OBSEVATIONAL LIMITS ON THE SPINDOWN TORQUE OF ACCRETION POWERED STELLAR WINDS

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1. INTRODUCTION

Classical T Tauri stars (CTTS) are known to be magnetically active protostars showing clear observational signatures of accretion (circumstellar disks) and ejection (jets and outflows). The stellar magnetic fields measured by spectropolarimetric observations (up to a few kG; Johns-Krull 2007) can deeply affect the dynamics of the circumstellar region: truncating the disk and channeling the accretion flow along the magnetic surfaces down to the stellar surface; providing an acceleration mechanism for different types of outflows, stellar winds along the opened magnetospheric field lines (Matt & Pudritz 2008a) and ejections associated with the magnetic star–disk interaction (Shu et al. 1994; Ferreira et al. 2000; Romanova et al. 2009; Zanni 2009).

Their spin represents a controversial issue. While a wide range of rotation periods is observed among low-mass young stars (0.2 up to 20 days; Irwin & Bouvier 2009), around half of them slowly rotate, much below the breakup limit. Besides, many slow rotators show clear accretion signatures (Herbst et al. 2007), as in the case of CTTS, which have an average rotation period around ~8 days, corresponding to ~10% of their breakup speed (Bouvier et al. 1993). Since CTTS are accreting mass and angular momentum from the surrounding accretion disk and they are still contracting, they would be expected to noticeably spin up in a few million years: conversely, there are indications that solar-mass slow rotators are prone to keep their rotation period constant for ~5 Myr (Irwin & Bouvier 2009). Therefore, some mechanism must act to efficiently remove angular momentum from these slowly rotating stars.

Grounded on models originally developed to explain the period changes of pulsars (Ghosh & Lamb 1979), one of the most widespread scenarios foresees that a significant spin-down torque is provided along the magnetospheric field lines connecting the star and the disk region rotating slower than the star (Königl 1991). On the other hand, both analytical (Matt & Pudritz 2005a) and numerical models (Zanni & Ferreira 2009) have questioned the efficiency of this mechanism, due to the limited extent of the connected region and the weakness of the magnetic connection.

Matt & Pudritz (2005b) have therefore proposed that stellar winds could efficiently remove angular momentum directly from the star along the opened field lines of the magnetosphere. Besides, they suggested that these outflows could derive their energy directly from the accretion power. This would be also consistent with the fact that accreting CTTS seem to have on average longer rotation periods than their non-accreting counterparts (weak-lined T Tauri stars, WTTS), indicating a connection between spin-down and accretion (Lamm et al. 2005). It is commonly assumed that the accretion power is liberated in a shock due to the impact of the accretion streams with the stellar surface. While a fraction of the accretion energy can be converted (e.g., into Alfvén waves, Scheurwater & Kuijpers 1988) and possibly injected into the wind, the emission of the shocked material can explain the observed optical excess and UV continuum (Gullbring et al. 2000). Observations of the accretion shock luminosity can be therefore used to constrain the accretion energy which is available to power the stellar wind.

In this Letter, we try to estimate the spin-down efficiency and the energy requirement of accretion powered stellar winds (APSWs) compatible with measurements of magnetic fields and accretion luminosities of several CTTS. In Section 2, we describe a simple analytical APSW model and we apply it to a specific CTTS example in Section 3. We determine the stellar parameters which are compatible with a spin equilibrium situation in Section 4 and we summarize our conclusions in Section 5.
2. THE MODEL

In our analysis we will assume that, even in the case of a multipolar, complex stellar field, the dipolar component controls the dynamics of both accretion (i.e., the disk truncation radius) and ejection (i.e., the wind magnetic lever arm). This assumption is somewhat supported by the results of Matt & Pudritz (2008a): in the case of a multipolar field (quadrupolar in their case) with no dipolar component, the stellar wind torque is strongly suppressed. The same assumption is done evaluating the disk truncation radius: the dipolar component is the one which can be defined as $\Omega_\star \sqrt{G M_\star / R_\star}$, where $\Omega_\star = 2 \pi / P_\star$ is the angular speed of the star, $P_\star$ is its rotation period, and $R_\star$ is the wind average magnetic lever arm. For $R_\star$ we assume the Matt & Pudritz (2008a) approximation:

$$R_\star = K \left( \frac{B_\star^2 R_\star^2}{M_\text{wind} v_{\text{esc}}^2} \right)^m R_\star,$$  

