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CHAPTER VIII

COUPLING OF AN ACCRETION DISK CORONA WITH A WEAKLY MAGNETIC NEUTRON STAR AND A POSSIBLE RELATION WITH QPO.

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Summary

We propose a mechanism whereby rotating, weakly magnetized neutron stars, which are thought to be present in some low mass X-ray binaries, may lose energy to the accretion disk via an electrodynamical coupling between the magnetic field of the star and magnetic loops in a disk corona. This mechanism may solve some of the problems posed by the observations of these systems. In particular, the properties of the power spectra of some of the sources, in which quasi-periodic oscillations are observed, can be explained, and especially it can be understood why the low-frequency noise in some sources may be much reduced in power with respect to the power in the oscillations.

Key words: Neutron stars - disk corona - quasi-periodic oscillations.

1. Introduction

It is generally believed that a large majority of the low-mass X-ray binaries (LMXB) consist of a neutron star and a low mass companion. Roche-lobe overflow drives matter from the companion onto the neutron star via an accretion disk. Most LMXB do not show X-ray pulsations, which could imply that the magnetic field of the neutron star is too weak to channel the accretion flow, and the accretion disk is therefore thought to extend down to the surface of the neutron star (see for a review: Lewin and Joss, 1983). During recent years, however, this picture has changed somewhat for some of the LMXB, mainly due to the discovery of millisecond radio pulsars in binaries and of quasi-periodic oscillations (QPO) in some of the LMXB.

The discovery of millisecond radio pulsars in binaries has lead to the conclusion that some of these systems (i.e., the wide ones) are formed by the
later evolution of binaries consisting of a neutron star and a normal companion star, in which the companion was less massive than the neutron star, (Van den Heuvel, 1984). When in such a system the companion has evolved to the giant stage it begins to overflow its Roche lobe, which leads to an expansion of the orbit. This may explain the existence of several wide LMXB, such as Cyg X-2. During this mass transfer phase, angular momentum is gained by the neutron star and its rotation period decreases. Therefore, when the mass transfer stops one will observe the neutron star as a radio pulsar with a short rotation period in a wide binary. Assuming that the radio pulsar loses its rotational energy due to magnetic dipole radiation (Pacini, 1967, 1968; Gunn and Ostriker, 1969), the magnetic field strength of the neutron star can be inferred. The measured values are in the range of $10^{8.5}$ to $10^{11.0}$ G (Dewey et al., 1986). This, of course, implies that the magnetic fields of the neutron stars in the LMXB, from which these radio pulsars evolved, were at least as strong. In a recent study by Van den Heuvel et al. (1986) it was shown that the magnetic fields of neutron stars probably do not decay below $10^9$ G or if they do, the timescale is larger than $10^9$ yr.

The magnetospheric radius for spherical accretion is given by (Davidson and Ostriker, 1973; Lamb, Pethick and Pines, 1973)

$$r_a = (2.9 \times 10^8 \text{ cm}) \mu_3^{4/7} m^{1/7} R_6^{2/7} L_{37}^{-2/7}$$  \hspace{1cm} (1)

Here $\mu_3$ is the magnetic dipole moment of the neutron star in units of $10^{30}$ G cm$^3$, $m$ its mass in solar masses, $R_6$ its radius in units of $10^6$ cm and $L_{37}$ the total luminosity in units of $10^{37}$ erg/sec. For the values of the magnetic field quoted above this leads to

$$1.6 \times 10^6 \text{ cm} < r_a < 4.2 \times 10^7 \text{ cm}$$  \hspace{1cm} (2)

where we have used $L_{37} = 10$, $R_6 = 1$ and $m = 1.4$. The values in (2) are valid for spherical accretion and dipole geometry and they may be a factor of two smaller in the case of disk accretion (Ghosh and Lamb, 1979). But even then it is clear that for some neutron stars in LMXB $r_a > R$ and the disk does not reach the surface of the star.

