Magnetically Induced Disk Winds and Transport in the HL Tau Disk

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Abstract

The mechanism of angular momentum transport in protoplanetary disks is fundamental to understanding the distributions of gas and dust in the disks. The unprecedented ALMA observations taken toward HL Tau at high spatial resolution and subsequent radiative transfer modeling reveal that a high degree of dust settling is currently achieved in the outer part of the HL Tau disk. Previous observations, however, suggest a high disk accretion rate onto the central star. This configuration is not necessarily intuitive in the framework of the conventional viscous disk model, since efficient accretion generally requires a high level of turbulence, which can suppress dust settling considerably. We develop a simplified, semi-analytical disk model to examine under what condition these two properties can be realized in a single model. Recent, non-ideal MHD simulations are utilized to realistically model the angular momentum transport both radially via MHD turbulence and vertically via magnetically induced disk winds. We find that the HL Tau disk configuration can be reproduced well when disk winds are properly taken into account. While the resulting disk properties are likely consistent with other observational results, such an ideal situation can be established only if the plasma $\beta$ at the disk midplane is $\beta_0 \simeq 2 \times 10^2$ under the assumption of steady accretion. Equivalently, the vertical magnetic flux at 100 au is about 0.2 mG. More detailed modeling is needed to fully identify the origin of the disk accretion and quantitatively examine plausible mechanisms behind the observed gap structures in the HL Tau disk.

Key words: accretion, accretion disks – magnetic fields – magnetohydrodynamics (MHD) – protoplanetary disks – stars: individual (HL Tauri) – turbulence

1. Introduction

Protoplanetary disks are the birthplace of planetary systems, including the solar system. The spatial distributions of both gas and dust in the disks provide the initial conditions for planet formation. The time evolution of the disks dictates when and where planet formation begins, and eventually determines the final architecture of planetary systems both in mass and in orbital distance (e.g., Ida & Lin 2004; Mordasini et al. 2009, 2016; Benz et al. 2014; Hasegawa 2016). Such distributions of planets can be confronted with a wealth of observed exoplanetary populations (e.g., Udry & Santos 2007; Borucki et al. 2011; Mayor et al. 2011; Winn & Fabrycky 2015; Twicken et al. 2016).

It is widely perceived that magnetic fields threading disks can regulate disk evolution (e.g., Armitage 2011). While the ultimate mechanism behind the observed high disk accretion rate is still uncertain (e.g., Hartmann et al. 2016), magnetorotational instability (MRI) and the resulting MHD turbulence can act as a promising candidate to account for the observations (e.g., Balbus & Hawley 1998). This becomes possible because MRI can efficiently transport the disk angular momentum radially by coupling magnetic fields with differentially rotating ionized gas in disks. Supportive evidence for the presence of magnetic fields in disks may exist in the form of magnetized chondrules found in chondrites (Fu et al. 2014). The results of the lab experiments imply that magnetic fields should have played an important role in both the distribution and the growth of planet-forming materials in the solar nebula (e.g., Shu et al. 1996; Desch & Cuzzi 2000; Ciesla 2007; Hasegawa et al. 2016; Wang et al. 2017).

The MRI and the resulting MHD turbulence do not necessarily operate fully in protoplanetary disks, since the disks are generally dense and cold in most regions (e.g., Turner et al. 2014). Accordingly, non-ideal MHD effects become important in order to realistically model a coupling between charged materials and magnetic fields in disks and to reliably compute the transport of disk angular momentum. Ohmic resistivity has been investigated extensively over the past two decades. It is known that inclusion of ohmic resistivity leads to the realization of layered accretion (e.g., Gammie 1996; Armitage et al. 2001; Fleming & Stone 2003; Terquem 2008; Kretke & Lin 2010; Hasegawa & Takeuchi 2015). In this picture, the disk angular momentum is transported radially via MHD turbulence; a higher disk accretion rate is achieved in the MRI-active surface layers, while MRI is quenched and MRI-dead zones are present in the disk midplane, which leads to a lower disk accretion rate there. Compared with ohmic resistivity, other non-ideal MHD terms (ambipolar diffusion and the Hall term) have been less explored (e.g., Wardle 1999; Sano & Stone 2002). The situation has changed rapidly since the recognition that ambipolar diffusion can suppress MRI in the disk surfaces significantly (e.g., Bai & Stone 2011). In this case, the observed high disk accretion rate can be reproduced only when disks are threaded by strong magnetic fields (e.g., Simon et al. 2013a). This condition is needed because then magnetically induced disk winds can be launched from the disk surface and efficiently transport the disk angular momentum vertically (e.g., Suzuki & Inutsuka 2009; Bai 2013b; Simon et al. 2013a). Inclusion of different non-ideal MHD terms thus draws a very different picture of how protoplanetary disks evolve with time (e.g., Bai 2016; Suzuki et al. 2016); depending on the strength of the vertical magnetic fields, the disk angular momentum can possibly be transported radially via the MRI and the resulting MHD turbulence and/or vertically via magnetically induced disk winds.
Observations of protoplanetary disks serve as an important complementary approach to understanding disk evolution and planet formation there (e.g., Williams & Cieza 2011). The advent of ALMA with significantly upgraded sensitivity and angular resolution has indeed revolutionized our view of planet formation (e.g., ALMA Partnership et al. 2015a). The best example is the observation taken toward HL Tau as a part of the ALMA long-baseline science verification campaign (ALMA Partnership et al. 2015b). The astonishing observations reveal that nearly concentric, multiple gaps are present in the dust continuum emission in the disk, which may be a potential signature of planet formation in a young stellar object (YSO). It is remarkable that similar gap structures in disks have recently been reported for a number of other targets (e.g., Andrews et al. 2016; Isella et al. 2016; Nomura et al. 2016; Tsukagoshi et al. 2016; Fedele et al. 2017a). Soon after the public release of the HL Tau image, a number of mechanisms were proposed to explain the origin of gaps for the HL Tau disk. These include disk–planet interaction (e.g., Dipierro et al. 2015; Dong et al. 2015; Kanagawa et al. 2015; Akiyama et al. 2016; Jin et al. 2016), the effect of ice lines (Zhang et al. 2015), sintering effects for growing dust particles (Okuzumi et al. 2016), and secular gravitational instabilities (Takahashi & Inutsuka 2016).

While all the studies can (partially) reproduce the observed gap structure for the HL Tau disk, the predominant mechanism of such local structures has not yet been fully identified. This is probably because different mechanisms are examined in different disk models, and hence a direct comparison among them may not be trivial. Furthermore, it is evident from the Brγ observations that the HL Tau disk has a high accretion rate ($\sim 10^{-7} M_\odot$ yr$^{-1}$, Beck et al. 2010). This is intuitively expected since HL Tau is a relatively young system ($\lesssim 1$ Myr). Most importantly, radiative transfer modeling of the ALMA data indicates that millimeter-sized grains have settled greatly at $\sim 100$ au from the star, reducing their scale height to just $\sim 1$ au (Pinte et al. 2016). Turbulence at this location must then be weak. Further constraints on the strength of turbulence come from the widths of the millimeter-wave emission lines of several gas molecules (Hughes et al. 2011; Guilloteau et al. 2012; Flaherty et al. 2015, 2017b), though concerns have been raised about the measurements' sensitivity to uncertainties in the gas temperature structure and flux calibration (Teague et al. 2016). Weak turbulence may be incompatible with the conventional viscous disk model, because rapid accretion requires strong turbulence, which counteracts the grains’ settling. Thus, it is worth examining the HL Tau disk’s global configuration before addressing the origin of the observed gaps.