(1)

where $K = 2.11$ and $m = 0.223$, obtained for a stellar wind flowing along the open field lines of a dipolar magnetosphere. In Equation (1), $M_\star$ and $R_\star$ are the stellar mass and radius, respectively, $B_\star$ is the intensity of the dipolar component of the magnetosphere at the stellar equator, and $v_{\text{esc}} = \sqrt{2 G M_\star / R_\star}$ is the escape speed. Assuming a Keplerian disk rotation, we write the spin-up accretion torque as $J_{\text{acc}} = M_{\text{acc}} \sqrt{G M_\star R_\star}$, where $M_{\text{acc}}$ is the disk accretion rate and $R_\star$ is the disk truncation radius. We assume that the radius at which the disk gets truncated by the dipolar component of the magnetosphere is proportional to the Alfvén radius: $R_t = C \Psi^{2/7} R_\star$, where $\Psi = B_\star^2 R_\star^2 / M_{\text{acc}} v_{\text{esc}}^2$ is a dimensionless magnetization parameter. Different theoretical works limit the multiplying factor in the range $C \sim 0.5$ (Bessolaz et al. 2008); we assume an average value $C = 0.75$. Combining all the previous expressions, the wind outflow rate necessary to extract a fraction $f_1 = J_{\text{wind}} / J_{\text{acc}}$ of the accretion torque is

$$M_{\text{wind}} = M_{\text{acc}} \frac{f_1 C^{1/2} \Psi^{1/7} - 2 m}{K^2 \delta} \frac{1}{\sqrt{\delta}},$$  

(2)

A value $f_1 = 1$ corresponds to a spin equilibrium situation, while $\delta = 1$ represents a star rotating at breakup speed. Employing Equation (2), we define the mass ejection efficiency as $f_\text{M} = \dot{M}_{\text{wind}} / M_{\text{acc}}$.

How much energy is necessary to drive such a stellar wind? Since CTTS rotate well below their breakup speed ($\delta \ll 1$), magneto-centrifugal processes are not efficient enough to accelerate the wind at the stellar surface and an extra energy input is required to give the wind the initial drive. This energy input corresponds roughly to a specific energy of the order of the potential gravitational energy: $E_{\text{wind}} = 1/2 M_{\text{wind}} v_{\text{esc}}^2$. In the APSW scenario, it is supposed that the wind gets the driving power from the energy deposited by accretion onto the stellar surface. The total available accretion power can be defined as (Matt & Pudritz 2005b)

$$E_{\text{acc}} = 1/2 M_{\text{acc}} v_{\text{esc}}^2 \left[ 1 - \frac{R_t}{2 R_1} - \delta \left( \frac{R_t}{R_1} \right)^{1/2} \right],$$  

(3)

given by the sum of the change in potential energy and kinetic energy minus the work done by the accretion torque. Terms proportional to $\delta^2$ have been neglected. If we assume that the stellar wind consumes a fraction $f_\text{E} = E_{\text{wind}} / E_{\text{acc}}$ of the accretion power, the remaining accretion shock luminosity $L_{\text{UV}} = E_{\text{acc}} - E_{\text{wind}}$ is given by

$$L_{\text{UV}} = 1/2 M_{\text{acc}} v_{\text{esc}}^2 \left[ 1 - 0.5 C^{-1} \Psi^{-2/7} - \delta C^{1/2} \Psi^{1/7} \right] \left( \frac{f_\text{E} C^{1/2} \Psi^{1/7} - 2 m}{K^2 \delta} \right) \frac{1}{\sqrt{\delta}}.$$  

