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Modulation of the Neutron Star Boundary Layer Luminosity by Disk Oscillations

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ABSTRACT

Paczyński (1987) pointed out that any modulation of the rate at which mass is accreted by a neutron star from the innermost part of a relativistic accretion disk will lead to a modulation of the luminosity of the boundary layer. Following this reasoning, we demonstrate that variability of the boundary layer X-ray flux must necessarily show frequencies of certain global accretion disk oscillations. This theoretical paradigm – clock in the disk, modulation at the boundary layer – resolves one puzzling paradox in the neutron star quasi periodic oscillation (QPO) data.

Key words: *Stars: neutron – Stars: oscillations – Accretion, accretion disks*

1. Introduction

Several similarities in observed properties of QPO frequency behavior (stability, the 3/2 ratio) suggest a similar physical nature of QPOs in black hole and neutron star sources. However, this idea is seemingly in a direct and acute conflict with a well established observational fact, discussed in detail by Gilfanov *et al.* (2003) and Revnivtsev and Gilfanov (2005), that in the case of neutron star sources, the X-ray luminosity of QPOs comes not from the accretion disk, but from the boundary layer that forms on the surface of the neutron star. In black holes, of course, there is no such boundary layer, so the direct mechanism of X-ray modulation must be different in black holes and in neutron stars.

In the present paper, we argue that the conflict could be naturally resolved in some versions of the Kluźniak–Abramowicz resonance model, and of other disk oscillation models of QPOs. Specifically, we show that the time varying X-ray flux that originates at the neutron star boundary layer, must necessarily have a memory of frequencies of global accretion disk oscillations modes. This has already been anticipated by Nowak and Wagoner (1993), in their discussion of disk p-modes. Neutron star and black hole accretion disks are very similar, and thus may have similar modes of oscillations, *i.e.*, a similar QPO clock.

1.1. Mode Excitation and Resonance

In the black hole case, excitation is most probably due to the MRI turbulence (Nowak and Wagoner 1993, Abramowicz 2005, Vio *et al.* 2006). In the case of neutron stars the excitation may be due to the neutron star spin (Lee *et al.* 2004, Kluźniak *et al.* 2004a, Pétri 2005a,b), a much stronger effect than excitation by turbulence. Detailed analysis of frequency–frequency correlations of kHz QPOs in neutron stars supports the notion of a resonance between two modes of oscillation (Kluźniak and Abramowicz 2000, Abramowicz and Kluźniak 2001, Abramowicz *et al.* 2003, Rebusco 2004, 2005, Horák 2004, Abramowicz *et al.* 2005a,b).

1.2. Modulation of X-rays

Bursa *et al.* (2004) have shown that for oscillating black hole accretion disks, the photon flux observed at infinity may be modulated, and enhanced, (mostly) by the general relativistic effect of gravitational lensing. In this paper we will argue that oscillations of neutron star accretion disks may modulate the X-ray flux *via* another general relativistic effect, the overflow of a "Roche lobe" near the innermost stable circular orbit (ISCO). For an oscillating disk this overflow occurs at a variable rate and results in a modulation of the luminosity of the boundary layer at the rigid surface of the neutron star.

For radial oscillations of an accretion torus, the variation of $\dot{M}$ has been found numerically by Zanotti *et al.* (2003), see also Fragile (2004). Here, we calculate $\dot{M}$ variation for a vertically oscillating disk. We will consider only an axisymmetric oscillation to allow an analytic treatment of the problem. Of particular interest is a rigid vertical oscillation mode of a torus, with the eigenfrequency equal to the epicyclic vertical frequency, $\omega = \omega_z$, recently identified by Kluźniak and Abramowicz (2002), Lee *et al.* (2004), Rubio-Herrera and Lee (2005a,b), Abramowicz *et al.* (2006).

Accretion disk oscillations are crucial for arguments presented here. Our general discussion is valid for all oscillations that influence the position of innermost part of the disk, *i.e.*, the part which is located closely to ISCO (Wagoner *et al.* 2001, Zanotti *et al.* 2003, Kluźniak *et al.* 2004a,b).

