IIb. ACCRETION POWERED PULSARS

Theoretical Considerations

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ACCRETION ONTO MAGNETIZED NEUTRON STARS: POLAR CAP FLOW AND CENTRIFUGALLY DRIVEN WINDS

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ABSTRACT. Some basic concepts of accretion onto the polar caps of magnetized neutron stars are reviewed. Preliminary results of new, multidimensional, time–dependent calculations of polar cap flow are outlined, and are used to suggest the possible observability of fluctuations in the X–ray intensity of accretion powered pulsars on time scales of 10 – 100 msec. The possible relevance of such fluctuations to Quasi–Periodic oscillations is suggested. Basic concepts of the interaction between a disk and the magnetosphere of a neutron star are also discussed. Some recent work on the disk–magnetosphere interaction is outlined, leading to the suggestion that a neutron star can lose angular momentum by driving some or all of the mass in the disk off as a centrifugally driven wind. The relevance of such mass loss to the orbital evolution of the binary is pointed out.

1. INTRODUCTION

I was asked to review accretion onto neutron stars. I don't know how to do this without writing a monograph. Therefore, I have chosen to discuss some basic concepts and recent work of immediate interest to me. These are the structure of accretion flow onto the polar caps of magnetized neutron stars, widely thought to be of relevance to the observable emission of X–rays from accreting neutron stars, and the interaction between a thin disk and the rotating magnetosphere of a neutron star, a subject of relevance to the rotational evolution of these objects.

2. POLAR CAP ACCRETION ONTO MAGNETIZED NEUTRON STARS

2.1. The Eddington limit and Optically Thick Accretion

I'll begin by discussing a few basic concepts. Most of the discussion in this section is review of well–known ideas, going back to Basko and Sunyaev (1976). Spherical, optically thin flow of a fully ionized plasma with cosmic composition onto a gravitating body cannot occur when the luminosity exceeds the Eddington luminosity $L_{ED} = 4\pi GM_{*} mc/\sigma_T = 1.3 \times 10^{38}$ erg/s, since the radiation pressure then exceeds gravity (eg., Shapiro and Teukolsky

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1983). Here M_* is the mass of the star, taken to be 1 $M_\odot$, m is the mean molecular weight, taken to be 61% of a proton's mass, and σ_T is the Thomson cross section.

Accretion onto a magnetic pole modifies the magnitude and meaning of this limit. If one imagines all the radiation to be coming out of the stellar surface through a polar cap area $\pi R_*^2 \theta_c^2 \sim 3\ (\theta_c/0.1)^2$ km^2 and assumes the radiation flux is directed anti–parallel to the plasma flowing in along the magnetic field B, then the Eddington luminosity for that pole is $L_{ED}^{(eff)} = (\theta_c^2 H_{\parallel}/4) L_{ED} \sim 3 \times 10^{35}\ (\theta\ /0.1)^2 H$ erg/s. $H_{\parallel} \geq 1$ is the ratio of the Thomson cross section to the magnetically modified cross section for co–herent scattering of photons travelling along B; for photon energies exceeding $11.6 B_{12}$ keV, $H_{\parallel} = 1$, while $H_{\parallel} \approx 0.06(B_{12}/T_{10})^2$ for $B_{12}/T_{10} \gg 4.5$; $H_{\parallel}$ departs substantially from unity only for $B_{12}/T_{10} > 3$ (Arons, Klein and Lea 1986, hereafter called AKL). Here, $B_{12} = B/10^{12}$ Gauss, $T_{10} = T/10$ keV with T the photon temperature. $L_{ED}^{(eff)}$ is determined mainly by the magnetic requirement that all the radiation come from a smaller area; in very strong fields, the reduction of the opacity pushes $L^{(eff)}$ up again, but in general, reduction of the area is the dominant effect.

If this effective Eddington limit meant the same thing as in spherical free fall, one would wonder how the accretion of plasma onto magnetic polar caps could be the only going model for the X–ray emission from the accre–tion powered X–ray pulsars, and also for the magnetized stars suspected of being in low mass X–ray binaries, since most of these objects have X–ray luminosities much in excess of $L_{ED}^{(eff)}$. However, one can readily show that the scattering optical depth from the surface to infinity, measured along an al–most radial dipole magnetic field through freely falling plasma is $1.5[c/v_{ff}(R_*)][L_{cap}/L_{ED}^{(eff)}] \gg 1$ when the polar cap accretion luminosity ex–ceeds $L_{ED}^{(eff)}$. Thus, locally super–Eddington flow is no longer optically thin, and the correct interpretation of $L_{ED}^{(eff)}$ is as a fiducial luminosity, above which scattering in an optically thick medium causes the radiation to be redirected out of the accretion column into a fan beam emission pattern. When accre–tion occurs from a thin disk, one suspects the plasma falls in the form of a hollow cone flowing along B. Radiation emitted from the outside of this op–tically thick cone forms a fan beam, while radiation emitted into the inner parts of the cone flows out as a broad or narrow pencil beam, with the opening angle determined by the distribution of scattering optical depth as a function of magnetic azimuth around the flow cone. The proportion of fan beam to pencil beam emission also depends on this optical depth distribution. For simplicity, I will discuss only the simplest case of optically thick accret–ing plasma falling in all over the polar cap, as may be a respectable model when plasma enters the magnetosphere with low angular momentum (Arons and Lea 1980; Burnard et al 1983; see also Arons et al 1984, Lamb 1984).