That some of the neutron stars in LMXB may have magnetic fields in the order of $10^8$ to $10^{10}$ G, and periods in the millisecond range was suggested by Alpar and Shaham (1985), by interpreting the quasi-periodic oscillations, discovered in GX5-1 (Van der Klis et al., 1985), in terms of a beat frequency model. In this model the observed frequency of the peak in the power spectrum is
due to a beat between the Kepler frequency at the inner edge of the accretion disk (taken to be the magnetospheric radius, $r_a$) and the rotation frequency of the neutron star. This gives a relation between the luminosity (if the assumption is made that the mass accretion rate is proportional to the luminosity) of the source and the beat frequency, that fits well to the data of GX5-1 and which predicts a rotation period of the neutron star of the order of 10 ms and a magnetic field of the order of $6 \times 10^9$ G, leading to a magnetospheric radius of $r_a \sim 60$ km. (see also Van der Klis et al. 1985).

So far, models have assumed that in order to produce such a beat frequency, an asymmetry in the stellar magnetic field must interact with structures in the inner region of the disk. Lamb et al. (1985) and Berman and Stollman (1986a) suggested that the magnetic field acts as a gating mechanism for the accretion flow from these structures. In the model put forward by Lamb et al. (1985) these structures are thought of as physical blobs created by some form of instability and are rotating in the inner region of the disk with the Kepler frequency. Berman and Stollman (1986a) argued that the interaction of the field-asymmetry with these blobs probably could not provide the degree of modulation observed, but rather that the structures must be regions of reduced conductivity.

After the discovery of QPO in GX5-1 more sources have been discovered, which show quasi-periodic oscillations; nl. Sco X-1 (Middleditch and Friedhorsky, 1985; Van der Klis and Jansen, 1985; Van der Klis et al. 1986 ), Cyg X-2 (Hasinger et al., 1986), GX349+2 (Lewin et al., 1985), GX17+2 (Stella et al., 1985a), 4U1820−30 (Stella et al., 1985b), GX3+1 (Lewin et al., 1986) and MXB1730−335 (the Rapid Burster, Stella et al., 1985c). Some of these sources show quite a different behaviour from GX5-1, which seems difficult to explain within the context of a beat frequency model (Lewin, 1986; Van der Klis, 1986). However, other theories which have been developed so far to explain the QPO and which do not use the idea of a beat frequency (e.g. Hameury et al., 1985; Boyle et al., 1986) have also problems explaining one or more properties of the observations (e.g., Lewin, 1986).

In this paper we assume that the neutron stars in some of the LMXB have magnetic fields, which are capable of holding the accretion disk at some distance from the stars surface. Within this context we develop a model for the interaction of the magnetic field of the neutron star and magnetic structures in an accretion-disk-corona (ADC). The existence of those coronae in low-mass X-ray sources can be inferred from the analysis of X-ray lightcurves of some of these sources (White and Mason, 1984). It seems plausible that these ADCs are magnetically structured with loops, which are anchored in the disk (e.g., Ionson and Kuperus, 1984).
In this paper we will show that energy can be transferred from the neutron star (probably rotational energy) to the disk-corona-system. That energy transport like this may be important in LMXB was suggested by Priedhorsky (1986), who pointed out that some of the problems in our understanding of X-ray spectra from Sco X-1 can be explained if energy can be transferred from the neutron star to the disk via the magnetic field.

Furthermore, we will show that the interaction between the magnetic field of the neutron star and the coronal loops may provide a mechanism that can explain the existence of QPO and which at the same time provides a solution to some of the problems that are posed by models, which are based on the gating of the accretion, such as the near absence of power in the low-frequency noise with respect to the power in the QPO, as, for example, observed in Sco X-1.

2. Interaction of a magnetic neutron star with a disk coronal loop

It was shown by Lightman and Eardly (1974) that the inner region of a "standard" accretion disk (cf. Shakura and Sunyaev, 1973), where radiation pressure dominates gas pressure, is secularly unstable to clumping into rings. The radius, $r_1$, at which the gas pressure and the radiation pressure are approximately equal is given by (Shapiro and Teukolsky, 1983)

$$r_1 \sim 1.2 \times 10^7 \alpha^{2/21} \dot{M}^{1/3} \frac{16}{21} \text{cm}$$

where $m$ is the mass of the neutron star in solar masses, $\dot{M}_{17}$ the accretion rate in $10^{17}$ g/sec and $\alpha$ is the well known non-dimensional viscosity parameter (Shakura and Sunyaev, 1973). In the case of the bright LMXB the value of $r_1$ is approximately $r_1 \sim 7.5 \times 10^7$ cm. Comparing this to the range of $r_a$ given in equation (2), we notice that the disks in the bright LMXB do have an inner region in which the gas pressure is dominated by the radiation pressure, and which may, therefore, be unstable as mentioned above.

