

MAGNETIC FIELD EVOLUTION OF MILLISECOND
PULSARS BEYOND THE X-RAY DEATH LINE

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Dedicated to my family.

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Abstract

Under the relativistic plasma diffusion theory (Kebede, 2002, 1996), we investigate the magnetic field evolution of millisecond pulsars beyond the X-ray death line by taking the competitive nature of magnetic field revival due to mass accretion and its decay due to one of the cooling sources called photon emission leading to pulsar field evolution at about threshold magnetic field for X-ray productions into account. Our study shows that the magnetic field is no more monotonous for old pulsars. This evolution of the magnetic field is applied to address the issue of the 35 day cycle of X-ray radiations of Hercules X-1.

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Introduction

At a present there is no as such a self-consistent theory for the origin of Neutron Star (NS) magnetic fields. The current understanding is that it can either be a fossil remnant (the standard picture) or it may also be generated by surface thermal processes soon after the formation of the NS [1, 2]. Models for field decay may follow either of these venues even though both of them have their own share of weakness. Very recently, however, separated charges have been suggested as a more likely source for NS magnetic fields [3, 4, 5]. In particular [5] has shown that charge diffusion driven by the huge plasma density gradient inherent to NS matter results in separated charges large enough to generate surface magnetic fields which, very early in the life time of a typical NS are at least two orders of magnitude stronger than those normally observed from young pulsars such as the Crab or Vela pulsars ($\sim 10^{12}\text{G}$). Recently pulsars data indicate that there are a good number of young pulsars (age $10^3 - 10^4$ years) having surface magnetic fields greater than 10^{14}G .

The author to the new model has shown [4] that magnetic and viscous coupling between the various parts of a neutron star will, according to this model, lead to dynamic properties of the magnetic moment which are totally foreign to the standard model. These extra dynamical features have been successfully utilized to address such worrisome issues in pulsar astronomy today as the paucity of long period pulsars, missing pulsars and delayed pulsar onset [6, 7, 8].

Investigation indicates that no serious magnetic decay occurs in a NS life time [9].

This raises question of how magnetic moments of the so called millisecond pulsars (MPs) in low-mass X-ray Binaries (LMXBs) decay from a typical value of $\sim 10^{30} Gcm^3$ down to $10^{25-26} Gcm^3$ in a time window of about 10^{7-8} years [10]. Obviously, Ohmic decay alone cannot be responsible for the indicated low magnetic moment. In fact, whether or not magnetic moments to pulsars undergo Ohmic decay is still an open question [11]. Whatever little effect there could be Ohmic decay, it is further suppressed by space-time curvature scenarios [12]. The problem of low magnetic moments of MPs has recently prompted a series of semi-empirical models for magnetic moment decay [13] including such approaches as the one involving mass-accretion induced decay scenarios [14]. The effects of mass-accretion induced decay scenarios, however, are expect to be over-shadowed by post-accretion in cease of the magnetic moment due to re-diffusion of the buried fields toward the surface [15]. In contrast, however, the new model [4, 5] suggests that NS magnetic fields are internal temperature and rotational frequency dependent indicating that pulsars fields decay primarily due to neutrino and photon emissions which according to standard calculations are responsible for NS cooling.

Most importantly it also hints to the fact that pulsars fields could go through a scenario of field revival whenever the magnetic pressure is weak enough to admit mass accretion onto the NS surface. This implies a possible competition between mass accretion and cooling sources leading to a unique scenario of pulsar field evolution at about the threshold field for X-ray production.

In this thesis I investigate the nature of the evolution of pulsar magnetic fields past the X-ray death line and apply it to a known X-ray source named Hercules X-1 which is known to have not a well understood 35 day cycle in its X-ray spectrum.

This thesis consists of four chapters. The first chapter is about neutron star cooling due to neutrino and photon emissions. Here we will derive an expression for the magnetic field decay due to photon emission during mass accretion. The second chapter

discusses about accretion, disc accretion and its formation. In this chapter we will calculate the magnetic field at which the neutron star admits accretion and also the time evolution of magnetic field due to mass accretion based on [5] model. Chapter three is devoted to magnetic field evolution of millisecond pulsars past the X-ray death line which is our ultimate goal. Based on the first two chapters we will investigate the competition between the magnetic fields found in chapters one and two. It will be shown that this field evolution fluctuates about the threshold field required for the production of X-ray radiation. In chapter four we made our conclusion based on the results found in the previous chapters.

Chapter 1

Neutron Stars Cooling Due to Neutrino and Photon Emissions

1.1 Neutron Star Cooling

The cooling of neutron stars offers a unique diagnostic of the physics of the interior. Because the emergent spectrum is nearly blackbody, the observed flux can be modelled to estimate the temperature, the intervening column of absorbing material, and the ratio of radius to distance. Astrometric measurement of the distance then allows an estimate of the photospheric radius, which is sensitive dependent on the high density equation of state. A simple model that adequately accounts for observations from the optical through the X-ray bands includes a power law component identified with non-thermal emission from the active magnetosphere, and two blackbody components: a harder component identified with the heated polar caps, and a softer component identified with emission from the entire photosphere. A blackbody radius as small as $R_\infty = 8\text{km}$ would pose serious challenges to standard neutron star models (e.g., Heiselberg & Pandharipande 2000), though it would be acceptable for a low mass strange star (Alcock, Fahri, & Olinto 1986). However, a more likely explanation for this anomalous result could be failure of the simple blackbody modelling [27].

When the iron-nickel core ($M_{\text{core}} \simeq 1.5M_\odot$) of a massive star becomes unstable and

collapses to nuclear density, roughly 10% of its gravitational energy, about 10^{53} erg is released. Most of this energy is emitted almost immediately in the form of neutrinos (see [35] and reference therein). Therefore, the internal temperature of a neutron star drops rapidly from about 10MeV to about 1MeV, and the matter is to all extents and purposes in its ground state. The Fermi energies of the neutrons, protons, and electrons are much larger. Neutrino emission dominates the further cooling of the neutron star, until the interior temperature falls to about 10^8 k which roughly corresponds to a surface temperature of about 10^6 k. Only then does the photoemission also begin to play an important role.

At the beginning of the collapse, the neutrinos are produced mainly by electron capture. Electron capture is a process whereby a proton in the nucleus absorbs an orbiting electron converts to a neutron. The mean free path of these neutrinos soon becomes so short that they remain trapped during the continued collapse and neutrino trapping and a highly degenerate neutrino gas is formed. The cooling rate of a not-too-old star (with an age $t < 10^5$ yr) is determined by the neutrino luminosity of its inner layers [16]. The surface temperature as a function of time of the neutron stars depends on a number of interesting aspects of the physics of neutron stars. The equation of state at high densities is particularly important, as is the possible existence of a pion condensate in the central region of a neutron star. Magnetic fields and the possible existence of a superfluid state of the nucleons also play some role.

Observation of the Crab pulsar during lunar eclipses have shown that the effective surface temperature T_s is less than 3.0×10^6 k. Since the age of the Crab is known, this upper limit provides an important restriction on models of neutron stars.

The Einstein observatory has made it possible to determine also upper limits for the thermal emission temperature of other suspected neutron stars contained in young supernova remnants [17].