In this paper, we build such a global view, focusing especially on the disk accretion rate and vertical mixing of dust particles by disk turbulence. We develop a semi-analytical disk model, making use of recent non-ideal MHD simulations. The key ingredient in our model is to properly include the effects of both MHD turbulence and disk winds simultaneously. We will show below that in the framework of magnetically driven disk accretion, the two observational features (a high disk accretion rate and efficient dust settling) can be reproduced in single-disk models when the plasma β at the disk midplane is $\beta_0 \approx 2 \times 10^4$ under the assumption of steady accretion.

This paper is organized as follows. In Section 2, we develop a simplified, but physically motivated disk model, using the results of recent non-ideal MHD simulations. In Section 3, we present our results and examine how the observed properties of the HL Tau disk can be reproduced in our model. We also perform a parameter study and consider applications of our results. These include the global configuration of magnetic fields and the gas-to-dust ratio. In Section 4, we take into account other observational features of the HL Tau disk. We also discuss limitations of our models and other potential mechanisms of disk accretion for the HL Tau system. Our conclusions are provided in Section 5.

### 2. Disk Model

#### 2.1. Magnetically Driven Disk Accretion

When magnetic fields play an important role in transporting angular momentum in disks, the disk accretion rate ($\dot{M}$) can be written as (e.g., Fromang et al. 2013; Suzuki et al. 2016)

$$\dot{M} = \frac{4\pi}{r\Omega} \left[ \frac{\partial}{\partial r} \int_{-H_*}^{H_*} dz \left( \rho v_z \delta v_\phi - \frac{B_x B_z}{4\pi} \right) \right] + r^2 \left( \frac{\rho v_z \delta v_\phi - \frac{B_x B_z}{4\pi}}{z=-H_*} \right) \left( \frac{\rho v_z \delta v_\phi - \frac{B_x B_z}{4\pi}}{z=-H_*} \right),$$

where the cylindrical coordinate system is adopted. The conventional notation is used here: $r$ is the disk radius from the central star, $\Omega = \sqrt{GM_\star/r^3}$ is the orbital frequency, $M_\star$ is the mass of the central star, $\rho$ is the volume density of the gas, and $v_\phi$ and $B_z$ are the $i$th components of the gas velocity and the magnetic field, respectively. Also, $\delta v_\phi$ is computed as $v_\phi - \Omega r$, and time-averaged values are utilized for some physical quantities, as denoted by an overbar. In addition, it is assumed in the equation that disk surfaces are located at $z = \pm H_*$ from the midplane, and disk winds will be launched from these surfaces. As discussed below (see Section 2.5), $H_* \approx 2H_T$, where $H_T = c_s/\Omega$ is the pressure scale height and $c_s$ is the sound speed. This equation is obtained from a vertical integration of the azimuthal component of the MHD momentum equation over $|z| < H_*$, and is applicable to both nonsteady and steady disk accretion.

We introduce the following notation for the normalized accretion stresses, which are written as

$$W_\phi \equiv \int_{-H_*}^{H_*} \left( \rho v_z \delta v_\phi - \frac{B_x B_z}{4\pi} \right) dz \frac{\Sigma_g c_s^2}{\Sigma_g c_s^2},$$

and

$$W_0 \equiv \frac{\rho v_z \delta v_\phi - \frac{B_x B_z}{4\pi}}{2\rho_{mid} c_s^2},$$

where $\Sigma_g$ is the gas surface density, and $\rho_{mid} = \Sigma_g \Omega/(\sqrt{2\pi c_s})$ is the gas volume density at the disk midplane. The vertically isothermal assumption is adopted for $c_s$. Here we follow the notation used in Simon et al. (2013a): our $W_\phi$ and $W_0$ correspond to their $\alpha_{\text{turb}}$ and $W_{0\text{turb}}$, respectively. Note that the normalization constants are different for the above two equations. This is because $W_\phi$ is the accretion stress integrated over the vertical direction, while the wind stress ($W_0$) can be regarded as the excess of the angular momentum flux coming into/out of the disk surfaces. This difference is properly taken into account when computing the disk accretion rate (see a
factor of $2r/(\sqrt{\pi}H_g)$ in Equation (9)). Then, Equation (1) can be rewritten as

$$M = \frac{4\pi}{\Omega r} \left( r^2 \frac{\partial}{\partial r} (2^\frac{3}{2} \Sigma_S W_{r0}) + 4\sqrt{2}\pi r \Sigma_S c_s W_{r0} \right).$$  

(4)

Thus, magnetic fields can regulate disk evolution by transporting the disk angular momentum both radially via MHD turbulence ($W_{r0}$) and vertically via magnetically induced disk winds ($W_{z0}$).

### 2.2. Conventional Viscous Disk Models

The classical 1D viscous disk model is still widely adopted in the literature to model protoplanetary disks. In this picture, it is assumed that the radial angular momentum transport is determined by a kinematic viscosity ($\nu$) without specifying its origin, and the disk accretion rate can be written as (e.g., Pringle 1981)

$$M = \frac{6\pi}{\Omega r} \left( r^2 \Omega \Sigma_S \right).$$  

(5)

We can obtain a relation between the accretion stress ($W_{r0}$) and viscosity ($\nu$) by comparing Equation (5) with the first term on the right of Equation (4):

$$W_{r0} = \frac{3}{2} \frac{\nu \Omega}{c_s^2}.$$  

(6)

It follows that the viscous $\alpha$-parameter ($\alpha_{SS}$) can be expressed as a function of $W_{r0}$ (Shakura & Sunyaev 1973):

$$\alpha_{SS} = \frac{2}{3} W_{r0},$$  

(7)

where $\alpha_{SS} \equiv \nu \Omega / c_s^2$. Note that the factor of 2/3 is often neglected in the literature.

Since Equation (5) considers only the radial angular momentum transport, it would work well for modeling the accretion rate and the resulting disk structure when turbulence plays a dominant role in disk evolution. In other words, disk winds should not transport the angular momentum efficiently either due to weak vertical magnetic flux or due to an unfavorable geometrical configuration of the magnetic field.

### 2.3. Steady-state Solutions

We explore steady disk accretion ($M = \text{const.}$) using Equation (4). While a more general solution can be obtained by integrating Equation (4) along the radial direction, we derive here an approximate but analytic solution by examining two limiting cases of disk accretion.

First, we consider the limit where disk accretion is dominated by the internal disk stress ($W_{r0}$). In this limit, the steady-state condition ($M = \text{const.}$) can be realized when the product $c_s^2 \Sigma_S W_{r0}$ is proportional to $\Omega$ ($\propto r^{-3/2}$). Then, $\partial(r^2 \Sigma_S W_{r0}) / \partial r = (1/2) r c_s^2 \Sigma_S W_{r0}$. Accordingly, we find that the disk accretion rate can be written as

$$\dot{M} \approx \frac{2\pi \Sigma_S c_s^2 W_{r0}}{\Omega}.$$  

(8)

The other limit is that disk accretion is regulated predominantly by the wind stress, that is, $\dot{M} \approx 4\sqrt{2}\pi r \Sigma_S c_s W_{r0}$.

