PROCEEDINGS OF SCIENCE

Quasi-spherical accretion in X-ray pulsars

K. Postnov[*a], **N. Shakura**[a], **A. Kochetkova**[a], **L. Hjalmarsdotter**[a]
[a]*Sternberg Astronomical Institute, 13, Universitetskij pr., 119992 Moscow, Russia*
E-mail: kpostnov@gmail.com, nikolai.shakura@gmail.com, apostnova@mail.com, astrogirl@telia.com

Quasi-spherical accretion in wind-fed X-ray pulsars is discussed. At X-ray luminosities $< 4 \times 10^{36}$ erg/s, a hot convective shell is formed around the neutron star magnetosphere, and subsonic settling accretion regime sets in. In this regime, accretion rate onto neutron star is determined by the ability of plasma to enter magnetosphere via Rayleigh-Taylor instability. A gas-dynamic theory of settling accretion is constructed taking into account anisotropic turbulence. The angular momentum can be transferred through the quasi-static shell via large-scale convective motions initiating turbulence cascade. The angular velocity distribution in the shell is found depending on the turbulent viscosity prescription. Comparison with observations of long-period X-ray wind-fed pulsars shows that an almost iso-angular-momentum distribution is most likely realized in their shells. The theory explains long-term spin-down in wind- fed accreting pulsars (e.g. GX 1+4) and properties of short-term torque-luminosity correlations. The theory can be applied to slowly rotating low-luminosity X-ray pulsars and non-stationary accretion phenomena observed in some SFXTs.

A scientific workshop in Sardinia "The extreme and variable high X-ray sky",
September 19-23, 2011
Chia Laguna, Sardegna, Italy,

*Speaker.

1. Introduction

Accreting magnetized neutron stars (NS) are observed as X-ray pulsars (XPSR). The observed characteristics of XPSRs are their X-ray luminosity L_x, pulse period P^* (or angular frequency $\Omega^* = 2\pi/P^*$) and its first time derivative $\dot{P}^*$ (or $\dot{\omega}^*$), the pulse profile at different energy bands, and X-ray spectrum. The timing properties (i.e., the temporal behavior of $\dot{\omega}^*$) and their correlation with variations of L_x provide important clues on the interaction mechanism of accreting plasma with the rotating NS magnetosphere [1]. Spectral features (especially, cyclotron lines) allow direct probing the strength of the near-surface magnetic field of NSs [2]. Their dependence on variable X-ray luminosity in persistent (Her X-1, [3]) and transient (e.g. 4U0115+63, [4]) XPSRs provides information on the accretion column properties (its geometry etc.).

XPSRs can be persistent (Her X-1, Cen X-3, Vela X-1, millisecond XPSRs, etc.) or transient (A0535+26, 4U0115+63, etc.), and are found in both high-mass and low-mass X-ray binaries (HMXB: Cen X-3, Vela X-1, transients, LMXB: Her X-1, 4U1626-67, millisecond XPSRs). The matter from the secondary star can be supplied to NS via Roche lobe overflow (the most frequent case in LMXBs) or can be captured from stellar wind of the optical companion (in HMXB). In the first case, the specific angular momentum of accreting matter j_m is usually high enought for an accretion disk to be formed around the NS magnetosphere, i.e. $j_m > j_K(R_A) = \sqrt{GMR_A}$, where R_A is the characteristic radius of magnetopshere (the Alfven radius), M is the NS mass, G is the Newton gravity constant.

2. Two regimes of quasi-spherical accretion

In the case of accretion from the wind considered here, the accretion flow geometry can be more complicated: the condition for the disk formation $j_m > j_K(R_A)$ can be satisfied or not, again depending on the size of the magnetosphere R_A and on the specific angular momentum of gravitationally captured matter from the stellar wind. The last quantity can be expressed through the gravitational capture radius $R_B = 2GM/v^2$(the Bondi radius), where $v^2 = v_w^2 + v_{orb}^2$ is the relative velocity of wind (moving with velocity v_w)and the NS (moving with orbital velocity v_{orb} through the wind) as $j_m = k\Omega_b R_B^2$. Here Ω_b is the orbital angular frequency, k is the numerical coefficient which can be positive (prograde j_m) or negative (retrograde j_m) [5], depending on properties of (generally inhomgeneous) stellar wind.