2. The Disk Structure Close to ISCO

Many models of neutron stars predict a neutron star radius that is smaller than the ISCO radius, accordingly the neutron star accretion disk may be decoupled from the stellar surface, *i.e.*, it could be very similar in its properties to a black hole accretion disk (Kluźniak and Wagoner 1985). A powerful analytic model for accretion flows near ISCO was worked out in terms of the flow equipotential structure by Paczyński and collaborators thirty years ago in Warsaw (see *e.g.*, Jaroszyński *et al.* 1980). The Warsaw model accurately describes the flow hydrodynamics for standard Shakura-Sunyaev thin disks, slim disks, adafs, ion tori, and thick Polish doughnuts. It is in impressive agreement with more recent numerical simulations of accretion flows, as directly checked, *e.g.*, by Igumenshchev *et al.* (1996). The model was formulated for a general hydrodynamical* flow, and with full and exact use of Einstein's strong gravity.

Here, for simplicity of presentation, we recall the Warsaw model assuming constant specific angular momentum $\ell(r,z,\varphi) = \ell_0 =$ const and a barytropic equation of state, $P = P(\rho)$. Here P is pressure and ρ is density. It would be convenient to introduce enthalpy h, by the standard definition $\mathrm{d}h = \mathrm{d}P/\rho$. We also use here Paczyński's pseudo-Newtonian model for Einstein's stationary and spherically symmetric strong gravity. In the standard cylindrical coordinates r, z, φ the Paczyński and Wiita (1980) pseudo-potential is given by

$$\Phi(R) = -\frac{GM}{R - R_{\mathrm{G}}}, \quad R^2 \equiv r^2 + z^2, \quad R_{\mathrm{G}} \equiv \frac{2GM}{c^2} \tag{1}$$

here M denotes the mass of the gravity source: in application to the problem in hand – the mass of the central neutron star. This pseudo-potential reproduces well the behavior of Keplerian angular momentum in spherical strong gravity,

$$\ell_{\mathrm{K}} = \frac{r}{r - R_{\mathrm{G}}} \sqrt{GMr} \tag{2}$$

in particular it exactly reproduces the location of the minimum of Keplerian distribution at ISCO, the innermost stable circular orbit at $r = r_{\mathrm{MS}} = 3R_{\mathrm{G}}$ (see the dotted line in the upper panel of Fig. 1).

The Euler equation for stationary, purely azimuthal flows,

$$\frac{\ell^2}{r^3}\vec{r} = \frac{\nabla P}{\rho} + \nabla\Phi \tag{3}$$

may be trivially integrated in the case considered here, *i.e.*, for constant angular momentum, $\ell = \ell_0$, and barytropic equation of state, $P = P(\rho)$. The integration

*Recently, Komissarov (2006) generalized the model by adding magnetic field into consideration. He demonstrated that the equipotential structure of the flow admits the Roche lobe. This means that, not surprisingly, the mechanism of the accretion rate modulation discussed in the present paper operates also when magnetic field is present.

