Global 3-D Radiation Magnetohydrodynamic Simulations for FU Ori’s Accretion Disk and Observational Signatures of Magnetic Fields

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ABSTRACT
FU Orionis systems are outbursting protoplanetary disks where the accretion disks outshine the central stars and strong disk winds are launched. Magnetic fields in these accretion disks have previously been detected through their Zeeman effects in spectropolarimetry observations. We carry out global radiation ideal MHD simulations to study FU Ori’s inner accretion disk. We find that (1) when the disk is threaded by vertical magnetic fields, most accretion occurs in the magnetically dominated atmosphere at z∼R, similar to the “surface accretion” mechanism in previous locally-isothermal MHD simulations. (2) A moderate disk wind is launched in these vertical field simulations with terminal speeds of ∼300-500 km/s and a mass loss rate of 1-10% disk accretion rate. Both the speed and loss rate are consistent with observations. Disk wind fails to be launched in net toroidal field simulations. (3) The disk photosphere at the unit optical depth can be either in the wind launching region or the accreting surface region, depending on the accretion rates and the disk radii. Magnetic fields have drastically different directions and magnitudes between these two regions. Our fiducial model agrees with previous optical Zeeman observations regarding both the field directions and magnitudes. On the other hand, simulations indicate that future Zeeman observations at near-IR wavelengths or towards other FU Orionis systems may reveal very different magnetic field structures. (4) Due to energy loss by the disk wind, the disk photosphere temperature is lower than that predicted by the thin disk theory, and the previously inferred disk accretion rate may be lower than the real accretion rate by a factor of ∼2-3.

Key words: accretion, accretion disks - astroparticle physics - dynamo - magnetohydrodynamics (MHD) - instabilities - turbulence

1 INTRODUCTION

Accretion disks have been observed in a wide range of astrophysical systems, ranging from around low mass stars (Hartmann 1998) to around compact objects and supermassive black holes (Begelman et al. 1984). The accretion process not only helps to build the central object, but the released radiation energy allows us to identify and study the central object (e.g. X-ray binaries). The high resolution M87 image by the Event Horizon Telescope (Event Horizon Telescope Collaboration et al. 2019) is an excellent example that we can constrain the properties of black holes by studying its surrounding accretion disks.

The leading theory to explain the accretion process involves magnetic fields, especially for sufficiently ionized disks1. Magnetic fields can drive turbulence through the magnetorotational instability (MRI; Balbus & Hawley 1991, 1998) or/and launch disk winds through the magnetocentrifugal effect in non-relativistic disks (Blandford & Payne 1982). The strengths of both MRI turbulence and disk winds depend on the field strength. Normally turbulence and wind are more prominent in systems having stronger magnetic fields (Hawley et al. 1995).

Despite the importance of magnetic fields, the observational evidences for magnetic fields in accretion disks remain

1 In poorly ionized disks where the non-ideal MHD effects become important, hydrodynamical processes may also play an important role in disk accretion (Turner et al. 2014).
to be scarce. The collimated jets/outflows provide some indirect evidence of magnetic fields since the confinement of jets may require the presence of magnetic fields (Pudritz et al. 2007; Frank et al. 2014). Another indirect evidence is from magnetic field measurements from meteorites. Paleomagnetic measurements by Fu et al. (2014) suggest that Semarkona meteorites were magnetized to 0.54 G in the solar nebulae.

The most direct evidence of magnetic fields in accretion disks comes from Zeeman splitting of atomic or molecular lines. Current Zeeman measurements of molecular lines using ALMA (Vlemmings et al. 2019) have only placed upper limits on the field strength (< 30 mG). So far, the only direct detection of magnetic fields in accretion disks is the detection of Zeeman splitting of atomic lines coming from the inner disk of FU Ori (Donati et al. 2005).

FU Ori is the prototype of FU Orionis systems: a small but remarkable class of variable young stellar objects that undergo outbursts in optical light of 5 magnitudes or more (Herbig 1977). While the outburst has a fast rise time (∼ 10−4 M⊙ yr−1) to ∼ 10−3 M⊙ yr−1 (Hartmann & Kenyon 1996). The strong accretion is accompanied by the strong jets (Class I-II rates) to RG (Hartmann & Kenyon 1998; Kraus et al. 2016; Köppél et al. 2017; Hillenbrand et al. 2018), the occurrence rate of these objects among young stars is still illusive (Hillenbrand & Findl甚至 2015; Scholz et al. 2013) with rates ranging from less than 1 outburst per young star to more than tens of outbursts per young star.

Such intense outbursts are due to the sudden increase of the protostellar disk’s accretion rates from ∼ 10−3 M⊙ yr−1 (Class I-II rates) to ∼ 10−2 M⊙ yr−1 (Hartmann & Kenyon 1996). The strong accretion is accompanied by the strong disk wind (Calvet et al. 1993; Millner et al. 2019). Although the outburst triggering mechanism is not clear2, the inner disk (∼ 1 au) during the outburst is hot enough (∼ 6000 K, Zhu et al. 2007) to be sufficiently ionized and MRI should operate in this disk. Since this disk with ∼ 100 L⊙ is much brighter than the central star and all the light we see is from this accretion disk, FU Orionis systems are ideal places to study accretion physics.

Taking advantage of many atomic lines available in these systems, Donati et al. (2005) have used the high-resolution spectropolarimeter to detect signals of Zeeman splitting in FU Ori. By splitting the circular polarization signal into symmetric and antisymmetric components, they constrain the magnetic fields in both the azimuthal and radial directions. Assuming that the disk’s rotational axis is 60° inclined with respect to our line of sight, their best fit model suggests that the vertical component of the fields is ∼ 1 kG at 0.05 au and points towards the observer, while the azimuthal component (about half as strong) points in a direction opposite to the orbital rotation.

In spite of these stringent observational constraints, theoretical work still lacks behind and its connection with observations has not been established. High numerical resolution is necessary for capturing MRI, while a large simulation domain is needed to study the disk wind. Only recently, with the newly developed Athena++ code which has both mesh-refinement and the special polar boundary condition, we can simulate the whole 4π sphere around the central object with enough resolution to capture MRI (Zhu & Stone 2018). Besides magnetic fields, radiative transfer is also crucial for understanding FU Ori’s inner accretion disk. Thermal instability was suggested to explain FU Ori’s outburst (Bell & Lin 1994). Although local shearing box MHD simulations with radiative transfer (Hirose et al. 2014) do not support the thermal instability theory for FU Ori outbursts (Hirose 2015), the disk’s thermal structure is still important for both the accretion physics (Zhu et al. 2009b) and the boundary layer physics (Kley & Lin 1999). Furthermore, radiative transfer is also important for making connections with observations (e.g. understanding the physical condition at the disk’s photosphere).

Thus, in this work, we include radiative transfer in the global MHD disk simulations to study the accretion structure of FU Ori’s inner disk. We will also compare our simulations with previous Zeeman magnetic field observations and disk wind observations. In Section 2, the theoretical framework for energy transport in accretion disks is presented. We will describe our numerical method in Section 3. The results are presented in Section 4. After connecting with observations and a short discussion in Section 5, the paper is concluded in Section 6.

2 THEORETICAL FRAMEWORK

Angular momentum transport and energy transport are two important aspects of the accretion disk theory. Angular momentum transport is essential for mass buildup of the central object, while energy transport is crucial for revealing disk properties from observations. Previously in Zhu & Stone (2018), we have done detailed analyses on angular momentum transport for disks threaded by net vertical magnetic fields. In this work, we will focus on energy transport of accretion disks threaded by net magnetic fields.

The fluid equations with both magnetic and radiation fields are

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

\[
\frac{\partial \rho \mathbf{v}}{\partial t} + \nabla \cdot (\rho \mathbf{v} \mathbf{v} - \mathbf{BB} + \mathbf{P} + \mathbf{\sigma}) = -\mathbf{S}_r(\mathbf{P}) + \mathbf{F}
\]

\[
\frac{\partial E}{\partial t} + \nabla \cdot [(E + P^*) \mathbf{v} - \mathbf{B} \cdot \mathbf{v} + \mathbf{\sigma} \cdot \mathbf{v}] = -c_s \mathbf{S}_r(\mathbf{E}) + \mathbf{F} \cdot \mathbf{v}
\]

\[
\frac{\partial \mathbf{B}}{\partial t} - \nabla \times (\mathbf{v} \times \mathbf{B}) = 0,
\]

where \( E = E_\gamma + p \frac{v^2}{2} + B^2/2 \) is the total gas energy density, \( E_\gamma = P/(\gamma - 1) \) is the internal energy, \( P^* = (P + B^2/2) \) is the pressure tensor (with \( I \) the unit tensor), and \( \mathbf{F} \) is the external force (e.g. gravity). We also include the dissipation tensor \( \mathbf{\sigma} \) in the equations. Although dissipation is not explicitly added in the simulations, shock dissipation is implicitly included in the Riemann solver, and dissipation terms are important for the energy analysis. The radiation equations are

\[
\frac{\partial E_r}{\partial t} + \nabla \cdot \mathbf{F}_r = c_s \mathbf{S}_r(\mathbf{E})
\]

\[
\frac{1}{c_s^2} \frac{\partial \mathbf{F}_r}{\partial t} + \nabla \cdot \mathbf{P}_r = \mathbf{S}_r(\mathbf{P})
\]

where the radiation flux \( \mathbf{F}_r \) and the radiation energy density
$E_c$ are Eulerian variables, and they are related to the comoving flux $\mathbf{F}_{c,0}$ through $\mathbf{F}_{c,0} = \mathbf{F}_c - (\mathbf{v} E_c + \mathbf{v} \cdot \mathbf{P}_c)$. The radiation pressure tensor $\mathbf{P}_c$ is related to the energy density though a variable Eddington tensor $\mathbf{P}_c = fE_c$. The source terms $cS_c(E)$ and $\mathbf{S}_c(P)$ are given in Jiang et al. (2013).