There can be two different regimes of quasi-spherical accretion[1]. The captured matter heated up in the bow shock at $\sim R_B$ to high temperatures $k_BT \sim m_p v^2$. If the characteristic cooling time of plasma t_{cool} is smaller than the free-fall time $t_{ff} = R_B/\sqrt{2GM/R_B}$, the gas falls supersonically toward the magnetosphere with the formation of a shock. This regime is usually considered in connection with bright XPSRs [7], [8]. The role of X-ray photons generated near the NS surface is two-fold: first, they can heat up plasma in the bow-shock zone via photoionization, and second, they cool down the hot plasma near the magnetosphere (with $k_BT \sim GM/R_A$) by Compton processes thus allowing matter to enter the magnetosphere via the Rayleigh-Taylor instability [9].

In the free-fall accretion regime, the X-ray luminosity (the mass accretion rate $\dot{M}$) is determined by the rate of gravitational capture of stellar wind at R_B. The accretion torque applied to NS

[1] See e.g. [6] for a recent review of previous studies of wind accretion

due to plasma-magnetopshere interaction is always of the same sign as the specific angular momentum of the gravitationally captured stellar wind, i.e. the NS can spin-up or spin down depending on the prograde or retrograde j_m.

If the relative wind velocity v at R_B is slow ($\lesssim 80$ km/s), the photoionization heating of plasma is important, but the radiation cooling time of plasma is shorter than the free-fall time, so the free-fall accretion regime is realized. If the wind velocity is larger than $\sim 80-100$ km/s, the post-shock temperature is higher than 5×10^5 K (the maximum photoionization temperature for a photon temperature of several keV); for $L_x \lesssim 4\times10^{36}$ erg/s, the plasma radiative cooling time is longer than the free-fall time, so a hot quasi-spherical shell is formed above the magnetosphere with temperature determined by hydrostatic equilibrium [10], [11]. Accretion of matter through such a shell is subsonic, so no shock is formed above the magnetosphere. The accretion rate $\dot M$ is now determined by the ability of hot plasma to enter magnetosphere. This is the settling accretion regime.

3. Settling accretion regime: theory

Theory of settling accretion regime was elaborated in [11]. In this regime, the accreting matter subsonically settles down onto the rotating magnetosphere forming an extended quasi-static shell. This shell mediates the angular momentum transfer to/from the rotating NS magnetosphere by voscous stresses due to large-scale convective motions and turbulence. The settling regime of accretion can be realized for moderate accretion rates $\dot M < \dot M_* \simeq 4\times10^{16}$ g/s. At higher accretion rates, a free-fall gap above the neutron star magnetosphere appears due to rapid Compton cooling, and accretion becomes highly non-stationary.

Mass accretion rate through the hot shell is determined by mean velocity of matter entering the magnetosphere, $u(R_A)=f(u)\sqrt{2GM/R_A}$. The dimensionless factor $f(u)$ is determined by the Compton cooling of plasma above magnetosphere and the critical temperature for Rayleigh-Taylor instability to develop [9], and is found to be

$$f(u)\approx 0.3\dot M_{16}^{4/11}\mu_{30}^{-1/11}\,, \tag{3.1}$$

where $\dot M_{16}=\dot M/[10^{16}\text{g/s}]$ and $\mu_{30}=\mu/[10^{30}\text{G}\,\text{cm}^3]$ is the NS magnetic moment. The definition of the Alfven radius in this case is different from the value used for disk accretion:

$$R_A\approx 10^9[\text{cm}]\left(\frac{\mu_{30}^3}{\dot M_{16}}\right)^{2/11}\,. \tag{3.2}$$

As the X-ray luminosity increases, Compton cooling occurs faster, and when $f(u)\to\approx 0.5$, the sonic point appears in the accretion flow above the magnetosphere, and the accretion becomes highly non-stationary. Thus the condition $f(u)=0.5$ determines the critical X-ray luminosity for the steady settling accretion: $L_x<4\times10^{36}\mu_{30}^{1/4}$ erg/s.

