ARE THERE MAGNETARS IN HIGH-MASS X-RAY BINARIES?
THE CASE OF SUPERGIANT FAST X-RAY TRANSIENTS
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ABSTRACT
In this paper we survey the theory of wind accretion in high-mass X-ray binaries hosting a magnetic neutron star and a supergiant companion. We concentrate on the different types of interaction between the inflowing wind matter and the neutron star magnetosphere that are relevant when accretion of matter onto the neutron star surface is largely inhibited; these include inhibition through the centrifugal and magnetic barriers. Expanding on earlier work, we calculate the expected luminosity for each regime and derive the conditions under which transition from one regime to another can take place. We show that very large luminosity swings (\(\sim 10^4\) or more on timescales as short as hours) can result from transitions across different regimes. The activity displayed by supergiant fast X-ray transients, a recently discovered class of high-mass X-ray binaries in our galaxy, has often been interpreted in terms of direct accretion onto a neutron star immersed in an extremely clumpy stellar wind. We show here that the transitions across the magnetic and/or centrifugal barriers can explain the variability properties of these sources as a result of relatively modest variations in the stellar wind velocity and/or density. According to this interpretation we expect that supergiant fast X-ray transients which display very large luminosity swings and host a slowly spinning neutron star are characterized by magnetar-like fields, irrespective of whether the magnetic or the centrifugal barrier applies. Supergiant fast X-ray transients might thus provide a new opportunity to detect and study magnetars in binary systems.

Subject headings: accretion, accretion disks — stars: neutron — supergiants — X-rays: binaries — X-rays: stars

1. INTRODUCTION

High-mass X-ray binaries (HMXBs) consist of a collapsed object, usually a magnetic neutron star (NS), that accretes matter from an OB companion star. Mass transfer takes place because of the intense stellar wind from the OB star, part of which is captured by the collapsed object (e.g., Verbunt & van den Heuvel 1995). Only in some short orbital period systems, the early-type star, often a supergiant, fills its Roche lobe and leads to mass transfer through Roche lobe overflow (Tauris & van den Heuvel 2006). Persistent HMXBs accrete all the time and in most cases display X-ray luminosities in the \(10^{35} - 10^{38}\) erg s\(^{-1}\) range. Many HMXBs are transient systems that remain at low X-ray luminosity levels (\(10^{12} - 10^{13}\) erg s\(^{-1}\)) most of the time and undergo outbursts lasting from weeks to months. During these outbursts they display properties nearly identical to those of persistent HMXBs. Transient systems usually comprise a Be star donor and relatively long, moderately eccentric orbits, such that the star sits deep in its Roche lobe and stellar wind capture is the only mechanism through which mass transfer takes place. The occurrence of the outbursts is likely associated with variations in the stellar wind of the Be star, such as shell ejection episodes, or buildup of matter around the resonant orbits in the slow equatorial wind component (van den Heuvel & Rappaport 1987). However, there are characteristics of the outbursts that are difficult to interpret if accretion onto the NS surface takes place unimpeded also in quiescence; these are (1) the large outburst-to-quiescence X-ray luminosity swing (factor of \(\sim 10^3\) larger) and (2) the presence in a given source of low-luminosity (type I) outbursts recurring close to periastron and, at different times, of high-luminosity (type II) outbursts that last for several orbital cycles and display little (if any) X-ray flux variations associated to the orbital phase. These characteristics of Be transients can be explained if the accretion rate (and thus X-ray luminosity) variations that are produced by the stellar wind alone could be amplified by some “gating” mechanism. Since most Be star HMXB transients contain relatively fast spinning X-ray pulsars, such a mechanism has been identified with the centrifugal barrier that results from the rotation of the neutron star magnetosphere (Stella et al. 1986).

About 10 transient systems have recently been discovered which display sporadic outbursts lasting from minutes to hours (i.e., much shorter than those of Be star transients) and reach peak luminosities of \(\sim 10^{36} - 10^{37}\) erg s\(^{-1}\). These systems spend long time intervals in quiescence, with X-ray luminosities down to \(\sim 10^{-5}\) times lower than those in outburst; in spite of their association with OB supergiant companions, their behavior is thus at variance with other persistent and transient HMXBs. They define a new class of HMXBs, collectively termed supergiant fast X-ray transients (SFXTs). An overview of the properties of SFXTs is given in § 2.

If accretion onto the collapsed object of SFXTs takes place both in quiescence and outburst, then the corresponding X-ray luminosity swing, typically a factor of \(\sim 10^4 - 10^5\), would require wind inhomogeneities with a very large density and/or velocity contrast (according to the standard wind accretion, the mass capture rate onto the NS scales like \(M_w v_w^{-4}\), with \(M_w\) the mass-loss rate and \(v_w^{-4}\) the wind velocity of the supergiant star; Davidson & Ostriker 1973). Several authors (In’t Zand 2005; Leyder et al. 2007; Walter & Zurita Heras 2007) suggested the presence of dense clumps in the wind of the OB companions in order to attain the luminosity variations of SFXTs. While some observations provide evidence for a clumpy wind, the characteristics of such inhomogeneities are still poorly known. Numerical simulations suggest that clumps may originate from small-scale perturbations in the radiation-driven wind (Dessart & Owocki 2003; Prinja et al. 2005). Models involving accretion of clumps are still being actively pursued (Negueruela et al. 2008; Walter & Zurita Heras 2007).

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Heras 2007). The requirement on the density and/or velocity contrasts in the wind can be eased if there is a barrier that remains closed during quiescence, halting most of the accretion flow, and opens up in outbursts, leading to direct accretion (this is similar to the case of Be star transients). This paper develops and discusses gated accretion models for SFXTs.

In § 3, after reviewing the theory of wind accretion in HMXBs, we describe the different regimes of a rotating magnetic NS immersed in the stellar wind from its companion. In § 4 we discuss the conditions under which transitions across regimes can take place in response to variations in the wind parameters. As compared to previous works addressing gating mechanisms in transient accreting magnetic NSs (notably Stella et al. 1986), we present here a more comprehensive treatment of the different physical processes that have been discussed in this context.

In § 5 we present an application to the different states of two SFXTs and discuss in turn the possibility that the outbursts are driven by a centrifugal or a magnetic barrier. The latter mechanism involves a magnetosphere extending beyond the accretion radius (Bondi 1952), which prevents most of the inflowing matter from accreting; it is discussed here for the first time in relation to the activity of transient binary sources. Unless the stellar wind were extremely clumpy (as envisaged in other SFXT models), the onset of a barrier inhibiting direct accretion would be required to explain the activity of SFXTs. In this case we conclude that if SFXTs host slowly rotating NSs (spin periods of several hundreds to thousands of seconds), then they must possess magnetar-like fields ($\sim 10^{14} - 10^{15}$ G), independent of whether the centrifugal or magnetic barrier operates. We summarize our discussion and conclusions in §§ 6 and 7.

2. THE OBSERVED PROPERTIES OF SFXTs

SFXTs are observed to exhibit sporadic outbursts, lasting from minutes to hours, with peak X-ray luminosities between $\sim 10^{36}$ and $10^{37}$ erg s$^{-1}$ (see, e.g., González-Riestra et al. 2004; Sidoli et al. 2005; Grebnev & Sunyaev 2005; Lutovinov et al. 2005; Sguera et al. 2005, 2006, 2007; Masetti et al. 2006; Götz et al. 2007). No firm orbital period measurement has been obtained yet. A recent list of confirmed ($\sim 5$) and candidate ($\sim 6$) SFXTs is given by Walter & Zurita Heras (2007).

Between outbursts, SFXTs remain in quiescence with luminosities in the range $\sim 10^{31} - 10^{33}$ erg s$^{-1}$ (González-Riestra et al. 2004; In’t Zand 2005; Smith et al. 2006; Kneeba & Campina 2006). In some cases, very high peak-to-quiescence X-ray luminosity swings (factor of $\sim 10^4 - 10^5$) were seen on timescales comparable to the outburst duration. Some SFXTs also showed flares-like activity at intermediate luminosity levels (e.g., González-Riestra et al. 2004). In the case of IGR J17544–2619, two states of intermediate luminosity were observed: one before the onset of the outburst and the other immediately after, with X-ray luminosities $\sim 3$ and $\sim 1$ decades below the value reached at the peak of the outburst, respectively (In’t Zand 2005).

The sporadic character of SFXT outbursts, as observed with INTEGRAL, suggested that the duty cycle of these sources (the fraction of time spent in a high-luminosity state) is small ($\approx 0.02 - 0.002$; Walter & Zurita Heras 2007). However, recent observations carried out with the very sensitive X-ray telescopes on board XMM-Newton and Chandra revealed that some SFXTs display flares around a luminosity of $\gtrsim 10^{34} - 10^{35}$ erg s$^{-1}$ (i.e., well below the INTEGRAL limiting sensitivity) for a large fraction of the time (González-Riestra et al. 2004; In’t Zand 2005; Tomrick et al. 2006). Therefore, the indication is that the active phase of SFXT sources (as opposed to true quiescence) lasts longer than previously thought, and the duty cycles are of order $\sim 0.1$ or higher.

