Binary interactions with high accretion rates onto main sequence stars

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Abstract Energetic outflows from main sequence stars accreting mass at very high rates might account for the powering of some eruptive objects, such as merging main sequence stars, major eruptions of luminous blue variables, e.g., the Great Eruption of Eta Carinae, and other intermediate luminosity optical transients (ILOTs; red novae; red transients). These powerful outflows could potentially also supply the extra energy required in the common envelope process and in the grazing envelope evolution of binary systems. We propose that a massive outflow/jets mediated by magnetic fields might remove energy and angular momentum from the accretion disk to allow such high accretion rate flows. By examining the possible activity of the magnetic fields of accretion disks, we conclude that indeed main sequence stars might accrete mass at very high rates, up to $\approx 10^{-2} \, M_\odot \, \mathrm{yr}^{-1}$ for solar type stars, and up to $\approx 1 \, M_\odot \, \mathrm{yr}^{-1}$ for very massive stars. We speculate that magnetic fields amplified in such extreme conditions might lead to the formation of massive bipolar outflows that can remove most of the disk's energy and angular momentum. It is this energy and angular momentum removal that allows the very high mass accretion rate onto main sequence stars.

Key words: stars: AGB and post-AGB — stars: jets — (stars:) binaries (including multiple): close

1 INTRODUCTION

Evidence has been acquired in recent years on the need to allow compact objects to accrete mass at rates exceeding the Eddington limit by about one to two orders of magnitude. Extreme cases of super-Eddington accretion rates are the magnetized neutron star M82X-11 that accretes from a massive companion at a rate of ≈ 100 times the Eddington limit (Bachetti et al. 2014), and two active galactic nuclei (AGNs) with $\approx 200 - 300$ times the Eddington limit (Du et al. 2015). Other examples include accretion onto a main sequence star during mergers of stars (Soker & Kashi 2011), during periastron passages of main sequence stars near asymptotic giant branch (AGB) stars (Staff et al. 2016), and outbursts of some massive stars that might be powered by accretion onto a main sequence star (e.g., Kashi & Soker 2010b,a; Kashi et al. 2013); a general name for these types of objects is intermediate luminosity optical transients (ILOTs; also termed red novae or optical transients). In many cases, super-Eddington or sub-Eddington outflow kinetic power is larger than the radiation luminosity, e.g., AGN HS 0810+2554 (Chartas et al. 2014). Chartas et al. (2014) suggest that this implies that magnetic activity is likely behind the acceleration of the outflow.

Despite the observational hints for high mass accretion rates onto main sequence stars, it is commonly assumed that main sequence stars cannot accrete mass at high rates, above $\approx 10^{-3} \, M_\odot \, \mathrm{yr}^{-1}$ (e.g., Hjellming & Taam 1991).

Our main theme in the present study is to introduce the idea that main sequence stars might accrete mass at very high rates. This has implications on transient events, on common envelope evolution and on the newly proposed grazing envelope evolution (GEE). We approach this challenge by building on strong magnetic activity in the accretion disk, including an efficient dynamo and magnetic field reconnection. We only deal with main sequence stars, and with accretion rates of up to $\approx 0.01 \, M_\odot \, \mathrm{yr}^{-1}$ for low mass main sequence stars, and up to $\approx 1 \, M_\odot \, \mathrm{yr}^{-1}$ for very high mass main sequence stars. In the merger model for ILOT V838 Mon (Soker & Tylenda 2003) for example, a companion of $\approx 0.3 \, M_\odot$ was destroyed on a star with $\approx 6 \, M_\odot$. It is assumed that several percent of the companion mass was accreted within several months; the rest inflated a large envelope. This leads to an accretion rate of $\approx 0.01 - 0.1 \, M_\odot \, \mathrm{yr}^{-1}$. We note that Nandez et al. (2014) simulated a merger process with parameters appropriate for ILOT V1309 Sco and found the formation of an accretion disk in the synchronized case. They did not find jets, as expected if magnetic fields are not included. As well, in the non-synchronized case, the merger proceeds without the formation of an accretion disk. Namely, it is possible, according to their results, that a merger and high accretion rate do not necessarily require an accretion disk.

