \left( \frac{B_z^2}{4\pi \rho R H \Omega_K^2} + \frac{\xi_{\text{mz}} B_z^2}{4\pi \rho R^2 \Omega_K^2} \right),
$$

(13)

and $\xi_{\text{mz}}$ is defined as

$$
\frac{\partial B_z}{\partial R} = -\xi_{\text{mz}}(R) \frac{B_z}{R},
$$

(14)

which can be calculated with the balance between the advection and diffusion of the fields in the accretion flow (see Section 2.2). We find that this formula can reproduce the numerical results quite well. We note that Equation (12) reduces to $H = c_s/\Omega_K$ in the absence of magnetic field.

We define dimensionless quantities by

$$
r = \frac{R}{R_g}, \quad \tilde{H} = \frac{H}{R}, \quad \tilde{v}_R = \frac{v_R}{c}, \quad \tilde{c}_s = \frac{c_s}{c}, \quad \tilde{j} = \frac{j}{R_g c}, \quad \tilde{\Omega} = \frac{\Omega}{R_g^2 c}, \quad \tilde{\Sigma} = \frac{\Sigma}{\Sigma(R_{out})}, \quad \tilde{B}_z = \frac{B_z}{B_0}, \quad \tilde{B}_R = \frac{B_R}{B_0}, \quad \rho_0 = \rho(R_{out}) \frac{B_0^2}{8\pi},
$$

(15)

where $B_0$ is the strength of the putative external imposed homogeneous vertical magnetic fields and $R_{out}$ is the outer radius of the accretion disk.

Differentiating Equation (12), we obtain

$$
\frac{1}{\tilde{H}} \frac{d\tilde{H}}{dR} = \frac{2c_s}{f_3 f_4 f_i H \Omega_K^2 (2\tilde{H} + f_1 R)} \frac{d\tilde{c}_s}{dR} - \frac{R f_1}{f_3 f_4 f_i v_R (2\tilde{H} + f_1 R)} \frac{dv_R}{dR}
$$

$$
- \frac{1}{f_4} \left[ \frac{2c_s^2}{f_3 f_4 f_i H (2\tilde{H} + f_1 R)} - \frac{2f_i R}{f_3 f_4 (2\tilde{H} + f_1 R) \Omega_K^2} \frac{d\Omega_K}{dR} \right] \frac{1}{f_3 f_4 R H \Omega_K^2 (2\tilde{H} + f_1 R)} - \frac{2c_s^2}{f_3 f_4 R H \Omega_K^2 (2\tilde{H} + f_1 R)} - \frac{2c_s^2}{f_3 f_4 R H \Omega_K^2 (2\tilde{H} + f_1 R)},
$$

(16)
The radial momentum equation is

\[ v_R = \frac{dR}{dt} \]

where

\[ f_2 = \frac{H}{8\pi(1-e^{-1/2})^2(2H+f_1R)\rho R \Omega_K^3} \left( 2\tilde{\xi}_m B_R^2 B_z + \frac{\xi_B^2 B_R B_z}{R} + \frac{\xi_B s_{\tilde{\xi}_m} B_R^2 B_z}{R} \right), \quad (17) \]

\[ f_3 = 1 - \frac{B_R^2 B_z}{8\pi(1-e^{-1/2})(2H+f_1R)\rho H \Omega_K^2}, \quad (18) \]

and \( \tilde{\xi}_m \) is defined as

\[ \frac{\partial B_R^2}{\partial R} = -\tilde{\xi}_m(R) \frac{B_R^2}{R}. \quad (20) \]

Substituting Equation (16) into the energy Equation (9), we finally obtain

\[ \frac{1}{c_s} \frac{dR}{dt} = \frac{1}{f_3 f_5} \left[ \frac{2c_s^2}{f_3 H \Omega_K^2(2H+f_1R) - \frac{2f_1 R}{f_3(2H+f_1R)}} \right] \frac{1}{\Omega_K} \frac{d \Omega_K}{dR} \]

\[ - \frac{f_1 R}{f_3 f_5(2H+f_1R)} \left( \frac{1}{f_3 f_5} \frac{dv_R}{dR} + \frac{fv_R}{f_4 f_5} R^2 H (j - j_m)^2 \right) \]

\[ + \frac{c_s^2}{f_3 f_4 f_5 R H \Omega_K^2(2H+f_1R)} + \frac{(f_2 R - f_3 H - f_5 f_4 H)(2H+f_1R) - f_1 RH}{f_3 f_4 f_5 R H \Omega_K^2(2H+f_1R)}, \quad (21) \]

where

\[ f_5 = \frac{2}{\gamma - 1} + \frac{2c_s^2}{f_3 f_5 H \Omega_K^2(2H+f_1R)}. \quad (22) \]