Second, we combine the solutions of the above two limits to consider a more general situation. For this case, the steady disk accretion rate can be approximately given as

$$\dot{M} \approx \frac{2\pi \Sigma_S c_s^2 W_{r0}}{\Omega} + 4\sqrt{2}\pi r \Sigma_S c_s W_{r0} = \frac{2\pi \Sigma_S c_s^2}{\Omega} \left( W_{r0} + \frac{2r}{\sqrt{\pi} H_g} W_{z0} \right).$$  

(9)

Once $\dot{M}$ is obtained, we can also compute both the surface density profile ($\Sigma_g$) and the effective $\alpha$-parameter ($\alpha_{SS,\text{eff}}$) for disks with $M = \text{const.}$:

$$\Sigma_g = \frac{M \Omega}{2 \pi c_s^2} \left( W_{r0} + \frac{2r}{\sqrt{\pi} H_g} W_{z0} \right)^{-1}$$  

(10)

and

$$\alpha_{SS,\text{eff}} = \frac{2}{3} \left( W_{r0} + \frac{2r}{\sqrt{\pi} H_g} W_{z0} \right).$$  

(11)

Note that $\alpha_{SS,\text{eff}}$ is derived from the comparison between Equation (9) and the disk accretion rate in the 1D viscous model, which can be given as (see Equation (5))

$$\dot{M} = 3\pi \alpha_{SS} c_s^2 \Sigma_g / H_g.$$  

(12)

Thus, disk winds provide an additional contribution to $M$, $\Sigma_g$, and the $\alpha$-parameter. Furthermore, since the contribution is scaled by $r / H_g$, they can play a dominant role in the evolution when disks are threaded by magnetic fields that have a corresponding plasma $\beta_0$ of $10^3$ or lower at the disk midplane (see Figure 1).

### 2.4. Dust Settling

Dust settling is one of the key processes from which the level of disk turbulence in protoplanetary disks can be estimated (e.g., Mulders & Dominik 2012). This is because the degree of dust settling is regulated both by the size of dust particles and by disk turbulence (e.g., Dubrulle et al. 1995; Youdin & Lithwick 2007). Here we briefly summarize the formulation of dust settling that is used in this paper.

The vertical diffusion coefficient for small dust particles can be defined as (e.g., Zhu et al. 2015)

$$D_z = \frac{1}{2} \frac{d \langle z^2 \rangle}{dt},$$  

(13)

where the brackets denote the ensemble average taken over the trajectories of a large number of particles. Then, the normalized diffusion coefficient can be written as

$$\alpha_D \equiv \frac{D_z}{c_s H_g}.$$  

(14)

In general, dust particles with a radius $a$ ($< 1$–$10$ mm) in protoplanetary disks can reside in the so-called Epstein regime where $a$ is smaller than the mean free path of gas particles. For this case, the Stokes number that is the dimensionless stopping time can be given as $St = \pi \rho_a a / (2 \Sigma_g)$ for such particles (e.g., Birnstiel et al. 2012). Finally, the scale height ($H_D$) of the dust particles can be written as (e.g., Youdin & Lithwick 2007)

$$H_D = \left( 1 + \frac{St}{\alpha_D} \right)^{-1/2} H_g.$$  

(15)
This is the outcome in which dust settling equilibrates with the vertical stirring due to disk turbulence. Note that $\alpha_D$ need not to be the same as $\alpha_{SS}$ and $\alpha_{SS,eff}$.

2.5. Fitting Formulae for $W_{r\phi}$, $W_{z\phi}$, and $\alpha_D$

Based on the above formalism, three quantities ($W_{r\phi}$, $W_{z\phi}$, and $\alpha_D$) are needed to uniquely determine the spatial distributions of both gas and dust for a given disk temperature profile (Equations (9) and (15)).

Here, we make use of recent non-ideal MHD simulations to model these quantities as a function of the plasma $\beta$, where

$$\beta = \frac{8\pi \rho_{mg} c_s^2}{B_z^2}. \quad (16)$$

More specifically, the results of simulations obtained by Simon et al. (2013a) and Zhu et al. (2015) are utilized, which are both appropriate for the outer part of protoplanetary disks. Simon et al. (2013a) perform vertically stratified, 3D MHD simulations with ambipolar diffusion using a local shearing box. In these simulations, the values of $W_{r\phi}$ and $W_{z\phi}$ are measured at the disk height of the wind base ($z = \pm z_{bw}$) for three different values of the initial $\beta$ ($\beta_0 = 10^3$, $10^4$, and $10^5$, see Table 1 of Simon et al. 2013a). They find that $z_{bw} \approx 2 H_g$ for a wide range of $\beta_0$. In this paper, it is assumed that $H_g = z_{bw}$ (see Equation (1)). For the actual fitting, we follow the approach taken by Armitage et al. (2013), where simple fitting formulae are derived to reproduce the simulation results:

$$\log W_{r\phi} = -2.2 + 0.5 \tan^{-1} \left( \frac{4.4 - \log \beta_0}{0.5} \right), \quad (17)$$

and

$$\log W_{z\phi} = 1.25 - \log \beta_0. \quad (18)$$

While the formula for $W_{z\phi}$ is identical to that used in Armitage et al. (2013), our formula for $W_{r\phi}$ is slightly different from theirs because here we consider the internal disk stress averaged over $|z| \leq z_{bw}$ rather than $|z| \leq 4 H_g$.

For $\alpha_D$, the results of Zhu et al. (2015) are adopted, wherein the authors perform both unstratified and stratified 3D, non-ideal MHD simulations using a local shearing box with Lagrangian particles. In these simulations, ambipolar diffusion is taken into account, and the values of $\alpha_D$ are measured over $|z| \leq H_g$ for different values of the initial $\beta$. Based on the results of their unstratified simulations (see their Table 3), we find that $\alpha_D$ is roughly proportional to $1/\beta_0$. Mathematically, the resultant fitting formula can be expressed as

$$\log \alpha_D = 1.1 - \log \beta_0. \quad (19)$$

Figure 1 shows the fitting profiles for $W_{r\phi}$, $W_{z\phi}$, and $\alpha_D$ in the left panel and the resulting $\alpha_{SS,eff}$ and $\alpha_{SS}$ in the right panel. As an example, we adopt $r/H_g = 10$ when computing $\alpha_{SS,eff}$ in this figure (see Equation (11)). In addition, the numerical results utilized for fitting are also plotted; the blue squares and the black triangles are adopted from Simon et al. (2013a). The red circles and the yellow circle are from the results of numerical simulations done by Zhu et al. (2015), respectively. Figure 1 shows that our formulae can fit the results of the numerical simulations very well (see the left panel). While it is beyond the scope of this paper, it is interesting that $\alpha_D$ tends to follow $W_{z\phi}$ more than $W_{r\phi}$. Some discussion about the ratio between $W_{r\phi}$ and $\alpha_D$, which is also known as the Schmidt number, can be found in Zhu et al. (2015). Also, our fitting suggests (right panel) that MHD turbulence is significant for the accretion stress ($\alpha_{SS,eff}$) when $\beta_0$ is large (see the turbulence-dominated regime). As $\beta_0$ decreases, disk winds play a more crucial role in disk evolution (see the wind-dominated regime).