It is commonly agreed upon that jets are produced at the center of accretion disks and that magnetic fields are a crucial ingredient in their launching process (e.g., Livio 2011; Pudritz et al. 2012). Many models for the formation

of astrophysical massive outflows are based on the operation of large scale magnetic fields within and around the accretion disk (e.g., Zanni & Ferreira 2013 and Narayan et al. 2014 and references therein). By large scale we refer to magnetic fields with coherence scale larger than the radius of the disk at the considered location. In the "X-wind mechanism" introduced by Shu et al. (1988) and Shu et al. (1991), the jets are launched from a narrow region in the magnetopause of the stellar field. A different group of models also uses open radial magnetic field lines and is based on an outflow from an extended disk region, but they do not rely on the stellar magnetic field (e.g., Konigl & Pudritz 2000; Shu et al. 2000; Ferreira 2002; Krasnopolsky et al. 2003; Ferreira & Casse 2004).

There are several arguments that point to problems with models for launching jets that are only based on large scale ordered magnetic fields.

(1) Precessing jets. In several young stellar objects (YSOs), the jets precess on a time scale of 100 years (e.g. Ybarra et al. 2006). A large scale magnetic field cannot change its symmetry axis in such a short time.

(2) Collimated jets in planetary nebulae (PNe). There is a highly collimated clumpy double-jet in the PN Hen 2-90 (Sahai & Nyman 2000) with very similar properties to jets from YSOs that form Herbig-Haro objects (dense blobs along jets launched by YSOs). The source of the accreted mass in PNe is thought to be a companion star. In such a system, large scale magnetic fields are not expected.

(3) No jets in DQ Her systems (intermediate polars). DQ Her systems are cataclysmic variables where the magnetic field of the accreting white dwarf is thought to truncate the accretion disk at its inner boundary. This magnetic field geometry is the basis for some jet-launching models in YSOs (e.g., Shu et al. 1991). However, no jets are observed in intermediate polars.

(4) Some of these models do not take into account the strong dynamo that is expected to operate in accretion disks. In a few cases, e.g., Hujeirat et al. (2003), the modification of the large scale magnetic field in the disk is considered. Such dynamos substantially modify the structure of the magnetic field in the accretion disk and lead to a behavior closer to that in active stars.

(5) Ferreira (2013) summarizes accretion models of classical T Tauri stars (TTS), and concludes that steady state models, such as X-winds or Accretion Powered Stellar Winds, cannot spin down a protostar. The balance between the accretion torque (working so as to spin up the star) and the magnetic torque (working so as to slow the star's rotation) should lead to a spin up on a time scale smaller than the disk lifetime, whereas observations show that the rotation period of TTSs is almost constant, at roughly 10% of their break-up speed. Instead, unsteady models, such as Magnetospheric Ejection, are the best candidates for removal of extra energy from the disk. The model suggested in this study is a type of unsteady model.

In the present paper, we take the view that the very high accretion rates onto main sequence stars in ILOTs, in the range of $\approx 10^{-4} - 1\,M_\odot\,\mathrm{yr}^{-1}$ (Kashi & Soker 2010a), require the formation of thick accretion disks. The accretion rates in our model are high, and the disks are thick enough so as to sustain turbulent eddies for long durations. Those eddies, combined with the differential rotation of the disk's plasma, generate the dynamo magnetic fields. Strong fields will eject mass and momentum out of the disk without the need for inward diffusion.

In Section 2 we build such accretion disks that we will use in Section 3 where we study their magnetic field properties and activity. We will take the magnetic dynamo to be similar in many respects to that of a stellar dynamo. The merit of comparison to the dynamo operating in the Sun is demonstrated in a recent paper by Bugli et al. (2014). They study the dynamo inside a thick disk around a Kerr black hole (BH) and find the evolution of the magnetic field to qualitatively occur in the same fashion as in the Sun. Our summary is in Section 4.

2 BUILDING THE DISK

2.1 A Thick Disk

Here, we construct an accretion disk with a high accretion rate where part of the energy is carried away by jets. High accretion rate disks were studied before, e.g., Abramowicz et al. (1980) for an accreting BH, as well as accretion disks where heat is removed to heat the corona (e.g., Svensson & Zdziarski 1994), or to an outflow (e.g., Kuncic & Bicknell 2004). To present a comprehensive picture we repeat the construction of the disk presented in these studies, and present it in a way that will serve our goals.