To study the energy budget, it is also helpful to write the equation for the internal energy of the gas. The kinetic and magnetic energy equation is

$$\frac{\partial}{\partial t} \left( \frac{\rho v^2}{2} + B^2 \right) + \nabla \cdot \left[ \mathbf{v} \left( \frac{\rho v^2}{2} \right) - \mathbf{B} (\mathbf{B} \cdot \mathbf{v}) + (P^* + \sigma) \right] v = -\mathbf{v} \cdot \mathbf{S}_v(P) + F \cdot v,$$

(4)

so that the internal energy is

$$\frac{\partial E_i}{\partial t} + \nabla \cdot (E_i \mathbf{v}) + \nabla \cdot (\mathbf{P} \cdot \mathbf{v}) = -cS_c(E) + \mathbf{v} \cdot \mathbf{S}_c(P),$$

(5)

which suggests that the change of the internal energy is due to the $P \mathbf{dv}$ work, the dissipation, and radiative transport.

We can use either the equation for the total energy (Equation 1) or the equation for the internal energy (Equation 5) to derive the disk luminosity. The equation for the total energy can be written as

$$\frac{\partial E}{\partial t} + \nabla \cdot \mathbf{A} = -Q_{\text{cool}} + \mathbf{F} \cdot \mathbf{v},$$

(6)

where $\mathbf{A} = (E + P^*) \mathbf{v} - \mathbf{B} (\mathbf{B} \cdot \mathbf{v})$, and $Q_{\text{cool}}$ is the radiative cooling rate. $\mathbf{A}$ can also be rewritten as

$$\mathbf{A} = \left( \frac{\gamma}{\gamma - 1} P + \frac{1}{\gamma} \rho v^2 \right) \mathbf{v} + \mathbf{B} \times (\mathbf{v} \times \mathbf{B})$$

(7)

using vector identities.

We will first review the thin disk theory under the cylindrical coordinate system and then we will write similar equations under the spherical-polar coordinate system which is more suitable to our simulations. The perturbed equation for the angular momentum under the cylindrical coordinate system can be written as

$$\frac{\partial (\rho \delta v_\phi)}{\partial t} = \frac{1}{R^2} \frac{\partial (R^2 \langle T_{R\phi} \rangle)}{\partial r} - \frac{\partial (\rho v_r \rho v_K)}{R} \frac{\partial v_N}{\partial r} \frac{\partial R}{\partial \phi} \frac{\partial v_K}{\partial \phi},$$

(8)

where

$$T_{R\phi} \equiv \rho v_r \delta v_\phi - B_r B_\phi,$$

$$T_{\phi\phi} \equiv \rho v_\phi \delta v_\phi - B_\phi B_\phi,$$

and $\langle \rangle$ denotes that the quantity has been averaged in the azimuthal ($\phi$) direction. Assuming a steady state, we have

$$\frac{\dot{M}}{2\pi} \frac{\partial R v_K}{\partial R} = \frac{\partial R^2 \langle T_{R\phi} \rangle}{\partial R} + R^2 \frac{\partial T_{\phi\phi}}{\partial \phi} + R^2 \langle \rho v_\phi \rangle \frac{\partial v_N}{\partial \phi} \frac{\partial v_K}{\partial \phi},$$

(9)

where $\dot{M} \equiv -2\pi R \langle \rho v_r \rangle$. Thus, the accretion is driven by the $T_{R\phi}$ stress within the disk or the $T_{\phi\phi}$ stress at the disk surface. If we assume that $\dot{M}$ is a constant along $R$, we have

$$\langle T_{R\phi} \rangle = \frac{\dot{M} v_K}{2\pi R} \frac{C}{R^2} \frac{1}{R^2} \int R^2 \left( \frac{\partial T_{\phi\phi}}{\partial \phi} + \langle \rho v_\phi \rangle \frac{\partial v_N}{\partial \phi} \frac{\partial v_K}{\partial \phi} \right) dR.$$

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The energy Equation (Equation 6) under the cylindrical coordinate is

$$\frac{\partial (E)}{\partial t} = -\frac{1}{R} \frac{\partial (R \langle A_r \rangle)}{\partial R} - \frac{\partial (A_z)}{\partial z} - \langle Q_{\text{cool}} \rangle + \langle \mathbf{F} \cdot \mathbf{v} \rangle,$$

(11)

where the leading term in $A_R$ (after removing the second-order terms) is

$$A_R = \frac{\gamma}{\gamma - 1} P v_R + \frac{1}{2} \rho \rho v_r^2 + v_K T_{R\phi},$$

(12)

and the leading term in $A_z$ is

$$A_z = \frac{\gamma}{\gamma - 1} P v_z + \frac{1}{2} \rho \rho v_z^2 + v_K T_{\phi\phi}.$$

(13)

If we ignore the pressure term in $A_R$, assume $v_z \sim 0$ in $A_z$, and assume a steady state, we have

$$\langle Q_{\text{cool}} \rangle = -\frac{1}{R} \frac{\partial (\langle -\frac{1}{8} M v_K^2 + R v_N T_{R\phi} \rangle)}{\partial R} - \frac{\partial (v_K T_{R\phi})}{\partial z} + \langle \mathbf{F} \cdot \mathbf{v} \rangle.$$  

(14)

If we plug in $T_{R\phi}$, ignore the $T_{\phi\phi}$ term, replace $\mathbf{F}$ with the gravitational force, and only consider the disk midplane, we have

$$2\pi \langle Q_{\text{cool}} \rangle = -\frac{1}{2} \dot{M} v_K^2 + \dot{M} v_N^2 - \frac{3}{2} C v_K^2 + \frac{M v_N^2}{R^2},$$

(15)

where the first term on the right is due to the radial derivative of the Keplerian kinetic energy flux, the second and third terms on the right are due to the radial derivative of the $R - \phi$ stress, and the last term on the right is the release of the gravitational potential energy. With the traditional zero stress inner boundary condition ($C = M R v_m v_N$, $\sigma = 0$), the cooling rate is

$$\langle Q_{\text{cool}} \rangle = \frac{3 M v_K^2}{4\pi R^2} \left( 1 - \left( \frac{R_m}{R} \right)^{1/2} \right),$$

(16)

or

$$\sigma T_{\phi\phi}^* = \frac{3 GM M}{8\pi R^2} \left( 1 - \left( \frac{R_m}{R} \right)^{1/2} \right).$$

(17)

After the vertical integration, this cooling rate is what we normally use in the thin disk approximation. If we integrate over the whole disk starting from $R_m$, the total cooling rate is half the release rate of the gravitational potential energy ($GM M/2R_m$). On the other hand, far away from the central star ($R \gg R_m$), the cooling rate ($3 M v_K^2/4\pi R^2$) is actually higher than the energy release rate from the gravitational contraction ($M v_K^2/2\pi R^2$). The additional $M v_N^2/4\pi R^2$ energy release is due to the energy transport in the radial direction. We note that the same equation can also be derived using the internal energy equation but with an additional step to derive the dissipation term.

On the other hand, our simulated disks are very thick and the disk photosphere roughly follows the radial direction in the spherical grids. Thus, we want to derive similar equations for the spherical-polar grid so that we can study energy transport in our simulations. The perturbed angular momentum equation under the spherical-polar coordinate is

$$\frac{\partial (\rho \delta v_\theta)}{\partial t} = \frac{1}{r^2} \frac{\partial (r^2 \langle T_{R\theta} \rangle)}{\partial r} - \frac{\partial (\rho \rho v_r)}{r} \frac{\partial v_N}{\partial r} \frac{\partial v_K}{\partial \theta} - \frac{1}{\sin^2 \theta} \frac{\partial (\sin^2 \theta \langle T_{\phi\phi} \rangle)}{\partial \theta} + \langle \rho v_\phi \rangle \frac{\partial \sin^2 \theta \frac{v_K}{r}}{\partial \theta},$$

(18)

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where
\[ T_{\phi} \equiv \rho v_j \delta v_0 \nabla_0 \cdot \vec{B}_R \]
\[ T_{\phi} \equiv \rho v_j \delta v_0 \nabla_0 \cdot \vec{B}_R \, . \] (20)

Assuming a steady state, we have
\[ \frac{\dot{M}}{M} \frac{\partial r v}{\partial r} = \frac{\partial (r^2 (\dot{r} T_{\phi}))}{\partial r} + \frac{r^2}{\sin^2 \theta} \frac{\partial (\sin^2 \theta \dot{r} T_{\phi})}{\partial \theta} + \frac{r^2 (\rho v_j) \frac{\partial (\sin \theta v_K)}{\partial \theta}}{\sin \theta} \] (21)

where \( \dot{M} = -r^2 (\rho v_j) \). Note that this \( \dot{M} \) definition is different from the \( M \) definition in the cylindrical grid. If we assume that \( \dot{M} \) is a constant along \( r \), we can integrate the equation to derive
\[ \langle \dot{r} T_{\phi} \rangle = \frac{\dot{M}_K}{r^2} - C \frac{r^2}{r^3} \int \frac{r^2}{\sin^2 \theta} \frac{\partial (\sin^2 \theta \dot{r} T_{\phi})}{\partial \theta} \, d\theta \] (22)

The energy equation (Equation 6) under the spherical-polar coordinate is
\[ \frac{\partial (E_{\phi})}{\partial t} = -\frac{1}{r^2} \frac{\partial (r^2 (A_{\phi}))}{\partial r} - \frac{1}{\sin \theta} \frac{\partial (\sin \theta (A_{\phi}))}{\partial \theta} - \langle Q_{\text{cool}} \rangle + \langle \mathbf{F} \cdot \mathbf{V} \rangle. \] (24)

The leading term in \( A_{\phi} \) is
\[ A_{\phi} = \left( \frac{\gamma}{\gamma - 1} \right) P + \frac{1}{2} \rho v_K^2 + \rho v_K \rho v_j \nabla \cdot \nabla_0 \cdot \vec{B}_R - \frac{v_K}{v_j} B_R \, B_R \] (25)
or
\[ A_{\phi} = \frac{\gamma}{\gamma - 1} P v_j + \frac{1}{2} \rho v_j v_j^2 + v_K T_{\phi} \] (26)
The leading term in \( A_{\phi} \) is
\[ A_{\phi} = \frac{\gamma}{\gamma - 1} P v_j + \frac{1}{2} \rho v_j v_j^2 + v_K T_{\phi} \] (27)