Note that care is needed in using these fitting formulae for the regimes of $\beta_0 > 10^3$ and $\beta_0 < 10^5$. For the case of $\beta_0 > 10^3$, our $W_{r\phi}$ approaches the constant value of $\sim 10^{-3}$, and both $W_{z\phi}$ and $\alpha_D$ decrease linearly with increasing $\beta_0$. The behavior of $W_{r\phi}$ is supported by non-ideal MHD simulations from Simon et al. (2013b), which show that the value of $W_{r\phi}$ is of the order of $10^{-3}$ for the case of no vertical magnetic fields...
(equivalently $\beta_0 = \infty$). For $W_{\varphi b}$ and $\alpha_D$, there is currently no simulation to examine these behaviors. In the regime of $\beta_0 < 10^3$, our formulae suggest that $W_{\varphi b}$ saturates to $\approx 10^{-2}$ while both $W_{\varphi b}$ and $\alpha_D$ keep increasing. As with the above case, the results of non-ideal MHD simulation are not available in the literature, so the behavior of these three quantities cannot be examined. (Even ideal MHD simulations are limited for the case of $\beta_0 < 10^3$, e.g., Stone et al. 1996; Bai 2013a.) Likely $W_{\varphi b}$ drops to zero as $\beta_0$ decreases, since strong vertical magnetic fields can fully quench MRI and hence suppress MHD turbulence (e.g., Königl et al. 2010).

Thus, both the accretion stress ($\alpha_{SS,eff}$) and the dust diffusion coefficient ($\alpha_D$) are evaluated simultaneously for a given value of $\beta_0$, yielding reasonable values for $10^5 > \beta_0 > 10^3$. This makes it possible to examine in our framework how the global structure of disks is determined as a function of $\beta_0$ and how different the vertical distribution of dust particles is for different values of $\beta_0$.

2.6. Model Parameters for the Disk around HL Tau

We finally summarize the key properties of the disk around HL Tau. These properties are derived either only from observations or from comparisons between observations and radiative transfer modeling.

We assume that $M_* = 1 M_\odot$ and $M = 1 \times 10^{-7} M_\odot$ yr$^{-1}$ for the HL Tau system. These two values are chosen because they are in the range suggested by observations: $0.5 \lesssim M_*/M_\odot \lesssim 1.3$ (Kwon et al. 2011; ALMA Partnership et al. 2015b) and $0.9 \times 10^{-7} \lesssim M_*/(M_\odot$ yr$^{-1}) \lesssim 2 \times 10^{-6}$ (Hayashi et al. 1993; Beck et al. 2010). For the dust surface density distribution ($\Sigma_d$) and the disk temperature ($T_{mid}$), we follow Kwon et al. (2011). In their model, similarity solutions are adopted for $\Sigma_d$, and a simple power-law profile is used for $T_{mid}$, which are given as

$$\Sigma_d = \rho_{d0} \sqrt{2\pi H_d} \left( \frac{r}{r_c} \right)^{-q} \exp \left[ -\left( \frac{r}{r_c} \right)^{7/2-p-q/2} \right],$$

$$T_{mid} = T_0 (r/r_0)^{-q},$$

where $r_0 = 10$ au, $q = 0.43$, and $T_0 = 62$ K. They find that the observed surface brightness of the disk around HL Tau can be reproduced very well when $\rho_{d0} \approx 2.2 \times 10^{-15}$ g cm$^{-3}$, $r_c = 78.9$ au, and $p \approx 1$, which are confirmed by their latest work (Kwon et al. 2015). Note that the above model parameters are obtained from a comparison with the observations at modest resolution made by CARMA (Kwon et al. 2011, 2015).

The recent studies suggest that, except for gap regions, Kwon’s disk model works well both for ALMA observations with very high spatial resolutions (Akiyama et al. 2016; Pinte et al. 2016) and for JVLA observations with comparable resolutions (Carrasco-González et al. 2016).

Another key result is that the midplane layer rich in millimeter-sized dust particles is surprisingly thin, with a scale height of $H_d \approx 1$ au at $r \approx 100$ au. This follows from the fact that the gaps between the bright rings are clearly visible along the projected minor axis, despite the system’s inclination to our line of sight. The resulting estimate of scale height comes from detailed radiative transfer calculations (Pinte et al. 2016).

In the following, we use the above parameters to compute the disk structure around HL Tau.

3. Results and Applications

We present the results that are derived from the magnetically driven disk accretion model. We also conduct a parameter study to examine how our results are maintained for a wide range of model parameters. Furthermore, we discuss the applications of our results such as the global structure of magnetic fields and the gas-to-dust ratio in the outer part of the HL Tau disk.

3.1. Results from Disk Wind Models

We first present the gas and dust structures coming from the magnetically driven disk accretion model. Substituting the prescriptions for the turbulent and wind stresses in Equation (9) yields the surface density profile ($\Sigma_g$) for a given value of the midplane plasma beta ($\beta_0$) of the net vertical flux. In turn, substituting the prescription for the turbulent diffusion into Equation (15) gives the scale height ($H_d$) of the millimeter-sized particles.

Figure 2 shows the results. The left panel has gas surface density profiles with distance from the star for three values of $\beta_0$, while the right panel has the corresponding scale heights for gas and dust particles of millimeter and centimeter sizes. In our formalism, self-gravity is neglected. Accordingly, the model breaks down in the shaded region, where the Toomre parameter ($Q = c_s \Omega/(\pi G \Sigma_g)$) is less than two, and self-gravity drives spiral waves or outright collapse (e.g., Vorobyov & Basu 2007; Shi & Chiang 2014). Our results show that $\Sigma_g$ becomes an increasing function of $\beta_0$ when the disk accretion rate is given. This occurs because $\alpha_{SS,eff}$ becomes smaller for a larger value of $\beta_0$.

In order to achieve a certain value of $M$, $\Sigma_g$ should increase (see Equation (9)). As $\beta_0$ decreases (i.e., the case of stronger magnetic fields), disk winds can contribute more to the angular momentum transport in the disk. This leads to a lower value of $\Sigma_g$ even if the disk accretion rate is high. We find that a relatively smaller value of $\beta_0 (\lesssim 10^4)$ is needed to prevent gravitational instability (GI) from operating in the disk (also see Lizano & Galli 2015). For $H_d$, our results show that both $H_d(1 \text{ mm})$ and $H_d(1 \text{ cm})$ are smaller than $H_T$. This is a simple reflection of dust settling, where turbulent mixing in the vertical direction is not significant enough for large dust particles to catch up with the vertical distribution of gas (see Equation (15)). We also find that particles of the same size have the roughly same value of $H_d$ even if $\beta_0$ varies. This arises because the dependence on $\beta_0$ cancels out in Equation (15) since $S_T \propto \Sigma_{g0}^{-1} \propto \alpha_{SS,eff} \propto W_{\varphi b} \propto \beta_0^{-1}$, while $\alpha_D \propto \beta_0^{-1}$.