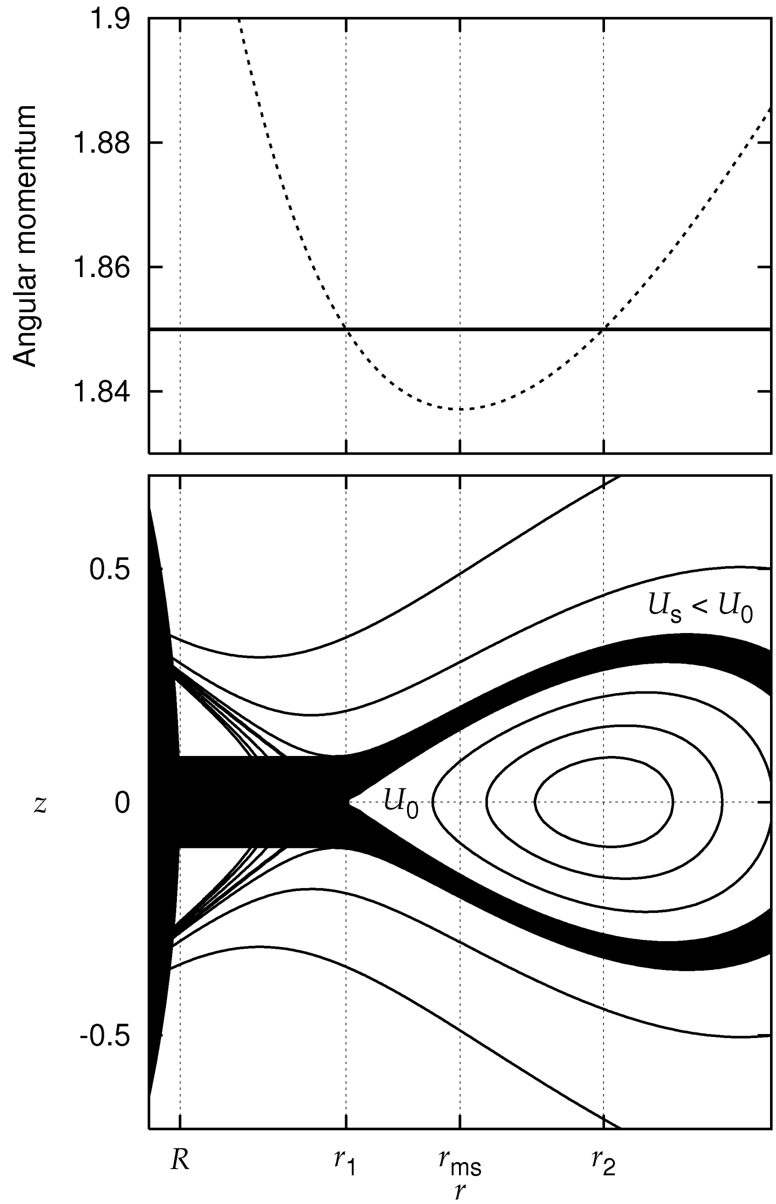

Fig. 1. *Top*: the Keplerian angular momentum (the dotted line) and the angular momentum in the flow (the solid line). *Bottom*: the equipotential surfaces (solid lines) and distribution of fluid (shaded region) in a meridional cross-section of an axially symmetric configuration.

yields,

$$h(P) + \mathcal{U}(r,z) = \text{const} \equiv \mathcal{U}_S, \tag{4}$$
$$\mathcal{U}(r,z) \equiv \Phi(r,z) + \frac{\ell_0^2}{2r^2}.$$

Eq. (4) is just the Bernoulli equation with the effective potential $\mathcal{U}(r,z)$, which is explicitly known, of course. Fig. 1 illustrated the obvious:

a) Circles $r = r_1$ and $r = r_2$, where $\nabla P = 0$, correspond to $\ell_0 = \ell_K$. At the $r = r_1$ circle, called the cusp, the pressure has a saddle point. The equipotential $\mathcal{U}(r,z) =$

$\mathcal{U}_0$, called the Roche lobe, crosses there itself. At the circle $r = r_2$, the pressure has its maximum.

b) In equilibrium described by the Bernoulli equation (Eq. 4), surfaces of constant enthalpy, pressure and density coincide with surfaces of constant effective potential, $\mathcal{U}(r,z) = \text{const}$. The surface of the disk is given by $P = \text{const} = 0$. Thus, equilibria may only exist if the disk surface corresponds to one of the equipotentials inside the Roche lobe, *i.e.*, in the region indicated by light shading. Equilibrium in the region $r \le r_1$ is not possible. If the fluid distribution overflows the Roche lobe, *i.e.*, the surface of the disk for $r \gg r_1$ coincides with $\mathcal{U}(r,z) = \mathcal{U}_S < \mathcal{U}_0$, the disk will suffer a dynamical mass loss, with accretion rate $\dot{M}$.