However, it has been demonstrated by Ionson and Kuperus (1984) that the inner parts of an accretion disk can be stabilized by the formation of a corona. This corona is likely to be magnetically structured with loops of a characteristic length $\lambda \sim h$, where $h$ is the thickness of the disk (Galeev et al., 1979; Coroniti, 1981; Burm, 1986). The stabilisation essentially occurs because part of the accretion energy is vented into the overlying corona, thus reducing the radiation pressure in the inner disk. The reason that such a disk corona is likely to be magnetically structured is that the magnetic field carried by the accreting matter is amplified by the differential rotation of the disk and the increasing magnetic pressure causes the field to become buoyant in...
the strong gravitational field.

An estimate of the strength of the field in the disk-corona system can be made by an equipartition argument. If we set the energy density of the magnetic field in the disk equal to the kinetic energy density of the Keplerian motion we find $B_d^2/8\pi \sim \rho v_K^2/2$. The strength of the stellar field at the magnetospheric radius can be estimated from the equation which determines the magnetospheric radius: i.e., $\dot{M} v_K r/\delta \sim r^2 B^2$ (Ghosh and Lamb, 1979). Here $\delta$ is the radial thickness of the boundary layer. The mass accretion rate is given by $\dot{M} = 4\pi r h \rho v_r$. Here $h$ is the half thickness of the disk and $v_r$ is the radial velocity in the disk. Substitution leads to the following expression for the magnetic field of the star at the magnetospheric radius: $B^2/8\pi \sim v_r v_K \rho h/2\delta$. Since the thickness of the disk, $h$, is probably not larger than the thickness of the boundary layer, and since the radial velocity in the disk is much smaller than the Keplerian velocity we find that the field strength in the disk is larger than the local stellar field. And since the coronal field may be somewhat smaller than the disk field we believe that the coronal field can be larger than or at least equal to the local stellar field.

Let us now consider the interaction of a magnetic loop in an accretion-disk-corona with the magnetic field of a rotating neutron star. The loop is anchored in the disk and hence is rotating with the Kepler angular velocity $Q_K(r)$, while the neutron star magnetic field is, up to the corotation radius, supposed to rotate with the neutron star angular velocity $Q_s$. We assume that a field asymmetry is present, which is, for example, due to a displacement of the aligned magnetic axis from the rotation axis, or due to a inclination of the magnetic axis with respect to the spin axis (Curtis Michel and Dessler, 1981; Lamb et al., 1985; Berman and Stollman, 1986a). In a coordinate system fixed to the loop, the loop experiences an oscillating stellar field with a frequency equal to the beat frequency $Q_B$ given by

$$Q_B = Q_K - Q_s$$

The magnetic field strength at radius $r$ and time $t$ is thus given by

$$B_s(r,t) = B_0(r) + B_1(r) \cos Q_B t$$

Here we have assumed the simplest possible form of the field asymmetry. A more general form is given by

$$B_s(r) = B_0(r) + \sum_{i=1}^{n} B_i(r) \{\cos(iQ_B t + \phi_i)\}$$
(Lamb et al., 1985). This magnetic field interacts with the current system inside the loop, thereby supplying an "external" power source for the loop.

In order to estimate the power transferred from this external source to the loop we make a number of simplifying assumptions. First, we assume the loop to be a stretched, cylindrical filament with length $l$ and diameter $a$, located at distance $r$ from the center of the neutron star and having a helical magnetic field $\mathbf{B}_f$ and a helical internal current distribution $\mathbf{j}_i$. Secondly, we assume that every part of the loop experiences the same external magnetic field, which we assume to be perpendicular to the axis of the stretched loop (see Fig. 1).