One might expect to find a neutron star near the center of the remnant of every

historical supernova, but this is definitely the case only for the Crab nebula. The observations force us to conclude either that neutron stars are not formed in other supernova explosions or that these have cooled off so fast that their present temperatures are below the given limits. We remark that as a result of interstellar X-ray absorption, these limits have an uncertainty of about 50% [18].

1.2 Magnetic Field Decay Due to Photon Emission

Neutron stars are born in supernova explosions at very high internal temperatures greater than 10^{11}k and eventually cool down to various cooling mechanisms. The internal evolution of a neutron star is indicative of such internal properties of the interior as viscous dissipation processes at the superfluid-crust interface, transport coefficients of neutron matter etc

Standard calculations indicate that neutrino emission is the dominant cooling mechanism for the first 10^5 years [16]. Neutrino producing reactions are the following type



where $\bar{\nu}_e$ means an antineutrino



Photon emission overtakes neutrinos only when the internal temperature falls to about 10^8k with a surface temperature of about 10^6k . The two temperature are interrelated by

$$T_s \sim (10T)^{2/3} \quad (1.2.3)$$

where T_s is the observed surface temperature.

The cooling rate due to photon emission can be calculated from

$$\frac{dT}{dt} = \frac{\int g_{00}^{-1} n \epsilon_\gamma dv_p}{\int n C_\nu dv_p} \quad (1.2.4)$$

where dv_p is the proper volume element, g_{00} is the metric tensor, n is the baryon density, and ϵ_γ is the photon emissivity per baryon defined by

$$\frac{L_\gamma}{4\pi R^2} = \epsilon_\gamma \quad (1.2.5)$$

L_γ being the photon luminosity.

$$L_\gamma \sim \sigma T^4 \quad (1.2.6)$$

and σ is the Steffen-Boltzmann constant. Also C_v is the specific heat per baryon given by

$$C_v = \frac{\pi^2(\chi^2 + 1)^{1/2}}{\chi^2} \left(\frac{k_\beta^2 T}{M_\beta c^2} \right) \quad (1.2.7)$$

where k_β is the Boltzmann constant, m_β is the baryon rest mass and

$$\chi = \frac{P_F}{m_\beta c^2} \quad (1.2.8)$$

where P_F is the Fermi-momentum of the degenerate baryon system.

The magnetic decay law calculated using the cooling rate indicated above shows that (Kebede, 2005-under peer review)

$$B_p = B_2^{(0)} \text{yr}^{3/2} \left[k 1.47 \times 10^{13} \left(\frac{M}{M_\odot} \right)^{1/2} T_2^{(0)-1} \right] t^{-3/2} \quad (1.2.9)$$

where t is in years, $B_2^{(0)} \simeq 10^{12}\text{G}$ and $T_2^{(0)} \simeq 10^6\text{k}$ are surface magnetic field and internal temperature respectively at $\sim 10^5$ yr and M_\odot is the mass of the sun. Temperature variations known to occur at certain spots on the surface of a neutron star (in a binary) such as the polar caps heated as a result of the trickling accreted gas flow guided by the surface magnetic fields are of no great consequences to our model. This is essentially because these spots cover very small areas as compared to the surface area of the neutron star and the suspected heat transfer to the rest of the neutron star surface is constrained to a relaxation time of $\sim 10^4$ years [19, 32]. We show that past the threshold value of $\sim 10^7\text{G}$ at which the ram pressure exceeds the

magnetic pressure the magnetic field passes through short periodic fluctuations that will not permit constant transfer of heat energy from the polar region to the rest of the neutron star surface.

In fact, the suggested heat transfer [19] is not practical essentially because the process is highly temperature gradient dependent indicating that the generated heat is automatically dissipated to the colder and highly conducting plasma in contact to the very cold interstellar medium having little chance to be transferred to the other parts of the neutron star surface which is naturally not.

Inserting the above indicated values for $B_2^{(0)}$, $T_2^{(0)}$ and $(\frac{M}{M_\odot}) = 1.4$ into eq.(1.2.9), we get

$$B_p \simeq 10^{19} t^{-3/2} \quad (1.2.10)$$

where B_p is in Gauss

If t is in days, then one finds

$$B_p \simeq 10^{23} t^{3/2} \quad (1.2.11)$$

Taking logarithms on both sides, this takes the form

$$\log_{10} B_p \simeq 23 - \frac{3}{2} \log_{10} t \quad (1.2.12)$$

Accretion starts at about $t = 10^8 \text{ yrs} \simeq 4 \times 10^{10} \text{ days}$ (see section 2.3), therefore transforming $\log_{10} t$ to $\log_{10} t + \log_{10}(4 \times 10^{10})$, eq.(1.2.12) results in

$$\log_{10} B_p \simeq 7 - \frac{3}{2} \log_{10} t \quad (1.2.13)$$

Chapter 2

Accretion in Binary Systems

Introduction

In this chapter we will discuss about accretion in neutron star, disc accretion and its formation. We will also calculate the magnetic field at which the neutron stars admits accretion. Finally we will investigate the magnetic field revival due to accretion based on [4, 5] model.

2.1 Disc Accretion in X-ray Binary System

By accretion, we mean the accumulation of diffuse gas or matter onto some object under the influence of gravity. Neutron stars in accreting pulsar systems are often called "X-ray pulsars" or "accretion -powered" pulsars. Accreting X-ray neutron stars provide a very interesting contrast to the spin-down of isolated neutron stars. These binary pulsars contain a rotating high-magnetic-field neutron star in orbit with a stellar companion which transfers mass to the neutron star via Roche-lobe Overflow or a stellar wind as a result of which the neutron stars are being spun up by the transferred angular momentum. If the accreted matter is close enough to the neutron star, the neutron star's gravitational attraction makes it fall onto the accreting star. However, since the neutron star is tiny, the matter has too much angular momentum to fall on

the star directly and therefore orbits around the star until it has rid itself of most of its angular momentum. This leads to the formation of accretion discs. Accretion will be inhibited by a centrifugal barrier if the neutron stars magnetosphere rotates faster than the Kepler frequency at the magnetosphere. Hence the radius of the inner edge of the accretion disc should be less than the co-rotating radius, otherwise accretion into the star will cease [34].

The accreted matter in a binary system is in a plasma state, which makes it an excellent conductor of electricity [20]. At the accretion rate typical of a binary, a magnetic field affects the motion of matter over several thousands of kilometers, i.e. over distances several hundreds of times greater than the size of the neutron star itself. When affected by a magnetic field matter moves anisotropically, which means that a neutron star will radiate anisotropically too. The stars own rotation will cause the radiation to fluctuate!

Zeldovich's pupil Viktory Shvartsman boldly suggested that the X-ray pulsars should be found in binary systems. His argument was that if a radio pulsar (for instance, the radio pulsar in the Crab Nebula) were placed in a binary system, the pulsar would radiate (eject) powerful fluxes of electromagnetic waves and relativistic particles [fig.2.1a]. The pressure produced by the ejected fluxes would be so high that all the matter in the stellar wind would be "swept out" making accretion impossible. This cannot go on forever. The pulsar will lose its rotational energy and will rotate more slowly. As this goes on, its luminosity and the pressure ejecting the stellar wind will lower. Then, there will come a moment a radio pulsar will "fade" to a degree at which the accreted matter attracted by its gravitation pull will rush to the neutron star's surface. Although the pulsar's radiation has "faded" accretion still impossible. The reason for this is the rapidly rotating magnetic field of the neutron star. Like enormous propeller, it scatters the matters and does not let it fall (fig.2.1b). This effect was latter called the propeller effected by Andrei Illarionov and Rashid Syunyeev.