In the same way, we can re-write the radial momentum equation as

\[ \frac{v_R^2 - v_{R,s}^2}{v_R} \frac{dR}{dt} = \frac{c_s^2(f_6 - f_5)}{f_4 f_5} \left[ \frac{2c_s^2}{f_3 H \Omega_K^2(2H+f_1R) - \frac{2f_1 R}{f_3(2H+f_1R)}} \right] \frac{1}{\Omega_K} \frac{d \Omega_K}{dR} \]

\[ - R \Omega_K^2 + \frac{j^2}{R^2} + \frac{c_s^2}{R} + \frac{B_R^2 B_z}{2\pi \Sigma} + \frac{\xi_B B_z^2 H}{2\pi \Sigma R} + \frac{(f_2 R - f_3 H)(2H+f_1R) - f_1 RH c_s^2}{f_3 f_4 f_5 R H (2H+f_1R)} \]

\[ - \frac{f_6 c_s^4}{f_3 f_4 f_5 R H \Omega_K^2(2H+f_1R)} + \frac{2c_s^4}{f_3 f_4 f_5 R H \Omega_K^2(2H+f_1R)} \]

\[ + \frac{c_s^4}{f_3 f_4 f_5 R H (2H+f_1R)} + \frac{f_6 f v_R}{f_3 f_4 f_5 R H (j - j_m)^2}. \quad (23) \]

where

\[ v_{R,s} = \left( 1 - \frac{f_1 R}{f_3 f_4(2H+f_1R)} - \frac{f_6}{f_5} \left[ 1 - \frac{f_1 R}{f_3 f_4(2H+f_1R)} \right] \right)^{1/2} c_s, \quad (24) \]

and

\[ f_6 = \frac{2c_s^2}{f_3 f_4 H \Omega_K^2(2H+f_1R)} - 2. \quad (25) \]

It is found that Equation (24) reduces to

\[ v_{R,s} = \left( \frac{2\gamma}{\gamma + 1} \right)^{1/2} c_s \quad (26) \]

in the absence of magnetic field. We re-write the equations in dimensionless form as follows.

The radial momentum equation is

\[ \frac{1}{\tilde{v}_R} \frac{d\tilde{v}_R}{dr} = \frac{1}{\tilde{v}_R^2 - \tilde{v}_{R,s}^2} \left\{ \frac{c_s^2(f_6 - f_5) \tilde{c}_s^2}{f_4 f_5} \left[ \frac{2\tilde{c}_s^2}{f_3 \tilde{\tilde{\xi}} \Omega_K^2 \tilde{\tilde{H}}(2\tilde{H} + f_1)} - \frac{2f_1}{f_3(2\tilde{H} + f_1)} \right] \left( \frac{1}{\tilde{\Omega}_K} \frac{d\tilde{\Omega}_K}{dr} - \tilde{\tilde{\Omega}}_K^2 \right) \right\} \]

\[ + \frac{j^2}{r^3} + \frac{c_s^2}{r} + \frac{2r \tilde{c}_s^2 \tilde{v}_R B_R^2 B_z}{r^2 \tilde{H} \tilde{\tilde{\xi}}_m \tilde{\tilde{H}}_m \tilde{\tilde{B}}_z^2} \]

\[ + \frac{2r \tilde{c}_s^2 \tilde{v}_R \tilde{\tilde{\xi}}_m \tilde{\tilde{B}}_z^2}{r^2 \tilde{H} \tilde{\tilde{\xi}}_m \tilde{\tilde{B}}_z^2} \left[ 1 - \frac{f_2 - f_3 \tilde{H}(2\tilde{H} + f_1) - f_1 \tilde{H}}{f_3 f_4 f_5 \tilde{H}(2\tilde{H} + f_1)} + \frac{f_6 f \tilde{v}_R}{f_3 f_4 f_5 \tilde{H}(j - j_m)^2} \right], \quad (27) \]
where