Thus, magnetically driven disk accretion can reproduce the high accretion rate observed for the HL Tau disk when the disk is threaded by a vertical magnetic field with $\beta_0 \lesssim 10^3$. Furthermore, this disk model provides a natural explanation for the observed, vertically thin layer of millimeter-sized dust particles.

3.2. Parameter Study

We now explore a broader parameter space, varying the particle size ($a$) and Stokes number ($St$) alongside the plasma beta. In this parameter study, we compute both $a$ and $St$, which

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5 Recently, Pinte et al. (2016) have found that the HCO$^+$ and CO line emission from the HL Tau disk are consistent with Keplerian motion around a star of $1.7 M_\odot$, although HL Tau is classified as a K-type star.
Figure 2. Modeled gas and dust distributions for three strengths of the net vertical magnetic field, specified using the midplane plasma beta ($\beta_0$). The left, $\Sigma_g$ is shown as a function of $r$. For comparison, the region where GI can operate is denoted by the shaded region. The value of $\Sigma_g$ increases with increasing $\beta_0$. A relatively strong magnetic field ($\beta_0 \lesssim 10^3$) is needed to suppress GI in the disk. On the right, the values of $H_g$, $H_d(1 \text{ mm})$, and $H_d(1 \text{ cm})$ are shown by the dotted, the solid, and the dashed lines, respectively. Dust particles of the same size experience the roughly same degree of dust settling for a wide range of $\beta_0$.

Figure 3. Radius ($a$) of the particles whose scale height ($H_d$) takes various values at $r = 100$ au, as a function of $\beta_0$. Each scale height is a different solid line. Vertical dashed lines denote the values of $\beta_0$ at which the Toomre $Q$ parameter is 1, 2, and 10. For the left panel, the accretion stress comes from turbulence alone ($\alpha_{SS}$, Equation (7)), while on the right, both turbulent and wind stresses are included ($\alpha_{SS,eff}$, Equation (11)). Because disk winds transport significant angular momentum when the magnetic fields are strong (i.e., when $\beta_0$ is smaller), the model in the right panel meets the constraints $H_d \lesssim 1$ au and $Q \gtrsim 2$ over a bigger region of the parameter space.

are needed to give a specified particle scale height ($H_d$, see Equation (15)), as functions of $\beta_0$ at $r = 100$ au. These calculations are carried out to investigate how the above results (see Figure 2) can be maintained for a wide range of $\beta_0$. In order to explicitly calibrate the importance of disk winds, we consider two cases here: $\alpha_{SS}$ is used in computing $a$ and $St$ in the first case (see Equation (7)), and $\alpha_{SS,eff}$ is adopted in the other case (see Equation (11)).

Figure 3 shows the results for $a$ with $\alpha_{SS}$ and $\alpha_{SS,eff}$ in the left and the right panels, respectively. The solid lines denote the resulting values of $a$ for different dust heights ($H_d = 0.3$ au, 1 au, and 3 au at $r = 100$ au), and the vertical dashed lines are for the values of $\beta_0$ at which the corresponding Toomre $Q$ parameter becomes 1, 2, and 10 at $r = 100$ au, respectively (also see Figure 1). Our results show that a parameter space that satisfies both $H_d \lesssim 1$ au and $Q \gtrsim 2$ at $r = 100$ au can expand when disk winds are taken into account (see the shaded regions in both panels). This is simply because only when MHD turbulence is considered is a lower value of $\beta_0$ required to achieve $\dot{M} = 10^{-7} \ M_\odot \ yr^{-1}$. This in turn pumps up the diffusion coefficient ($\alpha_{DS}$, see Figure 1). Then, a higher degree of dust settling can be established only for larger ($a \gtrsim 1$ cm) dust particles. Consequently, the shaded region shrinks in terms of both $a$ and $\beta_0$ (see the left panel). When disk winds are treated properly in the accretion stress ($\alpha_{SS,eff}$), however, a modest value of $\beta_0$ is good enough to satisfy both a higher value of $\dot{M}$ and a lower value of $H_d$. This is attributed to the fact that disk winds can remove the disk angular momentum considerably. As a result, when $\alpha_{SS,eff}$ is adopted, a parameter space becomes bigger, where both $H_d \lesssim 1$ au and $Q \gtrsim 2$ at $r = 100$ au are fulfilled simultaneously.

While Figure 3 shows that $a$ is relatively insensitive to the change in $\beta_0$ in the shaded region for the case of $\alpha_{SS,eff}$, it is interesting to examine whether a much tighter constraint can be obtained on the value of $\beta_0$. For this purpose, we also plot the resulting value of $St$ as a function of $\beta_0$ (see Figure 4). This is motivated by the calculations of dust growth, which show that when dust growth is limited by radial drift, the maximum size...
of dust particles is peaked around \( \alpha = 0.1 \) (e.g., Brauer et al. 2008; Birnstiel et al. 2010; Okuzumi et al. 2012). Since radial drift is one of the severest barriers to dust growth in protoplanetary disks, it is important to verify whether or not our results can satisfy the condition \( \alpha \approx 0.1 \), which may serve as an additional constraint.

Figure 4 shows the results. We first point out that the behavior of \( \alpha \) is the same with and without disk winds. This is simply because \( \alpha \) is a function of only \( \beta_0 \) when \( H_\alpha \) and \( H_g \) are given (Equation (15)). We then see that all three conditions \( (H_\alpha \lesssim 1 \text{ au} \text{ and } Q \gtrsim 2 \text{ at } r = 100, \text{ and } \alpha \approx 0.1) \) are realized simultaneously when disk winds are properly included (right panel). We therefore suggest that disk winds are necessary to account for the configuration of the HL Tau disk. In fact, we find that the wind stress contributes to the value of \( \alpha \) by more than 50\% across the entire region of the disk in our model (see Figure 5). It is important to emphasize, however, that only a small region of parameter space meets all the conditions, even for the case \( \alpha \) (see Figure 4).

Thus, our results suggest that in order to reproduce the disk configuration observed toward HL Tau, the presence of disk winds is inevitable and the initial vertical magnetic field should be relatively strong \( (\beta_0 \approx 2 \times 10^5) \).

### 3.3. The Global Structure of Vertical Magnetic Fields

The above results allow us to obtain the best value of \( \beta_0 \). In the following, we will make use of this value to investigate the resulting disk properties such as vertical magnetic fields and the gas-to-dust ratio. In this section, we compute the global structure of the initial, vertical magnetic field \( B_{z,0} \). This becomes possible because \( B_{z,0} \) can be derived from the resulting \( \Sigma_g \) and \( \beta_0 \) for a given value of \( c_s \) (see Equation (16)).