Optical identifications of SFXTs show that these sources are associated with OB supergiant companion stars (see, e.g., Walter & Zurita Heras 2007 and reference therein). The SFXT OB companions typically have a mass of $M_\ast \sim 30 M_\odot$, luminosity of $\log (L_\ast / L_\odot) \sim 5 - 6$, mass-loss rate of $M_\ast = 10^{-7}$ to $10^{-5} M_\odot$ yr$^{-1}$, and wind velocity of $1000 - 2000$ km s$^{-1}$. Note that isolated OB stars with $M_\ast > 5 M_\odot$ and $\log (L_\ast / L_\odot) \sim 5 - 6$ are persistent soft X-ray sources with luminosity around $\sim 10^{32}$ erg s$^{-1}$ (Cassinelli et al. 1981; Berghoefer et al. 1997).

It is widely believed that SFXTs contain sporadically accreting NSs. Only a little is known about their spin period. A coherent periodicity was detected at 4.7 and 228 s in AX J1841.0–0536 and IGR J16465–4507, respectively (Bamba et al. 2001; Lutovinov et al. 2005). However, IGR J16465–4507 showed only a factor of $\sim 100$ luminosity swing between quiescence and outburst, so it is unclear whether the source should be considered as a (transient) “intermediate system”; Walter & Zurita Heras 2007). The nature of the companion star in AX J1841.0–0536 is still debated (Halpern et al. 2004; Nespoli et al. 2007). Therefore, these two sources might not belong to the SFXT class. On the contrary, in the prototypical SFXTs XTE J1739–302 (Sguera et al. 2006) and IGR J16479–4514 (Walter et al. 2006) some evidence has been reported for periodicities in the $\sim 1000 - 2000$ s range. We assume in the following that SFXTs host a rotating magnetic NS.

3. STELLAR WIND ACCRETION

We investigate here the conditions under which a magnetized NS can accrete matter from the wind of a massive companion. In the theory of wind accretion in HMXBs, the following radii are defined (see, e.g., Illarionov & Sunyaev 1975; Stella et al. 1986):

1. The accretion radius $R_a$ is the distance at which the inflowing matter is gravitationally focused toward the NS (Bondi 1952). It is usually expressed as

$$R_a = 2GM_\text{NS}/v_w^2 = 3.7 \times 10^{10} r_g^{-2} \text{ cm,}$$

where $r_g$ is the wind velocity in units of 1000 km s$^{-1}$ and we assumed that the orbital velocity of the star is negligible (Frank et al. 2002). Throughout the paper we fix the NS radius and mass at $R_{\text{NS}} = 10^6$ cm and $M_{\text{NS}} = 1.4 M_\odot$, respectively. The fraction $M_{\text{capt}}/M_w$ of the stellar wind mass-loss rate ($M_w$) captured by the NS depends on $R_a$ through (Frank et al. 2002)

$$M_{\text{capt}}/M_w \simeq R_a^2/(4a^2) = 2 \times 10^{-5} r_g^{-4} a_{10}^{-2}. \quad (2)$$

$^4$ Only IGR J11215–5952 and AX J1749.1–2733 showed recurrent flaring activity, with periodicities of $\sim 165$ and $\sim 185$ days, respectively. These are interpreted as outbursts from two systems with unusually long orbital periods ($\geq 100$ days; Grebnev & Sunyaev 2005; Sidoli et al. 2007; Zurita Heras et al. 2007). Thus, Walter & Zurita Heras (2007) excluded these sources from their SFXT list.
Here $a = 4.2 \times 10^{12} \, a_{10, d} \, \text{cm}$ is the orbital separation, $a_{10, d} = P_{10, d}^{2/3} M_{30}^{1/3}$, $P_{10, d}$ is the binary orbital period in units of 10 days, and $M_{30}$ is the total mass in units of 30 M$_\odot$ (we assumed circular orbits).

2. The magnetospheric radius $R_M$ at which the pressure of the NS magnetic field [$\mu^2/(8\pi R_N^6)$, with $\mu$ the NS magnetic moment] balances the ram pressure of the inflowing matter ($\rho_w v_w^2$). In the case in which $R_M > R_a$, the magnetospheric radius is given by (Davies & Pringle 1981)

$$R_M = 3.3 \times 10^{10} M_{-6}^{-1/6} v_w^{-1/6} d_{10, d}^{1/3} \, \text{cm}. \quad (3)$$

Here we assumed a nonmagnetized spherically symmetric wind (Elsner & Lamb 1977), with density $\rho_w(R_M) \sim M_w/(4\pi a^2 v_w)$, a dipolar NS magnetic field with $\mu_{33} = \mu/10^{13} \, \text{G cm}^2$, and $M_{-6} = M_w/10^{-6} M_\odot \, \text{yr}^{-1}$. In the following sections we discuss the range of applicability of equation (3), and the regimes in which a different prescription for $R_M$ should be used.

3. The corotation radius $R_{co}$ at which the NS angular velocity equals the Keplerian angular velocity, i.e.,

$$R_{co} = 1.7 \times 10^{10} P_{3, s}^{2/3} \, \text{cm}. \quad (4)$$

Here $P_{s, s}$ is the NS spin period in units of $10^3$ s.

Changes in the relative position of these radii result into transitions across different regimes for the NS (Illarionov & Sunyaev 1975; Stella et al. 1986). Below we discuss these regimes individually, and provide a schematic representation of each regime in Figure 1. Being determined primarily by the spin of the NS, the corotation radius can change only over evolutionary timescales. Illarionov & Sunyaev (1975) summarize the different regimes experienced by a spinning down NS since its birth, from the initial radio pulsar (i.e., rotation powered) stage to the regime in which mass accretion onto the NS surface can take place. On the other hand, the accretion radius and magnetospheric radius depend on the wind parameters (see eqs. [1] and [3]), which can vary on a wide range of timescales (from hours to months). Therefore, variations in the wind parameters can cause the NS to undergo transitions across different regimes on comparably short timescales, thus opening the possibility of explaining the properties of some classes of highly variable X-ray sources through them. In particular, the transition from the accretion regime to the propeller regime (and vice versa), across the so-called centrifugal barrier, was identified as a likely mechanism responsible for the pronounced activity of Be X-ray pulsar transient systems (Stella et al. 1986). Below we summarize the different regimes of a magnetic rotating NS, subject to a varying stellar wind, with special attention to the condition under which a “magnetic” (as opposed to “centrifugal”) barrier inhibits accretion onto the NS (Harding & Leventhal 1992; Rutledge 2001; Mori & Ruderman 2003; Toropina et al. 2006, 2001). As is clear in the following, new motivation for investigating the magnetic inhibition of accretion comes from the discovery of magnetars, neutron stars with extremely high magnetic fields ($\sim 10^{14}$–$10^{15}$ G; Duncan & Thompson 1992).

### 3.1. Outside the Accretion Radius—the Magnetic Inhibition of Accretion: $R_M > R_a$

We consider here the case in which the magnetospheric radius is larger than the accretion radius. In systems with $R_M > R_a$ the mass flow from the companion star interacts directly with the NS magnetosphere without significant gravitational focusing, forming a bow shock at $R_M$ (Harding & Leventhal 1992; Toropina et al. 2001). A region of shocked gas surrounds the NS magnetosphere with density $\rho_{ps} \simeq 4\rho_w$ and velocity $v_{ps} \simeq v_w/4$ (the subscript “ps” stands for postshock). These are only rough estimates because the shock is very close to the magnetopause and it does not satisfy the standard Rankine-Hugoniot conditions (Toropina et al. 2001). At least in the front part of the shock, i.e., in the region around the stagnation point, the whole kinetic energy of the inflowing matter is converted into...
thermal energy, and the expected temperature of the heated gas is \( T \approx n_{mp} v_w^2 / (3k) \approx 4 \times 10^7 \, \text{v}_w^2 \, \text{K} \). Thus, the power released in this region is of order

\[
L_{\text{shock}} \approx \frac{\pi}{2} R_M^2 \rho_w v_w^3 = 4.7 \times 10^{29} R_{10}^2 \, \text{v}_w^2 \, \text{M}_{-6}^2 \, \text{erg s}^{-1}
\]  

\( (R_{10} \text{ is the magnetospheric radius in units of } 10^{10} \text{ cm}) \), and is mainly radiated in the X-ray band (Toropina et al. 2006). Below we distinguish two different regimes of magnetic inhibition of accretion.

### 3.1.1. The Super-Keplerian Magnetic Inhibition Regime: \( R_M > R_a, R_{\text{co}} \)

In this “super-Keplerian” magnetic inhibition regime the magnetospheric radius is larger than both the accretion and corotation radii \( (R_M > R_a, R_{\text{co}}) \). Matter that is shocked and halted close to \( R_M \) cannot proceed further inward, due to the rotational drag of the NS magnetosphere which is locally super-Keplerian. Since magnetospheric rotation is also supersonic, the interaction between the NS magnetic field and matter at \( R_M \) results in rotational energy dissipation and, thus, NS spin-down. In order to derive an upper limit on the contribution of this dissipation to the overall luminosity, we assume that the above interaction is anelastic (e.g., Perna et al. 2006), i.e., that matter at \( R_M \) is forced to corotate. This process releases energy at a rate

\[
L_{\text{shock}} \approx 2 \pi R_M^2 \rho_w v_w (R_M \Omega)^2 \approx 3.7 \times 10^{29} R_{10}^4 \, \text{M}_{-6}^2 \, a_{10}^{-2} \, \text{erg s}^{-1},
\]

which adds to the shock luminosity (eq. [5]).