Svensson & Zdziarski (1994) modified the standard Shakura-Sunyaev model for accretion disks (Shakura & Sunyaev 1973; Frank et al. 2002) by allowing a major fraction, f, of the power released from the accreted disk matter to be transported and dissipated in the corona. We replace their $(1 - f)$ factor by the parameter ε that is the fraction of energy liberated by the accretion process that goes to radiation

$$\frac{4\sigma T_c^4}{3\tau} = \varepsilon \frac{3GM\dot{M}}{8\pi R^3}\left[1 - \left(\frac{R_*}{R}\right)^{\frac{1}{2}}\right]. \qquad (1)$$

In our study, the rest of the energy is channeled to the outflowing gas, but the contributions to the pressure from this channel, e.g., magnetic fields, are neglected. The other symbols have their usual meanings: M is the mass of the star, $\dot{M}$ is the mass accretion rate, R is the radial coordinate in the disk plane, R_* is the stellar radius, τ is the optical depth, and σ is the Stefan-Boltzmann constant. We study accretion disks with very high mass accretion rates of $\approx 10^{-4} - 1\,M_\odot\,\mathrm{yr}^{-1}$ in which the primary star is a main sequence star with a radius of $R_* \approx 10^{11}$ cm. This can cause the opacity to be dominated by electron scattering processes rather than Kramers opacity induced by

free-free emission, and also the pressure to be dominated by radiation pressure rather than gas pressure.

In the gas pressure-dominated regime we only consider gas pressure, starting with the Kramers opacity. We will use a molecular weight of $\mu = 0.615$ throughout the paper. When the parameter ε is included, we find the following expressions for the radial dependance of the disk height H, the central temperature T_c, and the radial velocity of the disk v_R, respectively

$$H/R = 0.45\varepsilon^{1/20}\alpha^{-1/10}m^{-3/8}R_{12}^{1/8}\left(f^4\dot{M}_{24}\right)^{3/20}, \tag{2}$$

$$T_c = 2\times10^5\varepsilon^{1/10}\alpha^{-1/5}\left(\frac{m}{R_{12}^3}\right)^{1/4}\left(f^4\dot{M}_{24}\right)^{3/10}\text{K}, \tag{3}$$

and

$$v_R = -3.5\times10^6\varepsilon^{1/10}\alpha^{4/5}m^{-1/4}R_{12}^{-1/4}\dot{M}_{24}^{3/10}f^{-14/5}\text{cm s}^{-1}, \tag{4}$$

where α is the viscosity parameter in the disk which is assumed to be radius-independent, and we define the dimensionless variables

$$m \equiv \frac{M}{M_\odot}, \quad R_{12} \equiv \frac{R}{10^{12}\text{ cm}},$$

$$\dot{M}_{24} \equiv \frac{\dot{M}}{10^{24}\text{ g s}^{-1}}, \quad f^4 \equiv 1 - \sqrt{\frac{R_*}{R}}. \tag{5}$$

For these typical values the disk is no longer thin, and the radial velocity has the same order of magnitude as the Keplerian velocity. In addition, the high temperature suggests that electron scattering processes and radiation pressure might be important.

Regarding the pressure to be radiation pressure and taking electron scattering opacity give the following disk properties

$$H/R = 1.65\varepsilon\left(f^4\dot{M}_{24}\right)R_{12}^{-1}, \tag{6}$$

$$T_c = 1.3\times10^5\varepsilon^{-1/4}\alpha^{-1/4}\left(\frac{m}{R_{12}^3}\right)^{1/8}\text{K}, \tag{7}$$

and

$$v_R = -4.65\times10^7\varepsilon^2\alpha m^{1/2}R_{12}^{-5/2}\dot{M}_{24}^2f^4\text{cm s}^{-1}. \tag{8}$$

The properties are more sensitive to the parameter ε, which can strongly affect the thickness and the radial velocity of the disk. As noted in previous works (e.g., Socrates & Davis 2006), it is possible to build thin disks even for very high accretion rates by using $\varepsilon \ll 1$. Another important difference is that the disk height over radius is a decreasing function of radius where, in the gas pressure dominated case, it is an increasing function of radius.

In any case, in order not to require an extremely small value of ε, and demanding that H is not much larger than r, we find from the expressions for H/r that the accretion rate in our treatment is limited. To an order of magnitude, this limitation on the accretion rate is $\approx 10^{-2}M_\odot$ yr^{-1} for solar type stars, and up to $\approx 1M_\odot$ yr^{-1} for stars of 30 and more solar masses.

