Disk Accretion onto a Magnetized Young Star and Associated Jet Formation

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Abstract

We investigated disk accretion onto a magnetized young star by performing MHD numerical simulations. We considered the case in which the stellar magnetosphere truncates the accretion disk carrying the interstellar magnetic field, and the star is magnetized in the same direction as the disk. In such a case, the interface between the accretion disk and the magnetosphere contains a magnetically neutral ring in the equatorial plane. The numerical results show that the disk accretion drives the magnetic reconnection between the magnetospheric field and the disk magnetic field, which allows mass transfer from the disk to the magnetosphere. Most of the transferred mass accretes to the star along the reconnected magnetospheric field, while the rest of the mass is accelerated to the bipolar directions by a Lorentz force along the stellar open magnetic field. This “reconnection-driven” jet is further accelerated magneto-centrifugally due to stellar rotation, and corresponds to optical jets from young stars. The magnetic braking as a reaction of the magneto-centrifugal acceleration of the jet may explain the observed slow rotations of young stars in the disk-accretion stage.

Key words: Accretion, accretion disks — Astrophysical jets — Interstellar: medium — Magnetohydrodynamics — Stars: pre-main-sequence

1. Introduction

1.1. Jets from Young Stellar Objects and Models Proposed Thus-Far

Many observations suggest that well-collimated and high-velocity jets are ejected from the vicinity of young stellar objects (YSOs), which are considered to be in the disk-accretion stage. In some cases, we can resolve them spatially as “optical jets” in optical lines, or also recently by VLBI radio observations (Mundt, Fried 1983; Mundt et al. 1987; Reipurth, Heathcote 1993). On the other hand, spectroscopic observations also suggest that the morphology of spatially unresolved energetic winds from YSOs is bipolar rather than spherically symmetric (Appenzeller et al. 1984; Edwards et al. 1993a).

These collimated jets are often surrounded by more extended and larger scale outflows, high-velocity neutral winds (Lizano et al. 1988; Koo 1989; Masson et al. 1990), and molecular bipolar outflows (Snell et al. 1980; see review by Lada 1985) (figure 1a). For these extended flows, models concerning the stellar activities were proposed initially (De Campli 1981; Hartmann et al. 1982; Draine 1983; Hartmann, McGregor 1982). Recently, these extended flows are considered to come from the extended region of the surrounding accretion disk, and many people have developed steady magneto-centrifugal disk wind models and have applied them to this problem (Blandford, Payne 1982; Pudritz, Norman 1983, 1986; Königl 1989; Camenzind 1990; Wardle, Königl 1993). On the other hand, Uchida and Shibata (1985) proposed a model based on considering the interaction between the disk and the large-scale magnetic field penetrating it by using 2.5-dimensional MHD simulations. They showed that a supersonic spinning jet with a cylindrical shell shape is produced by both the toroidal magnetic-pressure gradient and the magneto-centrifugal effect (also cf. Shibata, Uchida 1986; Kudo, Shibata 1995; Lovelace et al. 1991).

As for the collimated jets considered to come from a region very close to the star, some people have proposed

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models in which the spherical stellar wind is collimated by the stellar open lines of force or by the surrounding outflows (Shu et al. 1988; Kwan, Tademaru 1988). However, there is growing evidence that these collimated jets are driven by disk accretion rather than stellar activities, for example, the correlation between the mass loss rate and the mass accretion rate (Lada 1985; Cabrit et al. 1990; Lavreault 1988; Edwards et al. 1993a). Furthermore, a comparison of Classical T Tauri Stars (CTTSs) with weak-line T Tauri stars indicates that while both have the same stellar characteristics, only CTTSs, which are considered to be in the disk-accretion stage, show evidence of jets (Edwards et al. 1993a). Therefore, their origins should be around the inner edge of the disk or the interface between the disk and the star. Torbett (1984) and Pringle (1989) considered the thermal pressure or the toroidal magnetic pressure amplified by velocity shear at the boundary layer as being the driving force of the jets.

Recently, another restriction to the model construction is given by VLBI and X-ray observations, which suggest that young stars have magnetospheres on the scale of several stellar radii and a strength of a few kilogauss (see review by Montmerle et al. 1993). A magnetosphere with such a field strength would truncate the accretion disk at several stellar radii, and the disk accretion would be changed to funnelled accretion, since such a truncation radius can be estimated as follows (e.g., Ghosh, Lamb 1979):

$$\frac{r_m}{R_*} \sim 0.5 \left( \frac{B^4 R_*^5}{G M_* M_{\text{acc}}^2} \right)^{1/7}. \quad (1)$$

(We obtain $\sim 6R_*$ for $r_m$; if we adopt the typical parameters in YSOs: $B_0 \sim 10^{3.6}$G, $R_* \sim 10^{11.4}$ cm, $M_* \sim 10^{3.3}$ g, $M_{\text{acc}} \sim 10^{-7}$M$_\odot$ yr$^{-1}$.) We should thus take into account the interaction between the disk and the magnetosphere around the truncation radius in order to consider the formation of collimated jets. Shu et al. (1994a, 1994b) have proposed a model of a magneto-centrifugally driven jet in this context. They considered the steady state where a non-magnetized accretion disk is truncated by the magnetosphere at the corotation radius, and formulated the dynamical problem of a mageto-centrifugally accelerated flow originating from there. They, following by Najita and Shu (1994) and Ostriker and Shu (1995), carried out an exact treatment of steady cold flows, and explained both the origins of the observed outflows and the angular-momentum problem in YSOs.