\[ \tilde{v}_{R,s} = \left\{ 1 - \frac{f_1}{f_3 f_4 (2 \tilde{H} + f_1)} \right\}^{1/2} \left( 1 - \frac{f_1}{f_3 f_4 (2 \tilde{H} + f_1)} \right) \]  

\[ \tilde{\Omega}_K = \frac{1}{\sqrt{2} r^{1/2} (r - 1)}; \quad \frac{1}{\tilde{\Omega}_K} \frac{d \tilde{\Omega}_K}{dr} = -\frac{3r - 1}{2r (r - 1)}. \]  

The functions \( f_1 \) - \( f_6 \) in the dimensionless form are given by

\[ f_1 = \frac{1}{1 - e^{-1/2}} \tilde{c}_s \left\{ \frac{2 c_s^2 H^2 \tilde{v}_R \tilde{B}_R^2 \tilde{B}_z^2}{H \tilde{c}_s \tilde{v}_{R,\text{out}} \tilde{\Omega}_K^2} \left( \frac{1}{\tilde{H} \tilde{c}_s^2} + \frac{\tilde{c}_s}{\kappa_0} \right) \right\}, \]  

\[ f_2 = -\frac{1}{1 - e^{-1/2}} \tilde{c}_s \tilde{v}_R \tilde{v}_{R,\text{out}} \tilde{\Omega}_K^2 \tilde{B}_R \tilde{B}_z \left( \frac{2 \tilde{c}_s \tilde{B}_R \tilde{B}_z}{\tilde{H} \kappa_0} + \tilde{c}_s \tilde{B}_R \tilde{B}_z + \tilde{c}_s \tilde{B}_R \tilde{B}_z \right), \]  

\[ f_3 = 1 - \frac{\tilde{c}_s \tilde{v}_R \tilde{B}_R \tilde{B}_z}{1 - e^{-1/2} \tilde{c}_s \tilde{v}_{R,\text{out}} \tilde{\Omega}_K^2 \kappa_0}, \]  

\[ f_4 = 1 + \frac{f_1}{f_3 (2 \tilde{H} + f_1)} \]  

\[ f_5 = \frac{2}{\gamma - 1} + \frac{2 c_s^2}{f_3 f_4 r^2 \tilde{\Omega}_K^2 (2 \tilde{H} + f_1)}, \]  

\[ f_6 = \frac{2 c_s^2}{f_3 f_4 r^2 \tilde{\Omega}_K^2 (2 \tilde{H} + f_1)} - 2. \]  

The dimensionless scale height of the disk is

\[ \tilde{H} = \frac{1}{2} \left( \frac{4 c_s^2}{r^2 \tilde{\Omega}_K^2} + f_1^2 \right)^{1/2} - \frac{1}{2} f_1. \]  

2.2. Magnetic Field Configuration of the ADAF

The advection/diffusion of large-scale (poloidal) magnetic fields is described by

\[ \frac{\partial}{\partial t} [R \psi(R, 0)] = -v_R \frac{\partial}{\partial R} [R \psi(R, 0)] - \frac{4 \pi \eta}{c} \frac{R}{2 \tilde{H}} \int_{-\tilde{H}}^{\tilde{H}} J_\phi(R, z_h) dz_h, \]  

where \( J_\phi(R, z_h) \) is the current density at \( z = z_h \) above/below the mid-plane of the disk, \( \psi(R, z) \) is the azimuthal component of the magnetic potential, and the vertical velocity \( v_z \) of the flow is neglected. Assuming the azimuthal current distribution in the vertical
direction to be homogeneous, we have

\[ J_{\phi}(R, z_h) = \frac{J^S_{\phi}(R)}{2H}, \]  