### 3.1.2. The Sub-Keplerian Magnetic Inhibition Regime: \( R_a < R_M < R_{\text{co}} \)

If \( R_a < R_M < R_{\text{co}} \) the magnetospheric drag is sub-Keplerian and matter can penetrate the NS magnetosphere. In this sub-Keplerian magnetic inhibition regime, the boundary between the inflowing matter and the magnetosphere is subject to the Kelvin-Helmholtz instability (KHI; Harding & Leventhal 1992). The mass inflow rate across \( R_M \) resulting from the KHI is approximately (Burnard et al. 1983)

\[
\dot{M}_{\text{KH}} \approx 2 \pi R_M^2 \rho_w v_{\text{conv}} = 2 \pi R_M^2 \rho_p v_{\text{sh}} \eta_{\text{KH}} (\rho_i / \rho_e)^{1/2} (1 + \rho_i / \rho_e)^{-1},
\]

where \( \eta_{\text{KH}} \sim 0.1 \) is an efficiency factor, \( v_{\text{sh}} \) is the shear velocity, and \( \rho_i \) and \( \rho_e \) are the density inside and outside the magnetospheric boundary at \( R_M \), respectively. Close to the stagnation point, virtually all the kinetic energy of the wind matter is converted into thermal energy (see also §3.1.1), and the shear velocity is thus dominated by the magnetosphere’s rotation, such that \( v_{\text{sh}} = v_{\text{rot}} = 2 \pi P_e^{-1} \dot{R}_M \).

Away from this region, the tangential component of the wind velocity with respect to the NS magnetic field lines increases up to \( v_{\text{ps}} \) and, in general, the shear velocity and rate at which plasma enters the magnetosphere due to the KHI, depend on both \( v_{\text{ps}} \) and \( v_{\text{rot}} \). For the aims of this paper, we adopt \( v_{\text{sh}} = \max (v_{\text{ps}}, v_{\text{rot}}) \), and use mass conservation across the KHI unstable layer to estimate the density ratio at \( R_M \) (Burnard et al. 1983). If matter crossing the unstable layer of height \( h_t \) is rapidly brought into corotation with the NS magnetosphere and free falls onto the NS, mass conservation implies

\[
R_M^2 \rho_i v_{\text{conv}} \approx R_M h_t \rho_i v_{\text{rot}} (R_M).
\]

The height \( h_t \) of the unstable layer, where matter and magnetic field coexist, is mostly determined by the largest wavelength of the KHI unstable mode (Burnard et al. 1983). A detailed analysis of this instability is beyond the scope of this paper. In equation (8) we conservatively use \( h_t \approx R_M \), and discuss in Appendix A the effect of smaller values of \( h_t \). Therefore, accretion of matter at a rate \( \dot{M}_{\text{KH}} \) onto the NS is expected to release a luminosity of

\[
L_{\text{KH}} \approx 3.5 \times 10^{34} \eta_{\text{KH}} R_M^2 \, a_{10}^{-2} \, \text{M}_{-6} (\rho_i / \rho_e)^{1/2} (1 + \rho_i / \rho_e)^{-1} \, \text{erg s}^{-1}
\]

if \( v_{\text{sh}} = v_{\text{ps}} \), or

\[
L_{\text{KH}} \approx 8.8 \times 10^{34} \eta_{\text{KH}} P_{-5}^{-1} R_M^2 \, a_{10}^{-2} \, \text{M}_{-6} (\rho_i / \rho_e)^{1/2} (1 + \rho_i / \rho_e)^{-1} \, \text{erg s}^{-1}
\]

if \( v_{\text{sh}} = v_{\text{rot}} \). The values of \( \rho_i / \rho_e \) that we use in equations (9) and (10) are derived numerically from equation (8).

In the sub-Keplerian magnetic inhibition regime, plasma penetration inside the NS magnetosphere is sustained also by Bohm diffusion (Ikhsanov & Pustil’nik 1996). This diffusion, being dependent on the temperature of the plasma (rather than velocity), is highest in the region close to the stagnation point, where the shock slows down the inflowing plasma most efficiently (see also §3.1.1). In accordance with Ikhsanov et al. (2001) the maximum inflow rate allowed by Bohm diffusion is

\[
\dot{M}_{\text{eff}} \approx 2 \pi R_M^2 \rho_p v_M = 4.5 \times 10^7 \zeta^{3/2} \, \text{M}_{-6} \, \text{a}_{10}^{-2} \, \text{R}_{16}^{11/4} \, \text{M}_{-6}^{-1/2} \, \text{erg s}^{-1},
\]

where \( v_M = (D_{\text{eff}} / t_f)^{1/2} \) is the diffusion velocity, \( D_{\text{eff}} = [\zeta c k T_i (R_M)] / [16 e B (R_M)] \) the diffusion coefficient, \( T_i (R_M) \approx m_p v_w^2 / (3k) \) the post-shock ion temperature, \( t_f = [p_{\text{ps}} (2GM)^{3/2}]^{1/2} \) the free-fall time, \( \zeta \approx 0.1 \) an efficiency factor, and \( m_p \) the proton mass. In the above equation we also approximated the density outside \( R_M \) around the stagnation point, with \( \rho_{pM} \) (although this might be underestimate by a

\footnote{This can be easily seen by comparing \( v_s \) with \( \Omega R_M \) for the values of the parameters used in this section.}
factor of a few; see Toropina et al. 2001). Over the whole range of parameters relevant to this work, the diffusion-induced mass accretion rate is orders of magnitude smaller than that due to the KHI. Similarly, the contribution to the total luminosity resulting from the shock and anelastic drag at the magnetospheric boundary can be neglected in this regime.

3.2. Inside the Accretion Radius: \( R_M < R_a \)

3.2.1. The Supersonic Propeller Regime: \( R_\infty < R_M < R_a \)

Once \( R_M \) is inside the accretion radius, matter flowing from the companion star is shocked adiabatically at \( R_a \) and halted at the NS magnetosphere. In the region between \( R_a \) and \( R_M \), this matter redistributes itself into an approximately spherical configuration (resembling an "atmosphere"), whose shape and properties are determined by the interaction between matter and NS magnetic field at \( R_M \). This scenario was considered previously by Davies et al. (1979) and Davies & Pringle (1981) and here we follow their treatment. These authors demonstrated that hydrostatic equilibrium ensues when radiative losses inside \( R_a \) are negligible (we discuss this approximation in Appendices B and C) and the atmosphere is stationary on dynamical timescales. Assuming a polytropic law of the form \( p \propto \rho^{1+\frac{1}{n}} \), the pressure and density of this atmosphere are

\[
p(R) = \rho(R) = \rho_p v_p^2 \left\{ 1 + \frac{1}{(1 + n)}[8 R_a / R]^{n+1} \right\},
\]

\[
\rho(R) = \rho_p \left\{ 1 + \frac{1}{(1 + n)}[8 R_a / R] \right\}^n.
\]

The value of the polytropic index \( n \) depends on the conditions at the inner boundary of the atmosphere, and in particular on the rate at which energy is deposited there.

When the rotational velocity of the NS magnetosphere at \( R_M \) is supersonic (see also § 3.1.1), the interaction with matter in the atmosphere leads to dissipation of some of the star’s rotational energy and thus spin-down. In the supersonic propeller regime, Davies & Pringle (1981) showed that turbulent motions are generated at \( R_M \) which convect this energy up through the atmosphere, until it is lost at its outer boundary. In this case \( n = 1/2 \). Accordingly, taking into account the structure of the surrounding atmosphere, the magnetospheric radius is given by

\[
R_M^{-\frac{6}{3}}[1 + 16 R_a / (3 R_M)]^{-3/2} = \frac{1}{2} \dot{M}_a a^{-2} \mu^{-2} v_w.
\]

This can be approximated by

\[
R_M \approx 2.3 \times 10^{10} d_{10}^{4/9} M_6^{-2/9} v_8^{4/9} \mu_{33}^{4/9} \text{ cm.}
\]

Matter that is shocked at \( \approx R_a \) reaches the magnetospheric boundary at \( R_M \), where the interaction with the NS magnetic field draws energy from NS rotation (see also § 3.1.1). According to Davies and Pringle (1981) this contributes

\[
L_{\text{ad}} = 2 \pi R_M^2 \rho(R) v_p^3 (R_M) \approx 5.4 \times 10^{31} \dot{M}_a a^{-2} v_8^{-1} R_{M10}^{1/2}[1 + 16 R_{a10} / (3 R_{M10})]^{1/2} \text{ erg s}^{-1}
\]

(16) to the total luminosity. In the above equation \( R_{a10} = 10^{-10} R_a \) and \( c_s(R_M) = v_a(R_M) = (2GM_{\text{NS}} / R_M)^{1/2} \) (Davies & Pringle 1981). In the supersonic propeller regime the energy released through the shock at \( R_a \),

\[
L_{\text{shock}} = \frac{9}{32} \pi R_M^2 \rho v_w \approx 2.6 \times 10^{26} R_{a10}^2 R_{M10}^{2/3} \dot{M}_a a^{-2} \mu^{-2} \text{ erg s}^{-1},
\]

(17)
is negligible.