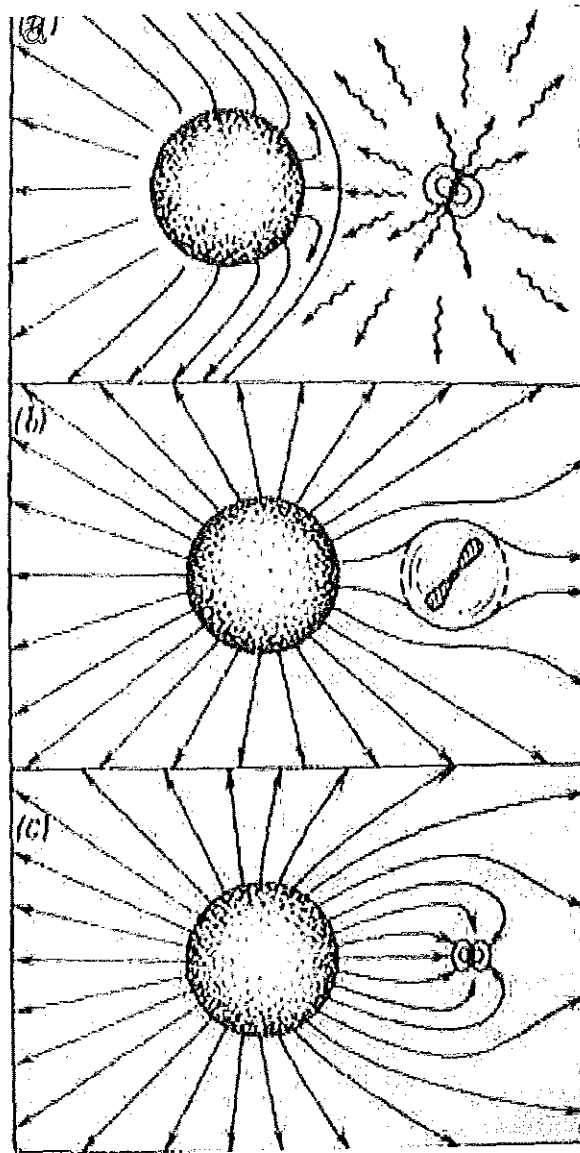


Figure 2.1: Three states of a neutron star in a binary system:
(a) an ejecting pulsar;
(b) a "propeller";
(c) an accreting neutron star

Whilst repelling matter, the pulsars continues decelerating until accretion is possible (fig.2.1c). The accreted matter rushes in to the neutron star's magnetic poles, and the result is a terrific outburst of energy. An X-ray pulsar powered by accretion flares up.

The dynamics of disc accretion is a more complex problem because the interaction between the magnetic field structure and the plasma in the accretion disc has to be understood [23]. Furthermore, the details of the interaction depend upon the angle between the magnetic axis of the compact star and the axis of the accretion disc. The radius within which the magnetic field of the star dominates the dynamics can

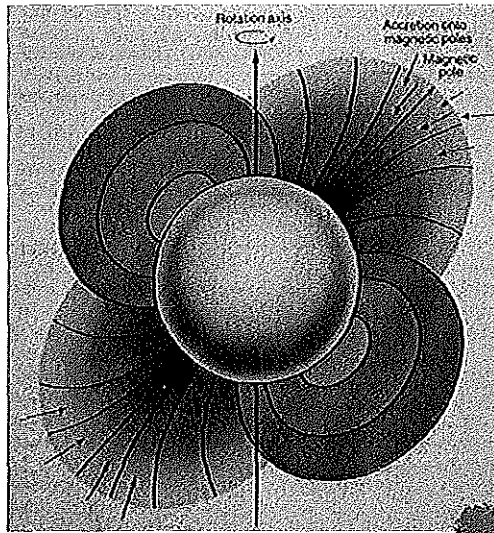


Figure 2.2: A schematic model illustrating how accretion onto a neutron star with a magnetic field nonaligned with the rotation axis can give rise to an accretion column, which, when observed at a distance, gives rise to the observation of X-ray pulses at the rotation frequency of the neutron star. The accreting matter is channelled down the magnetic field lines to the magnetic poles of the neutron star where its binding energy is deposited, resulting in strong heating of the plasma. The problems of radiative transfer under these conditions are highly complex.

be found by equating the torque exerted by the accretion disc on the magnetic field structure to the magnetic torque associated with the distorted magnetic field distribution at that radius. This is a non-trivial calculation, particularly when the effect of instabilities at the interface between the disc and the magnetic field structure have to be taken into account, but, in general, a result similar to that will be derived in the next section for the case of spherical accretion (fig. 2.2) is found. According to Frank, King and Raine (1992), a radius is obtained which is smaller than Aifvén radius [see eq.(2.2.11)] , typically being about half that value. This picture has a number of important consequences. First, it is apparent that, once again, accretion can only take place onto the poles of the compact star, but now the accretion disc feeds matter onto the poles of the neutron star via the dipole magnetic field, as illustrated in (fig.2.3). The net result is the formation of an accretion column with an opening angle which depends upon the angle of inclination between the axis of the magnetic field and the axis of the disc. The second important point is that the accretion disc exerts a torque upon the magnetosphere of the star, which, in turn, transmits the torque to the star itself. Thus, the process of accretion leads to the speeding up of the star, and this is observed in X-ray binary stars. It is also like to be the process which leads to the formation of the millisecond pulsars.

The case for associating the X-ray sources with accretion of matter from the primary onto the neutron star is wholly convincing, as can be appreciated from the following arguments.

First of all, if matter falls from the primary onto the secondary star, the kinetic energy which it gains is converted into thermal energy when the infalling matter hits the surface of the neutron star. The thermal energy released is just the gravitational binding energy of matter on the surface of the neutron star, so that about 5% of the rest mass energy of the infalling matter is available for emission as radiation. This represents a very efficient source of energy, at least five times the efficiency available

from nuclear energy sources. Secondly, let's now ask what the typical temperature of

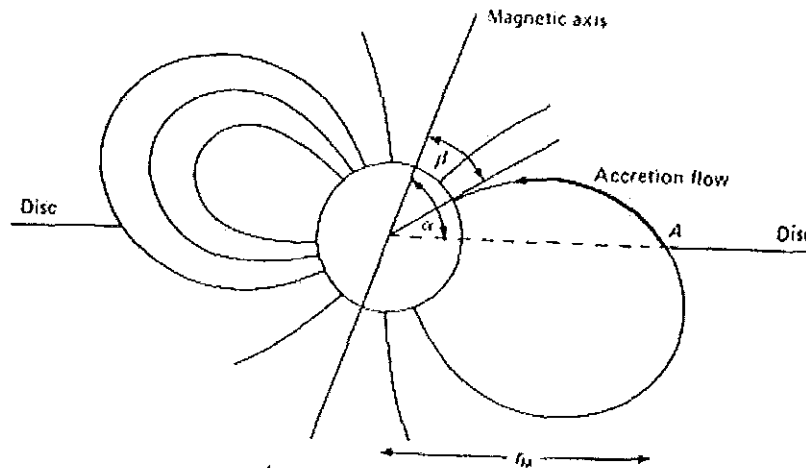


Figure 2.3: Illustrating the accretion of matter from an accretion disc onto the polar caps of a magnetized neutron star. (After J. Frank, A. King and D. Raine(1992). *Accretion power in astrophysics*, P.123. Cambridge: Cambridge University Press).

the radiating matter would have to be to account for the observed X-ray luminosities of binary sources. If we assume that this is emitted as blackbody radiation from the surface of a neutron star, then, equating the radiation flux of a blackbody at temperature T to this luminosity, the lower limit to the temperature of the emitting region is about 10^7 k. Thus, it is entirely natural that the radiation should be emitted in the X-ray waveband.