Accretion torques applied to NS in this regime are determined not only by the specific angular momentum of captured matter j_m (as is the case of the free-fall accretion), but also by the possibility to transfer angular momentum to/from the rotating magnetosphere through the shell by large-sclae convective motions. This means that the NS spin-down becomes possible even for prograde j_m. The plasma-magnetosphere interaction results in emerging of the toroidal magnetic field

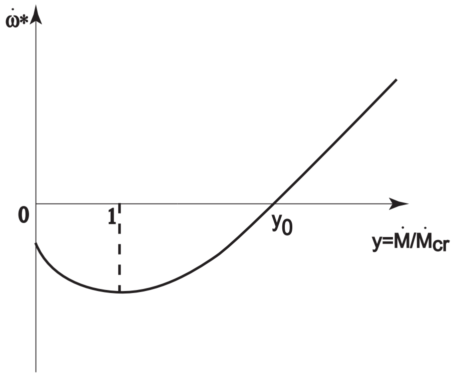

Figure 1: A schematic plot of $\dot{\omega}^*$ as a function of y. Torque reversal occurs at y_0. Sign of torque-luminosity correlations changes at $y = 1$.

$B_t = K_1 B_p(\omega_m - \omega^*)t_{inst}$, where ω_m is the angular frequency of matter at the Alfven radus, K_1 the dimensionless coupling coefficient which is different in different sources, t_{inst} is the characteristic time of RT instability. The gas-dynamical treatment of the problem of angular momentum transfer through the shell by viscous turbulence stresses [11] showed that $\omega_m \approx \Omega_b(R_A/R_B)^n$, where the index n depends on the character of turbulence in the shell. For example, in the case of isotropic near-sonic turbulence we obtain $n \simeq 3/2$, i.e. quasi-Keplerian rotation distribution. In the more likely case of strongly anisotropic turbulence (because of strong convection) we find $n \approx 2$, i.e. an iso-angular-momentum distribution. The NS spin evolution equation reads:

$$I\dot{\omega}^* = A\dot{M}^{\frac{3+2n}{11}} - B\dot{M}^{3/11}\,, \tag{3.3}$$

where I is the NS moment of inertia. The (independent of $\dot{M}$) coefficients A and B for the case $n = 2$ are (in CGS units):

$$A \approx 5.3\times 10^{31} K_1 \mu_{30}^{\frac{1}{11}} \left(\frac{v}{1000\text{km/s}}\right)^{-4} \left(\frac{P_b}{10\text{d}}\right)^{-1}, \qquad B \approx 5.4\times 10^{32} K_1 \mu_{30}^{\frac{13}{11}} \left(\frac{P^*}{100\text{s}}\right)^{-1} \tag{3.4}$$

The function $\dot{\omega}^*(\dot{M})$ reaches minimum at some $\dot{M}_{cr}$. By differentiating Eq. (3.3) with respect to $\dot{M}$ and equating to zero, we find $\dot{M}_{cr} = \left[\frac{B}{A}\frac{3}{(3+2n)}\right]^{\frac{11}{2n}}$. At $\dot{M} = \dot{M}_{cr}$ the value of $\dot{\omega}^*$ reaches an absolute minimum (see Fig. 1). In terms of the dimensionless mass accretion rate $y \equiv \frac{\dot{M}}{\dot{M}_{cr}}$ Eq. (3.3) can be written in the form

$$I\dot{\omega}^* = A\dot{M}_{cr}^{\frac{3+2n}{11}} y^{\frac{3+2n}{11}} \left(1 - \left(\frac{y_0}{y}\right)^{\frac{2n}{11}}\right), \tag{3.5}$$