In §4.2, we will measure the energy transport due to \( A_{\phi} \) and \( A_{\phi} \) from our simulations. On the other hand, in this section, we will continue the derivation by making several assumptions. If we ignore the pressure term in \( A_{\phi} \), assume \( v_j \sim 0 \) in \( A_{\phi} \), and assume a steady state, we have
\[ \langle Q_{\text{cool}} \rangle = -\frac{1}{r^2} \frac{\partial \left( -\frac{1}{2} \frac{\dot{M}_K^2 v_K^2 + r^2 v_K T_{\phi}}{\dot{r}} \right)}{\partial \theta} + \frac{1}{\sin \theta} \frac{\partial (\sin \theta (v_K T_{\phi}))}{\partial \theta} \] (28)

If we plug in \( T_{\phi} \) from Equation 23 and ignore \( T_{\phi} \) terms, we have
\[ \langle Q_{\text{cool}} \rangle = -\frac{1}{2} \frac{\dot{M}_K^2 v_K^2}{r^3} + \frac{\dot{M}_K^2 v_K^2}{r^3} - \frac{3}{2} C v_K^2 + \frac{\dot{M}_K^2 v_K^2}{r^3} \] (29)

Thus, if we can ignore the \( \theta \) direction energy advection/stress and the boundary \( C \) term, the cooling rate equals the release rate of the gravitational potential energy (the last term on the right) plus the radially advected energy (the first two terms on the right). As will be shown in Section 4.2, the energy advection in the \( \theta \) direction and the \( T_{\phi} \) stress cannot be ignored. Accordingly, the cooling rate is modified significantly.

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**Figure 1.** The Rosseland mean (solid black curves) and Planck mean (red dashed curves) opacities adopted in the simulations. Different curves represent opacities under different pressures (10^{-3} to 10^5 dyn cm^{-2}). Curves with overall lower values correspond to lower pressures.

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### 3 METHOD

We solve the magnetohydrodynamic (MHD) equations in the ideal MHD limit using Athena++ (Stone et al. 2019, in preparation). Athena++ is a newly developed grid-based code using a higher-order Godunov scheme for MHD and the constrained transport (CT) to conserve the divergence-free property for magnetic fields. Compared with its predecessor Athena (Gardiner & Stone 2005, 2008; Stone et al. 2008), Athena++ is highly optimized for speed and uses a flexible grid structure that enables mesh refinement, allowing global numerical simulations spanning a large radial range. Furthermore, the geometric source terms in curvilinear coordinates (e.g., in cylindrical and spherical-polar coordinates) are specifically implemented to conserve the angular momentum to machine precision. In this work, we adopt the second-order piecewise-linear method for the spatial reconstruction, the second-order Van-Leer method for the time integration, and the HLLC Riemann solver to calculate the flux.

The time-dependent radiative transfer equation has been solved explicitly and coupled with the MHD fluid equations using the radiation module of Jiang et al. (2014a). The general radiative transfer equation for the static fluid is
\[ \frac{1}{c} \frac{\partial I}{\partial t} + \mathbf{n} \cdot \nabla I = -\left( \sigma_{\nu,a} + \sigma_{\nu,f} \right) I_{\nu} + \mathbf{n} \cdot \sigma_{\nu,f} J_{\nu} \] (30)
where \( I_{\nu}(\mathbf{x}, t, \mathbf{n}) \) is the intensity at the position \( \mathbf{x} \), time \( t \) and along the direction of \( \mathbf{n} \), \( J_{\nu} = (4\pi)^{-1} \int \frac{d\Omega}{c} I_{\nu} \) and \( \mathbf{n} \cdot \sigma_{\nu,a} \) and \( \mathbf{n} \cdot \sigma_{\nu,f} \) are the absorption and effective scattering opacity at the frequency of \( \nu \). However, for a fluid that is moving at \( v \), additional correction terms on the order of \( (v/c) \) and \( (v/c)^2 \) need to be added (Jiang et al. 2014a). Jiang et al. (2019a) has adopted a mixed frame approach to solve the radiative transfer equation for moving fluid consistently. After integrating the radiative transfer equation over frequency, the equation becomes
\[ \frac{1}{c} \frac{\partial I}{\partial t} + \mathbf{n} \cdot \nabla I = S(I, \mathbf{n}) \] (31)

After carrying out the transport step in the lab frame, the
source terms on the right hand side are added. But instead of adding the source term $S(I, n)$ with all the $(v/c)$ and $(v/c)^2$ corrections to the intensity, the lab frame specific intensity $I(n)$ at angle $n$ is first transformed to the comoving frame intensity $I_0(n_0)$ via Lorentz transformation. Then the source terms in the comoving frame $(S_0(I_0, n_0))$ are added to $I_0(n_0)$,

$$S_0(I_0, n_0) = \sigma_{a,R} \left( \frac{a_r T^4}{4\pi} - I_0 \right) + \sigma_s (J_0 - I_0) + (\sigma_{a,P} - \sigma_{a,R}) \left( \frac{a_r T^4}{4\pi} - J_0 \right),$$

(32)

where $\sigma_{a,R} = \kappa_{a,R} \times \rho$, and $\sigma_{a,P} = \kappa_{a,P} \times \rho$, $\kappa_{a,R}$ and $\kappa_{a,P}$ are the Rosseland mean and Planck mean opacities. After this step to update $I_0(n_0)$, $I_0(n_0)$ are transformed back to the lab frame via Lorentz transformation. Then, the radiation momentum and energy source terms which are used in the fluid equations are calculated by the differences between the angular quadratures of $I(n)$ in the lab frame before and after adding the source terms.

For our particular FU Ori problem, we find that using the higher order PPM scheme (Colella & Woodward 1984) for the transport step is crucial for deriving the correct radiation fields in the extremely optically thick regime (see Section 3.2). Thus, the PPM scheme has been used in all our simulations for solving the radiative transfer equation. Since the characteristic speed in the transport step is the speed of light, solving this equation explicitly requires very small numerical timesteps. Thus, we adopt the reduced speed of light approach as in Zhang et al. (2018). We reduce the speed of light by a factor of 1000 in order to achieve a good timescale separation between the radiation transport and fluid dynamics. More discussions and tests on the reduced speed of light approach is in Section 3.2. We solve the radiative transfer equation along 80 rays in different directions. Integration of the specific intensity over angles yields various radiation quantities and source terms for the fluid equations.

The opacity that is adopted in the radiative transfer equation is generated in Zhu et al. (2007, 2009a). With this opacity, Zhu et al. (2007) find an excellent agreement between the synthetic spectral energy distributions and observations for FU Ori. This gives us great confidence to adopt it in this work for FU Ori hydrodynamical simulations. Both Rosseland mean and Planck mean opacities are shown in Figure 1. The dust opacity that is below $\sim 1500$ K is derived by the prescription in D'Alessio et al. (2001). The molecular, atomic, and ionized gas opacities have been calculated using the Opacity Distribution Function (ODF) method (Sbordone et al. 2004; Castelli & Kurucz 2004; Kurucz 2005) which is a statistical approach to handling line blanketing when millions of lines are present in a small wavelength range (Kurucz et al. 1974). More details on these opacities can be found in Zhu et al. (2009a) and Keith & Wardle (2014). On the other hand, we adopt a simple equation of state with a constant $\gamma = 5/3$ and $\mu = 1$ to avoid any complications due
to the change of $\gamma$ and $\mu$ or their interplay with the change of opacity at the same time.

Our grid setup is similar to Zhu & Stone (2018), where the whole 4$\pi$ sphere is covered by the spherical-polar $(r, \theta, \phi)$ grids with the special polar boundary condition in the $\theta$ direction (details in the appendix of Zhu & Stone 2018). The grid is uniformly spaced in $\ln(r)$, $\theta$, $\phi$ with $128 \times 64 \times 64$ grid cells in the domain of $[\ln(0.25), \ln(100)] \times [0, \pi] \times [0, 2\pi]$ at the root level. Two levels of mesh refinement have been adopted at the disk midplane. The outflow boundary condition which limits the inflow has been adopted at both the inner and outer radial boundaries.

The disk’s initial density, temperature, and velocity profiles are also similar to Zhu & Stone (2018) but with the midplane density slope of $p=2.125$, the temperature slope of $q=-3/4$, and $H/R=0.2$ at $R=1$. This is consistent with the structure of a viscously heated $\alpha$ disk. The initial disk scale height is thus resolved by 16 grids with two levels of mesh refinement. The density floor is also similar to Zhu & Stone (2018) except that an additional factor of $r_{\text{min}}/r$ was multiplied to Equation (10) of Zhu & Stone (2018) to further decrease the floor value at the disk atmosphere.

Simulations with both net vertical and net toroidal magnetic fields have been carried out. The net vertical field setup is similar to that in Zhu & Stone (2018) with a constant plasma $\beta$ at the disk midplane initially. In the net toroidal field simulations, magnetic fields are only present within 2 disk scale heights above and below the midplane initially, and the plasma $\beta$ is a constant anywhere within this region.

### 3.1 FU Ori Parameters and Simulation Runs

Our simulations adopt the disk parameters that are consistent with FU Ori observations. The detailed disk atmospheric modeling (Zhu et al. 2007, 2008) suggests that FU Ori’s inner accretion disk extends from $5 \ R_{\odot}$ to $\sim 1$ au with an accretion rate of $2.4 \times 10^{-3} \ M_{\odot} \ \text{yr}^{-1}$. The mass of the central star is $0.3 \ M_{\odot}$. The rotational axis of the disk is $55^\circ$ inclined with respect to our line of sight. Although these derived parameters are subject to change due to the recent Gaia distance measurement and ALMA disk inclination measurement for FU Ori (see Section 5.3), we will use these numbers as a guidance for our simulation parameters.