2.2 Relevant Time Scales

There are several relevant time scales. The dynamical time scale (or the Keplerian time divided by 2π) is given by

$$t_\phi \approx \frac{R}{v_\phi} \approx \Omega_K^{-1}. \tag{9}$$

This time scale is approximately equal to the sound crossing time along the disk height, which is about the dynamical time scale required to reach equilibrium along the z direction (perpendicular to the disk plane)

$$t_z \approx \frac{H}{c_s}, \tag{10}$$

where c_s is the sound speed. The viscous time scale is given by

$$t_{\text{visc}} \approx \frac{R^2}{\nu} \approx \frac{R}{v_R}, \tag{11}$$

where $\nu = \alpha c_s H$ is the turbulence viscosity coefficient. The thermal time scale is given by

$$t_{\text{th}} = \frac{\text{heat content per unit area}}{\text{dissipation rate per unit area}} \approx \frac{\Sigma c_s^2}{D(R)}, \tag{12}$$

where $D(R)$ is the dissipation due to differential viscous torque and Σ is the surface density of the disk.

In the radiation dominated regime it can be shown that

$$t_{\text{th}} = \varepsilon\frac{H}{c}\tau = \varepsilon\frac{t_{\text{diff}}}{3}, \tag{13}$$

where $t_{\text{diff}} = 3H\tau/c$ is the photon diffusion time from the disk's mid-plane outward. From Equation (13) we can interpret ε as the probability of losing energy (heat) via photon diffusion (and hence radiation) rather than outflow of kinetic energy. Approximately, as the diffusion time becomes longer, a larger fraction of the energy is transported by channels other than radiation. In our model these are magnetic fields and gas outflows. We note that when magnetic fields remove energy much faster than photons do, the disk might be geometrically thin, but still the radiation carries only a small fraction of the energy.

Let us dwell on this point. The viscosity in the disk and magnetic activity dissipate energy. But in our regime of very high accretion rate, the photon diffusion time is very long. Most of the dissipated energy that starts as thermal energy, rather than being emitted by the disk, will lead to internal motion, such as turbulence that amplifies magnetic fields, and to gas buoyancy outward. Near the surface, in particular due to magnetic field reconnection below the photosphere (see Sect. 3), the energy is channeled to winds or jets. The photons have no time to diffuse and most of the energy is carried by the outflow, i.e., kinetic energy. The acceleration occurs in many sporadic impulsive events resulting from magnetic field reconnection. Most of the reconnection events are much shorter than the dynamical time, and hence the mass and energy they eject merge into one continuous outflow.

For the treatment above to hold, the inflow time must be longer than one rotation time of the material in the disk

$$t_{\rm in} \gtrsim t_{\rm kep} = 2\pi t_\phi. \qquad (14)$$

This condition can be cast into a height over radius condition. With the definitions of time scales $t_{\rm in} = R/v_{\rm R}$ and $t_\phi = H/c_{\rm s}$, and using the relation $v_{\rm R} = (3\nu/2R) f^{-4}$, one gets

$$\frac{1}{2\pi} \gtrsim \alpha \frac{3}{2} \left(\frac{H}{R}\right)^2 f^{-4}. \qquad (15)$$

The area in the $\varepsilon - \dot{M}$ plane where this inequality holds is termed the "safe-zone." For other values of α, R_* and M, the safe-zone changes.

3 MAGNETIC ACTIVITY

3.1 Dynamo

The high accretion rate is associated with very large Reynolds flows and therefore the transition to turbulence is inevitable. The differential rotation in the disk is responsible for generation of the toroidal magnetic field, the Ω-effect and the combined action of helical turbulence and differential rotation generate the poloidal field, the α-effect and so forth. This is the basic premise of the α-Ω dynamo. We follow the treatment of the magnetic field amplification as described by Pudritz (1981a,b).

We study flow regimes with very high Reynolds numbers. Nath & Mukhopadhyay (2015) have shown recently that for high Reynolds numbers, as is the case for realistic astrophysical accretion disks, nonlinearity is achieved by transient growth (e.g., Umurhan & Regev 2004; Umurhan et al. 2007) rather than by magnetorotational instability (MRI). Nath & Mukhopadhyay (2015) also showed that when transient growth modes come into play, the magnetic fields increase much more than what the pure MRI model gives. For that, although the results obtained by Pudritz (1981a,b) might not be applicable when MRI dominates, in the regime studied here the transient growth dominates and we use the results by Pudritz (1981a,b).