3. Stationary "Roche Lobe" Overflow

An analytic formula for $\dot{M}$ for stationary flows was first calculated by Kozłowski *et al.* (1978), who used Einstein's theory. Here, we recall another derivation, done by Abramowicz (1981, 1985), as we later use the same assumptions and notation to calculate $\dot{M}$ for a non-stationary (oscillating) disk. In particular, we assume that the Roche lobe overflow is small (quadratic in disk thickness H),

$$h(r_1,z) = h^\star - \frac{1}{2}\kappa^2 z^2, \tag{5}$$

$$\kappa^2 \equiv -\left(\frac{\partial^2 h}{\partial z^2}\right)_L, \tag{6}$$

$$H = \frac{\sqrt{2h^\star}}{\kappa}, \tag{7}$$

that the equation of state is polytropic, and that the radial velocity v^r connected to the mass loss through the cusp equals the sound speed $c_{\rm s}$,

$$P = K\rho^{1+1/n}, \quad v^z \ll v^r = c_{\rm s} = \sqrt{\frac{h}{n}}. \tag{8}$$

In the above equations $h^\star = h(r_1,0)$ denotes a maximal value of the enthalpy on the cylinder $r = r_1$ and subscript L stands for the evaluation of the derivative in the point $[r_1,0]$. The local mass flux $\dot{m} = \rho v^r = h^{n+1/2}/K^n(1+n)^n n^{1/2}$, integrated vertically through the cusp thickness, and azimuthally, gives the desired total mass flux in terms of the enthalpy,

$$\dot{M} = \int_0^{2\pi} r_1\,{\rm d}\varphi \int_{-H}^{+H} \dot{m}\,{\rm d}z = (2\pi)^{3/2}\frac{r_1}{\kappa n^{1/2}}\left[\frac{1}{K(n+1)}\right]^n \frac{\Gamma(n+3/2)}{\Gamma(n+2)} h_0^{n+1}. \tag{9}$$

Here $\Gamma(x)$ is the Euler gamma function. From $v^2/2 + h + \mathcal{U} = \mathcal{U}_{\rm S}$ one gets $\Delta\mathcal{U} = \mathcal{U}_{\rm S} - \mathcal{U} = (1+1/2n)\,h$ and from

$$\kappa^2 \equiv -\left(\frac{\partial^2 h}{\partial z^2}\right)_L \left(\frac{n}{n+1/2}\right)\omega_z^2, \quad \omega_z^2 \equiv \left(\frac{\partial^2 \mathcal{U}}{\partial z^2}\right)_L, \tag{10}$$

we recover the result obtained by Abramowicz (1985),

$$\dot{M} = A(n)\frac{r_1}{\omega_z}\Delta \mathcal{U}^{n+1}, \tag{11}$$

$$A(n) \equiv (2\pi)^{3/2}\left[\frac{1}{K(n+1)}\right]^n \times \left[\frac{1}{n+1/2}\right]^{n+1/2}\frac{\Gamma(n+3/2)}{\Gamma(n+2)}. \tag{12}$$

4. Non-Stationary Roche Lobe Overflow

We will calculate $\dot{M}$ for non-stationary oscillating flows in a special but important case, assuming that the poloidal part of velocity $\vec{v} = (v^r, v^z)$ may be derived from a potential. In this case, the non-stationary version of the Bernoulli equation has the form,

$$\frac{\partial\chi}{\partial t} + \frac{v^2}{2} + h + \mathcal{U} = \mathcal{U}_{\rm S}, \quad \text{with} \quad \vec{v} = \nabla\chi. \tag{13}$$

We assume that oscillation is a small non-stationary perturbation to the stationary flow considered in Section 3,

$$\chi(\vec{r},t) = \chi_{(0)}(\vec{r}) + \varepsilon\chi_{(1)}(\vec{r},t) \tag{14}$$

where the subscript (0) refers to the stationary flow and the dimensionless parameter $\varepsilon \ll 1$ characterizes strength of the perturbation. From the definition $\vec{v} = \nabla\chi$ one derives,

$$v^2 = v^2_{(0)} + 2\varepsilon\vec{v}_{(0)}\cdot\vec{v}_{(1)} + \varepsilon^2 v^2_{(1)} = c_{\rm s}^2 + 2\varepsilon c_{\rm s}\frac{\partial\chi_{(1)}}{\partial r} + \varepsilon^2\left[\left(\frac{\partial\chi_{(1)}}{\partial r}\right)^2 + \left(\frac{\partial\chi_{(1)}}{\partial z}\right)^2\right]. \tag{15}$$