![Fig. 1: The magnetic field of the neutron star $B^*$, rotates with angular velocity $\Omega^*$ and interacts with a magnetic loop of length $l$, which is situated at distance $r_a$ (i.e. the magnetospheric radius of the disk) and which rotates with the Kepler angular velocity $\Omega_K$. Due to this interaction the loop oscillates with a velocity amplitude $v_0$.](image)

Apart from the internal current system, which is generated by the disk motions, there is an induced current $\mathbf{j}_i$, because of the periodic changes in the magnetic flux through the loop as a result of the rotating magnetic field of the neutron star (see eq. (5)). This induced current system consists of a component $\mathbf{j}_{i\parallel}$, parallel to the axis of the filament and thus perpendicular to $B_s$ and a
component \( j_{\perp} \) located in a plane perpendicular to \( j_{\|} \).

The Lorentz force acting on a volume element of the loop is then equal to

\[
\mathbf{f}_L = \frac{1}{c} \left( \mathbf{j}_1 \times \mathbf{B}_s + \mathbf{j}_2 \times \mathbf{B}_f \right)
\]

(6)

If we suppose that the filament is force free or pressure balanced we may neglect the term \( \frac{1}{c} (\mathbf{j}_f \times \mathbf{B}_f) \). If one integrates \( f_L \) over a circular slice of unitary thickness the term with \( \frac{1}{c} (\mathbf{j}_1 \times \mathbf{B}_f) \) integrates out because of the helical symmetry conditions. The same holds for the force component \( \frac{1}{c} (\mathbf{j}_1 + \mathbf{j}_f) \times \mathbf{B}_s \) so that the net force is given by

\[
F_L = \int f_L \, dS = -\frac{\pi a^2}{4c} j_{\|} B_s
\]

(7)

where \( j_{\|} = (\mathbf{j}_1 + \mathbf{j}_f)_{\|} \). The total force on the filament with length \( l \) is then

\[
F = \frac{\pi a^2 l}{4c} j_{\|} B_s
\]

(8)

If \( I = \int j_{\|} \, dS \) is the total current then

\[
F = \frac{I B_s l}{c}
\]

(9)

An upper estimate of \( j_{f\|} \) is made by \( j_{f\|} \approx c B_f/4\pi a \) which results in \( I_f/c \approx a B_f/16 \). The induced current \( I_\perp \) can be estimated from the induction equation \( L I = \Phi/c \), where the magnetic flux \( \Phi \approx B_1 l^2 \cos(Q_B t) \) is derived from the oscillating part of the magnetic field. Taking as an order of magnitude estimate for the selfinductance \( L \approx 4 c^{-2} l/\pi \) (Ionson, 1982) we find \( I/c \approx (\pi B/4) \cos(Q_B t) \). Hence using eq. (5) and (9) we obtain

\[
F \approx \frac{\pi}{4} B_1 B_0 l^2 \cos(Q_B t) + \frac{\pi}{4} B_1^2 l^2 \cos^2(Q_B t) + \frac{a l^2}{16} B_f B_0
\]

\[
+ \frac{a l^2}{16} B_f B_1 \cos(Q_B t)
\]

(10)

The force in eq. (10) consists of a constant component and a periodic component. The former one shifts the loop to another equilibrium position, while the latter one shakes the whole loop back and forth in the radial direction around this new equilibrium position. Below an estimate will be made of the importance of this constant force. It will be shown that it is of the same order as the force due to gravity at \( r_a \), implying that the loop will not be displaced over a large distance and will, therefore, not be destroyed.
It is plausible that through an interaction of the local stellar field and the coronal field the topology of the resultant field may change due to reconnection. If this process would develop faster than, for example, the beat period the field of the neutron star would reconnect with the coronal field before the loop does experience any significant oscillation. However, in order to treat this problem one should know the details of the plasma resistivity (presumably anomalous) in prescribed, little volumes of the loop. This is not possible, for our understanding of magnetic fields in accretion disks is poor. Therefore, at this moment we want to explicitly confine ourselves to the less dramatic case of an oscillatory filament.