(40)

where \( J^S_{\phi}(R) \) is the surface current density at \( R \) in the disk. As in the work by Lubow et al. (1994), the magnetic field potential \( \psi(R, z) = \psi_d(R, z) + \psi_{\infty}(R, z) \), where \( \psi_d(R, z) \) is contributed by the currents in the accretion flow, and \( \psi_{\infty}(R) = B_0 R / 2 \) is the external imposed homogeneous vertical field, which can be regarded as being contributed by the currents at infinity (see Lubow et al. 1994, for details). The potential \( \psi_d \) is related to \( J^S_{\phi}(R) \) with

\[
\psi_d(R, z) = \frac{1}{c} \int_{R_{\text{in}}}^{R_{\text{out}}} R' dR' \int_0^{2\pi} \cos \phi' d\phi' \int_{-H}^{H} \frac{J_{\phi}(R', z_h)}{[R'^2 + R^2 + (z - z_h)^2 - 2RR' \cos \phi']^{1/2}} d\phi. \]

(41)

Differentiating Equation (41), we have

\[
\frac{\partial}{\partial R}[R \psi_d(R, z)] = \frac{1}{2hc} \int_{R_{\text{in}}}^{R_{\text{out}}} J^S_{\phi}(R') R' dR' \int_0^{2\pi} \cos \phi' d\phi' \int_{-H}^{H} \frac{R'^2 + (z - z_h)^2 - RR' \cos \phi'}{[R'^2 + R^2 + (z - z_h)^2 - 2RR' \cos \phi']^{3/2}} d\phi. \]

(42)

For steady case, i.e., \( \partial/\partial t = 0 \), Equation (39) becomes

\[
- \frac{\partial}{\partial R}[R \psi_d(R, 0)] - \frac{2\pi \alpha c_s R}{c \nu_R} \mathcal{P}_m J^S_{\phi}(R) = B_0 R, \]

(43)

where the magnetic Prandtl number is defined as \( \mathcal{P}_m = \eta / \nu \). This equation can reduce to a set of linear algebraic equations, i.e.,

\[
- \sum_{j=1}^{n} P_{ij} J^S_{\phi}(R_j) \Delta R_j - \frac{2\pi \alpha c_s(R_j) R_j}{c \nu(R_j)} \mathcal{P}_m J^S_{\phi}(R_i) = B_0 R_i, \]

(44)

by using Equation (42), where \( J^S_{\phi}(R_j) \) is the surface current density of the ring at radius \( R_j \) in the accretion flow, \( \Delta R_j \) is the width of the ring, and

\[
P_{ij} = \frac{R_j}{2Hc} \int_0^{2\pi} \cos \phi' d\phi' \int_{-H}^{H} \frac{R_j^2 + z_h^2 - R_i R_j \cos \phi'}{[R_j^2 + R^2 + z_h^2 - 2R_i R_j \cos \phi']^{3/2}} d\phi. \]

(45)

The surface current density \( J^S_{\phi}(R) \) can be calculated by solving a set of linear algebraic equations provided the structure of the ADAF is given with the specified magnetic Prandtl number \( \mathcal{P}_m \) (see Lubow et al. 1994 for details), and therefore the configuration of the large-scale magnetic fields is available with the derived potential \( \psi \):

\[
B_R(R, z) = - \frac{\partial}{\partial z} \psi(R, z), \]

(46)

and

\[
B_z(R, z) = \frac{1}{R} \frac{\partial}{\partial R}[R \psi(R, z)]. \]

(47)

2.3. Boundary Conditions

In order to carry out the calculation of the structure of the ADAF, we have to specify the boundary conditions at the outer radius \( R_{\text{out}} \). There are six model parameters: viscosity parameter \( \alpha \), the outer radius of the accretion disk \( r_{\text{out}} \), the degree of advection \( f \), specific angular momentum \( j \) at \( r_{\text{out}} \), the temperature of the gas \( \Theta_{\text{out}} = c_s / R \Omega_k = \sqrt{2(r - 1) \zeta_0 / r^{1/2}} \) at \( r_{\text{out}} \), and \( \beta_0 = 8\pi p(R_{\text{out}})/B_0^2 \), which relates the external imposed vertical magnetic field strength to the gas pressure of the disk at \( r_{\text{out}} \). In all the calculations, we fix the ADAF parameters: \( \alpha = 0.2, f = 0.95, j(r_{\text{out}}) = j_k(r_{\text{out}}) \), and \( \Theta_{\text{out}} = 0.25 \), which are typical for ADAFs. We find that the final results are insensitive to the values of these parameters. We integrate Equations (27) and (31) inward from \( R = R_{\text{out}} \) numerically. The parameter \( j_k \), the specific angular momentum of the gas swallowed by the black hole, is tuned carefully till the derived global solution passes the sonic point smoothly.