Figure 6 shows the profile of \( B_{z,0} \) for our best case \( (\beta_0 = 2 \times 10^5) \). We find that the profile is characterized by \( B_{z,0} \propto r^{-1.25} \) and can be reproduced by the following calculation: as discussed above (see Figure 5), disk winds play the dominant role in transporting the disk angular momentum. Then, it can be assumed that \( \alpha \approx (r/H_g)W_{\alpha,0} \approx (r/H_g)\beta_0^{-1} \) (see Equations (11) and (18)). As a result, the disk accretion rate can be given as (see Equation (9))

\[
\dot{M} \sim \frac{\Sigma_g c_s^2}{\Omega} \frac{r}{H_g \beta_0} \sim \frac{r B_{z,0}^2}{\Omega}.
\]

Since steady disk accretion is currently considered \( (\dot{M} = \text{const.}) \), the resulting magnetic field profile becomes \( B_{z,0} \propto r^{-1.25} \).

This derivation is interesting in the sense that the \( \Sigma_g \) dependence cancels out when computing the disk accretion rate (see Equation (22)). In other words, the disk accretion rate is determined mostly by the initial \( B_{z,0} \). This finding is consistent with the previous results of ideal MHD simulations (Hawley...
et al. 1995) and non-ideal MHD simulations (Okuzumi & Hirose 2011; Bai 2013). Furthermore, our profile itself is comparable with the previous studies (Okuzumi et al. 2014); for instance, Okuzumi & Hirose (2011) perform resistive MHD simulations and find that \( B_{z,0} \propto r^{-1.24} \), while Bai (2013) performs resistive MHD simulations with ambipolar diffusion and finds that \( B_{z,0} \propto r^{-1.44} \) (see Bai 2014 for the correct formula).

The absolute magnitude of \( B_{z,0} \) is more uncertain. This is because the value of \( B_{z,0} \) should be controlled by the balance between the radial dragging and diffusion of magnetic fields (e.g., Lubow et al. 1994; Guilet & Ogilvie 2014; Okuzumi et al. 2014; Takeuchi & Okuzumi 2014). Here, we compare the above results with those obtained by Okuzumi et al. (2014). In their study, the upper limit of large-scale magnetic fields (\( B_{z,\text{max}} \)) is derived by considering the steady state, where radial inward advection of magnetic fields balances their outward diffusion. Mathematically, \( B_{z,\text{max}} \) can be given as (see their equation (48))

\[
B_{z,\text{max}} = 100 \left( \frac{r}{\text{au}} \right)^2 \left( \frac{r_{\text{out}}}{100 \text{au}} \right)^2 \left( \frac{B_{\infty}}{10 \mu G} \right) \text{mG},
\]

where \( B_{\infty} \) is the vertical magnetic field at \( r = r_{\text{out}} \). Figure 6 summarizes our comparison. We find that our \( B_{z,0} \) is much larger than the value of \( B_{z,\text{max}} \) in the outer part of the disk. Since the absolute magnitude of \( B_{z,\text{max}} \) is scaled by the magnetic field strength at the outer edge of disks, this difference arises either because the presence of the envelope may trigger more efficient, inward advection of magnetic fields or because disks in the stage of star formation may intrinsically possess a stronger field than those in the stage of planet formation. Note that Flock et al. (2015) perform global, non-ideal MHD simulations to model the detailed structure of protoplanetary disks. In these simulations, they adopt an initial \( B_z \) that is similar to our \( B_{z,0} \), and find that some observable structures such as gaps and rings are generated. It is obvious that more detailed calculations would be needed to fully address this point.

Thus, a coupling of our model with the ALMA observations of HL Tau has made it possible to discuss the global structure of magnetic fields threading disks in a self-consistent fashion.

3.4. The Gas and Dust Distributions in the Outer Part of the Disk

In this section, we consider the gas and dust distributions in the HL Tau disk.

To proceed, we compare our results with those of Kwon et al. (2011). As already described in Section 2.6, the surface density distribution of dust (\( \Sigma_d \)) for the HL Tau disk can be represented well by similarity solutions, except for the gap regions (see Equation (20), Kwon et al. 2015; Akiyama et al. 2016; Pinte et al. 2016). It should be noted that the dust thermal emissions observed by both ALMA and CARMA are optically thin only in the outer part of the HL Tau disk (ALMA Partnership et al. 2015b; Carrasco-González et al. 2016). Accordingly, we focus here on the surface density profile at \( r > 40 \text{au} \).

Figure 7 shows our results for the best case (\( \beta_3 = 2 \times 10^4 \), the red solid line), \( 10^2 \times \Sigma_d \) (the black, solid line), and \( \Sigma_d \) (the dashed line). The dust gaps observed by ALMA are also denoted by the dark, shaded regions. Notice that if the gas-to-dust ratio is 100, the measurements of Kwon et al. imply marginal gravitational instability at 60–80 au around the gaps (see \( 10^2 \times \Sigma_d \)). Our model avoids GI because of a slightly lower surface density. In other words, the gas-to-dust ratio may be lower than 100 there. In contrast, our model suggests that the gas-to-dust ratio may be higher than 100 beyond \( r = 100 \text{au} \) since \( \Sigma_g > 10^2 \times \Sigma_d \). Our results therefore imply that the gas-to-dust ratio may increase from the intermediate region of disks to the outer one. It is interesting that this
implication is indeed consistent with the trend observed in protoplanetary disks: larger dust particles tend to concentrate toward the central star there (e.g., Williams & Cieza 2011). Based on detailed radiative transfer modeling, Pinte et al. (2016) have also suggested that larger dust particles may be depleted at $r \approx 100$ au in the HL Tau disk, possibly due to either (efficient) radial drift of such particles or their slower growth.

Thus, the results of our modeling are useful in providing some implications for the gas-to-dust ratio in the HL Tau disk.

4. Discussion

We have so far focused on two of the key properties of the HL Tau disk: the high disk accretion rate ($\sim 10^{-7} M_{\odot}$ yr$^{-1}$) onto the central star and the vertical thin ($\sim 1$ au) dusty layer at $r = 100$ au. We have demonstrated above that magnetically driven disk accretion can reproduce these features naturally when disk winds are properly taken into account and the initial plasma $\beta$ is $\beta_0 \approx 2 \times 10^5$. It is important, however, that a number of other interesting features have been reported for the HL Tau disk. Here we examine how our results are consistent with these features. In addition, we discuss limitations of our model and other possible mechanisms of disk evolution for the HL Tau system.

4.1. Other Observational Features

In this section, we consider the observations of winds/outflow from the HL Tau disk, a polarization map for the disk, and the molecular line emission, and discuss how our disk model can account for these observational results.

It is well recognized that the HL Tau system exhibits large-scale ($\sim 10^4$ au) bipolar outflows (Monin et al. 1996). Observations trace the inner jet down to $\sim 100$ au from the disk (e.g., Pyo et al. 2006; Beck et al. 2010), and the ALMA observations with modest spatial resolution ($\sim 30$ au) are examined for characterizing the outward motion of molecular gas (e.g., Klaassen et al. 2016). Despite such efforts, identification of the launching points of disk winds/outflow has still not yet been done observationally for the HL Tau disk. Recently, Bjerkeli et al. (2016) have conducted ALMA observations toward TMC1A with high spatial resolution ($\sim 6$ au) and inferred that molecular ($^{13}$CO) outflows are very likely to originate from the Keplerian disk for this target. Interferometric observations with similar resolutions are demanded for the HL Tau system to specify the footpoint radius of outflows and to verify the importance of disk winds.