If \( B_1 < B_0, B_f < B_0 \) and \( a << l \) it is seen that the major component of the oscillating force is the first term in equation (10); i.e.

\[
F \sim B_1 B_0 \dot{\chi}^2 \cos(\Omega_B t) \tag{11}
\]

The loop is driven in forced oscillation by the varying neutron star magnetic field. In order to estimate the power consumption, \( P(t) \), we consider the whole loop as a forced oscillator with mass \( m \), repulsive elastic force \(-kx\), where \( x \) is the deviation from equilibrium, and damping force \(-\lambda \frac{dx}{dt}\). The equation of motion of the filament is then given by

\[
m \frac{d^2 x}{dt^2} + \lambda \frac{dx}{dt} + kx = F_0 \cos \Omega_B t, \tag{12}
\]

where \( F_0 \), using (11), is given by

\[
F_0 = B_1 B_0 \dot{\chi}^2 \tag{13}
\]

The solution of (12) consists of a homogeneous solution of a damped free oscillation with a characteristic frequency \( \omega_0 = (k/m)^{1/2} \) and a forced oscillation with frequency \( \Omega_B \) which is independent of the initial conditions. The free oscillations are damped so that finally the system is forced to oscillate with the frequency \( \Omega_B \). There is, however, a phase lag of the oscillation as compared to the forcing term. Therefore, the solution can be written as

\[
x = A \sin(\Omega_B t - \beta) \tag{14}
\]

where \( A \) is the amplitude, given by

\[
A = \frac{F_0 / m}{((\Omega_B^2 - \omega_0^2)^2 + (\frac{\lambda}{m} \Omega_B^2)^2)^{1/2}} \tag{15}
\]
and $\beta$ is the phase angle, given by

$$\tan \beta = \frac{m}{\lambda} \left( \frac{Q_B^2 \omega_0^2 - \omega^2}{Q_B^2} \right)$$  \hspace{1cm} (16)

The oscillation is damped and the amplitude grows until the damping force balances the driving force. The maximum amplitude occurs when there is amplitude resonance. In the damped system this occurs when the frequency $Q_B$ is slightly smaller than the characteristic frequency $\omega_0$ so that the denominator in (13) is minimal

$$Q_B^2 = \omega_0^2 - \frac{\lambda^2}{2m^2}$$  \hspace{1cm} (17)

From eq. (14) and (15) we can derive the velocity as a function of time

$$v = \frac{dx}{dt} = Q_B A \cos (Q_B t - \beta)$$  \hspace{1cm} (18)

where the velocity amplitude is given by

$$v_0 = \frac{F_0}{(mQ_B - \frac{k}{Q_B})^2 + \lambda^2}^{1/2}$$  \hspace{1cm} (19)

The maximum amplitude is reached when

$$Q_B = \left( \frac{k}{m} \right)^{1/2} = \omega_0$$  \hspace{1cm} (20)

Then the phase angle $\beta = 0$ and the velocity is in phase with the force. This is called energy resonance because the power delivered by the force is a maximum. The denominator in (19) is the impedance $Z$ given by

$$Z^2 = X^2 + R^2$$  \hspace{1cm} (21)

where the reactance $X$ is

$$X = mQ_B - \frac{k}{Q_B}$$  \hspace{1cm} (22)

and the resistance $R = \lambda$. From this it follows that $\tan \beta = X/R$ and $v_0 = F_0/Z$. The power produced by the force is given by

$$P(t) = F \cdot \dot{v} = \frac{F_0^2}{Z} \cos Q_B t \cos (Q_B t - \beta)$$  \hspace{1cm} (23)

Averaged over a period the absorbed power is

$$\langle P \rangle = \frac{F_0^2}{2Z} \cos \beta = \frac{1}{2} F_0 v_0 \cos \beta = \frac{F_0^2 R}{2Z^2} = \frac{1}{2} R(v_0)^2$$  \hspace{1cm} (24)
We observe that the damping should be known in order to estimate the velocity amplitude and the power. The damping of an oscillating filament in the solar corona has been studied by Kleczek and Kuperus (1969). They assumed that the main damping mechanism is acoustic radiation of the vibrating filament with an effective surface area $S$ into the ambient non-magnetic corona. They found good agreement with observed damping times of oscillating quiescent prominences after the excitation by a flare disturbance.