The dynamical structure of the ADAF is coupled with the magnetic field, i.e., the structure of the ADAF is affected by the large-scale field threading the flow and vice versa. In principle, the derived solution has to pass through several critical points (e.g., the slow/fast
Figure 1. Large-scale poloidal magnetic field configuration of an ADAF (the magnetic field strength $\beta_0 = 200$ is adopted). The dotted line is the scale height of the ADAF. The outer radius of the disk is assumed to be $R_{\text{out}} = 100 R_g$. The magnetic field configuration plotted in the upper panel is calculated for the magnetic Prandtl number $Pr_m = 1$, while the lower panel is for $Pr_m = 1.5$. Every magnetic field line corresponds to a certain value of the stream function $R\psi(R, z)$ with the lowest $[R\psi(R, z)]_{\text{min}} = 0.04 R_{\text{out}} \psi_{\infty}(R_{\text{out}}, 0)$ increasing by $\Delta R\psi(R, z) = 0.04 R_{\text{out}} \psi_{\infty}(R_{\text{out}}, 0)$ per line, where $\psi_{\infty}(R_{\text{out}}) = B_0 R_{\text{out}}/2$.

Figure 2. Same as Figure 1, but the outer radius of the disk $R_{\text{out}} = 1000 R_g$ and $\beta_0 = 20$ are adopted. Every magnetic field line corresponds to a certain value of the stream function $R\psi(R, z)$ with the lowest $[R\psi(R, z)]_{\text{min}} = 0.005 R_{\text{out}} \psi_{\infty}(R_{\text{out}}, 0)$ increasing by $\Delta R\psi(R, z) = 0.04 R_{\text{out}} \psi_{\infty}(R_{\text{out}}, 0)$ per line.

The configuration of the large-scale magnetic fields is plotted in Figures 1 and 2. The structure of ADAFs is altered due to the presence of such large-scale magnetic fields, which is plotted in Figures 3–6. It is found that the radial velocity of the accretion flow changes little except in the region of the accretion flow close to the black hole horizon (see Figures 3 and 4). The accretion flow is decelerated near the black hole by the magnetic field when $\beta_0$ is relatively small (see Figures 3–6), which implies that the accretion flow can be trapped by the magnetic field provided the external imposed field is strong enough. We find that the accretion flow may pass the magneto-sonic points and the Alfvén point. In this work, we first calculate the structure of the ADAF without magnetic field, and the magnetic field configuration/strength is then calculated based on the derived ADAF structure. The structure of the ADAF is re-calculated with this derived field configuration/strength. The calculations of the dynamical structure of the ADAF are decoupled with those of the field configuration, and the global solution of the ADAF is therefore only required to pass the sonic point. The calculations are iterated till the solutions converge. We find that the final solution is available usually after three or four iterations. This is because the dynamics of the ADAF is altered little by the magnetic field except in the inner edge of the flow.

3. RESULTS

The configuration of the large-scale magnetic fields is plotted in Figures 1 and 2. The structure of ADAFs is altered due to the presence of such large-scale magnetic fields, which is plotted in Figures 3–6. It is found that the radial velocity of the accretion flow changes little except in the region of the accretion flow close to the black hole horizon (see Figures 3 and 4). The accretion flow is decelerated near the black hole by the magnetic field when $\beta_0$ is relatively small (see Figures 3–6), which implies that the accretion flow can be trapped by the magnetic field provided the external imposed field is strong enough. We find that the accretion flow may
Figure 3. Structure of the accretion flows. The magnetic Prandtl number \( P_m = 1 \) is adopted in all the calculations. The colored lines represent the results with different magnetic field strengths, i.e., \( \beta_0 = 20 \) (red), 40 (green), 200 (blue), and 400 (yellow). The black lines represent the structure of the accretion flow without magnetic fields. Upper panel: the radial velocities (solid lines) and the sound speeds (dotted lines) as functions of radius. Middle panel: the scale heights of accretion flows as functions of radius. Lower panel: the specific angular momenta as functions of radius. The black dashed line represents the Keplerian specific angular momentum. (A color version of this figure is available in the online journal.)