Particles aligned in a magnetic field can produce polarized thermal continuum emission at submillimeter wavelengths (e.g., Bertrang & Wolf 2017). Mapping polarization vectors across the HL Tau disk with a modest spatial resolution of about 80 au, Stephens et al. (2014) inferred that the field’s strongest component is toroidal. A strong toroidal field was also found in synthetic maps of dust thermal polarization that are generated from global 3D MHD simulations (Bertrang et al. 2017). In contrast, Matsakos et al. (2016) found that MHD simulations with a large radial component better reproduced the observed polarization map for the HL Tau disk. The presence of magnetic fields in the radial direction would be viewed as evidence that magnetically induced disk winds are launched. Therefore, the polarization observation suggests that disk winds are very likely to play an important role in the HL Tau disk.

Note that the simulations of Matsakos et al. (2016) are carried out in the limit of ideal MHD and adopt a very low value of $\beta_0 (\sim 2.5)$. On the other hand, our results are based on non-ideal MHD simulations and find that $\beta_0 \sim 10^4$ for the base case. We thus cannot compare our results with those of Matsakos et al. (2016) directly. It is nonetheless interesting that both studies lead to the conclusion that disk winds would be crucial for the HL Tau disk.

Furthermore, Matsakos et al. (2016) have provided the following implication: given that the radial component becomes important only at the disk surface, small ($\lesssim 0.1$ mm) grains that are located in the upper layer of disks should contribute significantly to the observed polarized emission. This means that the maximum size of dust particles in the disk midplane may be too large to contribute to the polarized emission. Taking into account that the polarized emission at (sub)millimeter wavelengths can be generated by dust particles with a size of $\lesssim 0.1$ mm (e.g., Cho & Lazarian 2007), our results become compatible with their implication: the majority of dust in the midplane may be larger than 1 mm in size (see Figure 3). It should be pointed out that self-scattering of dust particles can also provide a reasonable explanation for the observed polarization map without the presence of magnetic fields and dust alignment (Kataoka et al. 2016; Yang et al. 2016; Tazaki et al. 2017).

We have so far discussed the observational features that can be used as a probe to detect magnetically induced disk winds. This is because disk winds are needed as an alternative process to transport angular momentum when MHD turbulence is suppressed in the outer part of the HL Tau disk. It would also be useful to directly infer the level of turbulence there, using other observables. Observations of the molecular line emission can play such a role. In general, it is very difficult to provide some meaningful constraints on the level of disk turbulence by detecting the molecular line emission. This is because of the need for spatially and spectrally resolved measurements to calibrate the width of nonthermal broadening, which can be less than a few per cent of the width caused by the thermal broadening. Recent advances in interferometric observations at (sub)millimeter wavelengths have enabled such observational studies to be conducted (e.g., Hughes et al. 2011; Guilloteau et al. 2012; Flaherty et al. 2015). For instance, Flaherty et al. (2017b) use the ALMA data taken toward HD 163296, and derive stringent upper limits on the nonthermal motion for the disk. They infer that the line width generated by turbulence is $\sim 5\%$ of the local sound speed. It is interesting that when MRI and the resulting MHD turbulence operate fully in the disk, these limits cannot be reproduced. In other words, turbulence should be weak in the outer part of the disk, which may be caused by non-ideal MHD effects. It is important to recognize, however, that the quality of the data is still not be high enough to fully address the strength of disk turbulence (e.g., Teague et al. 2016). In fact, the upper limit derived by Flaherty et al. (2017b) is still looser than the prediction of Pinte et al. (2016) that the velocity arising from turbulence is $\sim 1\%$ of the sound speed. More observations are demanded to tightly infer the level of turbulence for protoplanetary disks including the HL Tau system.

Thus, while the spatially and spectrally resolved observational data are needed to reliably examine the validity of magnetically driven disk accretion models for the HL Tau system, qualitative assessments suggest that a number of
important observational features can be explained by our model.

4.2. Limitations of Our Model

As discussed above, the observed properties of the HL Tau disk can be better understood by adopting magnetically driven disk accretion models. We must admit, however, that our disk model has been developed with a number of assumptions. Here, we discuss how these assumptions can affect our finding and what the limitations of our model would be.

First, our model relies heavily on the results of non-ideal MHD simulations. Specifically, the spatially integrated, temporally averaged values of $W_{r0}$, $W_{z0}$, and $\alpha_D$ are utilized to obtain the fitting formulae (see Figure 1). This approach can end up with either an overestimate or an underestimate of $H_d$ and $\Sigma_g$, depending on the spatial integration. For instance, the midplane value of $\alpha_D$ can be overestimated by integrating it up to a large value of the vertical height, which can puff up dusty layers in the vertical direction. On the other hand, the values of $W_{r0}$ and $W_{z0}$ at disk surfaces can be underestimated by smearing them out over the intermediate height, which can lead to a higher value of $\Sigma_g$. Numerical simulations are therefore needed to fully confirm the results of our semi-analytical model.

Second, our model does not include the mass loss from disk surfaces that can be triggered by disk winds. It is interesting that both numerical simulations and observations suggest that some degree of mass loss occurs from circumstellar disks. In fact, Klaassen et al. (2016) use the ALMA data on CO emission line and estimate that the mass loss rate caused by outflow from HL Tau can be as high as the disk accretion rate. Theoretically, both ideal and non-ideal MHD simulations confirm the mass loss due to disk winds (e.g., Suzuki & Inutsuka 2009; Bai 2013b; Simon et al. 2013a). Since the mass loss rate can be computed as the product of the gas density and sound speed at the launching point, the resulting value is quite sensitive to the location of the launching point. In other words, it is still hard to derive a reliable value of the mass loss rate both observationally and theoretically. Nonetheless, our models are very likely to overestimate the resulting value of $\Sigma_g$ due to neglect of the mass loss.

Third, we have assumed that the HL Tau disk is in a steady state (i.e., $M = \text{const.}$). While this assumption may be reasonable in the current calculations, it is intriguing to examine how the above results would be altered if this assumption is relaxed. This is motivated by the observation that some YSOs undergo so-called episodic accretion (see below, Hartmann et al. 2016). To proceed, we recompute the radius ($a$) of dust particles and the corresponding Stokes number ($St$) as was done in Section 3. We here reduce the value of $M$ from $10^{-7} M_\odot$ yr$^{-1}$ to $5 \times 10^{-8} M_\odot$ yr$^{-1}$. Note that when $M$ is increased by a factor of 2, even disk wind models cannot satisfy all the requirements simultaneously. Figures 8 and 9 show the results: the former is for $a$ and the latter is for $St$. In the left panels, the results are obtained taking into account only turbulence, while turbulence and disk winds are both included in the right panels. We find that if the accretion rate at the outer disk is lower than that at the inner disk by a factor of at least 2, then both the disk models with and without the effect of disk winds can satisfy all the requirements. Figures 8 and 9 show that the reduction in $M$ considerably expands the parameter space in which disk wind models can reproduce the configuration of the HL Tau disk (see the right panels). Based on this result, one may wonder whether disk winds may not be a necessary ingredient if the HL Tau disk is currently in nonsteady accretion. As discussed below, however, this is unlikely because the HL Tau configuration does not fit the basic properties of episodic accretion. Thus, whereas the assumption of steady-state accretion may not be fully valid, our finding that disk winds are important can still hold.