1.2. Situation Proposed in Our Model

Here we consider the case in which the accretion disk carries a large-scale interstellar magnetic field with the accreting mass (see discussion in section 5), and the magnetic moment of the stellar magnetosphere is in the same direction as the disk field. This would be the case when the interstellar magnetic field is brought into the star-forming core. The magnetosphere would be formed through magnetic reconnection in the equatorial plane when the stress on the magnetic field around the star is reduced, for example during the intermission of accretion (Uchida, Shibata 1984; Nakano, Umebayashi 1986; Montmerle et al. 1993).] In this case, a magnetically
neutral ring appears at the interface between the disk and the magnetosphere (figure 1b) (Uchida, Shibata 1984).
The truncation radius is estimated again by equation (1), because the magnetic field penetrating the disk would be passive to the dynamics of the disk. On the other hand, the corotation radius $r_c$, where the angular velocity of the magnetosphere equals that to the Keplerian disk, is defined as

$$\frac{r_c}{R_*} = \left( \frac{GM_*}{R_*^2 \Omega_*^2} \right)^{1/3} = \left( \frac{GM_*}{R_* v_*^2} \right)^{1/3},$$

and also becomes several stellar radii in the T Tauri stage. $v_*$ is the surface rotation velocity of the star, and is $\sim 20 \text{ km s}^{-1}$. Here, we consider the case in which the truncation radius $r_m$ is smaller than the corotation radius $r_c$, since, otherwise, the mass transferred from the disk to the magnetosphere would become super-Keplerian as a result of synchronization with the stellar rotation, and could not accrete onto the star. (In such a case, strong magnetic braking may operate, which slow down the stellar rotation until the corotation radius $r_c$ becomes larger than the truncation radius $r_m$.)

At the interface between the disk and the magnetosphere, magnetic reconnection between the disk magnetic field and the magnetospheric field can occur, which would control the mass transfer from the accretion disk to the magnetosphere, as proposed by Uchida and Shibata (1984). They adopted an MHD-equilibrium model of a stellar magnetic field loaded with the non-rotating accreting mass by Uchida and Low (1981), and performed 1-dimensional hydrodynamic simulations of mass accretion onto the star along the magnetospheric field. They showed that the accreting mass-carrying gas on the stellar surface at the polar crown, and reverse shocks propagating outwards accelerate the tail part of the accreting mass, and thus a high-velocity jet is formed.

In this study, we examined the interaction between the magnetosphere and a magnetized accretion disk by performing 2.5-dimensional MHD simulations, and propose a model for the formation of the collimated jet associated with disk accretion onto a magnetized young star. Although we based this simulation on Uchida and Shibata (1984), we newly took into account the rotation of the accreting mass and the change in the configuration of the magnetic field in a self-consistent way. For practical reasons, when performing our numerical simulations, we considered the case in which the star does not rotate, as a first step. (We will come back to the case of a rotating star in some of following papers.)

We describe the numerical method in section 2 and the results of the simulations in section 3. Discussions on the effect of the stellar rotation are given in section 4. In section 5 we discuss the application of our model to outflows and related problems in YSOs. Finally, we make some comments concerning the evolution of a magnetic field around young stars in section 6.

2. Numerical Method

2.1. Basic Equations

We adopt the following system of MHD equations:

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

$$\rho \left[ \frac{\partial}{\partial t} + (\mathbf{v} \cdot \nabla) \right] \mathbf{v} = -\nabla p + \rho \mathbf{g} + j \times \mathbf{B},$$

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

$$\frac{\rho^2}{\gamma - 1} \left[ \frac{\partial}{\partial t} + (\mathbf{v} \cdot \nabla) \right] \left( \frac{p}{\rho^2} \right) = \eta j^2,$$

$$\mathbf{j} = \frac{1}{4\pi} \nabla \times \mathbf{B}.$$  

Here, $\rho$, $p$, and $\mathbf{v}$ are the density, pressure, velocity of a fluid element, respectively. $\mathbf{B}$ is the magnetic field and $\mathbf{j}$ is the corresponding current density; $\mathbf{g}$ is the gravitational acceleration; $\eta$ is the resistivity, which is defined in the next subsection; $\gamma$ is the specific-heat ratio.

To make the above system of equations dimensionless, we introduce three characteristic quantities, the typical length $r_0$, the typical density $p_0$, and the typical velocity $v_0$ in the system. We normalize all of the variables by using these three quantities and their combinations, and obtain the following dimensionless variables (primed variables):

$$r' = r/r_0, \quad \rho' = \rho/p_0, \quad v' = v/v_0,$$

$$t' = t/(r_0/v_0), \quad p' = p/(p_0v_0^2), \quad \mathbf{B}' = \mathbf{B}/(p_0v_0^2)^{1/2},$$

$$\mathbf{g}' = \mathbf{g}/(v_0^2/r_0), \quad \eta' = \eta/(r_0v_0).$$  

Using these dimensionless variables, we obtain a dimensionless system of equations which takes the same form as the dimensional system of equations.