The third simple argument concerns the steady state X-ray luminosity of accreting compact objects. If the luminosity of the source were too great, the radiation pressure acting on the infalling gas would be sufficient to prevent the matter falling onto the surface of the compact object. In the simplest calculation, it is assumed that the radiation pressure acting on the infalling matter is due to Thomson scattering of the

emergent radiation by infalling electrons. If other sources of opacity are also important, these increase the radiation pressure and result in a lower value for the critical luminosity, above which accretion is suppressed. The resulting critical luminosity is known as the Eddington luminosity. This luminosity depends upon the mass of the gravitational body according to

$$L_{Edd} = 1.3 \times 10^{38} \left(\frac{M}{M_{\odot}} \right) \text{ergs}^{-1} \quad (2.1.1)$$

It is remarkable that the luminosity of most of the binary X-ray sources in the Galaxy and the Magellanic clouds are more or less consistent with this upper limit. Their luminosity extend up to about 10^{31} Watt, and the X-ray luminosity function cuts off sharply above this value. The calculation which leads to the Eddington limit is a very general one and does not depend strongly upon the details of the accretion process or the means by which the X-ray emission is produced.

These three arguments show how naturally accretion can account for properties of these X-ray sources. They also demonstrate the importance of accretion as a source of energy in astrophysics [23].

2.2 Magnetic Field at which the Neutron Star Admits Accretion

The compact stars possess magnetic fields, and these can strongly influence the accretion matter on them. We know the values of the magnetic fields not only from theoretical predictions but because they can actually be measured by observing synchrotron lines created by free electrons being accelerated in the field. When the accreted matter that has been captured by the neutron star (which is a high temperature plasma, consisting of ionized atoms and free electrons) comes close, it is trapped by the magnetic field. From then on, it can only move following the magnetic

field lines. They take it to the magnetic poles of the neutron star.

Now in this section we determine the magnetic field at which accretion is possible, (i.e. the ram pressure exceeds the magnetic pressure).

Let us assume that the magnetic field is dipolar, and hence that at a distance r from the center of the star the magnetic field strength is given by

$$B = B_s \left(\frac{R}{r} \right)^3 \quad (2.2.1)$$

where B_s is that typical magnetic field strength at the surface of the compact star having a stellar radius R .

Let us consider the case of spherical accretion onto the compact star, thus the magnetic pressure at radial distance r is

$$P_{mag} \simeq \frac{B^2}{2\mu} \simeq \frac{B_s^2}{2\mu} \left(\frac{R}{r} \right)^6 \quad (2.2.2)$$

where μ is magnetic permeability.

This magnetic pressure increases steeply as the matter approaches the stellar surface. The infalling matter can be considered to have fallen from rest at infinity, and so its infalling velocity V at radius r is

$$V = \left(\frac{2GM}{r} \right)^{\frac{1}{2}} \quad (2.2.3)$$

The pressure which this infalling gas exerts upon the magnetic field is known as the ram pressure, P_{ram} , and is just the rate at which momentum is transported inwards at radius r per unit area, that is,

$$P_{ram} = \rho V^2 \quad (2.2.4)$$

where ρ the density of the infalling matter.

In case of spherical accretion, we have from the equation of continuity the relation

$$4\pi r^2 \rho V = \dot{M} \quad (2.2.5)$$

where \dot{M} is the mass accretion rate.

The radius, r_M , where this ram pressure is balanced by the magnetic pressure of the compact star is known as Aifvén radius. Mathematically this now be expressed as

$$\frac{B_s^2 R^6}{2\mu r_M^6} = \frac{\dot{M}}{4\pi r_M^2} \left(\frac{2GM}{r_M} \right)^{\frac{1}{2}} \quad (2.2.6)$$

Then solving for r_M one obtains the Aifvén radius to be

$$r_M = (B_s R^3)^{4/7} \left(\frac{2\pi^2}{GM\dot{M}^2\mu^2} \right)^{1/7} \quad (2.2.7)$$

where $R = 10\text{km}$, $M = 1.4M_\odot = 1.4 \times 10^{33}\text{g}$, $B_s = 10^{12}\text{G}$. In order to obtain the value of r_M we first have to determine \dot{M} . The binding energy of an element of gas with mass m in the Keppler orbit which just grazes the surface of the primary is

$$E = \frac{GMm}{R}$$

Since the gas elements start at large distances from the star with negligible binding energy, the total luminosity in a steady state must be

$$L_{disc} = \frac{GM\dot{M}}{2R} = \frac{1}{2}L_{acc} \quad (2.2.8)$$

where L_{acc} is the accretion luminosity.

The other half of L_{acc} has still to be released very close to the star.

Plugging the value of M indicated above in eq.(2.1.1), we approximately get

$$L_{acc} = 1.3 \times 10^{38} \text{ergs}^{-1} \quad (2.2.9)$$

Using eq.(2.2.8) and eq.(2.2.9), we obtain

$$\dot{M} = 1.39 \times 10^{18} \text{gs}^{-1} \quad (2.2.10)$$

Making use of this, r_M found to be

$$r_M = 2 \times 10^7 \text{cm} \quad (2.2.11)$$

Substituting the values of B_s , R , and r_M back into eq.(2.2.1), the required magnetic field becomes

$$B = 10^7 G \quad (2.2.12)$$

2.3 Magnetic Field Revival Due to Mass Accretion

According to the new model [5] the magnetic field of a neutron star get revived as a result of mass accretion from its companion star. We investigate the resulting time evolution the magnetic field in this case.

To this end, consider an accreting millisecond pulsar in a binary system. Obviously its moment of inertia I and angular velocity Ω will be time dependent. Let us assume that its rotational kinetic energy is completely converted to emission as radiation.

So one can write as

$$\frac{d}{dt} \left(\frac{1}{2} I \Omega^2 \right) = \dot{E} \quad (2.3.1)$$

where \dot{E} is the total luminosity in a steady state situation.