In summary, it is obviously important to perform more detailed, numerical simulations to examine the validity of our results. Nonetheless, it can be concluded that our main results may be maintained even if some assumptions are relaxed.

4.3. Other Mechanisms of Disk Accretion

We finally discuss how other mechanisms could serve as a plausible process to account for the key features of the HL Tau disk. Here we examine episodic accretion and GI.

We first explore the possibility of episodic accretion. It is well known that some (perhaps most) YSOs experience
accretion outbursts during the class I and II stages (e.g., Hartmann et al. 2016). There are two famous types (e.g., Audard et al. 2014): one is the so-called FUor that is named from FU Ori (Herbig 1966, 1977), and the other is the EXor coined from EX Lupi (Herbig 1989). Both objects exhibit accretion outbursts, but the magnitude and the duration of the outbursts of FUors are much larger than those of EXors. It is still uncertain whether these two objects can be regarded as two distinct types and/or what would be the ultimate mechanism to trigger such outburst events. It is clear, however, that the disk mass should be piled up during the quiescent phase, and such masses are accreted onto the central star when an outburst occurs. Since FUors generally have accretion rates higher than $10^{-5} \, M_\odot \, yr^{-1}$ during outbursts, HL Tau may be better classified as an EXor if it is in the outburst phase. Note that outbursts can increase the accretion rate up to $10^{-3} - 10^{-6} \, M_\odot \, yr^{-1}$ for EXors. Also, a quiescent phase is unlikely to be applicable for the current HL Tau system. This is because, if this were the case, $M$ in the outer disk should be larger than that in the inner disk ($\approx 10^{-7} \, M_\odot \, yr^{-1}$), so that a pile-up of mass is established in the inner disk. We find that when $M > 10^{-7} \, M_\odot \, yr^{-1}$ in the outer disk, GI would be the main driver of disk accretion there (see Section 4.2). As discussed below, however, GI is unlikely to reproduce the observed features of the HL Tau disk.

We now consider how the current configuration of the HL Tau system can be (in)consistent with the basic properties of EXors. We begin by pointing out that EXor-like outbursts tend to occur on stars in the later stages of formation, when the circumstellar envelope is already dispersed (Sipos et al. 2009; Audard et al. 2014). The disappearance of envelopes is obviously inconsistent with the HL Tau case. Then we compute how much the disk mass can be piled up during the quiescent phase, and how long an outburst event can last. This calculation is motivated by the above results that if the accretion rate is $\lesssim 5 \times 10^{-8} \, M_\odot \, yr^{-1}$ in the outer disk, the MHD turbulence model can also reproduce the key features of the HL Tau disk (see the left panel of Figure 9). Utilizing this value with the assumption that the accretion rate in the outer disk is comparable between the outburst and quiescent phases, the mass ($\Delta m$) that can be accumulated during a quiescent phase ($\Delta t$) is written as $\Delta m = M_{\text{acc}} \Delta t$. The observational results suggest that the duration of outbursts and the interval between them are about one year and 10 years for EXors, respectively (Hartmann & Kenyon 1996; Audard et al. 2014). As a result, $\Delta m \lesssim 5 \times 10^{-7} \, M_\odot (= 5 \times 10^{-8} \, M_\odot \, yr^{-1} \times 10 \, yr)$. Furthermore, given that the current accretion rate is about $10^{-7} \, M_\odot \, yr^{-1}$, we can estimate that an outburst should continue for $\lesssim 5 \, yr$, if HL Tau is currently in an outburst phase. It is interesting that these numbers seem roughly consistent with the general properties of EXors. This means that the turbulence model might be able to account for the observed properties of the HL Tau disk, if HL Tau were an outbursting EXor. Finally, if HL Tau were an EXor, its outflows would show knots with spacing corresponding to the interval between outbursts. However, no such features are reported for the molecular outflow (Klaassen et al. 2016). Thus, it seems unlikely that HL Tau is an outbursting system.

Next, we discuss the possibility of GI. As already shown in Figure 7, a high value of the gas-to-dust ratio is needed for GI to operate in the HL Tau disk and to reproduce the observed high disk accretion rate. In addition, it is difficult to achieve a thin dusty layer in self-gravitating disks. This is because the rms vertical turbulent speed is comparable to the sound speed in such disks (Shi & Chiang 2014). Then the corresponding value of $\alpha_D$ becomes of the order of unity (see Equation (14)), which suggests that $H_d \approx H_g$ (also see Booth & Clarke 2016). Based on the above consideration, GI is unlikely to be currently the main driver of disk accretion for the HL Tau disk.

5. Conclusions and Summary

We have investigated how magnetically driven disk accretion could possibly account for a number of key features of the HL Tau disk. In particular, we have focused on the recent ALMA observations, which suggest a high degree of dust settling in the outer part of the disk. This observational finding is interesting, because in the framework of the standard 1D viscous disk model, such a configuration is not necessarily compatible with other observational results that suggest that the accretion rate onto the central star is high.

We have developed a simplified, but physically motivated disk model in which the recent non-ideal MHD simulations are utilized. Our model can thereby examine how the disk angular momentum is transported both radially via MHD turbulence and vertically via magnetically induced disk winds for given
values of the initial plasma beta ($\beta_0$) and of the sound speed ($c_s$). We find that these two features (a high disk accretion rate and a high degree of disk settling) can be reproduced well when magnetically induced disk winds are properly taken into account (see Figure 3). This becomes possible because disk winds can transport a significant amount of the disk angular momentum, so that a high disk accretion rate is achieved without a higher level of turbulence (see Figure 5). This naturally leads to efficient disk settling. We have also discussed that the framework of magnetically driven disk accretion would be useful for investigating other features of the HL Tau disk such as the global magnetic field configuration (see Figure 6) and the gas-to-dust ratio (see Figure 7).

Our results, however, show that the optimal configuration of the HL Tau disk is realized only in a very narrow parameter space under the assumption of steady accretion ($\beta_0 \approx 2 \times 10^2$, see Figure 4). In addition, our model relies heavily on the results of numerical simulations in which both the accretion and the wind stresses are vertically integrated (see Figure 1). As other limitations of our model, we have discussed neglect of mass loss due to disk winds and the assumption of steady-state accretion. Observations with higher spatial resolution for the disk gas and more self-consistent modeling would be required to fully identify the main driver of disk accretion for the HL Tau disk. For instance, a coupling of dust growth and non-ideal MHD effects would be important to further explore the origin of the observed gaps in both the gas and dust distributions for the disk. It is also worth noting that recent non-ideal MHD simulations can generate gap structures in gas disks (e.g., Moll 2012; Flock et al. 2015; Béthune et al. 2016, 2017; Suriano et al. 2017). While it is beyond the scope of this paper to identify the origin of dust gaps in the HL Tau disk, these results are promising for magnetically driven disk accretion models.

Thus, understanding magnetically driven disk accretion will be the fundamental step in drawing a better picture of disk evolution and the subsequent planet formation in the disk, and the HL Tau disk is one of the best examples in which such investigations will be performed.

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