We solve the dimensionless system of equations by two-step Lax-Wendroff scheme (Rubin, Burstein 1966) with an artificial viscosity (Lapidus 1966). We adopt a cylindrical coordinate system $(r, \phi, z)$, assuming axisymmetry around the z-axis, but allowing the $\phi$ components of the vector variables $\mathbf{v}$ and $\mathbf{B}$ (2.5-dimensional formulation). The numerical code is an extended version of that developed and tested by Shibata and Uchida (1985, 1986).

2.2. Initial Conditions

We consider the following simplified situation for the initial condition (figure 1b):

- The star is not rotating. It acts as a point source of the gravity, and has an isothermal hydrostatic corona. It possesses a magnetic dipole moment at
the center, which is perpendicular to the equatorial plane.

- The Keplerian disk, which is cool and dense compared to the corona, is rotating in the equatorial plane. It is penetrated by a uniform magnetic field parallel to its rotational axis. The direction of the uniform magnetic field is the same as the magnetic dipole moment of the star.

We assume that the star is located at the origin of the coordinates, and take the z-axis parallel to the rotational axis of the disk.

We take the coronal density and the coronal sound velocity at the point \((r_0, \phi, 0)\) as the typical density \(\rho_0\) and the typical velocity \(v_0\), respectively, where \(r_0\) is the typical length scale of the system. Then, the initial conditions we adopt can be expressed as follows:

\[
\rho' = \exp \left[ \frac{\gamma}{\epsilon} \left( \frac{1}{\sqrt{r'^2 + z'^2}} - 1 \right) \right] + \sum_{n=1}^{\infty} f_d(r', z'),
\]

\[
\rho'' = \frac{1}{\epsilon} \exp \left[ \frac{\gamma}{\epsilon} \left( \frac{1}{\sqrt{r''^2 + z''^2}} - 1 \right) \right],
\]

\[
u' = \left( 0, \frac{1}{\epsilon \sqrt{r'}} 0 \right) f_d(r', z'),
\]

\[
B' = \frac{8\pi}{\gamma \beta} \left( \frac{1}{r_n^3} - 1 \right)^{-1} \times \left[ \frac{3r'z'}{(r'^2 + z'^2)^{3/2}} + \frac{1}{r_n^3} \right],
\]

\[
g' = \left[ -\frac{1}{\epsilon} \frac{r'}{(r'^2 + z'^2)^{3/2}} + \frac{1}{\epsilon} \frac{r'}{(r'^2 + z'^2)^{3/2}} \right],
\]

\[
f_d(r', z') = \frac{1}{4} \left[ 1 + \tanh \left( \frac{r' - r_d}{\Delta} \right) \right] \times \left[ 1 - \tanh \left( \frac{z' - z_d}{\Delta} \right) \right].
\]

We defined dimensionless parameters as follows:

\[
\epsilon \equiv \frac{v_0^2}{GM_* r_0} = \left( \frac{v_0}{v_{K0}} \right)^2,
\]

\[
\beta \equiv \frac{8\pi \rho_0 v_0^2}{\gamma \left( B_o - B_o (R_*/r_0)^3 / 2 \right)^2} = \frac{p_0}{B_0^2 / 8\pi},
\]

\[
r_n' = \left( \frac{R_*}{r_0} \right)^{1/3}, \quad \Sigma' \equiv \frac{\Sigma}{\rho_0 r_0}, \quad R_*' = \frac{R_*}{r_0},
\]

\[r_d' = \frac{r_d}{r_0}, \quad z_d' = \frac{z_d}{r_0},
\]

where \(M_*, \ R_*, \) and \(B_0\) are the mass and the radius of the star, and the magnetic-field strength at the stellar pole, respectively, and \(\Sigma, \ r_d, \ z_d, \) and \(B_o\) are the surface density, the inner radius, the thickness of the disk, and the strength of the uniform magnetic field penetrating the disk, respectively. \(v_{K0}, \ p_0, \) and \(B_0\) are, respectively, the Keplerian velocity, the coronal pressure, and the magnetic field strength at the point \((r_0, \phi, 0)\).

\(\epsilon\) is the square of the ratio of the coronal sound velocity to the Keplerian velocity and \(\beta\) is the ratio of the coronal gas pressure to the magnetic pressure, where both are defined at the point \((r_0, \phi, 0)\). \(r_n'\) is the normalized radius of the magnetic neutral ring, and is related to the ratio of the strength of the dipolar magnetic field to that of the uniform magnetic field. \(f_d\) determines the configuration of the disk, and is unity inside the disk and zero outside the disk. \(\Delta\) is the width of the transition region between the disk and the corona, and is set to the value which allows the transition region containing about 10 grid points.