be disrupted by the magnetic field when the parameter \( \beta_0 \) is lower than a critical value \( \beta_{0,\text{crit}} \), provided all other disk parameters are fixed. The critical values of \( \beta_0 \) as functions of outer radius of the ADAF are plotted in Figure 7. The vertical structure of the accretion flow is altered in the presence of magnetic fields, i.e., the scale height decreases significantly in the inner region of the accretion flow (see the middle panels in Figures 3–6). In Figures 8–11, we compare the relative importance of the gas pressure and the magnetic pressure in the accretion flow with different magnetic field parameters (i.e., \( \beta_0 \) and \( P_m \)). The ratios of the magnetic field strength in the accretion flows to the external imposed field strength as functions of radius are plotted in Figure 12.

4. DISCUSSION

The configurations of the magnetic fields show that the advection of magnetic fields is very efficient if \( P_m \sim 1 \) is adopted, unlike the thin disk cases (Lubow et al. 1994). This is because the radial velocity of ADAFs is much higher than that of the thin disk. Such configurations may help launching outflows/jets from ADAFs. For simplicity, we have not considered magnetically driven outflows from ADAFs in this work. The radial velocity of the ADAF will increase if the angular momentum carried away by the outflows is properly taken into account, which will enhance the advection of the magnetic fields in the ADAF. We have not considered the azimuthal component of magnetic field \( B_\phi \) in the radial momentum equation (2), which may affect the dynamics of the accretion flow. Our results show that the dynamics of the accretion flow is dominantly determined by the gas pressure term except in the region very close to the black hole horizon. The previous MHD simulations showed that the magnetic stress never exceeds, and is usually much lower than, the gas pressure in the mid-plane of the accretion flow (e.g., Hawley & Krolik 2001; Igumenshchev et al. 2003), which implies that the main conclusions of this work will not be altered even if the azimuthal field component is included in the radial momentum equation.

The ratio of the magnetic pressure to gas pressure can reach \( \sim 0.5 \) in the inner edge of the ADAF (see Figures 8–11). The radial velocity and the temperature distribution have been altered little in the presence of large-scale magnetic fields, except in the inner region of the ADAF (\( \lesssim 3R_g \)). However, the ADAF is vertically pressured by the magnetic fields, which leads to the gas pressure significantly increased in the inner region of the ADAF. It is found that the magnetic pressure of the ADAF in the region close to the black hole horizon can be two orders of magnitude higher than the gas pressure of the ADAF without magnetic fields (see the dashed lines in the lower panels of Figures 8 and 9). In most previous works, the magnetic field strength is estimated with the gas/radiation pressure in the accretion disk on the so-called equipartition assumption, and the confinement of the disk by the magnetic fields in the vertical direction of the flow has not been considered (e.g., Moderski & Sikora 1996; Ghosh & Abramowicz 1997; Livio et al. 1999; Armitage & Natarajan 1999; Nemmen et al. 2007; Wu & Cao 2008; McNamara et al. 2009). This means that the maximal jet power estimated based on the normal accretion disk models without considering magnetically confinement in vertical direction is around two orders of magnitude underestimated. The assumption that the strength of the magnetic fields near the black hole horizon should not differ significantly in strength with the accretion disk may be incorrect, at least for ADAFs (e.g., Livio et al. 1999). Our present
Figure 4. Same as Figure 3, but the outer radius of the accretion flow $R_{\text{out}} = 1000R_g$ is adopted. The colored lines represent the results with different magnetic field strengths, i.e., $\beta_0 = 2$ (red), 4 (green), 20 (blue), and 40 (yellow).

(A color version of this figure is available in the online journal.)

Figure 5. Similar to Figure 3, the results with different values of the magnetic Prandtl number are compared ($\beta_0 = 200$ is adopted). The red lines represent the results calculated with $\mathcal{P}_m = 1$, while the green lines are for those with $\mathcal{P}_m = 1.5$.