Adopting their analysis we find as an estimate of the damping time

$$\tau_D \sim \frac{m}{\rho_C v^c_s S}$$  \hspace{1cm} (25)

where $\rho_C$ is the coronal density and $v^c_s$ is the coronal sound velocity. When the ambient medium is magnetically dominated one should use the fast mode velocity instead of $v_s^c$, which means in practice using the coronal Alfvén-speed, $v_A^c$.

From (12) we find that $\tau_D = 2\pi m/\lambda$ so that with $S = 2a\lambda$ we find, using (25)

$$\lambda \sim 4\pi a \rho_C v_A$$  \hspace{1cm} (26)

It should be noted that in this derivation we assume the acoustic damping to be larger than any other internal damping by viscosity and resistivity. If these were to be larger, $\lambda$ would be larger and so would be $R$, causing an increase in the power consumption, provided the system is out of resonance. Therefore our analysis yields a lower limit of the power consumption from the neutron star.

The Alfvén speed is given by $v_A^c = (B^2/4\pi\rho_c^c)^{1/2}$ and the magnetic field in the corona can be of the order of $2 \times 10^7$ G, where we have assumed a dipole field from the star with a value of $\sim 5 \times 10^9$ G at the pole and a typical magnetospheric radius of $\sim 60$ km (cf. eq (1)). In order to get an estimate for the damping constant, $R = \lambda$, we express $\lambda$ in units of $10^6$ cm, $a$ in units of $10^5$ cm and $\rho_c$ in units of $10^{-5}$ g/cm$^3$, where we have assumed the density in the corona to be about an order of magnitude smaller than in the disk (cf. e.g., Shapiro and Teukolsky, 1983). From eq. (26) we then find

$$R = \lambda \sim 2.2 \times 10^{16} a \rho_c^{1/2} \text{g/sec}$$  \hspace{1cm} (27)

An estimate of $v_0$ can be made in the following way. Notice that $v_0 = F_0/Z = B_0 B_0 a^2 / Z$. The value of $Z$ is given by $Z = (X^2 + R^2)^{1/2} = (\lambda^2 + m^2 c_B^2 - 2\pi k + k^2 / Q_B^2)^{1/2}$. The mass of the filament can be written as $m = \pi a^2 \rho_c$ and is of the order $3 \times 10^{11}$ g. The beat frequency $\Omega_B$ is determined by the rotation period of the neutron star, which is of the order of 10 ms, and by the radius of the
magnetosphere, \( r_a \), which is of the order of 60 km. We then find, using eq. (4), \( \Omega_B \sim 300 \text{ sec}^{-1} \). The value of \( k \) is determined by the eigenfrequency of the loop, \( \omega_0 = (k/m)^{1/2} \). We may estimate \( \omega_0 \) by using the Alfven crossing time for the loop: i.e., \( \omega_0 = \frac{k}{2} \) (Isonson, 1982), which leads to a value of the order \( \omega_0 \sim 5 \times 10^3 \text{ s}^{-1} \). Using the value we found for \( m \), this gives \( k \sim 7.5 \times 10^{18} \text{ g/sec}^2 \). The quantities found so far can be used to approximate \( Z \). We find

\[ Z \sim 3.3 \times 10^{16} \text{ g/sec} \]  

Furthermore, we assume \( B_1 \sim 0.1 B_0 \sim 2 \times 10^6 \text{ G} \). We then find

\[ v_0 \sim 1.2 \times 10^9 \text{ cm/sec} \]  

From this we may estimate the power in the oscillations, using eq. (24).

\[ \langle P \rangle = \frac{1}{2} R (v_0)^2 \sim 2 \times 10^{34} \text{ erg/sec} \]  

Using some of the quantities derived above we can now compare the constant force on the filament (see eq. 10) with the gravitational force at the magnetospheric radius. Their ratio is given by \( \frac{(a/16)B_fB_0/m_\chi(GM/r_a^2)}{R} \) and is of order 1, implying that the filament will probably not be displaced over a large distance.