Carrying out the differentiation and rearranging, we get

$$\frac{\Omega d\Omega}{\dot{E} - \frac{1}{2} \dot{I} \Omega^2} = \frac{dt}{I} \quad (2.3.2)$$

where the dots represent differentiation with respect to time. Hence, follows

$$\int_{\Omega_{min}}^{\Omega} \frac{\Omega d\Omega}{\dot{E} - \frac{1}{2} \dot{I} \Omega^2} = \frac{1}{\dot{I}} \int_{I_0}^I \frac{dI}{I}$$

After Integration, we obtain

$$\ln \left[\frac{\dot{E} - \frac{1}{2} \dot{I} \Omega^2}{\dot{E} - \frac{1}{2} \dot{I} \Omega_{min}^2} \right] = \ln \left[\frac{I_0}{I} \right]$$

Which solves Ω as

$$\Omega = \Omega_{min} \left[\frac{2\dot{E}}{\dot{I} \Omega_{min}^2} - \left(\frac{2\dot{E}}{\dot{I} \Omega_{min}^2} - 1 \right) \left(\frac{1}{1 + \frac{I}{I_0}} \right) \right]^{1/2} \quad (2.3.3)$$

where

$$\begin{aligned}
I &= I_o + \dot{I}t, \\
\dot{E} &= \frac{GM\dot{M}}{2R} = 6.5 \times 10^{37} \text{ ergs}^{-1}, \\
\dot{I} &= \dot{M}R^2 = 1.39 \times 10^{30} \text{ gcm}^2\text{s}^{-1}, \\
I_o &= M_o R^2 = 1.4M_\odot R^2 = 1.4 \times 10^{45} \text{ gcm}^2, \\
\Omega_{min} &= 6\text{s}^{-1}
\end{aligned}$$

By substituting all the constants with their corresponding values back into eq.(2.3.3), we see that at an accretion rate $1.39 \times 10^{18} \text{gs}^{-1}$ it would take only 10^8 years to spin up a neutron star to an angular frequency of 3696s^{-1} .

We can also see that from eq.(2.3.3) the angular frequency of an accreting neutron star is increasing with time. This spin-up can occur only if the Keplerian angular velocity,

$$\Omega_k(r) = \left(\frac{GM}{r^3} \right)^{1/2},$$

of the matter in the disc is greater than the angular velocity of the star Ω is the same sense as the disc rotation. Spin-up ceases when these angular velocities are equal.

In the new model for neutron star magnetic fields developed by Kebede, 2002 the surface magnetic fields are generated by the of spinning separated charges which came about as a result of plasma diffusion driven by density gradients inherent to the standard neutron matter.

According to this model, the surface magnetic fields at the poles is given by

$$B = \frac{2|Q|\Omega}{3cR} \quad (2.3.4)$$

where Ω is the spin frequency, $|Q|$ is the magnitude of the positive or negative charge. Substituting eq.(2.3.3) into eq.(2.3.4), we find the time evolution of magnetic field

due to accretion to be

$$B_a = \frac{2|Q|\Omega_{min}}{3cR} \left[\frac{2\dot{E}}{\dot{I}\Omega_{min}^2} - \left(\frac{2\dot{E}}{\dot{I}\Omega_{min}^2} - 1 \right) \left(\frac{1}{1 + \frac{\dot{t}}{I_0}} \right) \right]^{1/2} \quad (2.3.5)$$

As we find earlier, the surface magnetic field of an old neutron star ($\sim 10^8$ yrs. old) is about 10^7 G. Using this in eq.(2.3.5), we find

$$|Q| = 5.67 \times 10^{22} esu \quad (2.3.6)$$

Chapter 3

Magnetic Field Evolution of Millisecond Pulsars

3.1 Computation of Magnetic Field

In this section we will investigate the combined effects of the magnetic field decay due to photon emission and its revival due to mass accretion. Let us denote the sum of eq.(1.2.13) and eq.(2.3.5) by $2 \log_{10} B_s(t)$. Hence, one may write

$$2 \log_{10} B_r(t) = \log_{10} B_a(t) + \log_{10} B_p(t)$$

Solving for $B_r(t)$, we find

$$B_r(t) = \left(B_a(t) B_p(t) \right)^{1/2} \quad (3.1.1)$$

We write the total magnetic field as

$$B(t) = W_a(t) B_a(t) + W_p(t) B_p(t) \quad (3.1.2)$$

where $W_a(t)$ and $W_p(t)$ are the weights attached to the magnetic field due to accretion and photon emission, respectively. The weights are subjected to the constraints

$$W_a(t) + W_p(t) = 1 \quad (3.1.3)$$

and are defined by

$$W_a(t) = \left[1 - \left(\frac{B_r(t)}{10^7} \right)^{10^{-3}} \right] \quad (3.1.4)$$

and

$$W_p(t) = 1 - W_a(t) = \left[\frac{B_r(t)}{10^7} \right]^{10^{-3}} \quad (3.1.5)$$

The weights are so defined to reflect the competitive edges of each process on either side of the magnetic field at which mass accretion begins. In the next section it will be shown that this magnetic field determines the threshold for the production of X-ray. Inserting eq.(3.1.3) and eq.(3.1.4) back into eq.(3.1.2), we find

$$B(t) = \left[1 - \left(\frac{B_r(t)}{10^7} \right)^{10^{-3}} \right] B_a(t) + \left[\frac{B_r(t)}{10^7} \right]^{10^{-3}} B_p(t) \quad (3.1.6)$$

where $B_a(t)$ and $B_p(t)$ are given by eq.(2.3.5) and eq.(1.2.15) respectively.

3.2 Estimation of the threshold magnetic field for the production of X-ray radiation

The X-ray waveband is in the range

$$3 \times 10^{16} \leq \nu \leq 3 \times 10^{19} Hz;$$

$$10 \geq \lambda \geq 0.001 nm;$$

$$0.1 \leq E \leq 100 keV.$$

As in the case of the far-ultraviolet waveband, the atmosphere is opaque to X-ray because of photoelectric absorption by the atoms which make up the molecular gases of the atmosphere. The detectors of these energetic photons begin to resemble the detectors used in particle physics experiments. Proportional counters and scintillation detectors are used as well as other devices such as Charge Coupled Devices (CCDs)

in which the total energy deposited by the X-ray on entering the detector can be measured. X-ray astronomy is wholly carried out from above the atmosphere. The photons are of such high energy that they behave like particles and the telescopes for high energy X-rays are essentially collimators in which the resolution of the telescope is determined by the geometric design of the collimator. At low X-ray energies, $0.1 < E < 1\text{keV}$, grazing incidence optics can still be used to image the X-rays to a focal plane, but at higher energies the grazing incidence angles are so small that enormously long telescopes would be needed to focus the X-ray image [22].