The perturbed enthalpy profile can be approximated by an expansion

$$h = h_{(0)} + \varepsilon h_{(1)} + \varepsilon^2 h_{(2)} + o(\varepsilon^3). \tag{16}$$

By substituting Eqs. (15) and (16) into the Bernoulli Eq. (13) and equating coefficients of same powers of ε, we obtain

$$h_{(1)} = -\frac{\partial\chi_{(1)}}{\partial t} - \left(\frac{\bar{h}}{n}\right)^{1/2}\frac{\partial\chi_{(1)}}{\partial r}, \tag{17}$$

$$h_{(2)} = -\frac{1}{2}\left[\left(\frac{\partial\chi_{(1)}}{\partial r}\right)^2 + \left(\frac{\partial\chi_{(1)}}{\partial z}\right)^2\right]. \tag{18}$$

This way, we expressed perturbations of fluid quantities in terms of a perturbation $\chi_{(1)}$. To progress further, one must know the function $\chi_{(1)} = \chi_{(1)}(t,r,z)$ that describes oscillations. Finding it in general is a difficult global problem. We do not attempt to solve it here. Instead, we describe oscillations by an ansatz,

$$\chi_{(1)} = z v_z \cos\omega t, \quad v_z = \text{const} \tag{19}$$

which models a vertical "epicyclic" oscillation with frequency ω that rigidly moves the fluid up and down across the equatorial symmetry plane.

The quantity $\varepsilon v_z = \mathrm{const}$ can be interpreted as the amplitude of the vertical velocity. Eqs. (17) and (18) give

$$h_{(1)} = z v_z \omega \sin \omega t, \quad h_{(2)} = -\frac{1}{2} v_z^2 \cos^2 \omega t. \tag{20}$$

The vertical profile of the enthalpy at $r = r_1$ reads

$$h(r_1, z, t) = h^\star - \kappa^2 z^2 + \varepsilon z v_z \omega \sin \omega t - \frac{1}{2} \varepsilon^2 v_z^2 \cos^2 \omega t + o(\varepsilon^3) \tag{21}$$

that is quadratic in the variable z (see the left panel of Fig. 2). The position of the enthalpy maximum on the cylinder $r = r_1$ is shifted from $z = 0$ to height $\delta z(t)$ given as

$$\delta z(t) = \delta Z \sin \omega t, \quad \delta Z = \varepsilon \frac{\omega v_z}{\kappa^2}. \tag{22}$$

We can interpret δZ as the amplitude of the oscillations. Also the value of enthalpy in the maximum differs from the stationary case by

$$\delta h^\star \equiv h(r_1, \delta z) - h^\star = \frac{1}{2} \kappa^2 \left[\delta z^2 - \frac{\kappa^2}{\omega^2} (\delta Z^2 - \delta z^2) \right] + o(\varepsilon^3). \tag{23}$$