(4)

The accretion luminosity is usually employed to estimate the accretion rate assuming that all the accretion power is radiated at the accretion shock. This translates into a simplified formula like $M_{\text{obs}} = k L_{\text{UV}} R_\star / G M_\star$, with $k$ a numerical factor of order unity (e.g., $k = 1.25$; Gullbring et al. 1998), reflecting the uncertainty on the position of the truncation radius. This determines a possible discrepancy between the observed accretion rate $M_{\text{obs}}$ and the real one $M_{\text{acc}}$.

We try now to apply this simple model to some CTTS observations. We selected a sample of stars (see Table 1) for which spectropolarimetric observations are available, so that a magnetic topology reconstruction is possible: V2129 Oph (Donati et al. 2007, 2010b), BP Tau (Donati et al. 2008), CV Cha, CR Cha (Hussain et al. 2009), and AA Tau (Donati et al. 2010a). In the case of V2129 Oph, BP Tau, and AA Tau the

### Table 1

| Object Name | $M_\star (M_\odot)$ | $R_\star (R_\odot)$ | $B_\star$ (G) | $P_\star$ (days) | $\Omega_\star$ | $L_{\text{UV}} (L_\odot)$ | $M_{\text{obs}} (M_\odot \text{yr}^{-1})$ | Ref. |
|-------------|-------------------|-------------------|--------------|-----------------|--------------|-----------------|---------------------------------|-----|
| BP Tau      | (a) 0.7           | 1.95              | 600          | 7.6             | 0.05         | 0.179           | $2 \times 10^{-8}$               | 1.3 |
|             | (b) 0.023         | 2.5              | 900          | 6.5             | 0.06         | 0.143           | $2 \times 10^{-9}$               | 3.5 |
| V2129 Oph   | (a) 1.35          | 2.4               | 175          | 6.5             | 0.05         | 0.01            | $6.3 \times 10^{-10}$            | 6   |
|             | (b) 0.014         | 2.5              | 450          | 4.4             | 0.07         | 0.01            | $3 \times 10^{-8}$               | 4   |
| CV Cha      | (a) 2.0           | 2.5               | 300          | 4.4             | 0.14         | 0.02            | $10^{-9}$                        | 4   |
| CR Cha      | (a) 1.9           | 2.5               | 200          | 2.3             | 0.05         | 0.023           | $2.8 \times 10^{-9}$             | 1.5 |
| AA Tau      | (a) 0.7           | 2                 | 1500         | 8.2             | 0.05         | 0.005           | $6.3 \times 10^{-10}$            | 5   |

References. (1) Gullbring et al. 1998; (2) Donati et al. 2007; (3) Donati et al. 2008; (4) Hussain et al. 2009; (5) Donati et al. 2010a; (6) Donati et al. 2010b.
measurements of the dipolar component are available; in the case of CV Cha and CR Cha, we take maximum observed magnetic field as the maximum of the dipolar component, since the authors did not provide the intensity of the multipole. In Table 1, we report the intensity of the dipolar field at the stellar equator, which corresponds to half of its maximum, located at the pole. For V2129 Oph, BP Tau, and AA Tau two values of the accretion luminosity are provided: Donati and collaborators have recently re-estimated this quantity, giving luminosities which are systematically around one order of magnitude smaller than previous estimates (see, e.g., Gullbring et al. 1998, for BP Tau and AA Tau). We also take into account a Taurus–Auriga sample from Johns-Krull (2007): for these stars, the actual intensity of their dipolar component is unknown.