The length unit ($R = 1$) in the simulation is chosen as 0.1 au so that the whole domain extends from $5 \ R_{\odot}$ to $10$ au. The density unit is chosen as $10^{-8} \ g/\text{cm}^3$ with the midplane density of $10^{-7} \ g/\text{cm}^3$ at 0.1 au. The time unit is chosen as $1/\Omega$ at 0.1 au around a $0.3 \ M_{\odot}$ star. In this paper, we use $T_0$ to represent the orbital period ($2\pi/\Omega$) at 0.1 au around a $0.3 \ M_{\odot}$ star, which is 21 days.

Three main simulations have been carried out: 1) the disk that is initially threaded by net vertical fields with the strength of $\beta_0 = 1000$ at the disk midplane, labeled as V1000, 2) the disk that is threaded by vertical fields with $\beta_0 = 10^4$, labeled as V1e4, 3) the disk that is initially threaded by net toroidal fields with the strength of $\beta_0=100$, labeled as T100. We run these simulations to $T \sim 60 \ T_0$, which is equivalent to $\sim 3$ years. This time is equivalent to 500 innermost orbits in the simulation, and the disk at $R = 1$ has reached to a steady state as shown below.

#### 3.2 Code Tests

Although the radiative transfer scheme has been tested extensively (e.g. Jiang et al. 2014a, 2019a), we still need to test if the scheme is applicable to our particular FU Ori disk setup. Thus, we set up a 1-D plane-parallel atmosphere with a density profile of

$$\rho = \rho_0 e^{-z^2/2H^2},$$

(33)

to represent the disk vertical density structure at $R=1$ in our 3-D FU Ori simulations. $H$ is chosen as 0.02 au, and $\rho_0$ is chosen as $10^{-3} \ g/\text{cm}^3$. All other parameters are the same as our 3-D FU Ori simulations. To maintain this density structure, we don’t update the density and velocity during the run, and only allow the disk temperature to change. Only two rays have been used in this setup so that we can use two-stream approximation to calculate the analytical solution.

To represent the viscous heating in the accretion disk, we manually include a heating source term with the heating rate that is proportional to the disk local density ($\rho$) as

$$\frac{dE}{dt} = C \times \rho.$$

(34)

We have done three tests, two of which are steady state tests with a constant $C$ and one of which is the increasing heat test where $C$ suddenly increases at some time.

In the steady state tests, two different values of $C$ (0.0002316 and 0.02316 in the code unit) have been used to test if the disk can reach to the correct temperature in both low and high temperature regimes. The lower heating rate only heats the disk to $T \sim 10^3 \ K$, when the opacity is dominated by the dust and molecular opacity (the upper panels in Figure 2). The higher heating rate heats the disk sufficiently to test if the scheme is applicable to our particular FU Ori disk setup.
Figure 4. The poloidal slice of the temperature (the upper half) and density (the lower half) from the V1000 case at 50 $T_0$. This illustrated region represents FU Ori disk within 0.5 au from the central star. For the upper half of the image, the disk’s photosphere is illustrated with the iso-surface having $\rho R \times 0.1au = 10$.

to $T \sim 10^4$ K, when the opacity is dominated by the free-free and bound-free opacities (the lower panels in Figure 2).

These steady state tests show that we can accurately simulate the disk thermal structure, but also reveal the limitation of our setup. The black crosses in Figure 2 are results from simulations with 160 grids from -0.1 to 0.1 au (the same resolution as our 3-D simulations), while the red curves are from simulations with 1600 grids in the same domain range. The blue curves in the middle panel are the analytical solutions of this problem solved with the two-stream approximation:

$$T^4(\tau) = \frac{3}{4} T_{eff}^4 \left( \tau \left( 1 - \frac{\tau}{\tau_{tot}} \right) + \sqrt{\frac{1}{3}} \right),$$  (35)

where $\sigma T_{eff}^4$ is the flux emerging from half of the disk and $\tau_{tot}$ is the total optical depth from both sides of the disk. Clearly, when the opacity is low (e.g. the upper panels), the simulations with different resolutions agree with the analytical solution very well, even if the opacity has sharp changes among grids. On the other hand, when the opacity is high (e.g. the bottom panels), the optical depth can jump more than one order of magnitude from one grid to another grid. As expected, this jump leads to large errors in the calculations. Unfortunately, even with 10 times higher resolution (red curves in the lower panels), we still cannot recover the analytical solution accurately. One way to overcome this problem in future is using adaptive mesh-refinement for those grid cells having high optical depths. Overall, this test shows that, with our current setup, we may underestimate the temperature of some extremely optically thick grid cells by a factor of 2.

Since FU Ori’s disk temperature can change dramatically before and during the outburst, we also need to test if the code can capture the time evolution of the disk’s temperature accurately. Especially, our adoption of the reduced speed of light approach may delay the escape of the radiation energy. This is a particular concern when the disk is very optically thick (Skinner & Ostriker 2013) since the diffusion timescale $L \tau/c$ can now be longer than the dynamical timescale. For a typical size scale of 0.1 au and an optical depth of 1000, the radiation diffusion timescale is $\sim 1$ day. Naively, we would think that decreasing the speed of light by 1000 will increase the diffusion timescale to 1000 days, which is even longer than the total simulation timescale. On the other hand, it can be shown that the formulation in Zhang et al. (2018) guarantees that the radiative diffusion flux is the correct flux when the thermal energy of the gas dominates over the radiation energy. Thus, we should expect...
a correct diffusion timescale for our setup where the thermal energy of the gas always dominates. However, one could also argue that the optically thick region is joined by the optically thin region so that the disk will still cool/heat slower with the reduced speed of light approach.

To resolve these concerns, we carry out a test with a suddenly increased heating rate. We fix the absorption opacity to be 0.1 cm$^2$/g in this test. Initially, the disk is heated at the same heating rate as the above steady state test for a period of 2 $T_0$ so that the disk reaches to a steady state. Then, we suddenly increase the heating rate by a factor of 100 and watch the subsequent disk evolution. As shown in Figure 3, the reduced speed of light approach indeed slows down the heating of the disk. On the other hand, the temperature structure at 0.5 $T_0$ after the heating event for the disk using the reduced speed of light approach (the black dashed curve) overlaps with the temperature structure at 0.1 $T_0$ after the heating event for the disk using the normal speed of light (the red dotted curve). Thus, the reduced speed of light increases the diffusion timescale by a factor of $\sim 5$. This is larger than 1, but it is also much smaller than 1000 so that the diffusion timescale is still much smaller than the simulation timescale. Nevertheless, since the reduced speed of light approach increases the diffusion timescale to $\sim T_0$, we cannot trust short timescale variations of the radiation field in the simulations, and we can only study the state when the disk is relatively steady for the orbital timescale. Thus, in this paper, we only focus on the disk at the steady state with a constant accretion rate instead of discussing the outburst stage when the disk suddenly brightens by orders of magnitude within a short period of time.

4 RESULTS

The temperature and density structures of our fiducial model (V1000) at 50 $T_0$ are shown in Figure 4. We can see that the disk atmosphere at $z \sim R$ still has a significant density, which is similar to the disk structure in Zhu & Stone (2018). With the radiative transfer in our simulations, we can now study the disk’s temperature structure. The disk’s temperature is quite high ($\gtrsim$5000 K) close to the
central star (≲0.15 au). There is a sharp temperature jump around 0.15 au, indicating that the inner disk is at the upper branch of the equilibrium “S” curve which is dominated by the bound-free and free-free opacity while the outer disk is at the lower branch of the equilibrium “S” curve (≲2500 K) which is dominated by the molecular opacity. We also use $\rho R \times 0.1$ au ~10 to illustrate the disk’s photosphere. Clearly, the photosphere is hotter at the inner disk than the outer disk, and the photosphere is not smooth having noticeable structures. Due to these large scale structures at the photosphere, we expect that FU Ori has short timescale variations which has been implied by observations (Kenyon et al. 2000; Powell et al. 2012; Siwak et al. 2013).

After running for 50 $T_0$, our fiducial model has reached to a steady state within $R \sim 0.5$ au, i.e., the inner factor of ~20 in radius, as evident in Figure 5. From the mass accretion rate panel (the upper right panel), we can see that, at a later time, a larger disk region is accreting inwards since the outer disk region takes more time for MRI to grow. At 50 $T_0$, the region within 0.5 au, i.e., the inner factor of ~20 in radius, accretes inwards at a steady rate. Such constant accretion rates are also consistent with the stress shown in the middle left panel. The vertically integrated $R \theta$ stress follows $R^{-1.5}$ and this leads to a constant accretion rate based on Equation 10. Such accretion and stress structures are very similar to the global MHD simulations with the locally isothermal equation of state (compared with Figure 3 in Zhu & Stone 2018).

However, other quantities shown in Figure 5 are drastically different from those in Figure 3 of Zhu & Stone (2018). For example, the surface density in Figure 5 is almost flat, which is different from $R^{-0.8}$ in Zhu & Stone (2018). The midplane $\alpha$ is also flat compared with $R^{0.5}$ in Zhu & Stone (2018). Such differences are likely due to the temperature structure at the midplane. In the viscous heating dominated disk, the midplane temperature follows $R^{-3/4}$ (the lower left panel), while, in the locally isothermal simulations, the midplane temperature follows $R^{-1/2}$. Another evidence that the midplane temperature affects the $\alpha$ profile is that, at $R \sim 0.15$ AU where the midplane temperature jumps down, the $\alpha_{\text{total,mid}}$ there jumps up so that the total stress $T_{\text{total}}$ is still smooth. It is quite surprising that the accretion and stress profiles are smooth despite the jump of disk temperature. Considering that most stress is from the magnetic stress, this implies that the global disk accretion structure is mainly controlled by the global geometry of magnetic fields.
Figure 7. Density, velocities, temperature, mass flux, opacity, and optical depth along the $z$ direction at $R = 0.1$ au at 50 \( T_0 \). The quantities have been averaged azimuthally. The dashed curves in the velocity panels show the velocity components in the spherical-polar coordinates ($V_r$ and $-V_\theta$). The yellow curves are from the initial condition. The yellow shaded region labels the surface accreting region. Note the fast inward flow at the disk surface.