The dynamo action is described by mean field theory, in which the magnetic field and the velocity are decomposed into mean and fluctuating components,

$$\boldsymbol{B_t} = \boldsymbol{B} + \boldsymbol{b}, \qquad \boldsymbol{U_t} = \boldsymbol{U} + \boldsymbol{u}. \qquad (16)$$

The induction equation for the fluctuating component of the magnetic field is described by the so called "first order smoothing theory" as

$$\frac{\partial \boldsymbol{b}}{\partial t} = \boldsymbol{\nabla} \times (\boldsymbol{U} \times \boldsymbol{b} + \boldsymbol{u} \times \boldsymbol{B}) + \eta \nabla^2 \boldsymbol{b}. \qquad (17)$$

The last term comes from ohmic dissipation of the fluctuating component of the current, given by $\boldsymbol{j} = (c/4\pi)\nabla \times \boldsymbol{b}$, and reduces the magnetic field. The first term on the right hand side might lead to convective amplification of

the magnetic field, the "α-dynamo." From the assumption that the ohmic dissipation rate of fluctuating magnetic energy per unit volume $\dot{e}_{\rm Oh} = (\eta/l_u^2)b^2/4\pi$ equals the rate at which turbulence pumps energy back $\dot{e}_{\rm tr} = (\eta_T/l_u^2)B^2/4\pi$, where l_u is the largest turbulence scale, Pudritz (1981a,b) derived the magnitude of the fluctuating magnetic field

$$b^2 = \frac{\eta_T}{\eta}B^2. \qquad (18)$$

The magnetic diffusivity η is given by the expression

$$\eta \equiv c^2/4\pi\sigma_{\rm s} = 2.95 \times 10^{11} \ln \Lambda \, T^{-3/2} \ \ {\rm cm^2 \ s^{-1}} , \quad (19)$$

where $\sigma_{\rm s}$ is the Spitzer conductivity, T is the temperature, and $\ln \Lambda$ is of order 10 (see Huba 2013). η_T is the turbulent diffusivity for the mean magnetic field B,

$$\eta_T = M_t^2 H^2/t_K . \qquad (20)$$

Here H is disk height, t_K is the Keplerian orbital period and M_t is the turbulent Mach number defined as

$$M_t \equiv u/c_{\rm s} , \qquad (21)$$

with u being the rms turbulent speed and $c_{\rm s}$ the speed of sound.

Substituting Equation (2) into (20), and (3) into (19) we find the ratio

$$\begin{aligned} b^2/B^2 &= \eta_T/\eta \\ &= 3.2 \times 10^{13} \left(\frac{M_t}{0.5}\right)^2 \left(\frac{\varepsilon}{0.1}\right)^{1/4} \\ &\times \left(\frac{\alpha}{0.1}\right)^{-1/2} \left(\frac{m}{R_{12}^3}\right)^{1/8} \left(f^4 \dot{M}_{24}\right)^{3/4} . \end{aligned} \quad (22)$$

There are several large uncertainties in the derivation of Equation (22). Nonetheless, its implication is clear: the fluctuating magnetic field dominates. For any initial large scale, of the order of the disk radius or more, the magnetic field becomes negligible in the disk. At these high accretion rates large scale magnetic fields are not expected to play a significant role in the ejection of jets or disk winds. We should look for the role played by the fluctuating magnetic field, which is amplified by the dynamo, in launching the jets and winds.

The ratio of the fluctuating magnetic pressure to the total pressure is defined by

$$\lambda \equiv b^2/4\pi P, \qquad (23)$$

so that the mean magnetic field is

$$B = \lambda^{1/2}\sqrt{4\pi P \, \eta/\eta_T}. \qquad (24)$$

Taking $P = \rho c_{\rm s}^2$, Pudritz & Fahlman (1982) estimated a very large value λ of

$$10^{7/2} \leq \lambda \leq 10^4. \qquad (25)$$

As we are discussing accretion into main sequence stars where we apply the disk structure that neglects magnetic

pressure (Sect. 2), we limit the ratio of the magnetic to gas pressure to $\lambda \lesssim 1$. As well, we take a conservative approach in this preliminary study. We derive the following value for the magnetic field

$$|b| = 2 \times 10^4 \sqrt{\lambda} \left(\frac{\varepsilon}{0.1}\right)^{-1/40} \left(\frac{\alpha}{0.1}\right)^{-9/20}$$
$$\times \left(\frac{m}{R_{12}^3}\right)^{7/16} \left(f^4 \dot{M}_{24}\right)^{17/40} \text{G} \,. \quad (26)$$

Two comments are needed here regarding the numerical value given in Equation (26).