3.2.2. The Subsonic Propeller Regime: \( R_M < R_a, R_\infty, \dot{M}_a < \dot{M}_\text{lim} \)

The break down of the supersonic propeller regime occurs when \( R_M < R_\infty \), i.e., when the magnetosphere rotation is no longer supersonic with respect to the surrounding material. The structure of the atmosphere changes and the transition to the subsonic propeller regime takes place. Since the rotation of the magnetosphere is subsonic, the atmosphere is roughly adiabatic \( (n = 3/2) \), and the magnetospheric radius is given by (Davies & Pringle 1981)

\[
R_M^{-\frac{6}{3}}[1 + 16 R_a / (5 R_M)]^{-5/2} = \frac{1}{2} \dot{M}_a a^{-2} \mu^{-2} v_w.
\]

This can be approximated by

\[
R_M \approx 2 \times 10^{10} d_{10}^{4/7} M_6^{-2/7} v_8^{8/7} \mu_{33}^{4/7} \text{ cm.}
\]

(19)

In the subsonic propeller regime, the centrifugal barrier does not operate because \( R_M < R_\infty \), but the energy input at the base of the atmosphere (due to NS rotational energy dissipation) is still too high for matter to penetrate the magnetosphere at a rate \( \dot{M}_\text{capt} \) (Davies & Pringle 1981). Nevertheless, a fraction of the matter inflow at \( R_a \) is expected to accrete onto the NS, due to the KHI and Bohm diffusion.\(^8\)

\(^8\) To our knowledge this is the first application of the KHI to the subsonic propeller regime.
Based on the discussion in § 3.1.2, we estimate the accretion luminosity of this matter by using equations (7) and (11) (we approximate here the surface of interaction between matter and magnetic field with $4\pi R_{K}^{2}$). This gives

$$L_{\text{diff}} \simeq G M_{\text{NS}} \dot{M}_{\text{diff}} / R_{\text{NS}} = 4.5 \times 10^{30} \dot{M}_{-6} a_{10 d}^{-2} R_{M_{10}}^{9/4} M_{-6} a_{10 d}^{-2} v_{k}^{-1} [1 + 16 R_{a10}/(5 R_{M10})]^{3/2} \text{ erg s}^{-1}$$

(20)

and

$$L_{\text{KH}} \simeq G M_{\text{NS}} \dot{M}_{\text{KH}} / R_{\text{NS}}$$

$$= 1.8 \times 10^{35} \eta_{K} P_{6}^{-3} R_{M_{10}}^{5/2} M_{-6} a_{10 d}^{-2} v_{k}^{-1} [1 + 16 R_{a10}/(5 R_{M10})]^{3/2} (\rho_{i}/\rho_{e})^{1/2} (1 + \rho_{i}/\rho_{e})^{-1} \text{ erg s}^{-1}, \quad (21)$$

for the accretion luminosity arising from matter entering the magnetosphere through Bohm diffusion and KHI, respectively. For the range of parameters of interest here, equations (5), (20), and (21) show that $L_{\text{KH}}$ dominates. The rotational energy dissipation at $R_{M}$ (see § 3.2.1) gives a small contribution with respect to equation (21) (Davies & Pringle 1981):

$$L_{\text{ad}} = 2 \pi R_{M}^{5} \rho (R_{M}) \Omega^{3} = 2.2 \times 10^{30} P_{6}^{-3} R_{M_{10}}^{5/2} M_{-6} a_{10 d}^{-2} [1 + 16 R_{a10}/(5 R_{M10})]^{3/2} \text{ erg s}^{-1}.$$ 

(22)

The subsonic propeller regime applies until the critical accretion rate

$$\dot{M}_{\text{sub}} = 2.8 \times 10^{30} P_{6}^{-3} a_{10 d}^{2} v_{k}^{4} R_{M_{10}}^{5/2} [1 + 16 R_{a10}/(5 R_{M10})]^{-3/2} \text{ erg s}^{-1}$$

(23)

is reached, at which the gas radiative cooling (bremsstrahlung) completely damps convective motions inside the atmosphere (see Appendix C). If this cooling takes place, direct accretion at a rate $\dot{M}_{\text{capt}}$ onto the NS surface is permitted.

3.2.3. The Direct Accretion Regime: $R_{M} < R_{a}$, $R_{co}$, $M_{w} > M_{\text{sub}}$

If $R_{M} < R_{a}$ and matter outside the magnetosphere cools efficiently, accretion onto the NS takes place at the same rate $\dot{M}_{\text{capt}}$ (see eq. [2]) at which it flows toward the magnetosphere. The corresponding luminosity is

$$L_{\text{acc}} = G M_{\text{NS}} \dot{M}_{\text{capt}} / R_{\text{NS}} = 2 \times 10^{35} R_{M_{10}}^{2} a_{10 d}^{4} v_{k}^{-4} \text{ erg s}^{-1} \simeq 2 \times 10^{35} M_{15} \text{ erg s}^{-1},$$

(24)

where $M_{15} = \dot{M}_{\text{capt}}/10^{15} \text{ g s}^{-1}$. This is the standard accretion regime; the system achieves the highest mass-to-luminosity conversion efficiency.

4. TRANSITIONS AND PATHS ACROSS DIFFERENT REGIMES

We explore here the conditions under which transitions across different regimes take place. As emphasized in § 3, these transitions occur when the relative positions of $R_{M}$, $R_{a}$, and $R_{co}$ change; we concentrate here on transitions that occur in response to variations in the stellar wind parameters. In the following, since $R_{M}$ depends only weakly on the orbital period and the total mass of the system, we fix $a_{10 d} = 1$ (we explain this choice in § 5), and investigate variations in the other four parameters: $\mu_{33}$, $P_{6}$, $v_{k}$, and $M_{-6}$.

The equations that define the conditions between different regimes are

$$R_{M} > R_{a} \Rightarrow \dot{M}_{-6} \leq 0.45 \mu_{33}^{2} v_{k}^{11} a_{10 d}^{2}$$

(25)

(or, equivalently, $M_{15} \leq 0.6 \mu_{33}^{2} v_{k}^{11} a_{10 d}^{2}$; see eqs. [1], [3], and [2]) for the magnetic barrier; and

$$R_{M} > R_{co} \Rightarrow P_{6} \leq 2.6 \dot{M}_{-6}^{-1/4} v_{k}^{-1/4} a_{10 d}^{1/2} \mu_{33}^{1/2}$$

(26)

(see eq. [3]) if $R_{M} > R_{a}$, or

$$R_{M} > R_{co} \Rightarrow P_{6} \leq 1.8 a_{10 d}^{2/3} \dot{M}_{-6}^{-1/3} v_{k}^{2/3} \mu_{33}^{2/3}$$

(27)

(see eq. [15]) if $R_{M} < R_{a}$ for the centrifugal barrier.

The equation that defines the transition from the subsonic propeller to the direct accretion regime is

$$P_{6} \geq 4.5 \dot{M}_{-6}^{-15/21} a_{10 d}^{30/21} v_{k}^{60/21} 16^{1/21} \mu_{33}^{16/21} \simeq 5.5 \dot{M}_{15}^{15/21} a_{10 d}^{16/21} \mu_{33}^{16/21} .$$

(28)

In Figure 2 the above equations are represented as lines separating different regimes. In panels a and b we fixed $\mu_{33}$ and $v_{k}$ and investigated the different regimes in the $P_{6}$-$\dot{M}_{-6}$ plane. In panels c, d, e, and f, instead, $P_{6}$, $\mu_{33}$ were fixed and the relevant regimes shown in the $v_{k}$-$\dot{M}_{-6}$ plane. Below we summarize the different regimes that a system attains in response to variations of $M_{w}$, in the different panels.

Figure 2a shows that, for a wind velocity of $v_{k} = 1.2$, a strongly magnetized NS ($\mu_{33} = 0.8$) can undergo a transition between the super-Keplerian and sub-Keplerian magnetic inhibition regimes in response to changes in the mass-loss rate only for typical spin periods $\gtrsim 2000$ s. When the mass-loss rate reaches $\dot{M}_{-6} \sim 2.4$, the direct accretion or subsonic propeller regime sets in, depending on whether the spin period is longer or shorter than $\sim 3700$ s. For spin periods $\lesssim 420$ s the direct accretion regime is not attained for the interval of

9 Here we used eqs. (15) and (19) for the magnetospheric radius in the supersonic and subsonic propeller regimes, respectively.
mass-loss rates considered in Figure 2 and only transitions between the super-Keplerian magnetic inhibition and supersonic propeller regimes are expected.

For lower magnetic fields ($\mu_{33} = 0.01$), Figure 2b shows that only the supersonic propeller, subsonic propeller, and direct accretion regimes can be attained. For spin periods in the range $\approx 60$–$230$ s transitions can occur between all these three regimes, whereas systems with spin periods longer than $\approx 1300$ s and shorter than $\approx 20$ s are expected to be in the direct accretion and the supersonic propeller regimes, respectively. For $\mu_{33} = 0.01$ transitions to the super-Keplerian and sub-Keplerian magnetic inhibition regimes cannot take place because the magnetospheric radius is too small to exceed the accretion radius.