3. AN ILLUSTRATIVE CASE: BP Tau

Taking as an example the specific case of BP Tau, we assume the following values: \( M_* = 0.7 \, M_\odot \), \( R_* = 1.95 \, R_\odot \), \( B = 600 \, G \), \( P_* = 7.6 \) days, \( \delta = 0.05 \) (Donati et al. 2008). Fixing these quantities, the accretion luminosity \( L_{\text{UV}} \) depends only on the parameters \( M_{\text{acc}} \) and \( f_1 \). In Figure 1 (left panel), we plot the accretion luminosity as a function of the “true” accretion rate (not the “observed” one) for three different values of \( f_1 \) (Equation (4)). The difference between these curves and the available accretion power (dashed line) gives the power consumed to drive the stellar wind: since the energy requirement of the wind \( f_E \) increases with the accretion rate, it is possible to find a maximum accretion luminosity consistent with a given \( f_1 \) value. Therefore, in the BP Tau case, a spin equilibrium situation (\( f_1 = 1 \)) is only compatible with an accretion luminosity smaller than \( L_{\text{UV}} \approx 0.027 \, L_\odot \). This is actually consistent with the recent estimate \( L_{\text{UV}} = 0.023 \, L_\odot \) (Donati et al. 2010a), which in fact allows a maximum spin-down efficiency \( f_1 = 1.06 \). However, the higher luminosity \( (L_{\text{UV}} = 0.179 \, L_\odot) \) estimated by Gullbring et al. (1998) is not compatible with a zero-torque condition, independently of the value of the accretion rate and it allows a maximum spin-down efficiency \( f_1 = 0.52 \): the stellar wind is consuming too much accretion power. Anyway, it is possible to see in the right panel of Figure 1 that, independently of the accretion luminosity, the properties of a stellar wind at maximum spin-down efficiency are very demanding having a mass flux \( f_M \approx 0.4-0.5 \) and consuming an important fraction of the accretion power \( (f_E \sim 0.6) \). For \( f_1 < 0.2 \) the wind characteristics are less tough.

A summary of the limits obtained applying the APSW model to the star sample of Table 1 is given in Table 2. For each star \((M_*, R_*, B_*, P_*, L_{\text{UV}})\) specified, we show the maximum spin-down efficiency \( f_1 \) and the corresponding mass outflow rate \( f_M \), energy requirement \( f_E \), accretion rate, and wind outflow rate. If a value \( f_1 > 1 \) is found, we show in parentheses the characteristics of the same star–wind system at spin equilibrium (\( f_1 = 1 \)).

4. SPIN EQUILIBRIUM

In Figure 2, we plot several \( L_{\text{UV}}-\delta \) solutions of Equation (4) at spin equilibrium (\( f_1 = 1 \)) for different
different accretion rates (Equation (3) with different linestyles correspond to the maximum accretion power available for spin equilibrium solutions for a given magnetic field. Black thin lines with $B^\star$ values (using different linestyles).

The figure shows that, for a given magnetic field $B^\star$, the accretion luminosity–$\delta$ configurations at spin equilibrium have an envelope depending on the normalized dipolar component required by a spin equilibrium. Linestyles correspond to the solutions of Equation (5):

$$f_{M,\text{max}} \approx 0.48 \left[ 1.55 - 0.55 \left( \frac{\delta}{0.05} \right)^{0.467} \right]$$

$$f_{E,\text{max}} \approx 0.63 \left[ 1.03 - 0.03 \left( \frac{\delta}{0.05} \right)^{1.04} \right].$$

These values represent the maximum efficiencies required by the spin equilibrium: for a dipolar component stronger than the intensity defined by Equation (5), both efficiencies can be lower (see, for example, the values between parentheses in Table 2). On the other hand, a stronger field determines at some point the transition to a propeller regime ($R_t > R_{co}$). Imposing $\Psi = (\delta^{-2/3}/C^{7/2})$ into Equations (2) and (3), we can therefore estimate the minimum ejection and energy efficiencies of a stellar wind at spin equilibrium in a non-propeller regime:

$$f_{M,\text{min}} \approx 0.15 \left( \frac{\delta}{0.05} \right)^{-0.53}$$

$$f_{E,\text{min}} \approx 0.18 \left( \frac{\delta}{0.05} \right)^{-0.53} \left[ 1.25 - 0.25 \left( \frac{\delta}{0.05} \right)^{2/3} \right]^{-1}. $$