(i) In solar spots the magnetic fields reach equipartition, between the magnetic and thermal pressure, for magnetic fields of $\approx 2000 \, \text{G}$. The magnetic fields in the accretion disk studied here are about one order of magnitude stronger. This implies that the magnetic activity in the accretion disk is extremely strong, with magnetic pressures about two orders of magnitude above the photospheric pressure of the accreting star. It should be noted that for the very high accretion rates used here $\gtrsim 10^{-2} \, M_\odot \, \text{yr}^{-1}$ and the low value of $\varepsilon \approx 0.1$, both the density and temperature inside the disk are much higher than the corresponding values on the stellar photosphere. The pressure inside the disk is therefore higher by more than two orders of magnitude than that on the stellar surface, and the strong magnetic field is still below the thermal pressure inside the disk.

(ii) The scaling in Equation (26) takes $\lambda = 1$. This is applicable to the case where the magnetic field reaches equipartition with the thermal pressure in the accretion disk, as is the situation in solar spots. Above the solar surface the magnetic pressure can be much larger than the thermal pressure. Similarly, it is quite possible that in the accretion disk we study here the magnetic pressure will be larger than the thermal pressure, and be rather equal to the ram pressure of the Keplerian motion in the disk ρv_ϕ^2. This is the case for $\lambda > 1$, as applied by Pudritz & Fahlman (1982). Such magnetic fields will make the proposed scenario even more feasible.

3.2 Magnetic Power

We now show that most of the energy could potentially be carried off by magnetic activity rather than by radiation even when the magnetic energy in the disk is smaller than the thermal energy.

We assume that most of the magnetic activity takes place just inward of radius R_B, not too far from the stellar surface. We neglect further amplification of the magnetic field inward of this radius. The mass accretion rate is given by $\dot{M} \approx 4\pi R_B H \rho v_R$. We define the ratio of the inward radial velocity to the Keplerian velocity $\delta = v_R/v_\phi < 1$, and the kinetic energy density in the disk $e_d \simeq \rho v_\phi^2/2$. The accretion power of material crossing the radius R_B inward, i.e., the energy flux through radius R_B, is given by

$$L_{\text{acc}} = \frac{1}{2} \dot{M} v_\phi^2 \approx 2\pi \delta H R_B v_\phi e_d. \quad (27)$$

We take the time to amplify and release magnetic energy to be the dynamical time given by Equation (10), $\tau_B \approx t_z \approx H/c_s \approx R_B/v_\phi$. With an energy density $e_B = b^2/8\pi$ the magnetic power within radius R_B is $L_B \approx 2\pi R_B^2 H e_B/\tau_B$. The ratio of magnetic energy density to kinetic energy density is then

$$\frac{e_B}{e_d} \approx \delta \frac{L_B}{L_{\text{acc}}}. \quad (28)$$

We recall that in deriving Equation (28) we assumed that there is a global steady state, and hence the amplification time of the magnetic field is equal to the time scale to remove magnetic energy by reconnection. Hence, for a given power, shorter amplification and reconnection time scales imply weaker magnetic fields. Alternatively, if the ratio of the magnetic energy density to the kinetic energy density is some fixed value, like unity, then the magnetic power increases with decreasing time scale.

As $\delta < 1$, the above simple estimate shows that the magnetic energy can be lower than the kinetic energy in the disk, and yet the magnetic activity might carry a large fraction of the accretion energy. For that, the magnetic energy should be strongly amplified by vigorous turbulence in the disk. We assume that the high accretion rate leads to this behavior in the disk.

3.3 Reconnection and Outflow

In this subsection we speculate that the magnetic activity might lead to an energetic outflow.

We define three general coherence scales of the magnetic fields. Small scale is defined as much smaller than the disk scale height $l_B \ll H$. Moderate scale for which the coherence scale is of the order, but smaller than, the disk scale hight, $0.3H \lesssim l_B < H$. Finally, the large scale fields have $l_B \gtrsim r > H$. In the present study magnetic fields with moderate scale, about equal to the turbulence scale, dominate. We give no role to large scale magnetic fields (unlike many other models for the formation of jets).