Figure 8. Similar to Figure 7 but for quantities that are related to magnetic fields. The dashed curves are B components in the spherical-polar coordinates ($B_r$ and $-B_\theta$). The blue curves in the $T_{R\phi}$ and $T_{\phi z}$ panels are magnetic stresses that are calculated using the mean fields, $-B_R \times \overrightarrow{B}_\phi$ and $-B_z \times \overrightarrow{B}_\phi$. The mean fields are azimuthally averaged before being used to calculate the stress.

and is insensitive to the disk local temperature. The magnetic fields at the midplane and $\theta = 0.78$ are shown in the lower right panel, and we can see that the field strengths change smoothly in the disk despite the temperature jump at $R \sim 0.15$ au.

4.1 Accretion Structure

The flow structure in MHD disks is tightly coupled with the magnetic field geometry. Magnetic fields determine the accretion structure while the accretion process drags and alters the magnetic fields. We plot the azimuthally averaged
temperature, density, and magnetic field structures for our fiducial run in Figure 6.

The velocity and magnetic field structures are remarkably similar to the “surface accretion” picture in locally isothermal disks with net vertical fields (Zhu & Stone 2018). Although we called such surface accretion as “coronal accretion” in Zhu & Stone (2018) following Beckwith et al. (2009), this accreting surface may not be hot (as shown in this work and Jiang et al. 2019b). Thus, in this work, we call this as “surface accretion” instead. The flow structure can be separated into three regions: the disk region which is dominated by MRI turbulence, the surface accreting region which is above the $\beta = 1$ surface and extends all the way to $z \sim R$, and the disk wind region (with $v_r > 0$) at $z \gg R$. The accretion flow mainly occurs at the surface, as shown in the middle panel of Figure 6 where the velocity streamlines are towards the star in the surface accreting region. Such surface inflow drags magnetic fields inwards so that the fields are pinched at the disk surface (the right panel of Figure 6). Due to the increase of the Keplerian rotation speed towards the inner disk, these dragged-in magnetic fields are sheared azimuthally, leading to fields with opposite $B_\phi$ between the lower and higher surface regions. Such surface accretion has been seen as early as Stone & Norman (1994) and recently in several simulations (Beckwith et al. 2009; Zhu & Stone 2018; Suriano et al. 2018; Takasao et al. 2018; Mishra et al. 2019; Jiang et al. 2019b). Analytical works by Guilet & Ogilvie (2012, 2013) have also seen such surface accretion when the turbulent viscosity and diffusivity are considered.

On the other hand, our radiation MHD simulations reveal new information on the disk thermal structure, especially the position of the disk photosphere. The left panel of Figure 6 shows that the thermal radiation field is very smooth except at the sharp jump $\sim 0.15$ au separating the two states that reside at the upper and lower branches of the “S” curve. If we integrate the Rosseland mean opacity along the $z$ direction (starting from $20^\circ$ off the axis to avoid the coarse grids at the pole), the derived $\tau_R = 1$ surface is plotted as the blue curves in all three panels. We can see that the $\tau_R = 1$ surface is at the wind base or upper surface accreting region at the inner disk ($\lesssim 0.07$ au) and within the lower surface accreting region at the outer disk ($\sim 0.07$ au). Thus, $B_\phi$ derived from the atomic lines at the photosphere could have opposite directions depending on where these lines are produced. This has important implications to the B field measurements of FU Ori, which will be discussed in greater details in Section 5.1.

To understand the disk’s accretion structure quantitatively, we plot the vertical profiles of various quantities at 0.1 au in Figure 7. The yellow shaded region is the surface accreting region. We see that the density flattens out in the surface accreting region, and the radial accretion velocity can reach 20 km/s there (the $v_R$ panel). Considering that
the Keplerian velocity is 50 km/s at 0.1 au, the surface inflow velocity is \( \sim 40\% \) of the Keplerian velocity. Due to the high speed, most disk mass is accreted through this surface accreting region despite its low density (the \( \rho v_r \) panel). The azimuthal velocity also deviates from the Keplerian velocity. In the surface accreting region, the lowest azimuthal velocity can reach to 60\% of the Keplerian velocity (the \( v_\phi \) panel). Such low azimuthal velocity and high radial velocity can be understood as magnetic breaking by the midplane so that the surface loses angular momentum and falls inwards. The midplane is very hot with a high opacity. Here at \( R = 0.1 \) au, the disk’s photosphere (\( \tau_R = 1 \)) is within the surface accreting region (the \( \tau \) panel).

The magnetic field structure at \( R = 0.1 \) au is shown in Figure 8. The surface inflow drags the initially vertical magnetic fields inwards, pinching the magnetic fields at the disk surface. The radial component of the magnetic fields in the net magnetic fields inwards, pinching the magnetic fields at the disk midplane. The radial component of the magnetic fields in the net inward accretion of the surface. In other words, the midplane is magnetically breaking the surface region. On the other hand, the internal \( T_{\phi} \) stress will only transfer angular momentum from the surface to the disk midplane, and thus it won’t lead to the overall disk accretion. The overall disk accretion is led by the \( T_R \) stress within the disk and the \( T_{\phi} \) stress at the disk atmosphere (e.g. the magnetocentrifugal wind). The detailed analysis on the surface accretion can be found in Zhu & Stone (2018). The accretion mechanisms are very similar. The only difference we notice by comparing Figure 8 in this work with Figure 7 in Zhu & Stone (2018) is that \( T_{\phi} \) plays a more important role in FU Ori simulations which are thicker than simulations in Zhu & Stone (2018). We have verified that the radiation viscosity is not important here. It is at least 5 orders of magnitude lower than the magnetic stress, which is different from the sub-eddington accretion disks around supermassive black holes (Jiang et al. 2019b).

Although it is mainly the magnetic field that determines the accretion process, the radiation pressure in FU Ori plays some role in supporting the disk. The lower panel of Figure 9 shows the force balance with various terms in the total momentum equation (Equation 1). In a steady state, the stress tensor divergence and the vertical gradient of the total pressure are balanced by the vertical component of the gravitational force and the radiation pressure force. For a slowly moving fluid, the radiation pressure force is \( -\sigma_r F_{r,\phi}/c \). Close to the disk midplane (the white region around \( z = 0 \)), it is mainly the gradient of the gas pressure (the red curve) that balances the vertical forces (the black curves). The magnetic pressure gradient (the blue curve) has the same strength as the radiation pressure (the black dotted curve, \( \sim 30\% \) of the gas pressure), and thus they balance each other. The stress tensor also plays some role in compressing the disk. In the surface accretion region, It is mainly the gradient of the magnetic pressure that balances the gravity. Both the radiation pressure and the gradient of the gas pressure are negligible in comparison. This again suggests that the surface accretion occurs in the magnetically dominated region.

4.2 Energy Budget

Angular momentum transport and energy transport are the two most important aspects of accretion disks. In Zhu & Stone (2018), we have done analyses on the angular momentum budget of accretion disks threaded by net vertical magnetic fields. With the radiative transfer included in this work, we will do similar analyses for the disk’s energy budget. The formulas are laid out in §2. Since the energy budget is related to the angular momentum budget, we will first repeat the angular momentum analysis as we did in Zhu & Stone (2018).

The angular momentum budget is shown in the upper panel of Figure 10. Four different terms in the angular momentum equation (Equation 19) are plotted. The \( m_{\theta} \) term is the radial gradient of the \( r-\theta \) stress (the first term on the right hand side of Equation 19). After the integration over a volume in the disk, this term represents the transport due to the internal stress exerted at the face that is perpendicular to the disk midplane, either from the turbulent stress or the stress due to the large scale organized magnetic fields. The \( m_{\phi} \) term is the \( \theta \) gradient of the \( r-\phi \) stress (the third term on the right hand side of Equation 19). After the integration over a volume, it is the stress that is exerted at the disk surface. That is normally due to the magnetocentrifugal disk wind. The other two terms (the \( m_r \) term, which is the
second term on the right hand side of Equation 19, and the $\dot{m}_r$ term, which is the forth term on the right hand side of Equation 19) are the momentum transport due to the radial and poloidal mass flux. In the thin disk theory, the poloidal mass flux is normally ignored so that the radial mass flux is balanced by the $m_r\phi$ and $m_\theta\phi$ terms during the steady state.

In Figure 10, these terms are integrated over $\theta$ from $\theta=0.59$ to 2.55 covering both the surface accreting region and the midplane region. Similar to the results in Zhu & Stone (2018), the wind stress ($m_\theta\phi$) plays a less important role in accretion than the $r-\phi$ stress. The $m_\theta\phi$ term is $\sim 1/4$ of the $m_r\phi$ term around $R \sim 1$. Thus, only 20% of accretion is due to the $\theta-\phi$ stress. On the other hand, this value is larger than 5% in the simulation of Zhu & Stone (2018). Considering that this disk is thicker than the disk in Zhu & Stone (2018), it implies that wind plays a more important role for accretion in thicker disks. Nevertheless, most accretion is still due to the internal $r-\phi$ stress within the disk.