In this simulation we localized the finite resistivity only around the magnetically neutral point, and, based on the anomalous resistivity model in Sato and Hayashi (1979), we assumed a function of resistivity as follows:

\[
\eta'(r', z', t') = \frac{1}{2} \left[ 1 + \tanh \left( j_{ns}(t') - j_{crit}^\prime \right) \Delta_j \right] \frac{1}{R_m} f_n(r', z', t'),
\]

where \(R_m\) is the magnetic Reynolds number scaled by \(v_0\) and \(r_0\), \(j_{ns}^\prime\) is the normalized neutral sheet current, and \(f_n\) is a localization factor of order unity. The above equation means that a finite resistivity \(R_m^{-1}\) is set around the neutral point when the neutral sheet current \(j_{ns}\) exceeds the threshold \(j_{crit}^\prime\), which is set to the typical value of the neutral sheet current in the case of no resistivity. We set \(R_m\) to a value of \(10^{-2}\) of the numerical magnetic Reynolds number. The factor \(\Delta_j\) softens the transition of the resistivity between 0 and \(R_m^{-1}\), and is set to 10\% of \(j_{crit}^\prime\).

### 2.3. Boundary Conditions

The total simulation region in the \(r'\)--\(z'\) plane is taken to be \(0 \leq r' \leq r_{max}'\) and \(0 \leq z' \leq z_{max}'\), where \(r_{max}' = 2.7\) and \(z_{max}' = 5.7\). The number of grid points in the \(r'\)-direction and in the \(z'\)-direction is 480 and 350, respectively.

We set free boundary conditions for each quantity \(Q\) at the outer boundaries \((r' = r_{max}'\) and \(z' = z_{max}'\) as follows:

\[
\frac{\partial Q}{\partial r} = 0 \quad (r' = r_{max}') , \quad \frac{\partial Q}{\partial z} = 0 \quad (z' = z_{max}'),
\]

\[\Delta Q = Q(t + \Delta t) - Q(t).
\]

We adopted symmetric boundary conditions at the equatorial plane \((z' = 0)\) and at the \(z'\)-axis \((r' = 0)\). We introduced a thin shell-like damping region near to the stellar surface \(r' = R_*'\) where the disturbances in all physical quantities except for the magnetic field are damped as
follows:

\[ Q(r, z, t) = Q(r, z, t) - h_Q(r, z) [Q(r, z, t) - Q(r, z, 0)]. \]  

(19)

The damping factor \( h_Q \) for quantity \( Q \) was defined such that it is equal to zero outside the damping region and to unity inside the damping region. This corresponds to the absorbing boundary condition for the mass flows. In this case, we would miss such a phenomena as the accreting mass crashing onto the stellar surface (Uchida, Shibata 1984). On the other hand, we could deal with the reflection of the Alfvén wave at the stellar surface, and thus could deal with the angular-momentum transfer between the star and the accreting mass through the magnetic stress.

3. Numerical Results

3.1. Physical Parameters

In this section, we discuss the result of the typical case, in which we chose the numerical parameters introduced in the previous section as follows: \( \gamma = 5/3, \varepsilon = 1.2, \beta = 1.3, r_0^* = 0.63, \Sigma' = 4.8, R_* = 0.10, r_d' = 0.79, z_d' = 0.04. \) (The numerical results for other parameters will be discussed in a separate paper.) These parameters correspond to the following physical values. The normalizing length, density, velocity, and time-scale, respectively, were \( r_0 \sim 10^{12.4} \text{ cm}, \rho_0 \sim 10^{-14.0} \text{ g cm}^{-3}, \nu_0 \sim 10^{6.9} \text{ cm s}^{-1}, \) and \( t_0 \sim 10^{5.5} \text{ s}. \) [We assumed a sizable corona by implicitly considering that magnetic heating mechanisms are in operation, as suggested by recent observations of protostellar objects by ASCA (Koyama et al. 1996).] The radius, mass of the central star and the magnetic field strength at the stellar pole, respectively, were \( R_* \sim 10^{11.4} \text{ cm}, M_* \sim 10^{33.3} \text{ g}, \) and \( B_* \sim 10^{3.6} \text{ G}. \)

The surface density of the disk and the strength of the uniform magnetic field penetrating the disk, respectively, were \( \Sigma \sim 10^{-0.9} \text{ g cm}^{-2} \) and \( B_0 \sim 10^{0.9} \text{ G}. \) The normalizing factor of mass accretion/loss rate was \( \rho_0 r_0 v_0^2 \sim 10^{-8.1} M_* \text{ yr}^{-1} \) [see equations (20),(21)]. Hereafter, we discuss the numerical results using dimensionless variables with the prime omitted.

Figure 2 shows the time evolution of the system in the poloidal plane \((r-z)\) plane. We can see the formation of a bipolar jet, together with accretion flow along the magnetospheric field from the disk to the star. In the following subsections we describe the details of these results.
3.2. Magnetic Reconnection Driven by Disk Accretion

The shear in the angular velocity between the disk and the corona generates the toroidal component from the magnetic field penetrating the disk, and the angular momentum is transferred from the disk to the corona by the magnetic stress. Since this magnetic braking continuously acts on the disk, the disk accretion continues, which would be much more enhanced than in the case without a magnetic field. Note that we do not give the accretion rate of the disk arbitrarily, and that it is self-consistently determined by the efficiency of the magnetic braking.