3.2.1 The Electromagnetic Field Tensor

Because the vector potential is the space like component of the four-potential, we would expect the magnetic induction field to be given by the the four-dimensional curl of the four potential. We introduce the scalar and vector potentials, ϕ and \mathbf{A} respectively, in order to simplify the evaluation of the vector fields \mathbf{E} and \mathbf{B} at a distance \mathbf{r} from the accelerated charge

$$\mathbf{E} = -\frac{\partial \mathbf{A}}{\partial t} - \nabla \phi \quad (3.2.1)$$

$$\mathbf{B} = \nabla \times \mathbf{A} \quad (3.2.2)$$

In tensorial language we write these as

$$E_i = \partial_i A_0 - \partial_0 A_i \quad (3.2.3)$$

$$B_{ij} = \partial_i A_j - \partial_j A_i \quad (3.2.4)$$

where

$$B_{12} = B_3, B_{23} = B_1, B_{31} = B_2, \dots \quad (3.2.5)$$

The obvious covariant generalization is to define the electromagnetic field tensor as

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \quad (3.2.6)$$

where

$$A_\mu = (\phi, -\mathbf{A}) \quad (3.2.7)$$

$$\mu = 0, 1, 2, 3$$

The antisymmetric tensor $F_{\mu\nu}$ is gauge invariant (under $A_\mu \rightarrow A_\mu + \frac{\partial\chi}{\partial x^\mu}$) and has the right number of components to generate the electromagnetic field:

$$F_{ij} = \partial_i A_j - \partial_j A_i = B_{ij} \quad (3.2.8)$$

and

$$F_{i0} = \partial_i A_0 - \partial_0 A_i = E_i \quad (3.2.9)$$

where $i, j = 1, 2, 3$

Maxwell equations are obtained from the somewhat less elegant third-rank tensor equation

$$\partial_\sigma F_{\mu\nu} + \partial_\mu F_{\nu\sigma} + \partial_\nu F_{\sigma\mu} = 0 \quad (3.2.10)$$

This expands as

$$\nabla \cdot \mathbf{B} = 0 \quad (3.2.11)$$

and

$$\nabla \times \mathbf{E} = -\frac{1}{c} \frac{\partial \mathbf{B}}{\partial t} \quad (3.2.12)$$

3.2.2 The Pulsar magnetosphere

The induced electric field of a rapidly rotating magnetized neutron star are so strong that the region surrounding the star cannot be vacuum, but must contain a substantial space charge. This was first pointed out by Goldreich and Julian [37]. We repeat here their argument.

Consider the simple case of an aligned dipole field, where an angular frequency Ω

is parallel to magnetic moment vector μ , and assume that the star is surrounded by vacuum and its spin axes is along the z-axis. In polar coordinates:

$$\Omega = \Omega(\cos\theta\hat{r} - \sin\theta\hat{\theta}) \quad (3.2.13)$$

$$\mu = \mu(\cos\theta\hat{r} - \sin\theta\hat{\theta}) \quad (3.2.14)$$

The corresponding dipolar magnetic field \mathbf{B} is

$$\mathbf{B} = \frac{2\mu}{r^3}(\cos\theta\hat{r} + \frac{1}{2}\sin\theta\hat{\theta}) \quad (3.2.15)$$

The magnitude of \mathbf{B} at the poles is thus

$$B_s = \frac{2\mu}{R^3} \quad (3.2.16)$$

where R is the stellar radius. For a given magnetic field \mathbf{B} , the electric field inside and outside of the neutron star surface are given by

$$\mathbf{E}^{in} = \frac{1}{c}(\Omega \times \mathbf{r}) \times \mathbf{B} \quad (3.2.17)$$

Using eq.(3.2.15) these expand to

$$\mathbf{E}^{in} = \frac{B_s R^3 \Omega}{cr^2} \left(\frac{1}{2} \sin^2\theta\hat{r} - \sin\theta\cos\theta\hat{\theta} \right) \quad (3.2.18)$$

and

$$\mathbf{E}^{ext} = -\nabla\phi \quad (3.2.19)$$

Where

$$\phi = C \frac{3\cos^2\theta - 1}{r^3} \quad (3.2.20)$$

where C is the constant to be determined.

The θ -component of the external electric field is

$$E_{\theta}^{ext} = -\frac{1}{r} \frac{\partial\phi}{\partial\theta} = 6C \frac{\sin\theta\cos\theta}{r^4} \quad (3.2.21)$$

From the continuity of the tangential component of \mathbf{E} at R then follows that

$$C = -\frac{B_s R^5 \Omega}{6c} \quad (3.2.22)$$

This eventually leads to

$$\begin{aligned} E_r^{ext} &= -\frac{B_s R^5 \Omega}{2c} \frac{3 \cos^2 \theta - 1}{r^4}, \\ E_\theta^{ext} &= -\frac{B_s R^5 \Omega}{cr^4} \sin \theta \cos \theta, \\ E_\varphi^{ext} &= 0 \end{aligned} \quad (3.2.23)$$

The normal component of the electric field is discontinuous, corresponding to a surface charge density

$$\sigma_s = \frac{1}{4\pi} (E_r^{ext} - E_r^{int}) = -\frac{B_s R \Omega}{4\pi c} \cos^2 \theta \quad (3.2.24)$$

The Lorentz invariant scalar product $\mathbf{E}^{ext} \cdot \mathbf{B}$ does not vanish. This quantity gives a measure of the force which a co-rotating charged particle feels in the direction of the magnetic field. From eq.(3.2.24) and eq.(3.2.15) we find

$$\mathbf{E}^{ext} \cdot \mathbf{B} = \frac{B_s^2 R \Omega}{c} \frac{R^7}{r^7} \cos^3 \theta \quad (3.2.25)$$

while the corresponding quantity vanishes inside the neutron star. Thus inside a thin transition layer at the surface of the star, there is a non-vanishing $\mathbf{E} \cdot \mathbf{B}$ of magnitude [we take $\frac{1}{2}$ of eq.(3.2.25)]

$$\mathbf{E} \cdot \mathbf{B} = -\frac{B_s^2 R \Omega}{2c} \cos^3 \theta \quad (3.2.26)$$

leading to an acceleration of the particle in the direction of the magnetic field:

$$a = \frac{e\mathbf{E} \cdot \mathbf{B}}{m_e B} \quad (3.2.27)$$

Let us compare this with the gravitational acceleration

$$g = \frac{GM}{R^2} \quad (3.2.28)$$

Thus,

$$\frac{a}{g} = \frac{eB_s R^3 \Omega}{m_e G c M} f(\theta) \quad (3.2.29)$$

where,

$$f(\theta) = \cos^3 \theta (3 \cos^2 \theta + 1)^{-\frac{1}{2}} \quad (3.2.30)$$

Typically the ratio $\frac{a}{g}$ is very large ($\sim 10^{13}$).

The electric fields parallel to \mathbf{B} are very strong. From eq.(3.2.26) we obtain

$$E_{||} \simeq \frac{\Omega R B_s}{c} \approx 6 \times 10^{10} B_{12} P^{-1} (\text{V cm}^{-1}) \quad (3.2.31)$$

P is measured in seconds, and $B_{12} = \frac{B}{10^{12}}$ G. Fields of this magnitude give rise to field emission and charge will flow from the star to fill the surrounding region.

In plasma filled magnetosphere, we have

$$\mathbf{E} = -\frac{1}{c} (\boldsymbol{\Omega} \times \mathbf{r}) \times \mathbf{B} \quad (3.2.32)$$

Corresponding to space charge density

$$\rho = \frac{1}{4\pi} \nabla \cdot \mathbf{E} = -\frac{1}{2\pi c} \boldsymbol{\Omega} \cdot \mathbf{B} \quad (3.2.33)$$

Then the plasma density becomes

$$n_e = -\frac{1}{2\pi c e} \boldsymbol{\Omega} \cdot \mathbf{B} \quad (3.2.34)$$

where e is the electric charge.