Equation (29) for the velocity amplitude of the filament would imply a displacement \( A \) (see eq. 14) of \( A = \frac{v_0}{\Omega_B} \sim 4 \times 10^6 \text{ cm} \). This value is, however, somewhat larger than the length of the loop. This problem can be solved by increasing the damping of the oscillation. In that case the velocity amplitude would decrease, reducing \( A \) but, however, also the average power. An interesting possibility is that the loop is in resonance.

For the loop to be in resonance \( \omega_0 \sim \Omega_B \sim 300 \text{ sec}^{-1} \), which would imply \( \lambda \sim 10^7 \text{ cm} \). In this case \( Z = R = \lambda \sim 2.2 	imes 10^{18} \text{ cm} \), if again we assume \( a \sim 0.1 \lambda \). For \( v_0 \) we then find \( v_0 \sim 2 \times 10^9 \text{ cm/sec} \), implying a displacement amplitude of \( 6 \times 10^6 \text{ cm} \). This value is now smaller than the length of the loop and does not create the same problem as mentioned above. In this case the power is given by

\[ \langle P \rangle \sim 4 \times 10^{36} \text{ erg/sec} \]  

If we, furthermore, would assume a somewhat stronger modulation field \( B_1 \) (e.g. \( B_1 \sim 4 \times 10^6 \text{ G} \)) this power could be as large as \( 1.5 \times 10^{37} \text{ erg/sec} \). This last estimate is probably an overestimate but if we compare the values with the
Eddington limit, $L_{\text{Edd}} \sim 10^{38}$ erg/sec, of a neutron star we see that one isolated coronal disk loop, powered by electrodynamic coupling of the neutron star magnetic field, may provide a few percent of the X-ray flux.

3. Discussion and conclusion

In section 2 we have shown that the neutron star may lose (rotational) energy to the disk-corona-system via the interaction of its magnetic field with magnetic structures (i.e. loops) in the corona, which are anchored in the disk. The energy that is transferred directly to the disk can be radiated thermally. However, the energy transferred to the corona cannot be radiated efficiently but is probably lost due to comptonisation of "soft" photons coming from both the neutron star and the underlying disk. In the simple picture of disk accretion half of the gravitational binding energy, which is liberated when the matter falls from "infinity" to the magnetospheric radius, $r_a$, can be radiated by the disk. Therefore

$$L_{\text{disk}} = \frac{GM^2}{2r_a}$$

(32)

Using $M = 1 \times 10^{18}$ g/sec and $r_a = 60$ km we find

$$L_{\text{disk}} \sim 1.5 \times 10^{37} \text{ erg/sec}$$

In the previous section we showed that the energy transferred to one loop can be of the order of $10^{34}$ to $10^{37}$ erg/sec. Since it seems plausible to assume that more than one loop is present, the luminosity coming from the disk may be increased by a substantial fraction. As pointed out by Priedhorsky (1986) this transfer of energy from the rotating neutron star to the disk may explain why in some LMXB, such as Sco X-1, the ratio between the luminosity coming from the disk and that coming from the star is so large.

Another point of interest is the fact that the mechanism discussed in this paper may give a natural explanation for the observed quasi-periodic oscillations in some LMXB, while in the mean time it may solve some of the problems posed by some other models for QPO. This can be seen in the following way. The instantaneous power transferred from the star to the loop is given by eq. (23), i.e.

$$P(t) = \frac{F^2}{2} \cos Q_B t \cos (Q_B t - \beta)$$

$$= \frac{F^2}{2z} [(\cos 2Q_B t) \cos \beta + (\sin 2Q_B t) \sin \beta + \cos \beta]$$

(33)

If, for reasons of simplicity, we assume that $\beta = 0$ (i.e., the system is in
resonance) then for one loop it holds that

\[ P_1(t) = \frac{F_{01}}{2\pi I}[\cos 2\Omega_B t + 1] \]

If we furthermore assume that there are \( N \) loops at different azimuthal positions, \( \phi_i \), on the disk, and, furthermore, that the energy transferred to these loops can be reradiated instantaneously, then the total luminosity from the system (star + disk) is given by

\[ L_{\text{tot}}(t) = L_0(t) + \sum_{i=1}^{N} A_i [\cos(2\Omega_B(t-t_i) + \phi_i) + 1]f(t-t_i) \quad (34) \]

where \( A_i = \frac{F_{01}^2}{2\pi I} \) and \( f(t) \) is an envelope function, representing the finite lifetime of each loop (Burw, 1986) As pointed out by Lamb et al. (1985) a luminosity as given by eq. (34) leads to a power spectrum with a broad peak centered around frequency \( f_p = \Omega_B/\pi \) and a low frequency component (i.e., red noise). The latter is due to the presence of the term \( A_i f(t-t_i) \) in eq. (34) or, in other words, the presence of an average contribution from the loop, which is different from zero, leads to a red noise component in the power spectrum.