Our situation in which the rotating disk is penetrated by magnetic field was similar to that of Uchida and Shibata (1985), in which the mass is swirled up from the surface of the disk by the action of the toroidal magnetic field (cf. Shibata, Uchida 1986). In our case, however, such outflow is not formed, because the generation of a toroidal magnetic field is not large enough to form that type of outflow, because the inertia of the rotation of the disk is relatively small in our parameter range.

Figure 3 shows the motion of some selected magnetic lines of force together with the disk mass frozen to them. The disk mass accreting inward presses the disk magnetic field onto the magnetosphere of the star, which causes magnetic reconnection between the magnetospheric field and the disk magnetic field. This is because the enhancement of the neutral sheet current leads to the set up of an anomalous resistivity based on our assumption. Thus, the magnetic reconnection is driven by disk accretion in this case. Figure 4 is an enlargement of the region around the neutral point at $t = 3.70$, and we can see the magnetic reconnection between the disk field and the magnetospheric field in an asymmetric “X”-configuration. Note that the magnetic reconnection occurs successively due to continuous disk accretion.

The magnetic reconnection changes the topology of the magnetic field. A line of force penetrating the disk and a closed line of force of the magnetosphere reconnect, and they change to two open lines of force rooted at the polar
crown of the star (figure 3). Figure 5 shows the magnetic reconnection of one pair of magnetic surfaces in three-dimensional space. We can see that lines of force penetrating the disk are squeezed and twisted by the disk accretion ($0.0 \leq t \leq 1.8$) and reconnected with the magnetospheric field ($t \sim 2.4$), and then relax ($3.0 \leq t \leq 4.2$). Note that the magnetic flux of the star does not change in this accretion process because of the topology change.

3.3. Transfer of Mass and Angular Momentum through Magnetic Reconnection

Through the magnetic reconnection, the rotating disk mass is transferred to the open line of force rooted to the star. Then, the mass and angular momentum are redistributed along this magnetic line of force. Figure 3 shows the transfer and redistribution of the mass through the magnetic reconnection.

Since we set the fixed boundary condition at the stellar surface, the star behaves as if it has infinite inertia when it interacts with the transferred disk mass via the magnetospheric field rooted to the star. Thus, the rotation of the disk mass is synchronized with the stellar rotation. Since the star is not rotating in this case, the rotating disk mass is braked and the greatest part of it falls toward the star along the opened magnetospheric field (figure 6). The rest of the disk mass (the disk surface part) is ejected as a jet (figure 7), as is discussed in the next subsection. In figure 4 we can see that the disk accretion flow divides into two flows: the accretion flow to the star and the jet flow. The mass-accretion rate was evaluated by integrating the mass flux over the cross section of the disk accretion $S_{\text{acc}}$ as

$$M_{\text{acc}} = \int_{S_{\text{acc}}} \rho v \cdot dS = 1.29 \times 10^{1},$$

and the mass-loss rate by the jet was evaluated by integrating the mass flux over the surface penetrated by the stellar open lines of force $S_{\text{open}}$ as

$$M_{\text{loss}} = \int_{S_{\text{open}}} \rho v \cdot dS = 1.28 \times 10^{0}.$$

Both values were time averages over the period of the quasi-steady jet formation ($3.5 \leq t \leq 4.5$) and normalized with $\rho v_{\infty} r_{\odot}^2$. Taking account of the fact that some of the coronal mass is also ejected in the jet, we can see that more than 90% of the disk mass accretes to the star, and that the rest is ejected in the jet.
Fig. 6. Accretion of the disk mass onto the star through magnetic reconnection. For the sake of clarity, only one fluid element of the disk mass (small circle) together with the line of force (solid curve) frozen to the element is followed. Each axis extends from 0 to 1 in the normalized scale. The numbers show the time sequence (1: \( t = 3.00 \), 2: \( t = 3.15 \), 3: \( t = 3.30 \), 4: \( t = 3.45 \), 5: \( t = 3.60 \), 6: \( t = 3.75 \), 7: \( t = 3.90 \), 8: \( t = 4.05 \), 9: \( t = 4.20 \), 10: \( t = 4.35 \), 11: \( t = 4.50 \)).

Fig. 7. Acceleration of the mass initially located at the disk surface through magnetic reconnection. The notations are the same as in figure 6.

As stated in the above, the disk mass splits into the accreting mass and the jet mass through magnetic reconnection (figure 3), and the rotations of both components are finally synchronized with that of the non-rotating star, that is, stopped by the non-rotating star. This means that the angular momentum brought to the magnetosphere by the disk mass is transferred to the star. This can be roughly confirmed in the numerical results. The angular momentum brought by the disk accretion per unit time is evaluated by integrating the angular-momentum flux over the surface \( S_{\text{acc}} \) as follows:

\[
\dot{L}_{\text{acc}} = \iiint_{S_{\text{acc}}} r (\rho v_\phi v - B_\phi B / 4\pi) \cdot dS = 2.80 \times 10^9. \tag{22}
\]

The angular momentum carried away from the star per unit time was evaluated by integrating the angular-momentum flux over the surface \( S_{\text{open}} \) as follows:

\[
\dot{L}_{\text{loss}} = \iiint_{S_{\text{open}}} r (\rho v_\phi v - B_\phi B / 4\pi) \cdot dS = 8.61 \times 10^{-2}. \tag{23}
\]