In Figure 2c ($\mu_{33} = 0.8$ and $P_{33} = 1$), transitions can occur virtually between all the regimes described in $\S$ 3. In particular, for $v_8$ in the range $0.9$–$1.5$ transitions are expected to take place between the super-Keplerian magnetic inhibition, the supersonic and subsonic propeller, and the direct accretion regimes, as the mass-loss rate increases from $\dot{M}_{\odot} \sim 0.1$ to $\dot{M}_{\odot} \sim 100$. For velocities $v_8 < 0.9$ transitions can occur only between the supersonic propeller, the subsonic propeller, and the direct accretion regimes, whereas transitions to the super-Keplerian and sub-Keplerian magnetic inhibition regimes are impeded by the fact that the accretion radius cannot be overtaken by $R_M$. On the contrary, for $v_8 > 1.5$ the magnetospheric radius is located beyond the accretion radius for any considered value of the
mass-loss rate, and thus transitions can take place only between the super-Keplerian and sub-Keplerian magnetic inhibition regimes. Similar considerations apply to Figures 2a, 2e, and 2f.

For a system with $\mu_{33} = 0.01$ and $P_{3} = 1$ (Fig. 2a), the super-Keplerian magnetic inhibition regime never occurs, because the magnetic field is too low and $R_{M} < R_{\text{co}}$ for any $0.1 < M_{-6} < 100$. Instead, the sub-Keplerian magnetic inhibition regime can be attained for high wind velocities ($v_{\infty} = 2$), because $R_{\infty} \propto v_{\infty}^{-2}$.

In Figure 2e, for $\mu_{33} = 0.8$ and $P_{3} = 0.1$, the magnetospheric radius is larger than the corotation radius for the entire range spanned by $M_{-6}$. Thus, accretion onto the NS does not take place (note that in the region below the dashed line accretion can occur even if $R_{M} \gtrsim R_{\text{co}}$; Davies & Pringle 1981).\(^{10}\)

Finally, in Figure 2f we show the transitions for a system with $\mu_{33} = 0.01$ and $P_{3} = 0.1$ s. In this case all the regimes described in § 3 are present in the figure, similar to the case of Figure 2c. However, the region corresponding to the subsonic propeller regime is larger, such that there is only a modest range of (fixed) velocities for which mass-loss rate variations in our chosen range ($0 < M_{-6} < 100$) can cause transitions through all regimes, from the super-Keplerian magnetic inhibition regime to the direct accretion regime. As we discuss below, this has important consequences for the expected luminosity variations.

We now compute the luminosity swings for some of the examples discussed above in a fashion similar to what was done in the context of centrifugally inhibited accretion in NS X-ray transients (Corbet 1996; Campana et al. 1998). Figure 3a applies to a system with $\mu_{33} = 0.8$, $P_{3} = 1$, and $v_{\infty} = 1.2$ (this corresponds to the $P_{3} = 1$ case of Fig. 2a). The bottom panel of this figure shows that, for $0.1 < M_{-6} < 100$, the magnetospheric radius crosses both the centrifugal ($R_{\text{co}}$) and magnetic ($R_{\nu}$) barriers. Correspondingly, the system moves from the super-Keplerian magnetic inhibition regime to the supersonic and subsonic propeller regimes, and, finally, to the direct accretion regime, giving rise to a six-decade luminosity swing from $\sim 10^{33}$ to $\sim 10^{37}$ erg s$^{-1}$. We note that a large part of this swing (about five decades) is attained across the transitions from the super-Keplerian magnetic inhibition to the direct accretion regimes, which take a mere factor of $\sim 5$ variation of $M_{\nu}$.

In the presence of a standard NS magnetic field ($10^{12}$ G), Figure 3b shows that such abrupt luminosity jumps are not expected for a very slowly rotating (1000 s) NS (the other system parameters are the same as those of Fig. 3a), since the magnetospheric radius is smaller than both $R_{\infty}$ and $R_{\text{co}}$ for any reasonable value of $M_{\nu}$. Therefore, the direct accretion regime applies, with the the luminosity proportional to $M_{\nu}$.

In Figure 4 we show the transitions for a system with $\mu_{33} = 0.01$ and $P_{3} = 0.1$. The wind velocity is $v_{\infty} = 1.2$ in Figure 4a, and $v_{\infty} = 2.2$ in Figure 4b (see also Fig. 2f). These two figures show that, for submagnetar fields, a 100 s spinning NS can undergo a transition across the magnetic barrier (besides the centrifugal barrier), for suitable parameters (a high wind velocity in the case at hand). Such transitions take place over a more extended interval of mass-loss rates. For instance, Figure 4b shows that an increase by a factor $\sim 100$ in the mass-loss rate is required, in this case, to achieve a factor $\sim 10^{5}$ luminosity swing comparable with the magnetar case of Figure 3a.

In order to illustrate further the role of the magnetic field, spin period, and wind velocity, we show in Figure 5 the way in which the transitions across regimes take place, by holding two of the above variables fixed and stepping the third variable. Figure 5a shows the effect of increasing the value of $v_{\infty}$ from 1 to 1.8 (in turn resulting in a decrease of the accretion radius), in a system with $\mu_{33} = 0.8$ and $P_{3} = 1$.

The behavior of the luminosity changes mainly because a different set of regimes is involved in each case. For $v_{\infty} = 1$ (solid line) the system passes through the super-Keplerian magnetic inhibition regime and the supersonic and subsonic propeller regimes, finally reaching the direct accretion regime at $M_{-6} \lesssim 6$. In case $v_{\infty} = 1.2$ (dotted line), the transition to the supersonic propeller shifts toward higher mass-loss rates, such that super-Keplerian magnetic inhibition applies up to $M_{-6} < 2$. Further increasing the wind velocity to $v_{\infty} = 1.4$ (dashed line), the system first undergoes a transition to the subsonic propeller regime at $M_{-6} \lesssim 10$, bypassing the supersonic propeller regime (this is because for $M_{-6} \gtrsim 10$ the magnetospheric radius defined by eq. [14] is smaller than the corotation radius). As the mass-loss rate increases further, the direct accretion regime sets in for $M_{-6} \lesssim 20$. In the case $v_{\infty} = 1.8$ (dot-dashed line), the corotation radius exceeds the accretion radius. The system is thus in the super-Keplerian magnetic inhibition regime for $M_{-6} < 20$, while for higher mass-loss rates the sub-Keplerian magnetic inhibition regime applies and the luminosity is dominated by accretion through the KHI. A transition to the direct accretion regime is expected for mass-loss rates $M_{-6} > 100$.

\(^{10}\) This is because, in this region, the mass flow rate is so high that eq. (B2) is violated (the supersonic propeller is no longer self-consistent), and convective motions are damped by radiative cooling.
For a wide range of wind parameters, accretion can take place, and a high persistent luminosity is released (the case corotation radius for the whole range of mass-loss rates spanned in the figure, and the system is always in the direct accretion regime. In the case $\mu_{33} = 0.1$ (dotted line), the system is in the subsonic propeller regime for $M_{-6} < 1$, while for higher mass inflow rates the direct accretion regime applies. For $\mu_{33} = 0.5$ (dashed line) the luminosity behavior becomes more complex, as the system goes through the super-Keplerian magnetic inhibition regime, the supersonic and subsonic propeller regimes, and eventually reaches the direct accretion regime at $M_{-6} \simeq 8$. For $\mu_{33} = 0.8$ the same sequence of transitions applies, with the entire luminosity swing taking place over a smaller interval of mass-loss rates.

Finally, in Figure 5c we show the effects of increasing the spin period in a system with $\mu_{33} = 0.1$ and $v_b = 1.2$. For the lowest spin period considered here ($P_{s3} = 0.01$, solid line), the system remains in the supersonic propeller regime for the whole range spanned by $M_{-6}$. In the case $P_{s3} = 0.5$ (dotted line) transitions occur from the supersonic propeller regime ($M_{-6} < 0.7$), to the subsonic propeller regime ($0.7 < M_{-6} < 4$), and then to the direct accretion regime ($M_{-6} > 4$). By further increasing $P_{s3}$ to a value of 1 (dashed line), the system goes through the subsonic propeller and the direct accretion regimes, while the supersonic propeller regime does not occur due to the longer spin period as compared to the previous case. For $P_{s3} = 6$ (dot-dashed line) the spin period is so high that the direct accretion regime applies for any reasonable value of $M_{-6}$.

The above results show that over a range of values of the key parameters $\mu_{33}$, $P_{s3}$, and $v_b$, large luminosity swings can be achieved with comparatively modest changes in the mass-loss rate, as the NS undergoes transitions from one regime to another. More generally, these transitions result from changes in the relative position of the accretion, magnetospheric, and corotation radii, reflecting short-term variations of the wind velocity and mass-loss rate, the only parameters that can vary on shorter timescales than secular. Therefore, transitions between different regimes take place once the source parameters are such that the NS straddles the centrifugal barrier (i.e., $R_M \simeq R_{co}$ when $R_a > R_{co}$) or the magnetic barrier (i.e., $R_M \simeq R_a$ when $R_a < R_{co}$) or both (i.e., $R_M \simeq R_{co} \simeq R_a$). The centrifugal barrier applies to relatively short spin periods. It is well known that the longer the spin period of a transient NS, the higher its magnetic field must be for the centrifugal barrier to operate (see eqs. [26] and [27]). In particular, for periods of hundreds of seconds, or longer, magnetic field strengths of $\geq 10^{13} - 10^{14}$ G are required. On the other hand, we have shown that NSs with even longer spin periods and magnetar-like fields are expected to undergo transitions across the magnetic barrier and thus are expected to have an inherently different “switch-off” mechanism than short spin period systems. A necessary condition for this is that $R_{co} > R_a$, which translates into $P_{s3} \geq 3v_b^{-1}$. The magnetic and centrifugal barriers set in (nearly) simultaneously (i.e., $R_a \simeq R_{co} \simeq R_M$) for $\mu_{33} \approx 0.3P_{s3}^{1/6}M_{15}^{1/2}$.