Numerically, this corresponds to a plasma density

$$n_e = 7 \times 10^2 B P^{-1} \text{ cm}^{-3} \quad (3.2.35)$$

where B is in Gauss and P is in second.

Since $\mathbf{E} \cdot \mathbf{B} = 0$, the magnetic field lines become equipotentials and the strong magnetic fields force the charged particles to corotate with the star in regions where magnetic

field lines form closed loops ($\frac{\mathbf{E} \times \mathbf{B}}{B^2}$ drift). Corotation can, however, not be maintained beyond the light cylinder of radius

$$R_L = \frac{c}{\Omega} \approx 5 \times 10^9 P[s] cm \quad (3.2.36)$$

where the tangential velocity equals to the velocity of light [12]. Let us work out the threshold magnetic field strength in the region of the source that would be required for the production electromagnetic radiation at X-ray frequency at an average energy of 0.1 keV.

To this end we write the mean free path for the electrons as:

$$\ell = \left(\frac{1}{n_e} \right)^{1/3} = \left(\frac{P}{7 \times 10^2 B} \right)^{1/3} \quad (3.2.37)$$

where ℓ is in meter. We assume this mean free path to be the average distance travelled by an electron between successive collisions. The kinetic energy gained by the electron along ℓ is given by

$$W = \Delta\phi e \quad (3.2.38)$$

where $\Delta\phi$ is the potential difference which may be written as

$$\Delta\phi = -\mathbf{E} \cdot \ell = E_{\parallel} \ell \quad (3.2.39)$$

and

$$E_{\parallel} = \frac{2\pi R B}{P c}$$

From this follows

$$W = \Delta\phi e = \frac{2\pi R e B^{2/3}}{c} \left(\frac{1}{7 \times 10^2 P^2} \right)^{1/3} \quad (3.2.40)$$

We assume the electron loses all of its kinetic energy after collision so that

$$W = h\nu \quad (3.2.41)$$

where ν is the frequency of the emitted photon.

Solving for $B(h\nu = 1.602 \times 10^{-10} \text{ erg})$, we obtain

$$B(h\nu = 1.602 \times 10^{-10} \text{ erg}) = 10^7 G \quad (3.2.42)$$

where $h\nu = 1.602 \times 10^{-10} \text{erg}$ is the average energy at X-ray frequency, P is the spin period of the pulsar which in this case is taken as a unique X-ray source named as Hercules X-1.

This magnetic field intensity is taken to be the minimum magnetic field intensity that would be required for the production of X-ray.

3.3 Application to Real Experimental Situation (e.g. Hercules X-1)

The X-ray pulsar Hercules X-1, discovered in 1972 with the UHURU X-ray satellite (Tananbaum 1972), is one of the best studied X-ray sources with its companion star Hz Herculis. Its radiation is periodic but has three different periods. This remarkable property has earned it the name a wonder clock. First, it pulsates with a period of 1.24 s. The X-ray light curve with this period is given in Fig.(3.2). Second, it has a period of 1.7 days. This light curve has one II-shaped minimum.

Physiologists say that cats sleep for 18 hours a day. In other words, they sleep for 75% of their lives. Hercules X-1 is like a cat. Every 35 days, Hercules X-1 is at "work" producing X-rays for 11 days, and for the next 24 days its X-rays are not observed from the earth. This is the third period.

The shortest period of 1.24 seconds is the period of the neutron star's rotation about its axis. The star's pulsation light curve in this case has two maxima.

The period of 1.7 days is, undoubtedly, the binary's orbital period. The short period of 1.24 s changes with the same period due to the Doppler effect. The 1.7-day change in the X-ray curve is the result of the eclipsing of the X-ray source by its companion. It is interesting that the eclipsing begins very abruptly. Recall that the duration of the decreasing part in a light curve is proportional to the size of the eclipsed.

The star that periodically obscures the X-ray pulsar was found by the Soviet astronomer Nikolai Kurochkin.

In 1967, Shklovsky emphasized that a very strong reflection effect must be inherent in X-ray binary systems. This phenomenon is basically the same as for normal stars, the only difference being that in an X-ray binary the star is warmed up by intercepting the pulsar's hard X-radiation. The inner face of Hz Herculis is strongly irradiated by X-rays from the neutron star, however the warped accretion disc shadows part of it. An observer near the accretion disc would see the following sight. At the "horizon", there is the optical component, i.e. the star Hz Herculis and a stream of gas following from it. As the disc approaches the neutron star bending its surface, which is covered with "ripple" due to Kelvin-Helmholtz instability. Eventually, the disc is broken up by the magnetic forces at a distance of several thousands of kilometers. The matter is then "frozen" into the magnetic field and flows down onto the neutron star's magnetic pole, which produce two X-ray beams.

It is not yet clear, though, why Hercules X-1 is as sleepy as a cat for 24 days. By the way, the fact that for us on earth the radiation disappears for 24 days every 35 days does not indicate that the pulsar takes a 24-days holiday. Even when the pulsar is "off", its optical light curve remains the same, that is, the reflection effect does not disappear. Something seems to be obscuring the neutron star from us, but this something cannot screen all the X-radiation falling onto its companion, Hz Herculis. The accretion disc around the neutron star, probably, serves as a screen. Since the binary's orbital plane lies actually in the line of sight, the disc which wobbles slightly with respect to this plane is believed to block the pulsar's radiation from us for some time. The accretion disc's shadow on the companion star is not large and has virtually no impact on the reflection effect.

The reasons given above for the anomalous behavior of Hercules X-1 are not totally convincing. For one reason, the eclipsing is short and for another the disc's wobble

is not periodic. Making use of the new model developed by Kebede (2002, 1996), we can address issue of Hercules X-1. According to this model this phenomenon is basically related to the fact that the magnetic field evolution of the neutron star beyond the X-ray death line (see eq.(3.1.6) and its graph fig. 4.1. As a result of the competing effects of magnetic field decay due to photon emission and magnetic field revival due to mass accretion, the magnetic field fluctuates about the threshold magnetic field (i.e. the minimum magnetic field that would be required for X-ray spectrum production, 10^7G) almost with 35 day cycle.