On the other hand, the disk wind seems to play a much more important role in the energy transport. $\text{Flux}_r$, $\text{Flux}_\theta$, $E_{\text{cool}}$, and $E_{\text{pot}}$ in the lower panel of Figure 10 are the four terms on the right hand side of Equation 24. The traditional thin disk theory (Equation 29) actually adds the disk energy by an amount that is equal to half the released gravitational energy. The energy gain/loss in the poloidal direction is normally ignored. Thus, the total cooling rate is 1.5 times the released gravitational potential energy. However, our particular simulation suggests that energy transport in the radial direction (the red curve) is small compared with the energy loss in the poloidal direction by the wind (the blue curve). The wind carries half of the gravitational potential energy (the green curve) so that only the rest half gravitational potential energy needs to be radiated away (the cyan curve). Thus, the cooling rate is

$$\langle Q_{\text{cool}} \rangle = \frac{\dot{M}_r^2}{2\pi},$$

(36)

which is roughly $1/3$ of the value in the thin disk theory. This cooling rate is plotted as the green dashed curve in the lower panel of Equation 24, and it agrees with simulations very well (even at the inner disk close to the boundary). Thus, the disk temperature in the simulation can be approximated by

$$\sigma T_{\text{eff}}^4 = \frac{G\dot{M}M}{8\pi R^3}.$$  

(37)

Based on our simulations, such temperature estimate indeed agrees with the measured temperature at the $\tau \sim 1$ surface. The disk vertical structure at $R=0.1$ au is shown in Figure 11. At $\tau_R = 1$ (the dotted line in the right panels),
the value calculated using Equation 37 (the blue curve in the temperature panel) agrees with the measured temperature very well.

However, except for the similar $T_{eff}$, the temperature structure along $z$ in simulations is very different from the temperature structure based on the analytical theory. First, the radiation flux in the $\theta$ direction deviates significantly from the flux in the $z$ direction when $\tau \lesssim 1$ (the bottom panels in Figure 11). This is because the radiation from the inner disk ($R < R_0$) is so strong that the flux measured in the optically thin region at $R_0$ consists of a significant contribution from the disk inside $R_0$. Thus, we use the mea-

Figure 13. Similar to Figure 5 but for the V1000 case at $t = 55T_0$ (black curves) and the T100 case at $t = 60T_0$ (red curves). In the temperature and $(B^2)/2P_{\text{mid,0}}$ panels, the solid curves are the midplane quantities and the dashed curves are the quantities along $r$ at $\theta = 1.1$. The black dotted line in the temperature panel is from Equation 37 with an accretion rate of $10^{-3}M_\odot\text{yr}^{-1}$.

Figure 14. Similar to Figures 7 and 8 but for the V1000 case (black curves) and the T100 case (red curves).
sured flux at $\tau \sim 1$ to represent the flux emitted by the local annulus at $R$. Second, the measured flux in the $\theta$ direction rises much slower from the midplane to the $\tau = 1$ surface than the models (red and blue solid curves) where the heating rate is proportional to the disk local density (Equation 35). The measured radiative flux only rises quickly beyond one disk scale height. This is due to: 1) energy transport by turbulence is as important as the radiative energy transport within the disk so that less temperature gradient is needed to radiate the thermal energy, as shown in the upper panel of Figure 9; 2) both heating and accretion processes are more efficient at high above the disk midplane. Even with the similar emergent flux, the midplane temperature of the $\alpha$ disk model is hotter than the measured midplane temperature by a factor of $\gtrsim 3$. This result is consistent with previous local radiation MHD simulations (Turner 2004; Hirose et al. 2006; Jiang et al. 2014b), suggesting that, towards the disk surface, MHD heating is more efficient compared with heating in viscous models. Third, the emergent flux at $\tau = 1$ is significantly lower than the flux (red curves) estimated based on the traditional accretion disk theory (Equation 17) using the measured disk accretion rate of $4 \times 10^{-4} M_\odot \text{yr}^{-1}$. This is mostly due to the energy lost in the poloidal direction as discussed above. Equation 37 which has accounted for the energy loss in the poloidal direction agrees with the measured $F_\tau$ at $\tau = 1$ much better. We note that Equation 37 only stands at the inner disk. As shown in the temperature panel of Figure 5, the measured disk temperature is higher than the dotted line beyond $R \sim 0.2$ au. This is probably due to the fact that the outer disk is irradiated by the inner disk so that it gets heated up.

4.3 Different Field Strengths and Geometries

Since the disk temperature structure is self-consistently determined by the radiative transfer process in these simulations, the only major disk parameters that we can vary are the initial field geometry and strength. Thus, we carry out two additional simulations (V1e4 and T100) to explore how a weaker field or a toroidal field can affect the disk accretion. The disk temperature, density, velocity, and magnetic field structures are shown in Figure 12. Although these two simulations have similar temperature structures, one major difference which is quite noticeable in the middle panels is that disk wind fails to be launched in the net toroidal field simulations. In T100, disk material high above the atmosphere falls to the disk (green curves) instead of leaving the disk. Furthermore, the surface accreting region in T100 is much thinner if it exists at all. In the right panels, V1e4 shows an extended surface accreting region with high $B_\phi$ and $B_z$ values due to the surface accretion mechanism, while T100 only shows a thin region at the disk surface with noticeable $B_\phi$ and very weak fields above that. There is no large-scale organized fields in T100 either. The disk is dominated by turbulent fields in T100.

This lack of surface accretion in net toroidal field simulations is also evident in Figure 13 where the radial profiles of various quantities are shown. In the $T_{total}$ and $\alpha$ panels, the two simulations have similar values at the disk midplane for both $T_{rad}$ and $\alpha$, while the vertically integrated $T_{rad}$ and $\alpha$ are significantly higher for V1e4. This indicates that V1e4 has a higher stress level at the disk atmosphere than that in T100. The magnetic field panel also shows that, while $B^2$ at the midplane is similar between two simulations, V1e4 has much stronger fields at the disk atmosphere. This leads to a higher accretion rate for V1e4 even though these two simulations have very similar turbulent levels at the disk midplane.

The vertical profiles of various quantities clearly show the difference of disk wind between the net vertical and toroidal field simulations (Figure 14). At the wind region above $Z \sim R$, V1e4 has a much higher density than T100. The outflow nature of this region in V1e4 is clearly shown in the velocity panels, while this region in T100 is falling back to the disk. The magnetic fields and stresses are also very weak in the wind region of T100. Although there are some hints of surface accretion for T100 at $z/0.1$ au~1 shown in the $\tau_R$ panel, the density there is more than 5 orders of magnitude lower than the disk midplane (the $\rho$ panel) so that the radial accretion of this surface is negligible in net toroidal field simulations.

5 DISCUSSION

After studying the disk structure, we will compare the simulations with existing observations regarding both the magnetic field and disk wind.

5.1 Comparison with Magnetic Field Zeeman Observations

Donati et al. (2005) use a high resolution spectropolarimeter to measure circularly polarized light (Stokes $V$) from thousands of spectral lines for FU Ori. The circular polarized light is produced by Zeeman splitting which depends on both the field geometry and strength. The measured polarization signal corresponds to the line-of-sight magnetic field of $\sim 32$ G. Together with some additional constraints on the disk parameters (e.g. 60° inclination) and theoretical disk wind models (Ferreira 1997), the detailed decomposition of the Stokes $V$ into antisymmetric and symmetric components has put a much more stringent constraint on the magnetic fields of FU Ori. To summarize: 1) comparing the polarized light with the unpolarized light reveals that strong magnetic fields occupy $\sim 20\%$ of the disk surface, and the magnetic plasma rotates $\sim 2-3$ times slower than the local Keplerian velocity; 2) the vertical component of the magnetic fields (leaving the disk surface) is pointing towards us with a strength of $\sim 1$ kG at 0.05 au; 3) the toroidal field in the disk points to a direction which is opposite to the disk’s orbital rotation with a strength of $\sim 500$ G at 0.05 au.

Although these measurements are consistent with previous resistive MHD simulations (Ferreira 1997) where the MRI turbulence is simplified by the resistivity parameters, we can now compare these observations directly with our first-principle radiation MHD simulations. We thus measure the magnetic field direction and strength at the $\tau_R = 1$ surface in our simulations. The magnetic fields at $R=0.05$ au and 0.1 au are shown in Figure 15. Please note the direction of the field in this figure. $a_z$ is a parameter that equals 1 if $B_z$ at the $\tau_R = 1$ surface is pointing in a direction that is leaving the disk midplane and it is -1 if $B_z$ is pointing towards the disk midplane. $V\hat{z}$ is the unit vector in the disk’s rotational direction. The reason that we express $B_z$ in this $B_z \cdot V\hat{z} a_z$ form is due to the facts that we can view the disk from either the top or bottom side of the disk in Figure 6 and the disk’s
$B_z$ can also be either aligned or anti-aligned with the angular momentum vector of the disk's rotation. Let's take the V1000 case as an example. As shown in the upper left panel of Figure 15, $\vec{B}_\phi \cdot \vec{V}_\phi a_z$ (the solid black curve) is negative. If we observe the disk downwards from the upper side of the disk in Figure 6, $B_z$ is pointing to us so that $a_z = 1$. In this case $\vec{B}_\phi \cdot \vec{V}_\phi$ is negative implying that $\vec{B}_\phi$ is in the opposite direction from the disk rotation. This can be seen in Figure 6 where $\vec{B}_\phi$ has negative values in the wind region. If we view the disk from the bottom and $\vec{B}_z$ is pointing towards the disk midplane, $a_z = -1$ so that $\vec{B}_\phi$ at the $\tau_R = 1$ surface on this side of the disk is in the same direction as the disk rotation (as shown with the positive $B_\phi$ values at the bottom side of the wind region in Figure 6). On the other hand, since we don’t know if the rotational axis of the disk is aligned or anti-aligned with the magnetic fields (e.g. both Sun and Earth have magnetic reversals), we can reverse the field direction in simulations and the disk velocity structure will be unchanged. In that case, if we look at the disk downwards from the upper side of Figure 6, $a_z = -1$ and $\vec{B}_\phi$ at the wind region will be positive (in the same direction as the disk rotation) so that $\vec{B}_\phi \cdot \vec{V}_\phi a_z$ is still negative.