Taking into account all the examples discussed in this section, we conclude the following:

1. Long spin period systems ($P_{s3} \geq 1$) require magnetar-like magnetic fields ($\mu_{33} \geq 0.1$) in order for a large luminosity swing ($\sim 10^5$) to arise from modest variations in the wind parameters (e.g., a factor $\sim 5$ in $M_{-6}$). These luminosity swings might result from transition across different regimes through both the centrifugal and magnetic barriers.
2. Shorter spin period systems ($P_{s3} \ll 1$) must possess lower magnetic fields ($\mu_{33} \ll 0.1$) for similar transitions to take place. Somewhat larger variations in the wind parameters are required in order to achieve similar luminosity swings to those of the long-period case, and transitions between different regimes occur in most cases through the centrifugal barrier.
3. Few or no transitions are expected for systems with either high magnetic fields and short spin periods, or systems with lower magnetic fields and long spin periods. In the first case the centrifugal barrier halts the inflowing matter at $R_N$ and accretion does not take place; such systems might thus be observed only at very low (X-ray) luminosity levels ($\geq 10^{32}-10^{33}$ erg s$^{-1}$). In the second case, $R_M < R_{co}$ for a wide range of wind parameters, accretion can take place, and a high persistent luminosity is released ($\geq 10^{32}-10^{33}$ erg s$^{-1}$).

5. APPLICATION TO SFXT SOURCES

In this section we propose that transitions across different regimes caused by relatively mild variations of the wind parameters are responsible for the outbursts of SFXTs. In consideration of the wide range of spin periods inferred for SFXTs, the outbursts of these sources are expected to result from the opening of the magnetic barrier in very long spin period systems and the centrifugal barrier in all other systems (see also Greben'ev & Sunyaev 2007; Sidoli et al. 2007). More crucially, we conclude that slowly spinning SFXTs should host magnetars, independent of which of the two mechanisms is responsible for their outbursts.

11 This point was already noted in Stella et al. (1986).
As a case study we consider IGR J17544–2619 (Sunyaev et al. 2003), a SFXT observed by Chandra during a complex transition to and from a ~1 hr long outburst, yielding the first detailed characterization of a SFXT light curve. The spin period of IGR J17544–2619 is presently unknown. In’t Zand (2005) showed that four different stages, with very different luminosity levels, could be singled out during the Chandra observation: (1) a quiescent state with \( L_X \approx 2 \times 10^{32} \) erg s\(^{-1}\), (2) a rise stage with \( L_X \approx 1.5 \times 10^{34} \) erg s\(^{-1}\), (3) the outburst peak with \( L_X \approx 4 \times 10^{37} \) erg s\(^{-1}\), and (4) a postoutburst stage (or “tail”) with \( L_X \approx 2 \times 10^{36} \) erg s\(^{-1}\) (see Fig. 6a; these luminosities are for a source distance of ~3.6 kpc; Rahoui et al. 2008). The maximum luminosity swing observed across these stages was a factor of \( \gtrsim 6.5 \times 10^4 \).

Motivated by the evidence for >1000 s periodicities in XTE J1739–302 and IGR J16479–4514 (see § 2), we discuss first the possibility that IGR J17544–2619 contains a very slowly spinning NS. We use \( \mu_{33} = 1, P_{33} = 1.3, \) and \( v_8 = 1.4 \), and show in Figure 6b the different regimes experienced by such a NS as a function of the mass-loss rate. For \( M_{\infty} < 20 \) the above values values give \( R_M > R_{\infty} \) and \( R_M > R_{\infty} \), such that super-Keplerian magnetic inhibition of accretion applies. The expected luminosity in this regime, \( \sim 10^{31} \) erg s\(^{-1}\), is likely outshined by the X-ray luminosity of the supergiant star (the companion star’s luminosity is not shown in Fig. 6, but it is typically of order \( \sim 10^{32} \) erg s\(^{-1}\); Cassinelli et al. 1981; Berghoefer et al. 1997). We conclude that the lowest emission state (quiescence) of IGR J17544–2619 can be explained in this way, with the companion star dominating the high-energy luminosity (In’t Zand 2005). The rise stage is in good agreement with the sub-Keplerian magnetic inhibition regime, where the luminosity (\( \sim 10^{34} \) erg s\(^{-1}\)) is dominated by accretion of matter onto the NS due to the KHI. The uncertainty in the value of \( h \) translates into an upper limit on the luminosity in this regime which is a factor of \( \sim 10^1 \) higher than that given above (see § 3 and Appendix A).

During the outburst peak the direct accretion regime must apply at a mass-loss rate of \( M_{\infty} = 500 \). In this interpretation direct accretion must also be at work in the outburst tail at \( M_{\infty} \sim 3 \), where a slight decrease in \( M_{\infty} \) would cause the magnetic barrier to close and the source to return to quiescence. According to this interpretation, if IGR J17544–2619 has a spin period of >1000 s, then it must host a magnetar.

Figure 6c shows an alternative interpretation of the IGR J17544–2619 light curve, where we fixed \( \mu_{33} = 0.08, P_{33} = 0.4, \) and \( v_8 = 1 \). For this somewhat faster spin (and lower magnetic field), the luminosity variation is mainly driven by a transition across the centrifugal barrier (as opposed to the magnetic barrier). In this case, the quiescent state corresponds to the supersonic propeller regime (\( M_{\infty} < 0.6 \)) and the rise stage to the subsonic propeller regime (\( 0.6 < M_{\infty} < 2 \)), while both the peak of the outburst and the tail take place in the direct accretion regime at \( M_{\infty} = 200 \) and \( M_{\infty} = 10 \), respectively.

Assuming an even faster NS spin period for IGR J17544–2619, a weaker magnetic field would be required. In Figure 6d we show the results obtained by adopting \( \mu_{33} = 0.001, P_{33} = 0.01, \) and \( v_8 = 2 \). The \( \sim 10^{30} \) erg s\(^{-1}\) luminosity in the subsonic propeller regime compares well with the luminosity in the rise stage. However, the luminosity of the supersonic propeller regime is now significantly higher than the quiescence luminosity of \( \sim 10^{32} \) erg s\(^{-1}\) (this is consequence of the higher value of \( M_{\infty} \) for which the supersonic propeller regime is attained in this interpretation). We note that the whole luminosity swing takes place for a wider range of mass-loss rates, and the outburst peak luminosity requires \( M_{\infty} \sim 3000 \), an extremely high value even for an OB supergiant.

Interpreting the properties of IGR J17544–2619 in terms of a NS with a spin periods <100 s is more difficult. For instance, for the subsonic propeller regime to set in, the mass-loss rate corresponding to the transition across \( R_M = R_{\infty} \) must be lower than the limit fixed by equation (23). If instead the transition takes place at higher mass-loss rate, the system goes directly from the supersonic propeller to the direct accretion regime (or vice versa), bypassing the subsonic propeller regime: therefore, the rise stage would remain unexplained. Since fast-rotating NSs require lower magnetic fields for direct accretion to take place while in outburst, equation (23) is satisfied only for very high wind velocities (\( v_8 > 2 \)). On the other hand, an increase by a factor of \( \sim 2 \) in the wind velocity (with respect to the longer spin period solutions) would give a substantially lower \( M_{\infty,\text{opt}} \), such that the subsonic and the direct accretion regime luminosities fall shortwards of the observed regime values (unless unrealistically high mass-loss rates are considered).

Based on the above discussion, we conclude that IGR J17544–2619 likely hosts a slowly rotating NS, with spin period >100 s. Whether the magnetic barrier or the subsonic barriers sets in, causing inhibition of accretion away from the outbursts, will depend on whether the spin period is longer or shorter than \( \sim 1000 \) s. We note that IGR J16418–4532, a ~1240 s pulsating source with a 3.7 day orbital period, displayed short-duration flares similar to those of SFXTs and thus might be considered a candidate for hosting an accreting NS with a magnetar-like field. However, there is no clear evidence yet that IGR J16418–4532 is a transient source, since the very low state revealed with SWIFT might well be due to an eclipse (Tomskick et al. 2004; Walter et al. 2006; Corbet et al. 2006).

As another example we discuss the case of IGR J16465–4507, a SFXT with a spin period of 228 s. The luminosity behavior of this source is still poorly known. An outburst at \( 5 \times 10^{36} \) erg s\(^{-1}\) was observed with INTEGRAL (assuming a distance of 12.5 kpc; Lutovinov et al. 2005; Smith 2004), which did not detect the source before the outburst down to a level of \( 5 \times 10^{35} \) erg s\(^{-1}\). About a week later,
state, the rise state, the tail, and the peak of the outburst, respectively. According to In’t Zand (2005) the count rates on the 
XMM-Newton revealed the source at \(5 \times 10^{34} \text{ erg s}^{-1}\) and discovered the 228 s pulsations (Lutovinov et al. 2005; Zurita Heras & Walter 2004). If the direct accretion regime applied all the way to the lowest luminosity level observed so far, then an upper limit of \(\dot{P}_{33} = 0.001\) would be obtained by imposing that the NS did not enter the subsonic propeller regime (see eq. [28]). On the other hand, if the luminosity measured by XMM-Newton signalled that the source entered the subsonic propeller regime, while direct accretion occurred only during the outburst detected by INTEGRAL, then a considerably higher magnetic field of \(\mu_{33} = 0.07\) would be required.