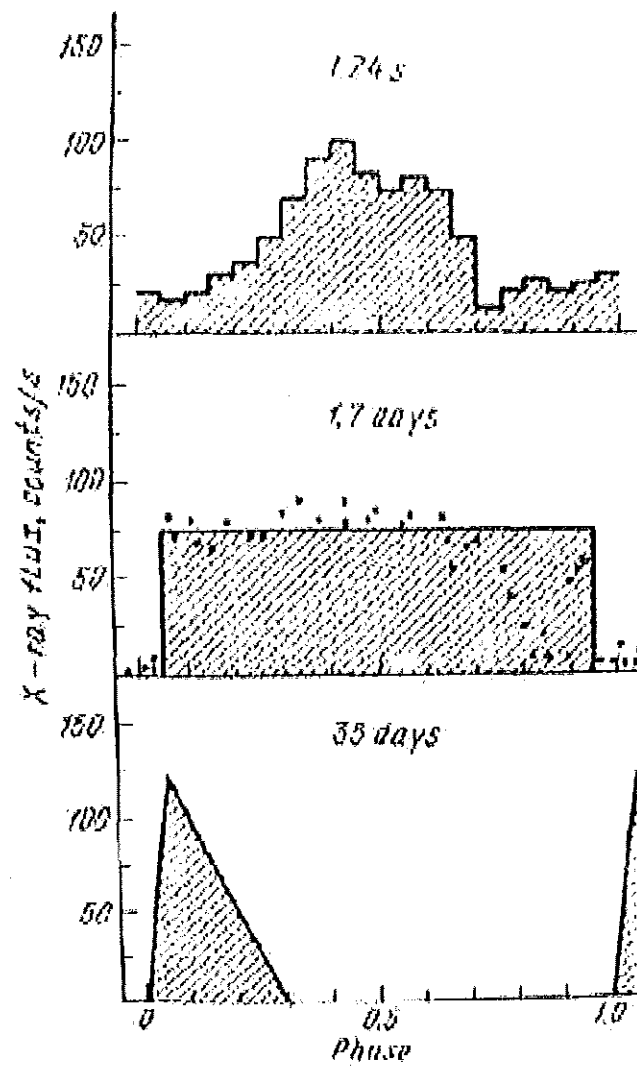


Figure 3.1: The three periods of the X-radiation from Hercules X-1

Chapter 4

Discussion and Conclusion

Considering the competitive nature of mass accretion which according to the new

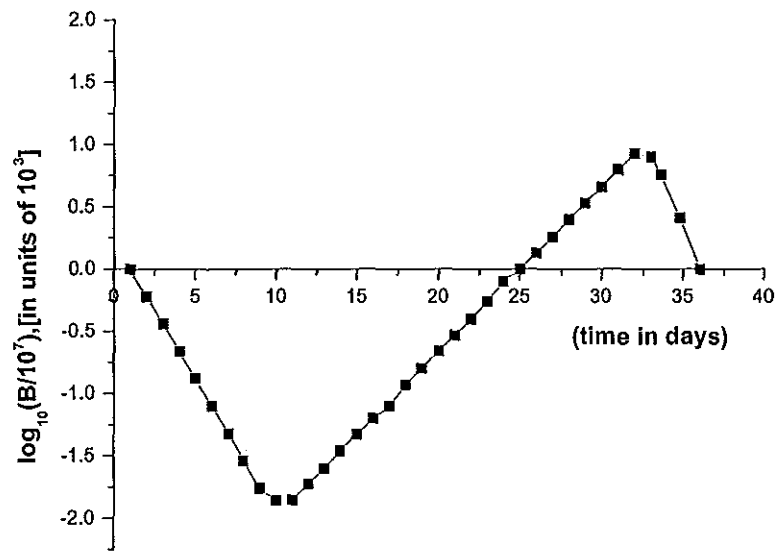


Figure 4.1: This graph illustrates the competition between magnetic field revival due to accretion and its decay due to photon emission and B is given in eq.(3.1.6).

model is responsible magnetic field revival and neutron star cooling via photon emission which results in field decay; we have presented convincing theoretical explanation for the 35 day cycle of X-ray radiation from Hercules X-1, an accreting NS in a binary (see fig. 4.1).

This is reminiscent of the so called X-ray transients (SXRTs) which just like Hercules X-1 are accreting NSs. SXRTs undergo [38] the periods of outburst activity for days (even months) super imposed with quiescent periods for months. The reason for their peculiar activity is not yet well understood but it is generally believed that their X-ray emission is regulated by accretion from discs around the NSs.

In quiescence, the accretion is switched off or suppressed for some reason standard theories fail to pin point. We have shown that neutron stars admit accreted mass from their companions below the threshold surface field of $\sim 10^7$ G. This is also the field which according to the standard calculations is able to induce electric fields than can lead to low X-ray emissions. Since the accreted mass transfers angular momentum to the star this would mean a revival of the magnetic field that has been decaying (weakening) until then due to photon emission. Obviously then this would lead to a repetitive pattern of active and quiescent periods in low X-ray production from such accreting old NSs.

Our model therefore holds that pulsars magnetic field evolution beyond "the X-ray death point" will not be monotonous but rather involve cyclic patterns not only in X-ray but perhaps in thermal emission as well. To the best of our knowledge no previous study has addressed this issue of magnetic field evolution beyond the X-ray death line.

Bibliography

- [1] Blandfor et al., 1983
- [2] Uprin et al., 1986
- [3] Cummings et al., 1992
- [4] L. W. Kebede, 1996
- [5] L. W. Kebede, *Astrophysics and Space Science* **282**: 131-140, 2002
- [6] Ostimerker & Gunn, 1996
- [7] Manchester & Taylor, 1997
- [8] Vivekanard & Naryana, 1981
- [9] Uprin et al., 1994
- [10] Burdier & D'Amico, 1997
- [11] Bhanttacharya & Srrinivasan, 1995
- [12] San gupta,1997
- [13] Konar & Bhattacharya, 1997

- [14] Uprin et al., 1998
- [15] Young & Chanmngan, 1995
- [16] K. P. Levenfish and D. J. Yakovlev, astro-ph/9608032 v1 7 Aug 1996
- [17] D.J. Helfad et al.: Nature 283:337(1980)
- [18] Norbert Straumann, General Relativity and Relativistic Astrophysics
- [19] Colpi et al., 2001
- [20] V. M. LIPUNOV, SCIENCE FOR EVERYONE IN THE WORLD OF BINARY STARS
- [21] Ya. B. Zel'dovich, and I. D. Novikov, Stars and Relativity
- [22] MALCOLM S. LONGAIR, High Energy Astrophysics, Second Edition,
VOLUME 1
- [23] MALCOLM S. LONGAIR, High Energy Asrtophysics, Second Edition,
VOLUME 2
- [24] JIHAN FRANK, ANDREW KING and DEREK RAINE, Accretion Power in Astrophysics, SECOND EDITION
- [25] L. ZHANG, K. S. CHENG, Z. J. JIANG, AND P. LUENG, THE ASTROPHYSICAL JOURNAL, 604:317-327, 2004 MARCH 20

- [26] BING ZHANG AND ALICE K. HARDING, THE ASTROPHYSICAL JOURNAL, 535: L51-L54, 2000 May 20
- [27] Walter F. Brisken, S. E. Thorsett, A. Golden, & W. M. Goss, astro-ph/0306232 v1 11 Jun 2003
- [28] I. A. SHOVKOVY and P. J. ELLIS, astro-ph/0303073 v1 10 Mar 2003
- [29] W. Becker & J. TtÜmper, Max-Plank-Institut für extraterr. Physik, Giessenbachstrasse 1, D-85740 Garching
- [30] Heisenberg & Pandharipade, 2000
- [31] Alock, Fahri, & Olinto, 1986
- [32] Yakovler et al., 2003
- [33] Ohanian, Gravitation and Space Time
- [34] Lars Bildsten, astro-ph/9801043 v1 7 Jan 1998
- [35] N. Straumann: Weak Interactions and Astrophysics, Proceedings of the GIFT Seminar on Electro-Weak Interactions, Peniscola (Spain), May 1980, to be published
- [36] Dipankar Bhattachary Raman Research Institute, Bangalore 560 080, India, J. Astrophys. Astr. (2002) 23, 67-72
- [37] P.Golgreich, W.H. Julian: Astrophys.J. 157, 869(1969)
- [38] Chen et al., 1997