Anyway, the upper limits on $f_M$ and $f_E$ are likely more robust than the lower ones. When the truncation radius approaches corotation, which is the condition used to derive Equation (7), it is not certain if our approximation for the truncation radius is still valid. Using the $C$ factor as a measure of this uncertainty, the $f_{M,\text{min}}$ value is more sensible to errors on the truncation position ($f_{M,\text{min}} \propto C^{2.82}$) than $f_{M,\text{max}}$ ($f_{M,\text{max}} \propto C^{0.22}$) in the fiducial values $C = 0.75, \delta = 0.05$. A field larger than Equation (5) limit would be also compatible with an $f_1 > 1$ situation, but it would require mass and energy efficiencies even greater than Equation (6) estimates (see Table 2).

5. DISCUSSION AND CONCLUSIONS

We applied the APSW model (Matt & Pudritz 2005b) to a sample of CTTS to verify if stellar winds are a viable mechanism to spin down the rotation of accreting and contracting protostars. According to this scenario, a fraction of the energy deposited by the magnetospheric accretion flow onto the surface of the star could be used to drive the stellar wind: we added the additional constraint that the same accretion energy must be used to power the emission of the accretion shock. In Section 3 we showed that, for a given spin-down efficiency $f_1 \neq 0$, a maximum accretion luminosity can be attained: when the accretion power and, consequently, the spin-up torque become too large, the stellar wind consumes too much energy and a smaller and smaller fraction is left to support the emission.

In Table 2, we showed that the stars in our sample characterized by a high accretion luminosity ($L_{UV} \geq 0.1 L_\odot$) impose severe limits on the accretion power which is available to drive the stellar wind so that the spin-down torque is not strong enough to achieve a spin equilibrium. This is consistent with Equation (5): in the range of parameters covered in our sample, a dipolar component of kG intensity is required by the wind to spin down a star characterized by such a high UV emission. In Figure 2, we also showed that an important fraction of our sample (around

$$B^\star (M_*/M_\odot)^{-1/4} (R_t/R^0)^{3/4}$$

values (using different shades of gray) and different $M_\text{acc}$ values (using different linestyles).

We also truncate the solutions whenever the truncation radius is equal to the corotation radius: $R_t = R_{co} = \delta^{-2/3} R^*$. In fact, when the disk truncation occurs very close to corotation, disk-locked solutions are in principle possible (Matt & Pudritz 2005a), while, for $R_t > R_{co}$, accretion becomes intermittent and highly variable (“propeller” regime; Ustyugova et al. 2006): in both cases the magnetic star–disk interaction can provide an efficient enough braking torque.

The figure shows that, for a given magnetic field $B^\star$, the accretion luminosity–$\delta$ configurations at spin equilibrium have an envelope depending on the normalized $B^\star$ value (solid lines): the solid curves therefore give an estimate of the minimum dipolar component compatible with a spin equilibrium condition. Applying the model to our entire star sample, it is possible to see that dipolar fields between 100 G and 3 kG are required to power an APSW at spin equilibrium compatible with the observed accretion luminosity. In many cases a field stronger than 1 kG is required, which has been currently observed only in the case of AA Tau.

By fitting the solid curves in Figure 2 we can define the minimum intensity of the dipolar component required by a spin equilibrium configuration, given the characteristic of a particular star ($L_{UV}$, $M^\star$, $R^\star$, $\delta$):

$$B^\star_{\text{min}} \approx 1025 \left[ \frac{L_{UV}}{0.1 L_\odot} \right]^{1/2} \left( \frac{M^\star}{M_\odot} \right)^{-1/4} \left( \frac{R^\star}{2 R^0} \right)^{-3/4} \left[ 0.933 \left( \frac{\delta}{0.05} \right)^{-1.5} + 0.067 \left( \frac{\delta}{0.05} \right)^{0.45} \right] G.$$ (5)