where the frequency derivative vanishes at $y = y_0$: $y_0 = \left(\frac{3+2n}{3}\right)^{\frac{11}{2n}}$. The qualitative behaviour of $\dot{\omega}^*$ as a function of y is shown in Fig. 1. So in the settling accretion regime transitions from spin-up ($y > y_0$) to spin-down ($y < y_0$) occurs with changing L_x. By varying Eq. (3.5) with respect to y we find:

$$I(\delta\dot{\omega}^*) = I\frac{\partial\dot{\omega}^*}{\partial y}(\delta y) = \frac{3+2n}{11} A\dot{M}_{cr}^{\frac{3+2n}{11}} y^{-\frac{8-2n}{11}} \left(1 - \frac{1}{y^{\frac{2n-1}{11}}}\right)(\delta y)\,. \tag{3.6}$$

We see that depending on whether $y > 1$ or $y < 1$, *correlated changes* of $\delta\dot{\omega}^*$ with X-ray flux should have different signs. Indeed, for GX 1+4 in [12], [13] a positive correlation of the observed δP with $\delta\dot{M}$ was found using *Fermi* data (see also [14] for BATSE data). This means that there is a negative correlation between $\delta\omega^*$ and $\delta\dot{M}$, suggesting $y < 1$ in this source.

4. Settling accretion regime: observations

First, the settling accretion regime explains long-term spin-down states as observed, for example, in GX 1+4 [13]. Another immediate application of Eq. (3.3) to observations is the estimation of the NS magnetic field in XPSRs with equilibrium spin period (i.e. where $\dot{\omega}^* = 0$ on average):

$$\mu_{30}^{(eq)} \approx 0.3 \left(\frac{P_*/100s}{P_b/10d}\right)^{\frac{11}{12}} \dot{M}_{16}^{\frac{1}{3}} \left(\frac{v}{1000\text{km/s}}\right)^{-\frac{11}{3}} \tag{4.1}$$

(here we assumed $n = 2$ and maximum anisotropic turbulence). In wind-fed pulsars Vela X-1 and GX 301-2, the NS magnetic field estimated in this way was found to be close to the value inferred from the cyclotron line measurements (see [11] for more detail).

This equation also implies that the equilibrium period of XPSRs in the settling accretion regime is

$$P_{eq} \approx 1000[\text{s}]\mu_{30}^{12/11}\dot{M}_{16}^{-4/11}\left(\frac{P_b}{10\text{d}}\right)\left(\frac{v}{1000\text{km/s}}\right)^4. \tag{4.2}$$

There are known several very slowly rotating XPSRs, including some in HMXB (4U 2206+54 with $P^* = 5550$ s), 2S 0114+65 with $P^* = 9600$ s) and some in LMXB (e.g. 3A 1954+319, $P^* = 19400$ s [15]). Assuming disk accretion in such pulsars would require incredibly high NS magnetic fields. But in 4U 2206+54 the NS field is measured to be normal [16] , $B \simeq 3 \times 10^{12}$ G (in 2S 0114+65 possibly too, with $B \sim 2.5 \times 10^{12}$ G [17]). X-ray luminosity in these pulsars is low, $\sim 10^{36}$ erg/s, which suggests the settling accretion regime in these pulsars.

Another possible implication of the settling accretion theory can be for non-stationary phenomena in XPSRs. For example, a dynamical instability of the shell on the time scale of the order of the free-fall time from the magnetosphere can appear due to increased Compton cooling and hence increased mass accretion rate in the shell. This may even result in a complete collapse of the shell resulting in an X-ray outburst with duration similar to the free-fall time scale of the entire shell ($\sim$ 1000 s). Such a transient behaviour is observed in supergiant fast X-ray transients (SFXTs) (see e.g. [18]). Slow X-ray pulsations are found in some of them (e.g. IGRJ16418-4532, $P^* \approx 1212$ s [19]). In this source, the regular pulsations distinctly seen at a small X-ray luminosity $L_x \sim 10^{34}$ erg/s were found to disappear when the average X-ray flux increased by more than an order of magnitude, and quasi-regular flares were observed on a time scale of 200-250 s. In such a low-luminosity source the magnetospheric radius should be very large, $R_A \simeq 3 \times 10^{10}$ cm (assuming the standard value for the NS magnetic field $\mu_{30} = 1$). So the observed flaring behavior can be the manifestation of a Rayleigh-Taylor instability from the magnetospheric radius which takes place on the dynamical time scale $\sim R_A^{3/2}/\sqrt{GM}$. Note that at small X-ray luminosities the radiation cooling (and not the Compton cooling) controls plasma entering the magnetosphere, so the critical luminosity for unstable accretion can be lower than 4×10^{36} erg/s found for the Compton cooling case.