(A color version of this figure is available in the online journal.)
calculations are based on the pseudo-Newtonian potential, which can only simulate the gravitational potential of non-rotating black holes. However, the gas in the accretion flow plunges rapidly to the black hole in the region of the flow with \( R \lesssim R_{ms} \) (\( R_{ms} \) is the radius of the marginal stable circular orbit of the hole) for a rotating black hole, which is similar to the cases calculated in this work for non-rotating black holes. We believe the conclusion that the magnetic fields are significantly strengthened near the black hole horizon due to very large radial velocity of the gas in the flow should still hold even for rapidly spinning black holes. This implies that the efficiency of the Blandford–Znajek mechanism (Blandford & Znajek 1977) is significantly underestimated in most previous works (e.g., Moderski & Sikora 1996; Ghosh & Abramowicz 1997; Livio et al. 1999; Nemmen et al. 2007; Wu & Cao 2008; McNamara et al. 2009). The detailed calculations for rotating black holes are beyond the scope of this paper, and will be reported in our future work.

The scale height of the ADAF is sensitive to the external imposed homogeneous vertical magnetic field strength. We find that it decreases with increasing field strength, i.e., low-\( \beta_0 \) case (see the middle panels in Figures 3 and 4). The scale height also depends on the value of the magnetic Prandtl number \( P_m \). It increases with the value of \( P_m \) (see Figures 5 and 6), because magnetic fields are less strong for a higher \( P_m \) (see Figures 10 and 11). The relative thickness of the disk \( H/R \)}
Figure 8. Magnetic pressure/gas pressure of the accretion flow, and the inclination of field lines at the disk surface as functions of radius. The magnetic Prandtl number $P_m = 1$ is adopted in all the calculations. The different colored lines represent the results with different magnetic field strength, i.e., $\beta_0 = 20$ (red), 40 (green), 200 (blue), and 400 (yellow). Upper panel: the magnetic field pressure/gas pressure of the accretion flow as functions of radius. The black line represents the gas pressure of the ADAF without magnetic fields as a function of radius. Middle panel: the inclination of the magnetic field lines at the disk scale height. Lower panel: the ratio of magnetic pressure to the gas pressure in the accretion flow (solid lines) and the ratio of magnetic pressure to the gas pressure in the accretion flow without magnetic fields (dashed lines).

(A color version of this figure is available in the online journal.)

Figure 9. Same as Figure 8, but the outer radius of the accretion flow $R_{out} = 1000R_g$ is adopted. The different colored lines represent the results with different magnetic field strength, i.e., $\beta_0 = 2$ (red), 4 (green), 20 (blue), and 40 (yellow).

(A color version of this figure is available in the online journal.)

becomes very small (i.e., geometrically thin) in the inner edge of the ADAF, due to strong magnetic pressure. This implies that the optical depth in the vertical direction will increase significantly, and the temperature difference between ions and electrons will decrease in the inner edge of the ADAF, though no detailed radiative processes are included in our present calculations. Its
implications on the observational features of the objects accreting at low rates can be explored, if the detailed radiative processes are considered and the black hole mass and the mass accretion rate need to be specified, which is beyond the scope of this work.