Our fiducial case (V1000) roughly reproduces the velocity and field geometries inferred from Donati et al. (2005). At R=0.05 au, the $\tau_R = 1$ surface is at $z \sim R$ which is the top of the surface accreting region or the bottom of the wind region (Figure 6). At $z \sim R$, the disk rotates with $\sim60\%$ of the midplane Keplerian velocity (the lower left panel of Figure 15), while the disk becomes Keplerian slightly deeper in the disk (the $V_\phi$ panel in Figure 7). At the $\tau_R = 1$ surface of R=0.05 au, the field strength is quite strong with $B_z \sim 150$ G. If $B_z$ is pointing to us, $V_\phi$ will be in a direction that is opposite to the disk rotation, which is consistent with observations. $B_\phi$ is half of $B_z$, which is also consistent with observations. At deeper regions in the disk, both $B_\phi$ and $B_z$ decreases significantly. In the surface accreting region and down towards the disk midplane, $B_\phi$ changes from negative to zero and to positive. Thus, the 20% covering factor from observations could be that 20% light comes from the strong $B$ and sub-Keplerian region, while the rest 80% comes from the deeper Keplerian and weaker $B$ region. The only difference between our simulations and the observations is that the field strength measured in simulations is weaker than the observed inferred kG strength by a factor of $\sim5$. On the other hand, we note that the first-order moment of the observed Zeeman signature is only $\sim32$ G. The kG strength is inferred from matching models considering the 60° inclination and the assumed filled geometry and filling factor. As will be shown in Section 5.3, the assumed inclination is too high compared with recent ALMA observations. Overall, the relatively good agreement regarding the field and velocity structure is very encouraging.

Our model also predicts that new observations by SpIROU at near-IR may reveal a different field structure.

Figure 15. Upper panels: the vertical (blue curves) and azimuthal (black curves) components of magnetic fields measured at the $\tau_R = 1$ surface at R=0.05 au (solid curves) and 0.1 au (dashed curves) for three simulations (from left to right panels). $a_z$ equals 1 if the $B_z$ field at the $\tau_R = 1$ surface is pointing in a direction that is leaving the disk midplane and equals -1 if the $B_z$ field is pointing towards the midplane. $\vec{V}_\phi$ is the unit vector in the disk’s rotational direction, and $\vec{B}_\phi$ is the projection of the magnetic field vector to the disk's rotational direction. Lower panels: the radial (blue curves) and azimuthal (black curves) velocity at the $\tau_R = 1$ surface at R=0.05 au (solid curves) and 0.1 au (dashed curves) for three simulations.
than earlier results using optical lines from Donati et al. (2005) since near-IR lines come from further out in the disk (e.g. 0.1 au). The simulation indicates that the $\tau_R = 1$ surface has very different field geometries and strengths at $R = 0.1$ au (the dashed curves in Figure 15) compared with those at $R = 0.05$ au. From Figure 6, we can see that, further away from the central star, the $\tau_R = 1$ surface is closer to the disk midplane due to the lower disk surface density there. The upper left panel in Figure 7 shows that both $B_z$ and $B_\phi$ at the $\tau_R = 1$ surface change their signs moving from 0.05 au to 0.1 au and the field strength gets a lot weaker. Furthermore, unlike at 0.05 au, $B_\phi$ is stronger than $B_z$ at the photosphere of 0.1 au since the photosphere is at the bottom of the surface accreting region and closer to the disk midplane.

The surface accreting regions in our other two simulations, V1e4 and T100, have much lower density so that the $\tau_R = 1$ surface is closer to the disk midplane even at $R = 0.05$ au (Figure 12). Thus, $B_\phi$ is always stronger than $B_z$ at the photosphere as shown in the right two panels of Figure 15. If $B_z$ is pointing towards us, $B_\phi$ will be in the same direction as the disk rotation in these cases.

Various possible scenarios for $B_z$ and $B_\phi$ measurements are summarized in Figure 16. Under the surface accretion picture, $B_z$ becomes quite strong at the upper surface/the base of the wind region at $R \sim z$, and $B_\phi$ changes sign there. Thus, if the disk has a very high density and the photosphere is only in the wind region or at the wind-base region (the thin dashed curve is the photosphere under this scenario), we are expecting to measure strong $B_z$ and $B_\phi$ at all disk radii. On the other hand, the disk normally has a lower density at the outer cooler region and the opacity there is lower, it is more likely that the photosphere changes from the wind-base region to the lower surface/disk region (e.g. V1000 case). In this case, the $B_z$ at the photosphere decreases dramatically at the outer disk and $B_\phi$ changes sign from the inner photosphere to the outer disk photosphere, indicating observations at different wavelengths may reveal different field and velocity geometries. For the third scenario that the photosphere is always closer to the disk (e.g. V1e4 and T100 cases), $B_z$ will be significantly smaller than $B_\phi$ at all radii and observations at different wavelengths may reveal similar field and velocity geometries. We note that the signs of various $B$ components can change depending on our viewing angle and the orientation between the fields and the rotational axis (as described in Figure 16).

5.2 Comparison with Disk Wind Observations

FU Ori shows evidence of strong winds in P Cygni profiles, especially in the Na I resonance lines (Bastian & Mundt 1985; Croswell et al. 1987). The blue-shifted line absorption implies a disk outflow with a typical velocity of 100-300 km/s and a mass loss rate of $\sim 10^{-3} M_\odot$ yr$^{-1}$ (Calvet et al. 1993). Recent work by Milliner et al. (2019) suggests that the wind may be turbulent.

We have plotted the gas radial velocity and mass loss rate at different poloidal directions in Figure 17. As long as the disk is threaded by net vertical fields, the magnetic fields accelerate the gas flow along the radial direction, reaching $\sim 400$ km/s terminal velocity. The integrated outflow rate at a distance $r$ from the central star is

$$M_{\text{wind}}(r) = \int_0^{2\pi} d\phi \int_0^\pi d\theta r^2 \sin(\theta) \rho v_r \sin(\theta) \rho v_r d\theta, \quad (38)$$

where $\int$ means that the quantities have been averaged over the azimuthal direction. The lower left panel of Figure 17 shows that $2\pi v^2 \sin(\theta) \rho v_r$ is around $10^{-3} M_\odot$ yr$^{-1}$. Thus, the integrated wind mass loss from the pole to 30$''$ (0.52 in Radian) away from the pole is $\sim 10^{-3} M_\odot$ yr$^{-1}$ * 0.52 ~ 10$^{-6} M_\odot$ yr$^{-1}$ which comes from both sides of the disk. Thus, our fiducial simulation can reproduce both the observed outflow velocity and outflow rate.

If the disk is threaded by net toroidal fields, wind can not be launched, as shown in the right panel of Figure 17. Thus, the existence of disk wind in FU Ori implies that the disk is threaded by net vertical magnetic fields.

5.3 New FU Ori Parameters

While we are preparing this manuscript, the distance to FU Ori is more precisely constrained by Gaia. The new distance is 416±9 pc (Gaia Collaboration et al. 2018) instead of 500 pc assumed in Zhu et al. (2007). The disk inclination is also better constrained to be 35$''$ by ALMA (Pérez et al. 2019) instead of 55$''$ assumed in Zhu et al. (2007). With these updated parameters, Pérez et al. (2019) derive that the central star mass is updated to be 0.6 $M_\odot$ instead of 0.3 $M_\odot$, and the disk accretion rate is $3.8 \times 10^{-5} M_\odot$ yr$^{-1}$ instead of $2.4 \times 10^{-4} M_\odot$ yr$^{-1}$. The disk accretion rate now is only 1/6 of the earlier estimate, since both the closer distance and more face-on configuration reduce the disk accretion rate estimate. In the Appendix and Figure A1, we have shown the SED fitting using the new parameters.

To be consistent with these new parameters, we have
Figure 17. The radial velocity (upper panels) and mass loss rate (lower panels) at 0.2 au and 1 au along the \( \theta \) direction in our three simulations (from left to right). The quantities have been averaged over both time (the last 2 \( T_0 \) close to the end of each simulation) and azimuthal direction.

carried out a simulation which is similar to the V1e4 case but with \( M_* = 0.6 M_\odot \). The results are shown in Figure 18. The overall “surface accretion” picture still stands. But due to the short duration of this simulation (only to 31.5 \( T_0 \)), the field structure at the surface accreting region is not fully established. The high disk accretion rate and the high central star mass releases a significantly amount of gravitational energy so that the disk is significantly hotter than the V1e4 case with \( M_* = 0.3 M_\odot \). The real FU Ori system may have weaker net vertical fields or a lower surface density than those we assumed in Figure 18.

6 CONCLUSIONS

We have carried out three-dimensional global ideal MHD simulations to study the inner outbursting disk of FU Ori. Since the accretion disk outshines the central star, the radiation field of the disk plays an important role in the disk accretion dynamics. The radiative transfer is also crucial for connecting with observations. Thus, we self-consistently solve the radiative transfer equations along with the fluid MHD equations. We have carried out simulations where the disk is threaded by either net vertical or net toroidal magnetic fields.

We find that, when the disk is threaded by net vertical fields, most accretion occurs in the magnetically dominated atmosphere at \( z \sim R \), very similar to the “surface accretion” mechanism in previous simulations with the simple locally isothermal equation of state. This implies that the “surface accretion” is a general feature for accretion disks threaded by net vertical fields. With radiative transfer, we can study the accretion disk’s temperature structure. The radiation pressure is \( \sim 30\% \) of the gas pressure at the inner disk (e.g. 0.1 au). The disk midplane has a sharp temperature transition at \( \sim 0.15 \) au separating the inner and outer disks which are at the higher and lower branches of the equilibrium “S” curve. But the accretion and stress profiles are smooth despite the jump of disk temperature. This implies that the global disk accretion structure is mainly controlled by the global geometry of magnetic fields and is insensitive to the disk local temperature.

Compared with the simulations for thinner disks in Zhu & Stone (2018), the simulations here have stronger disk wind. 20% of disk accretion is due to the wind stress, which is higher than 5% in Zhu & Stone (2018). The wind mass loss rate from the disk surface spanning one order of magnitude in radii is 1-10% of the disk accretion rate, which is also higher than 0.4% in Zhu & Stone (2018). Thus, the disk wind seems to be stronger in thicker disks. The mass loss rate of \( \sim 10^{-5} M_\odot \) yr\(^{-1} \) in our FU Ori simulations is consistent with observations. The wind’s terminal speed is \( \sim 300-500 \) km/s. This speed is also consistent with the observed wind speed and is several times the Keplerian speed at the launching point (\( V_K \) at the inner boundary is 100 km/s).
km/s). On the other hand, no disk wind is launched when the disk is threaded by net toroidal fields, implying that net vertical fields are crucial for launching the disk wind. The net toroidal field simulation also shows weaker accretion and smaller vertically integrated stresses due to the lack of the surface accretion at the disk surface.