As magnetic flux accumulates and tension increases due to the dynamo action, moderate-scale turbulence triggers local explosive events related to magnetic reconnection (Lazarian et al. 2015). Those events are sporadic and decay as the extra magnetic energy is released. At any given radius R, the largest turbulence scale l_u is described by $l_u = M_t \cdot H$ (Pudritz 1981a,b), and for highly turbulent Mach numbers such as $M_t \sim 0.5$ expected for the presently studied accretion flows (Sect. 2), the vertical scale can be populated by only two such large eddies. Large magnetic flux tubes will also buoy quite fast, relative to small flux tubes, outward from the disk, allowing the replenishing time of the magnetic fields to be of the order of the dynamical time.

There are two key processes resulting from the large flux tubes. These are magnetic torque and reconnection in scales not much smaller than the disk scale height. In accretion disk models where the large scale magnetic fields extend beyond the disk, there is a magnetic torque on the

disk from gas outside the disk, mainly a disk wind. When there is no such field, a torque acts within the disk. When the flux tubes are small relative to the radius, the torque is local and plays no large scale role. In our situation, the flux tubes are within the disk, but they are very large. Therefore, we speculate that they might transfer angular momentum and energy from the internal parts of the disk to near the surface of the disk. Near the surface, the magnetic field is likely to reconnect and accelerate the wind/jets. This process further reduces the amount of energy that is radiated from the disk, hence lowering the value of ε.

The vigourous turbulence facilitates magnetic field reconnection (Lazarian et al. 2015). We therefore propose that magnetic field reconnection will not only take place above the disk, like in the Sun, but also below the disk surface. The outcomes of the reconnection process for the large flux tubes might be as follows.

(1) *Mass outflow.* Like in the Sun, once reconnection occurs between two adjacent magnetic flux tubes, plasma flows along reconnected magnetic field lines, in two opposite directions. In our setting, we assume that a large fraction of the reconnection events takes place below the photosphere. Although initially a large fraction of the dissipated energy is thermal, before photons have time to diffuse out the gas is accelerated to form an outflow. Each reconnection event is like a small explosion, much like in supernova explosions where the diffusion time for photons is very long and most of the final energy is kinetic with only a small fraction emitted as radiation.

(2) *Huge magnetic arcs.* The reconnection process implies more than just direct mass ejection. The reconnection might also lead to the build up of huge magnetic arcs above the disk. Although the original flux tube emerging from the disk is expected to reside at about the same radial distance R, we speculate that the new arcs formed by reconnection can have their two footpoints at different radii, much like trans-equatorial magnetic loops in the Sun (Tadesse et al. 2014). We term these speculative structures trans-radial loops. The trans-radial loops experience very large shear, and might lead to further amplification of the magnetic energy. This process will be studied in the next paper in the series (see also Heyvaerts & Priest 1989).

(3) *Flares.* Reconnection of such huge arcs might lead to extremely energetic flares. Such an energetic flare was speculated to be the cause of the 2009 November outburst of Aquila X-1 (Soker 2010). de Gouveia Dal Pino et al. (2010), Khiali et al. (2015) and Kadowaki et al. (2015) discuss such rapid reconnection processes in relation to different properties of accretion disks associated with microquasars. Here, we aim at accretion disks around main sequence stars accreting at a very high rate. We expect the wind from the disk to obscure radiation from such flares.

4 SUMMARY

We describe a crude accretion flow setting that allows main sequence stars to accrete mass at very high rates. The key to allowing such high accretion rates is that most of the energy liberated in the accretion process, and the angular momentum carried by the inflowing gas, are removed by massive outflows. Svensson & Zdziarski (1994) used the removed energy to heat the corona, and Kuncic & Bicknell (2004) produced jets by magnetic torque, both in accretion disks onto BHs. Here, we studied accretion onto main sequence stars, and the removal of energy and angular momentum by jets that are formed from sporadic reconnection events of magnetic fields that are amplified by a strong dynamo. Although in previous papers from our group it was just assumed that main sequence stars might accrete at a very high rate, in the present study we have outlined the scenario for that. The main result is that the removal of angular momentum and energy cannot be achieved without magnetic fields that are efficiently amplified by the operation of a dynamo in the dense accretion disk.