It is interesting to find that the accretion flow is decelerated near the black hole by the magnetic field when $\beta_0$ is relatively small (see Figures 3 and 4). There is a critical value $\beta_{0,\text{crit}}$, below which the accretion flow is decelerated near the black hole. The value of $\beta_{0,\text{crit}}$ decreases with increasing outer radius $R_{\text{out}}$, and it decreases with increasing magnetic Prandtl number $\mathcal{P}_m$ (see Figure 7). This implies
that the accretion flow may be disrupted by the magnetic field in the inner region of the accretion flow, when the external imposed field is strong enough or the gas pressure of the flow is low at the outer radius, or both, which justifies the main assumption in the qualitative analysis on the magnetic arrested accretion disks (Narayan et al. 2003). In this case, the gas may accrete as magnetically confined blobs diffusing through field lines in the region very close to the black hole horizon (see Narayan et al. 2003, for a detailed discussion, and references therein), which is similar to the case in compact stars (e.g., Elsner & Lamb 1984; Kaisig et al. 1992). For an ADAF surrounding a black hole with a given external imposed homogeneous field, the value of \( \beta_{0,\text{crit}} \) corresponds to a certain accretion rate at \( R_{\text{out}} \), below which the accretion flow will be trapped by the magnetic field, for given temperature of the gas in the flow at \( R_{\text{out}} \). Our calculations show that the gas can alternatively be trapped by the magnetic field at outer radius of the accretion flow provided the magnetic pressure significantly exceeds the gas pressure, i.e., \( \beta_0 \ll 1 \) (see the green line in Figure 7 at large radii). We compare the amplification of the magnetic field advected in the accretion flows with different values of the disk parameters and find that the amplification of the field is predominantly determined by the value of \( \mathcal{P}_m \), while it is insensitive to the adopted value of \( \beta_0 \).

This implies that the strength of the field near the black hole horizon is mainly determined by the strength of the external imposed field and the outer radius of the ADAF, which is qualitatively consistent with the results for thin accretion disks in Lubow et al. (1994), though the advection of the field is rather inefficient in the thin disk case. This means that the strength of the magnetic field near the black hole can be estimated if the strength of the ordered field threading the ambient gas and the outer radius \( R_{\text{out}} \) of the ADAF are known. It is found that stronger magnetic field will be in the inner edge of the ADAF, if the gas starts to accrete from a larger radius (see Figure 12). This should be very useful for estimating the field strength near the black hole horizon in some low-luminosity AGNs, of which the properties of the ambient gas at the Bondi radius have been well measured with observations (e.g., Baganoff et al. 2003; Allen et al. 2006).

We note a series of works of numerical simulations on the magnetic fields in ADAFs (Igumenshchev & Narayan 2002; Igumenshchev et al. 2003; Igumenshchev 2006, 2008). They found that the initial uniform \( z \)-direction magnetic field can be efficiently dragged inward by the ADAF, and their derived magnetic field configuration is similar to that obtained in this work (see Figure 3 in Igumenshchev & Narayan 2002). The numerical simulations show that the magnetic field in the inner edge of the accretion can be strong enough to disrupt the accretion flow (Igumenshchev 2008), which agrees qualitatively with our calculations. Compared with these MHD simulations on ADAFs, our calculations explicitly indicate that the vertical scale height of the ADAF near the black hole is significantly reduced, and hence the field strength can be much higher than that estimated based on the conventional ADAF model.

In this work, the viscosity \( \alpha = 0.2 \) is fixed. The radial velocity of the ADAF in the outer region is roughly proportional to the kinetic viscosity \( \nu \), and the magnetic Prandtl number \( \mathcal{P}_m = \eta/\nu \), which implies that the advection/diffusion process in the accretion flow is almost independent of the adopted value of \( \alpha \). The dynamical structure of ADAFs is insensitive to the value of \( \mathcal{P}_m \) provided the accretion flow is advection dominated, i.e., \( 0.5 \lesssim f \lesssim 1 \). We also calculate the problem with different values of the disk parameters (e.g., \( f, \Theta_{\text{out}}, \) and \( \tilde{j}_{\text{out}} \)) and find that the main conclusions have not been altered.

The ADAF may connect to a thin accretion disk at a certain transition radius \( R_{\text{tr}} \). This is required by modeling on a variety of observations of AGNs/X-ray binaries (e.g., Esin et al. 1997; Quataert et al. 1999; Lu & Wang 2000; Cao 2003; Yuan & Narayan 2004; Xu & Cao 2009) and is also predicted by some theoretical model calculations (e.g., Abramowicz et al. 1995; Liu et al. 1999; RózaŃska & Czerny 2000; Spruit & Deufel 2002; Lu et al. 2004). Our calculations are also valid for the case that the inner ADAF connects to the outer cold thin disk at a certain radius. In this case, the advection of the external fields is quite inefficient in the outer thin disk due to its low radial velocity, and the field lines thread the disk almost vertically (see, e.g., Lubow et al. 1994), while these field lines can be efficiently dragged inward by the radial motion of the inner ADAF.
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