Note that Equation (5) gives a lower limit for the dipolar field intensity at the stellar equator, while the value at the stellar pole is twice larger. The dependency on $\delta$ has been obtained by fitting the solutions in the range $0.01 < \delta < 0.266$, where $\delta \approx 0.266$ corresponds to an extrema of the solid curves in Figure 2. The other dependencies are exact. We can also estimate the mass and energy efficiencies $f_M$ and $f_E$ of the stellar wind which correspond to the solutions of Equation (5):

$$f_{M,\text{max}} \approx 0.48 \left[ 1.55 - 0.55 \left( \frac{\delta}{0.05} \right)^{0.467} \right]$$

$$f_{E,\text{max}} \approx 0.63 \left[ 1.03 - 0.03 \left( \frac{\delta}{0.05} \right)^{1.04} \right].$$

Anyway, the upper limits on $f_M$ and $f_E$ are likely more robust than the lower ones. When the truncation radius approaches corotation, which is the condition used to derive Equation (7), it is not certain if our approximation for the truncation radius is still valid. Using the $C$ factor as a measure of this uncertainty, the $f_{M,\text{min}}$ value is more sensible to errors on the truncation position ($f_{M,\text{min}} \propto C^{2.82}$) than $f_{M,\text{max}}$ ($f_{M,\text{max}} \propto C^{0.22}$) in the fiducial values $C = 0.75, \delta = 0.05$. A field larger than Equation (5) limit would be also compatible with an $f_1 > 1$ situation, but it would require mass and energy efficiencies even greater than Equation (6) estimates (see Table 2).
50%) would require such a strong dipolar component to be compatible with a zero-torque condition: at the moment of writing a dipolar component of kG intensity has been measured only in the case of AA Tau.

Equation (5) and Figure 2 clearly show that lower UV luminosity and/or faster spinning stars require weaker fields, more consistent with the dipolar intensities currently measured, to be in spin equilibrium with an APSW. In fact, stars in the sample with a low accretion luminosity \( L_{\text{UV}} \ll 0.1 L_\odot \) are compatible with an \( f_1 \geq 1 \) situation, but the corresponding APSW at maximum spin-down efficiency is energetically very demanding, as confirmed by Equation (6). Some low-luminosity cases at spin equilibrium are less demanding, see, e.g., the CR Cha or AA Tau examples. Still, the mass fluxes would correspond roughly to the entire mass flux of T Tauri jets (1%–20%; Cabrit 2009): this would imply that stellar winds are the primary ejecting component of young stars, which seems unlikely (Ferreira et al. 2006). Besides, even for relatively low ejection rates \( f_M \ll 0.1 \) the energy input is still an issue. It is already known that the wind cannot be thermally driven (Matt & Putzitz 2007): a temperature close to virial (\( \sim 10^6 \) K) determines an emission that is too high and incompatible with observations (Dupree et al. 2005). Turbulent Alfvén waves represent another possible pressure source (DeCampli 1981). Furthermore, it has been suggested that the amplitude of the waves generated by the impact of the accretion streams onto the surface of the star is greater than interior convection-driven wave amplitude (Cranmer 2009). In this case, the accretion/ejection energy coupling is not easy to determine: recent models suggest anyway that the wind mass-loss rates due to this mechanism are generally very low \( 10^{-5} < f_M < 10^{-2}; \) Cranmer 2009). Besides, it is important to remark that in our sample, when the APSW ejection efficiency \( f_M \) at spin equilibrium becomes \( \ll 0.1 \), the star–disk system approaches a propeller regime (Cranmer 2008). The problem of the spin of accreting and contracting stars like T Tauri has still many open issues. It is likely that diverse mechanisms contribute at the same time with different degrees: stellar winds, magnetospheric star–disk angular momentum exchanges (Zanni & Ferreira 2009), and magnetospheric ejections driven by the star–disk interaction (Ferreira et al. 2000; Zanni 2009).

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