Here, again, both values are time averages over the period of the quasi-steady jet formation \((3.5 \leq t \leq 4.5)\), and normalized with \( \rho_0 v_0^2 r_0^3 \). Thus, almost all (97%) of the angular momentum brought to the magnetosphere by the disk accretion is absorbed by the star. Note that this result is due to the fact that the star is not rotating, and should thus be contrasted with the case of the rotating star discussed in sections 4 and 5.
3.4. Formation of the Fast Jet Associated with Magnetic Reconnection

The surface part of the disk is accelerated to the bipolar directions by the Lorentz force. This Lorentz force is associated with the reconnected lines of force relaxing and synchronizing with the stellar rotation in three-dimensional space (figure 5).

Figure 7 shows the motion of fluid particle initially located at the disk surface together with the lines of force frozen to the particle; figure 8 shows specific powers due to various forces acting on the particle in figure 7. Figure 8a shows the total power, the powers due to the Lorentz force, the pressure gradient, and the gravitational force. (Here, since we consider a "Lagrangian" fluid particle, the inertia forces, such as the centrifugal force, do not affect the energy change of the particle.) We can see that the particle is accelerated two times at \( t \sim 3.9 \) and \( t \sim 4.2 \), mainly by the Lorentz force. The particle acquires the local escape velocity by these two accelerations (figure 8c).

In order to see what happens at these times, we decompose the Lorentz force, \( \mathbf{j} \times \mathbf{B} \), into three components, \( \mathbf{j}_t \times \mathbf{B} \), \( \mathbf{j}_p \times \mathbf{B}_p \), and \( \mathbf{j}_p \times \mathbf{B}_p \) ("t" and "p" denote toroidal and poloidal, respectively), and examine the specific powers due to these three components. The result is shown in figure 8b. We can see that the first acceleration is mainly due to \( \mathbf{j}_t \times \mathbf{B}_p \), and that the second one is mainly due to \( \mathbf{j}_p \times \mathbf{B}_p \). The first acceleration is produced by the magnetic tension associated with the relaxation of the reconnected lines of force in the poloidal plane, which is a direct consequence of a topology change through magnetic reconnection (figure 7a: the time indices 6 to 8). The second acceleration in the toroidal direction leads to a centrifugal acceleration of the particle. This Lorentz force is associated with the relaxation of the lines of force distorted through synchronization with the stellar rotation (figure 7b: the time indices 9 to 11). Thus, the second acceleration is due to the angular-velocity difference between the magnetosphere and the inner edge of the disk, and, thus, is reduced in the case of a rotating star, while the first acceleration in the poloidal plane is also essential in such a case.

In our case, most of the magnetic energy liberated in the magnetic reconnection process is used to lift up the disk mass from the equatorial plane, which allows the disk mass to accrete to the star along the reconnected magnetospheric field; the rest of the magnetic energy is used to accelerate the surface part of the disk mass (figure 3). That is, only a part of the magnetic energy liberated in the magnetic reconnection is converted to the kinetic energy of the jet. This is because the fluid motion after the magnetic reconnection is strongly affected by the gravity and the asymmetric configuration of the magnetic field; these factors determine the ratio of the mass-loss rate to the accretion rate and the velocity of
Fig. 9. Schematic picture for explaining the angular-momentum transfer between the disk, the star, and the jet. Here, we consider the angular-momentum transfer along only one reconnected line of force for simplicity; however, note that this process occurs continuously. (i) shows the initial state just after the reconnection. (ii) shows the final state after the magneto-centrifugal acceleration of the disk mass up to the Alfvén radius.

the ejected mass. In our parameter range, the ratio of the mass-loss rate $M_{\text{loss}}$ to the accretion rate $M_{\text{acc}}$ is about 0.1, as shown in the previous subsection, and the velocity of the ejected mass is about 50% of the local escape velocity (figure 8c). This velocity is also about 50% of the local Alfvén velocity.

4. Effects of Stellar Rotation

Since we could not take the stellar rotation into account in the simulation due to numerical problems, here we briefly discuss the effects of stellar rotation on the numerical results, assuming that the features of the reconnection driven-jet which would be further accelerated magneto-centrifugally by the stellar rotation can be roughly described by a simplified cold-flow model.

Since the reconnected lines of force are synchronized with the stellar rotation, the mass ejected by the magnetic reconnection can be accelerated magneto-centrifugally, which leads to additional magnetic braking of the star. To see such an effect of the stellar rotation, we consider here the mass and angular momentum transfer along one reconnected line of force. This is justified because the mass and angular momentum are continuously transferred from the disk to the star through successive magnetic reconnection (figure 3). Furthermore, we adopt some simplifications. First, we assume that the inner edge of the disk, which coincides with the reconnection point, is inside the corotation radius, as discussed in section 1. Second, we follow the basic idea of a steady magneto-centrifugal disk wind model in the cold-gas limit.