The above discussion emphasizes the importance of determining, through extended high sensitivity observations, the luminosity at which transitions between different source states occur, in particular the lowest luminosity level for which direct accretion is still at work. In combination with the NS spin, this can be used to infer the NS magnetic field. Alternatively, accretion might take place unimpeded at all luminosity levels of SFXTs, a possibility which requires a very clumpy wind as envisaged in other scenarios (Negueruela et al. 2008; Zurita Heras et al. 2007). In this case the NS magnetic field can be considerably lower than discussed here. More extensive studies of these sources (and, by analogy, other SFXTs) are clearly required.

6. DISCUSSION

If the centrifugal barrier operates in IGR J17544—2619, then the activity of this source (and by extension that of SFXTs with similar properties) should parallel that of long spin period X-ray pulsar transients with Be star companions (Stella et al. 1986). One crucial difference between the two classes, namely, the duration of the outbursts, might well result from the presence in Be systems of an accretion disk mediating the flow of matter outside the NS magnetosphere. In fact, while in Be systems the star’s slow equatorial wind has enough angular momentum to form such a disk, the supergiant’s wind in SFXTs is fast and possesses only a little angular momentum relative to the NS. In the absence of a disk, variations in the wind parameters take effect on dynamical timescales, whereas in the presence of a disk they are smoothed out over viscous timescales.

If the spin period is sufficiently long in IGR J17544—2619 and other SFXTs, the onset of the magnetic barrier will inhibit accretion. While steady magnetic inhibition of accretion is familiar, e.g., from the Earth magnetosphere–solar wind interaction, transitions in and out of this regime have not yet been observed, to the best of our knowledge. Therefore, very long spin period SFXTs might provide the first opportunity to study transitions across such a magnetic barrier. Irrespective of whether the centrifugal or magnetic barrier operates in IGR J17544—2619, a long spin period would imply a high magnetic field, comparable to those inferred for magnetar candidates.

Scenarios involving magnetars with spin periods well above 1000 s have been considered in several studies (Rutledge 2001; Mori & Ruderman 2003; Toropina et al. 2001; Liu & Yan 2006; Zhang et al. 2004). Moreover, a few known accreting X-ray pulsars with unusually long spin periods have already been proposed as magnetar candidates (see, e.g., Ikhsanov 2007 and references therein). However, these sources display properties different from those of SFXTs: some are persistent sources, others display week- to month-long outbursts; the high spin-up measured in IGR J16358—4726 testifies that the accretion flow is likely mediated by a disk (Patel et al. 2007). Therefore, some of the features of the model discussed in this paper would not be applicable to these sources.

In the context of wind-fed HMXBs, Zhang et al. (2004) pointed out that magnetars might be hosted in binary systems with relatively short orbits (~1–100 days) and long pulse periods (\(10^7 \text{–} 10^8\) s). By using evolutionary calculations, these authors showed that magnetars would be easily spun-down to such long spin periods by the interaction with the wind of the companion star, in less than \(10^8\) yr. The systems we considered in this paper would likely result from a similar evolutionary path.

Once a system approaches the spin period that is required for direct accretion to occur, the short-timescale (~ hours) erratic variations that characterize a supergiant’s wind will cause transitions across different NS regimes. Insofar as the average wind properties evolve only secularly, the NS spin will then remain locked around such a period, alternating spin-up intervals during accretion and spin-down intervals during quiescence, when accretion is inhibited.

The relevant spin-up timescale during accretion intervals is approximately (see, e.g., Frank et al. 2002)

\[
\tau_{su} = -\frac{P_{spin}}{\dot{P}_{spin}} = \Omega I / (\dot{M}_{capl}) \approx 8 \times 10^2 I_{45} v_8^3 P_{10}\dot{P}_{3}\dot{M}_{17}^{-1} \text{ yr},
\]  
(29)
where \( l = 2\pi R_c^2/(4P_{\text{orb}}) \) is the specific angular momentum of wind matter at the accretion radius, \( I = 10^{45} I_{45} \) g cm\(^2\) is the NS moment of inertia, and \( M_{17} = M_{\text{cep}}/10^{17} \) g s\(^{-1}\) (this corresponds to an outburst luminosity of \( \sim 10^{37} \) erg s\(^{-1}\)).

A rough estimate of the spin-down timescale is obtained by assuming that most of the quiescence luminosity draws from the rotational energy of the NS (Davies & Pringle 1981); this gives

\[
\tau_{\text{dd}} = P_{\text{spin}}/l = \frac{L_X}{1.3 \times 10^2 I_{45} L_{31}^{-1} P_{31}^{-2}} \text{yr}, \tag{30}
\]

where \( L_{31} \) is the luminosity produced by the interaction between the magnetosphere and the wind in units of \( 10^{31} \) erg s\(^{-1}\). It is apparent that, for magnetar-like fields and long spin periods, the above timescales are much shorter than the lifetime of the supergiant’s strong wind phase. Therefore, secular changes in the wind parameters and/or the NS magnetic field will be easily tracked by the NS spin. Moreover, since spin-up and spin-down take place on comparable timescales, it is to be expected that spin-up intervals during outbursts are compensated for by spin-down during quiescent intervals of comparable duration. In other words, the accretion state of long spin period SFXTs would be expected to have a duty cycle of order \( \sim 0.5 \) or higher. As we discussed in § 2, evidence for high values of the duty cycle is gradually emerging from high-sensitivity observations of SFXTs.

On the other hand, spin-down may also occur during the accretion intervals as a result of velocity and density gradients in the supergiant’s wind that lead to temporary reversals of the angular momentum of the captured wind relative to the NS (note that persistent wind-accreting X-ray pulsars in HMXBs have long been known to alternate spin-up and spin-down intervals; see, e.g., Henrichs 1983). This would tend to favor spin-down, such that the NS spin might gradually evolve longward of the spin period that is required to inhibit direct accretion and a transient source becomes a persistent one. Interestingly, this might apply to 2S 0114+65, a persistent X-ray pulsar with a luminosity of \( \sim 10^{36} \) erg s\(^{-1}\), which displayed variations by a factor of \( \sim 10^4 \) (Gavriil et al. 2002).

We note that for SFXTs to host magnetars their dipole magnetic field must retain values in the \( 10^{13} \) G and \( 10^{15} \) G, respectively) are in the magnetar range. According to our proposed scenario the combination of a long spin period and a very large luminosity swing is indicative of the presence of a magnetar. This can be further corroborated through other magnetars signatures, such as, e.g., proton cyclotron features in the X-ray spectrum (Zane et al. 2001) or sporadic subsecond bursting activity such as that observed in AXPs and SGRs (Gavriil et al. 2002).

Finally, we remark on the orbital period of SFXTs: in all regimes described in § 3, the luminosity scales with \( L_X \propto P_{\text{spin}}^{-3} v_w^4 \). For our fiducial wind parameters, orbital periods of tens of days are required of the transition between low- and high-luminosity states to occur at \( \sim 10^{36} \) erg s\(^{-1}\), a typical luminosity for the onset of SFXT outbursts (Walter & Zurita Heras 2007). This is why throughout this paper we scaled our equations by \( P_{\text{orb}} = 10 \) days and used the same value in the examples of Figures 2–6. Negueruela et al. (2008) showed that orbital periods around \( \sim 10 \) days are compatible with the NSs being embedded in the clumpy wind from the supergiant companions, rather than in a quasi-continuous wind. In this case SFXT outburst durations might be associated with the transit time of a clump,

\[
\tau_{\text{out}} \approx a/v_w = 4.2 \times 10^4 a_{10} v_{8}^{-1} \text{s}, \tag{31}
\]

in reasonable agreement with the observed durations of individual flares (see § 2).

We conclude that clumpiness of the stellar wind, an often-used concept for interpreting SFXT activity, applies to the gating scenarios described here as well. The main advantage of introducing a gating mechanism rests with the possibility to model the very large luminosity swings of SFXTs with much milder density (or velocity) contrasts in the wind.

7. CONCLUSIONS

In this paper we reviewed the theory of wind accretion in HMXBs hosting a magnetic NS with a supergiant companion, and considered in some detail the interaction processes between the inflowing plasma and the magnetosphere that are expected to take place when direct accretion onto the NS surface is inhibited. We then applied this theory to SFXTs and showed that their large luminosity swings between quiescence and outburst (up to a factor of \( \sim 10^4 \)) can be attained in response to relatively modest variations of the wind parameters, provided the system undergoes transitions across different regimes. Expanding on earlier work, we found that such transitions can be driven mainly by either (1) a centrifugal barrier mechanism, which halts direct accretion when the NS rotation becomes super-Keplerian at the magnetospheric radius, a mechanism that has been discussed extensively in Be star X-ray transient pulsars; or (2) a magnetic barrier mechanism, when the magnetosphere extends beyond the accretion radius. Which mechanism and wind interaction regime apply will depend sensitively on the NS spin period and magnetic field, besides the velocity and mass-loss rate in the supergiant’s wind. In particular, the magnetic barrier mechanism requires long spin periods (\( \geq 1000 \)) coupled with magnetar-like fields (\( \geq 10^{14} \) G). On the other hand, magnetar-like fields would also be required if the centrifugal barrier set in at relatively high luminosities (\( \geq 10^{36} \) erg s\(^{-1}\)) in NSs with spin periods of hundreds seconds.