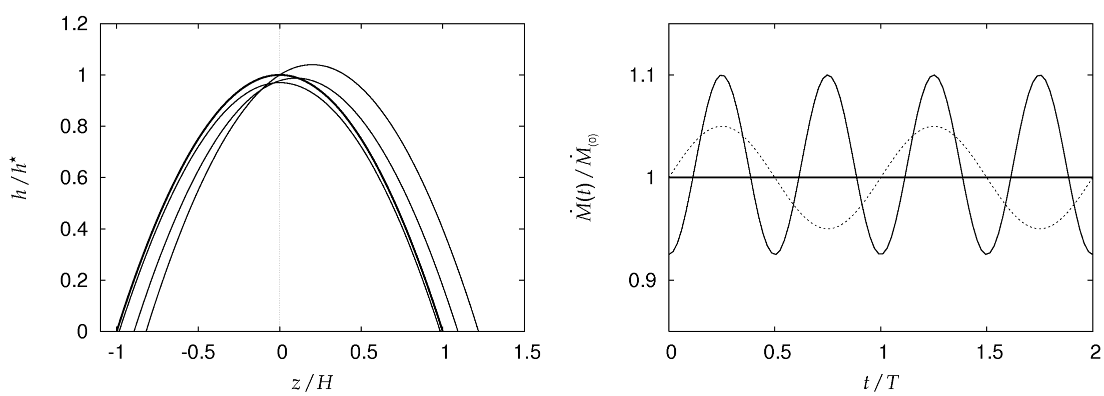

Fig. 2. *Left*: The vertical profiles of the enthalpy $h(r_1, z)$ on the cylinder $r = r_1$ during vertical disk oscillations (thin lines). The amplitude of oscillations is $\delta Z = 0.2H$ and we chose the polytropic index of the fluid $n = 1.5$. The figure captures profiles with the enthalpy maxima at $\delta z = 0, 0.1H$ and $0.2H$. The enthalpy profile for the unperturbed stationary disk is also shown (thick line).
Right: The modulated accretion rate from the oscillating disk (thin solid line). The accretion rate for the stationary disk is plotted by thick line. Time is rescaled by the period of oscillations. For reference we also plot the phase of disk oscillations (dotted line). The accretion rate is modulated with twice the frequency of oscillations if the vertical oscillation is axially symmetric.

According to Eq. (9) the instantaneous accretion rate depends on the maximal enthalpy as $\dot{M} \propto (h^\star)^{n+1}$. This relation can also be applied in the case of vertical oscillations because the z-dependence of enthalpy on the cylinder $r = r_1$ remains quadratic also in this case and the oscillations do not contribute to the radial velocity of accreted matter. Hence, using Eqs. (7), (10) and (23) and assuming that the frequency of oscillations equals to the local vertical epicyclic frequency, $\omega = \omega_z$, we arrive at our final result

$$\frac{\delta \dot{M}}{\dot{M}_{(0)}} = (n+1)\frac{\delta h^\star}{h^\star} = \frac{2-p}{2-2p}\left[(1+p)\frac{\delta z^2}{H^2} - p\frac{\delta Z^2}{H^2}\right] \tag{24}$$

where $\delta\dot{M} \equiv \dot{M} - \dot{M}_{(0)}$ and $p \equiv n/(n+1/2)$. The quadratic dependence of $\delta\dot{M}$ on perturbation implies that the frequency of the modulation of $\dot{M}$ is twice the frequency of the disk oscillation (see the right panel of Fig. 2).

5. Discussion

Neutron stars have rigid surfaces, black holes do not. Narayan forcefully advocated in a series of papers that this fundamental difference shows up in several observational properties of these two classes of strong gravity sources (see a summary of his arguments in Narayan 2003). We add to the Narayan list also the difference discussed by us here: the modulation mechanism of the X-ray luminosity in QPOs. While in both neutron and black hole sources the high frequency double peak QPOs probably originate in accretion disk oscillations, the modulation of the X-ray luminosity is very different in the two cases.

While gravitational lensing may directly modulate the X-ray flux at infinity of an oscillating black hole disk (Bursa *et al.* 2004), the variable flow rate leaving the disk near the ISCO is not an additional source of modulated luminosity in black holes. The inner mass flux disappears behind the black hole event horizon, and thus modulation of the X-ray luminosity may hardly be expected due to the Roche lobe overflow caused by disk oscillations in black hole systems.

The situation is quite different in neutron stars. A modulated rate of mass flow leaving the accretion disk and crossing the "relativistic accretion gap" (Kluźniak and Wagoner 1985) will give an enhanced luminosity modulation as the mass enters the boundary layer (Paczyński 1987). The accreting fluid releases more energy in the boundary layer than that emitted by the whole accretion disk. Sunyaev and Shakura (1986) estimated that 69% of the total luminosity if $R_\star = 3R_{\rm G}$, or even 86% if $R_\star = 1.5R_{\rm G}$ is radiated from the boundary layer. It is important that the boundary layer cools rapidly (Kluźniak and Wilson 1991, Kluźniak *et al.* 1990), hence the kHz modulation is not erased by the thermal capacity of the boundary layer reservoir.

We have shown that the mass accretion rate carries a modulation imprint of the vertically oscillating accretion disk, at double the mode frequency if the oscillation mode is antisymmetric with respect to the equatorial plane. This modulation frequency is in turn imprinted on the X-ray flux released in the boundary layer.

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