In the case of the cold-gas limit, magneto-centrifugal acceleration is possible only when the mass is brought over the “ridge” in the effective potential field with positive velocity. (In the disk-wind model, this condition is satisfied because the end of the line of force where the jet mass is supplied is on the “ridge” in the effective potential field.) In our case, although the jet mass is supplied at the reconnection point, which is inside the “ridge” of the potential, the initial velocity needed for the magneto-centrifugal acceleration is acquired during the process of ejection by the magnetic reconnection, as can be seen in the numerical results.

Then, the terminal velocity of the jet $v_\infty$ is roughly estimated as follows (Michel 1969):

$$v_\infty \sim r_A \Omega_* = \left( \frac{r_A}{R_*} \right) v_*, \quad \text{(24)}$$

where $r_A$ is the Alfvén radius up to which the magneto-centrifugal acceleration operates.

The angular-momentum transfer can be roughly described in the following simplified model (figure 9):  

(i) In the initial state, just after reconnection, suppose that a star having moment of inertia $I_*$ rotates with angular velocity $\Omega_*$, while the transferred mass $\Delta m_d$ rotates at the Keplerian angular velocity $\Omega_K(r_0)$ at the reconnection point ($r = r_0$).

(ii) In the final state, suppose that a part $\Delta m_a$ of the transferred mass accretes to the stellar surface ($r = R_*$), and the other part $\Delta m_j$ is magneto-centrifugally accelerated up to the Alfvén radius $r_A$. The star, the accreting mass $\Delta m_a$ and the jet mass
\( \Delta m_j \) are all synchronized with the stellar rotation \( \Omega_\ast + \Delta \Omega_\ast \), because the magnetic field anchored to the star is considered to be “rigid” up to the Alfvén radius.

Then, the angular-momentum conservation is written as

\[
I_\ast \Omega_\ast + \Delta m_d \frac{r_0^2}{R_\ast} \Omega_K (r_0) = (I_\ast + \Delta m_d r_0^2 + \Delta m_j) \frac{r_0^2}{R_\ast} \Omega_\ast + \Delta \Omega_\ast,
\]

where \( \Delta m_d = \Delta m_a + \Delta m_j \). We then have

\[
\Delta \Omega_\ast = \frac{\Delta m_d \frac{r_0^2}{R_\ast} \Omega_K (r_0) - \Delta m_a \frac{R_\ast}{R_\ast} \Omega_\ast - \Delta m_j \frac{r_0^2}{R_\ast} \Omega_\ast}{I_\ast + \Delta m_d \frac{r_0^2}{R_\ast} + \Delta m_a \frac{R_\ast}{R_\ast}}.
\]

Here, we do not consider the changes in the stellar radius \( R_\ast \) and that of the moment of inertia \( I_\ast \) of the star for simplicity. [\( \Delta m_d \frac{r_0^2}{R_\ast} \Omega_K (r_0) - \Delta m_a \frac{R_\ast}{R_\ast} \Omega_\ast \)] represents the angular momentum transferred from the disk to the star, and \( - \Delta m_j \frac{r_0^2}{R_\ast} \Omega_\ast \) represents the magnetic braking as a reaction of the magneto-centrifugal acceleration of the jet. If the Alfvén radius has an appropriate value, the former can be cancelled by the latter, and, thus, the steady rotation of the star \( \Delta \Omega_\ast = 0 \) becomes possible. From equation (26) such a value of the Alfvén radius can be written as

\[
\left( \frac{r_0}{R_\ast} \right)^2 = 1 + \frac{1}{f} \left[ \left( \frac{r_0}{R_\ast} \right)^{3/2} - \left( \frac{r_0}{R_\ast} \right)^{3/2} - 1 \right],
\]

where \( f \equiv \Delta m_j / \Delta m_d \) and \( r_c \) is the corotation radius [equation (2)].

Even when the stellar rotation contributes to the acceleration of the jet (as discussed in this section), the magnetic reconnection driven by the disk accretion still plays a vital role in starting up the ejection; in that sense the jet is “driven by the disk accretion,” which is consistent to what the observational results suggest. Namely, it is the magnetic reconnection that accelerates the mass initially, and thus determines the correlation between the mass loss rate and the accretion rate in the disk. Furthermore, the following magneto-centrifugal acceleration by the stellar rotation is subsequently made possible by this magnetic reconnection.

5. Outflows from Young Stellar Objects and Related Problems

In the situation considered in our model, two types of outflows having different origins may be formed in the vicinity of a young star. One is the collimated jet ejected by the magnetic reconnection from the interface between the disk and the magnetosphere as discussed in section 3 and section 4; the other is the magnetically accelerated outflows originating from the extended regions of the disk penetrated by large-scale magnetic field (Uchida, Shibata 1985; Shibata, Uchida 1986; Kudo, Shibata 1995). We consider that the former corresponds to optical jets, and that the latter corresponds to high-velocity neutral winds (figure 1b). Since Shibata and Uchida (1986) as well as Kudo and Shibata (1995) showed that the typical velocity of outflows from the disk penetrated by large-scale magnetic field is on the order of the Keplerian velocity at its origin, the high-velocity neutral winds (~ 200 km s\(^{-1}\)) may come from the inner part of the accretion disk (~ 0.01 AU). As for the molecular bipolar outflows, since the mass in the two lobes is large, most of material is likely to be gas of the molecular cloud core swept up or entrained by outflows from the disk. In the mechanism proposed in Uchida and Shibata (1985) the mass not only of the disk surface, but also on the way, are swept up; also, the mass initially existing in the hourglass-shaped region threaded by large-scale magnetic field is collected by the pinch of the toroidal field as the non-linear torsional Alfvén waves propagate.