5. Conclusion

At X-ray luminosities $< 4 \times 10^{36}$ erg/s wind-fed X-ray pulsars can be at the stage of subsonic settling accretion. In this regime, accretion rate onto NS is determined by the ability of plasma to enter magnetosphere via Rayliegh-Taylor instability. A gas-dynamic theory of settling accretion regime is constructed taking into account anisotropic turbulence [11]. The angular momentum can be transferred through the quasi-static shell via large-scale convective motions initiating turbulence cascade. The angular velocity distribution in the shell is found depending on the turbulent viscosity prescription. Comparison with observations of long-period X-ray wind-fed pulsars GX 301-2 ad Vela X-1 shows that an almost iso-angular-momentum distribution is most likely realized in their shells. The theory explains long-term spin-down in wind- fed accreting pulsars and properties of short-term torque-luminosity correlations. Long-period low-luminosity X-ray pulsars are most likely experiencing settling accretion too. Spectral and timing measurements can be used to further test this accretion regime.

References

[1] L. Bildsten L., D. Chakrabarty, J. Chiu, et al., *ApJS* **113** 367 (1997).

[2] R. Staubert, *Chinese J. Astron. Astrophys. Suppl.* **3** 270 (2003).

[3] R. Staubert, N.I. Shakura, K.A. Postnov, et al., *AA* **465** L25 (2007).

[4] S. Tsygankov, A. Lutovinov, E. Churazov and R. Sunyaev, *AstL* **33** 368 (2007).

[5] C. Ho, R.E. Taam, B.A. Fryxell, T. Matsuda and H. Koide, *MNRAS* **238** 1447 (1989).

[6] E. Bozzo, M. Falanga and L. Stella, *ApJ* **683** 1031 (2008).

[7] J. Arons and S. Lea, *ApJ* **207** 914 (1976).

[8] D.J. Burnard, J. Arons and S.M. Lea, *ApJ* **266** 175 (1983).

[9] R. F. Elsner and F. K. Lamb, *ApJ* **215** 897 (1977).

[10] M.E. Davies and J.E. Pringle, *MNRAS* **196** 209 (1981).

[11] N.I. Shakura, K.A. Postnov, A.Yu. Kochetkova and L. Hjalmarsdotter, *MNRAS*, in press (2011). [arXiv:1110.3701]

[12] A. González-Galán, E. Kuulkers, P. Kretschmar, et al., *PoS(8th INTEGRAL Workshop)* p. 016 [arXiv:1105.1907].

[13] A. González-Galán, E. Kuulkers, P. Kretschmar, et al., *AA* in press (2011).

[14] D. Chakrabarty, L. Bildsten, M.H. Finger, et al., *ApJ* **481** L101 (1997).

[15] D.M. Marcu, F. Fürst, K. Pottschmidt, et al., *ApJ* **742** L11 (2011).

[16] W. Wang, *MNRAS* **398** 1428 (2009).

[17] E.W. Bonning and M. Falanga, *AA* **436** L31 (2005).

[18] L. Ducci, L. Sidoli and A. Paizis, *MNRAS* **408** 1540 (2010).

[19] L. Sidoli, S. Mereghetti, V. Sguera and F. Pizzolato, *MNRAS* submitted; [arXiv:1110.5218]