Evidence has been found that the spin periods of a few SFXTs might be as long as \( 1000–2000 \) s. Motivated by this, we presented an interpretation of the activity of IGR J17544−2619 (whose spin period is unknown) in terms of the magnetic barrier by a 1300 s spinning NS and showed that the luminosity stages singled out in a Chandra observation of this source are well matched by the different regimes of wind-magnetosphere interaction expected in this case. We discussed also an interpretation of this source based on the centrifugal barrier and a slightly shorter spin period (400 s), which reproduced the luminosity stages comparably well. We emphasize that in both solutions the required magnetic field strength (\( \geq 10^{15} \) G and \( \geq 8 \times 10^{13} \) G, respectively) are in the magnetar range.
While the possibility that magnetars are hosted in binary systems with supergiant companions has been investigated by several authors (e.g., Zhang et al. 2004; Liu & Yan 2006), clear observational evidence for such extremely high magnetic field NS in binary systems is still missing. According to the present study, long spin period SFXTs might provide a new perspective for detecting and studying magnetars in binary systems.

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APPENDIX A

ON THE HEIGHT OF THE KHI UNSTABLE LAYER

In this section we expand on the approximation $h_t \sim R_M$, introduced previously in § 3.1.2. As discussed by Burnard et al. (1983), the height $h_t$ of the layer where matter and magnetic field coexist due to the KHI is mostly determined by the largest wavelength-unstable mode of the KHI itself. These authors suggested that

$$ h_t \sim h H_t, \quad (A1) $$

where $h$ is a factor of order $\sim 1$, and $H_t = R_M^2 k_b T(R_M)/(G M m_p)$ is the scale height of the magnetosheath (note that $H_t$ is roughly of the same order as $R_M \sim 10^{10}$ cm for $T \sim 10^8$ K). In the subsonic propeller regime, equation (A1) might be a reasonable assumption due to the presence of an extended atmosphere around the magnetospheric boundary; on the contrary, it cannot be used in the sub-Keplerian magnetic inhibition regime, where matter flowing toward the NS is shocked very close to the magnetospheric boundary. Despite all these uncertainties, we show below that the assumption $h_t = h R_M$, with $h \sim 1$, gives a conservative estimate of the mass accretion rate due to the KHI. Using equation (8) and the definition $v_{\text{conv}} = \eta_{\text{KH}} v_{\text{sh}} (\rho_i/\rho_e)^{1/2} (1 + \rho_i/\rho_e)^{-1}$, we derive the equations that define the ratio of the densities inside ($\rho_i$) and outside ($\rho_e$) the magnetosphere. These are

$$ (\rho_i/\rho_e)^{1/2} (1 + \rho_i/\rho_e) = 0.3 \eta_{\text{KH}} h^{-1} R_M^{1/2} P_s^{-1}, \quad (A2) $$

if $v_{\text{sh}} = v_{\text{rot}}$, and

$$ (\rho_i/\rho_e)^{1/2} (1 + \rho_i/\rho_e) = 0.1 \eta_{\text{KH}} h^{-1} R_M^{1/2} P_s, \quad (A3) $$

if $v_{\text{sh}} = v_{\text{ps}}$. Equations (A2) and (A3) show that $\rho_i/\rho_e$ is an increasing function of $h^{-1}$ (for fixed values of $v_{\text{sh}}$, $P_s$, and $R_M$). Therefore, being $v_{\text{conv}} \propto (\rho_i/\rho_e)^{1/2} (1 + \rho_i/\rho_e)^{-1}$ and $\dot{M}_{\text{KHI}} \propto v_{\text{conv}}$, the KHI rate of plasma entry inside the magnetosphere is also an increasing function of $h^{-1}$ (provided that $\rho_i \leq \rho_e$). Since the KHI unstable layer is located inside the NS magnetosphere, the maximum height attainable is $h_t = R_M$, and thus the approximation used in §§ 3.1.2 and 3.2.2 gives a lower limit on the mass flow rate controlled by the KHI in both the sub-Keplerian magnetic inhibition and the subsonic propeller regimes.

Note also that the above lower limit does not violate the stability condition of the quasi-static atmosphere in the subsonic propeller regime. In fact, following Khansanov (2001b) this atmosphere can be considered quasi-static if the relaxation timescale of the envelope is less than the drift timescale of the mass flow crossing the magnetospheric boundary. In our case matter penetration inside the NS magnetosphere in the subsonic propeller regime is mostly provided by the KHI, and the above condition translates into $\dot{M}_{\text{KHI}} \lesssim (R_M/R_{\text{co}})^{3/2} \dot{M}_{\text{capt}}$, which is satisfied for a wide range of parameters (see Figs. 3–5).

An upper limit to the KHI mass flow rate can be obtained by assuming $\rho_i/\rho_e = 1$ (a solution adopted, e.g., by Rutledge 2001). According to Burnard et al. (1983) the density $\rho_i$ can be increased until the thermal pressure inside the magnetosphere $P_i \propto \rho_i$ is comparable to the magnetic pressure $p_m = B^2(R_M)/(8\pi)$. When this limit is reached, an instability occurs on the lower surface of the unstable layer that increases $h_t$ until $P_i < p_m$ is restored. Using the postshocked gas temperature at $R_M$ (see § 3), it is shown that $\rho_i/\rho_e \sim 1$ does not violate the condition $P_i < p_m$ (at least in the sub-Keplerian magnetic inhibition regime). Therefore, even though we restricted ourselves to the lower limit $h = 1$, the upper limit on $\dot{M}_{\text{KHI}}$ might be attainable in some instances. The KHI would then provide matter penetration inside the magnetosphere at a rate $\sim \dot{M}_{\text{capt}}$, thus allowing almost all the captured matter to accrete onto the NS. A detailed calculation of KHI accretion in the sub-Keplerian magnetic inhibition and subsonic propeller regimes will be reported elsewhere.

APPENDIX B

RADIATIVE LOSSES IN THE SUPERSONIC PROPELLER

As shown by Davies & Pringle (1981), the treatment of the supersonic propeller regime (see § 3.2.1) is self-consistent only if the energy input at the base of the atmosphere, due to the supersonic rotating NS magnetosphere, is larger than radiative losses within the atmosphere itself. The range of the validity of this assumption can be determined by using the convective efficiency parameter

$$ \Gamma = \mathcal{M}_r^2 v_t(R) t_{\text{ps}}(R) R^{-1}, \quad (B1) $$

where $\mathcal{M}_r = v_t(R)/c_s(R)$ is the Mach number, $v_t$ and $c_s$ are the turbulent and sound velocity, and $t_{\text{ps}} = 2 \times 10^{11} T^{1/2} m_p \rho^{-1}(R) \text{s}$ is the bremsstrahlung cooling time. For most of the NS rotational energy dissipated at the magnetospheric radius to be convected away
through the atmosphere’s outer boundary, $\Gamma$ should be $\geq 1$ across the entire envelope. Taking into account that in the supersonic propeller regime $v_t \simeq c_s$, $c_s \sim v_{\infty}$, $T \sim T_{\infty}$, and $\rho$ is given by equation (13), one finds $\Gamma \propto R^{-3/2}$. Therefore, the said requirement is satisfied when $\Gamma(R_{\ast}) \gtrsim 1$, i.e.,

$$M_{-6} \leq 2.2 \times 10^{-2} v_t^2 a_{10}^2.$$  \hfill (B2)

For mass-loss rates larger than the above limit, radiative losses are not negligible and the treatment used in § 3.2.1 for the supersonic propeller regime is no longer self-consistent. We checked that the limit of equation (B2) is never exceeded in the cases of interest. We note that a similar value was also derived by Ikhsanov (2002), but his limit is a factor 10 larger than ours. This might be due to the fact that Ikhsanov (2002) used $\rho_{\ast}$ instead of $\rho_{\ast}(1 + 16/3)^{1/2}$ in the expression for the matter density at $R_{\ast}$.

**APPENDIX C**

**RADIATIVE LOSSES IN THE SUBSONIC PROPELLE**

A calculation similar to that in Appendix B in the subsonic propeller regime shows that $\Gamma \propto R^{1/2}$, and radiative losses are negligible if $\Gamma(R_M) \gtrsim 1$, i.e.,

$$M_{-6} \leq 2.8 \times 10^{-2} P_{-3} a_{10}^2 \omega_{8}^{5/2} [1 + 16 R_{10}/(5 R_{M10})]^{-3/2},$$  \hfill (C1)

which is the same value as that in equation (23) (a somewhat different value was obtained by Ikhsanov [2001a] assuming a density $\simeq \rho_{\ast}$ at $R_M$ instead of $\rho_{\ast}(1 + 16 R_M/(5 R_{\ast}))^{1/2}$).

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