Umebayashi and Nakano (1988) showed that in the middle part of the disk the ionization degree is low and the coupling between the gas and the magnetic field is weak near to the midplane of the disk. However, the local condition, whether the ionized gas in the surface layer continues to accrete toward the star or not, is independent of the frozen-in condition near to the midplane of the disk. On the other hand, the magnetically accelerated outflow from the inner region of the disk does not blow out the field lines of the middle part of the disk, because we consider the case in which the field lines are aligned in all regions of the disk (figure 1b). Thus, the magnetic field brought by disk accretion from the outer ionized part to this middle part will be advected inward by the accretion of the surface frozen-in layer. (In this region the magnetic field inside the disk will be current-free field determined by the motion of the frozen-in gas near the surface.) The magnetic field will be coupled to the whole part of the disk again when it comes to the inner part, and will then be brought towards the star by disk accretion, generating a “high-velocity neutral wind,” and eventually reconnecting with the magnetospheric field, and drives an “optical jet.”

YSOs, especially CTTSs, which are considered to be surrounded by the accretion disk, show evidence for rather slow rotation in spite of receiving a high specific angular momentum from the disk together with the mass. This riddle can be answered if the angular momentum can be extracted from the star-disk system by some mechanism (Hartmann, Stauffer 1989; Durisen et al. 1989; Edwards et al. 1993b; Königl 1991). In our model, the reconnection-driven jet, which is accelerated magneto-centrifugally by stellar rotation, is a candidate of such an extractor of the angular momentum. From equation (27), if the Alfvén radius is \( \sim 10 R_\ast \), the torque
given to the star by the disk accretion can be canceled by the negative torque by the magneto-centrifugal acceleration of the jet, and, thus, the rather slow rotation of the star may be made possible. [Here, we assumed that \( r_c \) and \( r_0 \) are several stellar radii, as discussed in section 1, and that \( f \sim 0.1 \) from the numerical results, which is consistent with observations (Edwards et al. 1993a).] Furthermore, if the Alfvén radius is \( \sim 10 r_\ast \), the terminal velocity of the jet is estimated to be \( \sim 200 \text{ km s}^{-1} \) from equation (24), which agrees with the typical velocity of optical jets. (We adopt a conventional value for the surface rotational velocity of CTTs, \( \nu_s \sim 20 \text{ km s}^{-1} \).)

Thus, our model consistently explains the jet-formation and angular-momentum problems of young stars in the disk-accretion stage.

In our model, the jet propagates along the stellar open field twisted by the stellar rotation. Uchida et al. (1992) and Todo et al. (1992, 1993) showed that when such a jet interacts with the ambient medium it experiences a helical-kink instability, which well explains the wiggled structure of some HH objects.

6. Discussion

Here, in this section, we comment on the evolution of the magnetic field around a young star and related observational features.

In our simulation, rather continuous magnetic reconnection is driven by the disk accretion, since we must adopt a low threshold of the neutral sheet current for the anomalously resistivity to be induced, because of the small numerical magnetic Reynolds number. When the threshold is sufficiently high, there is a possibility that the following two stages repeat alternatively. At first, accretion onto the star does not occur while disk accretion occurs, since magnetic reconnection does not occur readily in this case. (Note that in our model, the accretion to the star occurs only when the magnetic reconnection occurs.) Then, the accreting mass accumulates at the magnetospheric boundary, which leads to the build-up of a neutral sheet current (Stage I). When the neutral sheet current exceeds the threshold, an anomalously resistivity is induced, and magnetic reconnection occurs. Then, the accumulated mass accretes onto the star accompanying jet formation (Stage II). After all of the accumulated mass which stressed the magnetic field around the star accretes onto the star, the magnetosphere is reformed in the current-free space, and the state returns to Stage I. Such a cycle may correspond to a non-steady accretion process (FU Orionis phenomena) (see review by Hartmann et al. 1993; Lin et al. 1994) or the intermittence of mass ejection suggested by the multiple bow shocks in HH objects (Bührke et al. 1988; Reipurth 1989; Hartigan et al. 1990).

The collimation of the jet changes as the reconnection proceeds, because the reconnected lines of force in which the jet is confined are bundled outside the preceding reconnected lines of force. Multiple mass ejections with different collimations has actually been observed in HH 90/91 (Reipurth, Heathcote 1993).

Finally, when the disk accretion ceases, and, thus, no stress acts on the magnetic field around the star, the configuration of the magnetic field quickly changes to that of a current-free field, and the magnetosphere is reformed in empty space. Namely, the open lines of force of the star return to the closed states through magnetic reconnection, which reduces the angular-momentum loss through the open lines of force. Then, the star may well conserve its angular momentum and spin up due to a contraction in its evolution process, which is consistent with the observation that the rotation of stars after the disk-accretion stage is faster than that of stars in the disk-accretion stage (Edwards et al. 1993b).

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