The moderate disk wind also carries half of the accretion gravitational potential energy so that only the rest half of gravitational potential energy needs to be radiated away. The emergent flux is only \(\sim 1/3\) of the traditional value with the same disk accretion rate (comparing Equation 37 with Equation 18). Thus, the disk photosphere temperature is lower than that predicted by the thin \(\alpha\)-disk theory. Using the observed flux, the previously inferred disk accretion rate may be lower than the real disk accretion rate by a factor of \(\sim 2-3\). The disk midplane is also much cooler than that predicted by viscous models due to the energy transport by turbulence at the midplane and the efficient heating at the disk surface. With the surface accretion, the disk is heated up at the surface and the energy there can be more easily radiated away.

We have compared the magnetic fields at the photosphere in our simulations with Zeeman observations from Donati et al. (2005). The disk’s \(r_R = 1\) photosphere can be either in the wind launching region or the accreting surface region, depending on the accretion rates and the disk radii. Magnetic fields have drastically different directions and magnitudes between these two regions. It is very encouraging that the photosphere in our fiducial model, which is at the base of the wind launching region, agrees with previous Zeeman observations regarding both the field direction and magnitude. On the other hand, we suggest that the magnetic fields probed by future Zeeman splitting observations at different wavelengths (e.g. near-IR) or for different systems (e.g. with lower accretion rates) can be quite different from the existing measurements in Donati et al. (2005) since the photosphere can be deep into the surface accreting region.

Overall, we find excellent agreements between the first-principle MHD simulations having net vertical fields and existing observations regarding both the wind and magnetic field properties. This strongly supports that accretion disks in FU Orionis systems are threaded by net vertical magnetic fields and MHD processes are important for the accretion process. More comparisons between simulations and future observations will allow us to probe the 3-D structures of magnetic fields and gas flow in accretion systems.

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APPENDIX A: SED FITTING FOR FU ORI

With the updated FU Ori inclination, Pérez et al. (2019) use the disk atmospheric radiative transfer model (Zhu et al. 2007) to update FU Ori’s parameters. The best fit SED is shown in Figure A1.

REFERENCES

Armitage P. J., Livio M., Pringle J. E., 2001, MNRAS, 324, 705
Audard M., et al., 2014, in Beuther H., Klessen R. S., Dullemond C. P., Henning T., eds, Protostars and Planets VI. p. 387 ([arXiv:1401.3358]), doi:10.2458/azu_uapress_9780816531240-ch017
Bae J., Hartmann L., Zhu Z., Nelson R. P., 2014, ApJ, 795, 61
Balbus S. A., Hawley J. F., 1991, ApJ, 376, 214
Balbus S. A., Hawley J. F., 1998, Reviews of Modern Physics, 70, 1
Bastian U., Mundt R., 1985, A&A, 144, 57
Beckwith S., Hawley J. F., Kroll J. H., 2009, ApJ, 707, 428
Begelman M. C., Blandford R. D., Rees M. J., 1984, Reviews of Modern Physics, 56, 255
Bell K. R., Lin D. N. C., 1994, ApJ, 427, 987
Blandford R. D., Payne D. G., 1982, MNRAS, 199, 883
Bonelli E., Bastian P., 1992, ApJ, 401, L31
Calvet N., Hartmann L., Kenyon S. J., 1993, ApJ, 402, 623
Castelli F., Kurucz R. L., 2004, A&A, 419, 725

Colella P., Woodward P. R., 1984, Journal of Computational Physics, 54, 174
Connelley M. S., Reipurth B., 2018, ApJ, 861, 145
Croswell K., Hartmann L., Avrett E. H., 1987, ApJ, 312, 227
D’Alessio P., Calvet N., Hartmann L., 2001, ApJ, 553, 321
Donati J.-F., Paletou F., Bouvier J., Ferreira J., 2005, Nature, 438, 466
Event Horizon Telescope Collaboration et al., 2019, ApJ, 875, L1
Ferreira J., 1997, A&A, 319, 340
Frank A., et al., 2014, in Beuther H., Klessen R. S., Dullemond C. P., Henning T., eds, Protostars and Planets VI. p. 451 ([arXiv:1402.3553]), doi:10.2458/azu_uapress_9780816531240-ch020
Fu R. R., et al., 2014, Science, 346, 1089
Gaia Collaboration et al., 2018, A&A, 616, A1
Gardiner T. A., Stone J. M., 2003, Journal of Computational Physics, 205, 509
Gardiner T. A., Stone J. M., 2008, Journal of Computational Physics, 227, 4123
Guilet J., Ogilvie G. I., 2012, MNRAS, 424, 2097
Guilet J., Ogilvie G. I., 2013, MNRAS, 430, 822
Hartmann L., 1998, Accretion Processes in Star Formation
Hartmann L., Kenyon S. J., 1996, ARA&A, 34, 207
Hawley J. F., Gammie C. F., Balbus S. A., 1995, ApJ, 440, 742
Herbig G. H., 1977, ApJ, 217, 695
Hillenbrand L. A., Findeisen K. P., 2015, ApJ, 808, 68
Hillenbrand L. A., et al., 2018, ApJ, 869, 146
Hirose S., 2015, MNRAS, 448, 3105
Hirose S., Kroll J. H., Stone J. M., 2006, ApJ, 640, 901
Hirose S., Blaes O., Kroll J. H., Coleman M. S. B., Sano T., 2014, ApJ, 787, 1
Jiang Y.-F., Stone J. M., Davis S. W., 2013, ApJ, 778, 65
Jiang Y.-F., Stone J. M., Davis S. W., 2014a, ApJS, 213, 7
Jiang Y.-F., Stone J. M., Davis S. W., 2014b, ApJ, 784, 149
Jiang Y.-F., Stone J. M., Davis S. W., 2019a, ApJ, 880, 67
Jiang Y.-F., Blaes O., Stone J. M., Davis S. W., 2019b, ApJ, 885, 144
Kadom K., Vorobyov E., Regály Z., Ábrahám P., 2019, ApJ, 882, 96
Keith S. L., Wardle M., 2014, MNRAS, 440, 89
Kenyon S. J., Kokotilo E. A., Ibrahimov M. A., Mattei J. A., 2000, ApJ, 531, 1028
Kley W., Lin D. N. C., 1999, ApJ, 518, 833
Kóspál Á., Ábrahám P., Westhus C., Haas M., 2017, A&A, 597, L10
Kraus S., Caratti o Garatti A., Garcia-Lopez R., Kreplin A., Aarnio A., Naylor T., Weigelt G., 2016, MNRAS, 462, L61
Kurucz R. L., 2005, Memorie della Societa Astronomica Italiana Supplementi, 8, 14
Kurucz R. L., Peytremann E., Armitage P. J., Simon J. B., 2019, ApJ, 885, 144
Kurucz R. L., Peytremann E., Avrett E. H., 1974, Blanketed model atmospheres for early-type stars
Martin R. G., Lubow S. H., Livio M., Pringle J. E., 2012, MNRAS, 423, 2718
Millner K., Matthews J. H., Long K. S., Hartmann L., 2019, MNRAS, 483, 1663
Mishra B., Begelman M. C., Armitage P. J., Simon J. B., 2019, arXiv e-prints, p. arXiv:1907.08995
Pérez S., et al., 2019, arXiv e-prints, p. arXiv:1911.11282
Powell S. L., Irwin M., Bouvier J., Clarke C. J., 2012, MNRAS, 426, 3315
Pudritz R. E., Ouyed R., Fedt C., Brandenburg A., 2003, Protostars and Planets V, pp 277–294
Shobdane L., Bonifacio P., Castelli F., Kurucz R. L., 2004, Memorie della Societa Astronomica Italiana Supplementi, 5, 93
Scholz A., Froebrich D., Wood K., 2013, MNRAS, 430, 2910
Semkov E. H., Peneva S. P., Munari U., Milani A., Valisa P., 2010, A&A, 523, L3
Siwak M., et al., 2013, MNRAS, 432, 194
Skinner M. A., Ostriker E. C., 2013, ApJS, 206, 21
Stone J. M., Norman M. L., 1994, ApJ, 433, 746
Stone J. M., Gardiner T. A., Teuben P., Hawley J. F., Simon J. B., 2008, ApJS, 178, 137
Suriano S. S., Li Z.-Y., Krasnopolsky R., Shang H., 2018, MNRAS, 477, 1239
Takasao S., Tomida K., Iwasaki K., Suzuki T. K., 2018, ApJ, 857, 4
Turner N. J., 2004, ApJ, 605, L45
Turner N. J., Fromang S., Gammie C., Klahr H., Lesur G., Wardle M., Bai X.-N., 2014, Protostars and Planets VI, pp 411–432
Vlemmings W. H. T., et al., 2019, A&A, 624, L7
Vorobyov E. I., Basu S., 2006, ApJ, 650, 956
Zhang D., Davis S. W., Jiang Y.-F., Stone J. M., 2018, ApJ, 854, 110
Zhu Z., Stone J. M., 2018, ApJ, 857, 34
Zhu Z., Hartmann L., Calvet N., Hernandez J., Muzerolle J., Tannirkulam A.-K., 2007, ApJ, 669, 483
Zhu Z., Hartmann L., Calvet N., Hernandez J., Tannirkulam A.-K., D’Alessio P., 2008, ApJ, 684, 1281
Zhu Z., Hartmann L., Gammie C., 2009a, ApJ, 694, 1045
Zhu Z., Hartmann L., Gammie C., McKinney J. C., 2009b, ApJ, 701, 620

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