The picture presented so far can explain all the features of the observation of QPO and red noise in sources like GX5-1 and Cyg X-2, in the same way as the models of Lamb et al. (1985) or Berman and Stollman (1986a). A problem which is inherent in these models and in any beat frequency model, that is based on the modulation of the accretion, is that one does not expect any accretion and therefore QPO to be present when the magnetospheric radius is larger than the corotation radius. However, in some of the sources (e.g., Sco X-1 or the Rapid Burster) the intensity is sometimes anti-correlated with the observed QPO-frequency. Within the context of a beat-frequency model this would imply that the magnetospheric radius is larger than the corotation radius. This problem can be overcome if one introduces an arbitrary parameter, which relates the accretion rate, \( \dot{M} \), to the observed intensity, \( I \), (Lamb et al. 1985; Berman and Stollman, 1986b).

The model presented in this paper is a beat frequency model but is not based on the modulation of the accretion flow, and the mechanism presented does also work when the magnetospheric radius is larger than the corotation radius since there is still a coupling between the rotating field and the loops on the disks. This may lead to an anti-correlation between the frequency and accretion rate.

Another fundamental problem which has to be explained by any model of QPO, is the fact that in some sources (e.g., Sco X-1) the power in the red noise is
much less than that in the QPO-peak (van der Klis et al. 1986). This is very difficult to explain within the context of a beat frequency model, which is based on the modulation of the accretion flow. This can be seen in the following way. The accretion modulation from the blobs in the model of Lamb et al. (1985) or from the accretion sites in the model of Berman and Stollman (1986a) can be described by a function of the form of eq. (34), where the term $1$ is replaced by an arbitrary constant $B_i > 1$. The term $A_iB_if(t-t_i)$, which thus results in eq. (34), represents the nonmodulated or average accretion from each blob. This can never be zero, since then, due to the term $A_i\cos(2\Omega_B + \phi_i)$, the accretion luminosity could be negative. As pointed out above the term $A_iB_if(t-t_i)$ causes the red noise component in the power spectrum. The ratio, $r$, of the power in the red noise and the power in the QPO-peak is given by (Lamb et al., 1985)

$$\text{(35)}$$

$$r = 2 \left< B_i^2 \right> > 2$$

The model presented in this paper, in principle, can give $r = 0$. Consider again eq. (33). If we set $\beta = \frac{1}{2} \pi$ then

$$P_i(t) = \frac{F_{oi}^2}{2Z_i} \sin(2\Omega_B t)$$

and

$$L_{tot}(t) = L_0(t) + \sum_{i=1}^{N} A_i \sin(2\Omega_B t + \phi_i)f(t-t_i)$$

This expression for $L_{tot}$ leads to a power spectrum with again a broad peak centered on frequency $f_d = Q/\pi$, but with no red noise component, since the average contribution from each loop is zero. However, again it must be stressed that this conclusion is based on the assumption that the power transferred back and forth between the loops and the neutron star is instantaneously "translated" into radiation. In particular, in the case that the loops transfer energy to the star, we assume that this energy is extracted from the disk leading to a reduction of the flux coming from the disk. Whether this is a valid assumption must be investigated in more detail. However, this lies outside the scope of this paper.

Our conclusion is that the mechanism described in this paper, in which energy is transferred from the rotating, magnetized neutron star to the disk, can explain some of the observations of LMXB. In particular the model seems suited to describe the properties of the power spectra, observed from some sources, in which a broad peak (QPO) and a red noise component is present. Although the model is based on a beat frequency model, it does not seem to have some of the difficulties which are present in the beat frequency models that are
based on a modulation of the accretion flow.

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