As was done in the past (e.g., Svensson & Zdziarski 1994), we described a disk structure where only a fraction $\varepsilon \ll 1$ of the energy liberated by accretion is channeled to radiation (Sect. 2). From this crude structure we could constrain the properties of the disk around main sequence stars. For the parameters used here, we found a thick accretion disk dominated by thermal pressure (Eqs. (2) – (4)). To still be regarded as an accretion disk in the usual meaning, the inflow time should be longer than the orbital time.

Using the disk properties, we speculated on possible magnetic activities that might lead to rapid mass ejection from the disk. We presented the expected dynamo amplification of the magnetic field, and concluded that a very strong fluctuating magnetic field is expected in the disk (Eq. (22)). We then argued (Sect. 3.2) that it is possible for the magnetic power to be very large without altering the general disk structure. We also note that for geometrically thin disks, for which $\delta \ll 1$, the removal of energy by magnetic fields can be efficient even for $e_B \ll e_d$. Such a thin disk might be obtained if the disk is radiation dominated, as seen from Equation (6) for $\varepsilon \ll 1$. This outlined scenario for allowing high mass accretion rates onto main sequence stars, although speculative in some parts, is our main new result. We hope it will encourage relevant models of astrophysical systems, such as common envelope evolution, to consider main sequence stars that accrete mass at very high rates, as well as motivate numerical studies of the accretion flow. Such 3-dimensional numerical studies should include magnetic fields and have high spatial resolution to resolve the turbulence in the accretion disk.

Some key ingredients of the proposed scenario should be emphasized.

(1) In the proposed model magnetic fields with moderate coherence scales dominate. By moderate scale we refer to a scale of $0.3H \lesssim l_B < H$, where H is the disk scale height. Large-scale magnetic fields with $l_B \gg H$

play no role in our model (unlike in many other models of the formation of jets).

(2) Although magnetic energy is dissipated locally to heat by reconnection, in our setting most of this energy ends up as kinetic energy of the outflow. The reason is that because of the very high mass accretion and outflow rates, the flow is optically thick. Therefore, before photons have time to diffuse out the gas expands. Most of the thermal energy goes to adiabatic expansion rather than to radiation (much like in supernova explosions).

The outcomes of strong magnetic field reconnection below the photosphere, discussed in Section 3.3, lead us to propose the existence of trans-radial magnetic field lines, similar to those studied by Heyvaerts & Priest (1989), namely, huge magnetic arcs above the disk that have their two footpoints at different radii. These trans-radial magnetic fields might lead to further magnetic field amplification. We expect trans-radial loops to play a major role in the magnetic activity of the boundary layer, where the disk angular velocity substantially decreases to match the angular velocity of the star. The dynamics of trans-radial loops is the subject of a forthcoming paper.

Our main conclusion is that main sequence stars might accrete mass at very high rates, up to $\approx 10^{-2}\,M_\odot\,\mathrm{yr}^{-1}$ for solar type stars and up to $\approx 1\,M_\odot\,\mathrm{yr}^{-1}$ for very massive stars. Such accretion rates could potentially account for some outbursting objects that have their luminosity in the bulk region between novae and supernovae and that are thought to be powered by high accretion rates onto main sequence stars in binary systems (Kashi & Soker 2010a,a; Kashi et al. 2013). As well, such high accretion rates might take place in main sequence companions in eccentric orbits around AGB stars (Staff et al. 2016). A general name for these types of objects is ILOTs (also termed red novae or optical transients). Examples of specific objects include stellar mergers, such as V838 Mon (Soker & Tylenda 2003), the nineteenth century Great Eruption of Eta Carinae (Kashi & Soker 2010b), progenitors of some PNe (Soker & Kashi 2012), such as KjPn 8 (Boumis & Meaburn 2013), and the pre-explosion outbursts of SN 2009ip (Soker & Kashi 2013).

High accretion rates onto, and hence high outflow rates from, main sequence stars might supply extra energy to remove a common envelope where the secondary is a main sequence star. When the envelope removal by jets is sufficiently efficient as to remove the entire envelope residing outside the secondary star, no real common envelope phase is established. Instead, GEE commences (Soker 2015). The GEE might nicely explain the observation that the kinetic energy of the bipolar outflow from the binary system HD 101584 exceeds the released orbital energy (Olofsson et al. 2015).

We can summarize our study by stating that the possibility for main sequence stars to eject winds/jets at extremely high rates, as we argued here, opens a rich variety of processes that could potentially account for some properties of and some evolutionary puzzles in binary stellar systems.

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