2.2. *Physical Conditions in Polar Columns*

Over the years, quite of lot of data have been gathered by the satellite ex–periments equipped with proportional counters. A typical example of the continuum spectral fluxes of these objects is contained in the survey of the accretion powered pulsars published by White et al (1983). The parameters typical of polar cap flow, inferred by assuming densities on the order of free fall, "temperatures" comparable to the quasi–exponential roll–overs seen in the spectra, and photospheric areas and magnetic field strengths characteristic of the magnetospheric models, yield an impression of a very lively part of the universe. Typically, for photospheric emission area ~ 1 km^2, one finds "temperatures" ~ 10 keV, plasma densities $\sim 10^{21}$ cm^{-3} (higher, in regions where the flow is decelerated below free fall), free fall dynamic pressures $\sim 10^{17}$ dyne/cm^2, LTE radiation pressure $\sim 10^{17.5-18}$ dyne/cm^2, plasma pressure $\sim 10^{14}$ dyne/cm^2, and magnetic pressure $\sim 10^{22.5}B_{12}^2$ dyne/cm^2. Clearly, radiation pressure is enough to decelerate the flow to speeds below free–fall, while plasma pressure is unimportant, at least at the observable photosphere. If pressures everywhere in the column are comparable to these observationally constrained values, then perfect magnetic confinement in the directions across the magnetic field looks like a good idea[1].

I occasionally remind myself of this environment's lively character by looking up the X–ray exposure limits issued by the United States Occupational Health and Safety Administration, which tell me that even a few minutes exposure to the X–rays from these stars at a distance of one parsec would be enough to make me very sick indeed. Put differently, the X–ray flux at the source typical of these objects corresponds to roughly 300,000 nuclear wars[2]/m^2–sec. Examples like these remind one of the hostility of some parts of the universe to our rather fragile form of life.

2.3. Spectroscopic Problems

When one tries to advance beyond these order of magnitude views of polar columns, an immediate problem is encountered. The observed spectra of the brighter sources are not well represented by the simplest models. No model employing one value of density, temperature, with a single radiation process converting plasma energy into photons, has succeeded in giving a good account (meaning physically self–consistent) of the emergent spectra, with their flat power law form below a quasi–exponential roll–off. Furthermore, the columns clearly form scattering atmospheres, whose photospheric surfaces have radiation number density depleted well below the Planck level appropriate to the plasma temperature – the emergent fluxes might be modified black bodies, if the lines of sight through the column ever become optically thick to absorption at all. Some weak observational evidence that such modifications of the emission, due to the dominance of scattering in the atmosphere, really do play a role, can be gleaned by comparing the exponential roll–overs expected from complete local thermodynamic equilibrium models to the observed rollovers. Because freely falling, optically thick plasma traps the photons emitted from the higher parts of the decelerated mound of plasma contained within the free fall zone, the emission area of the fan beam turns

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1. Of course, anyone who has followed the efforts to achieve controlled fusion will realize how dangerous and unrealistic such a conclusion can be, based only on a simple comparison of magnetic to total gas pressure. In fact, polar cap flows may admit an interesting class of interchange flows which may allow plasma to "diffuse" anomalously fast across B (Arons et al 1985). For my discussion here, I follow the conventional astronomical discussions and assume perfect magnetic confinement.
2. Being a bit of a pessimist, I have assumed one war is a 10,000 megaton "exchange" (1 megaton $\approx 10^{22}$ ergs).

out to be comparable to the area of the polar cap, even when the de–celerated column has a height large compared to the polar cap's diameter (Arons and Klein 1987a,b); the emission is never that of a tall skinny cylin–der, in high luminosity flows. Then the effective temperature of the photons is expected to be $\sim 4\ (L_{cap}/L_{ED})^{1/4}(\theta_c/0.1)^{1/2}$ keV, corresponding to a quasi–exponential roll–over in the spectrum at energies $\sim 11\ (L_{cap}/L_{ED})^{1/4}(\theta_c/0.1)^{1/2}$ keV. Observed spectra have exponential rollovers usually factors of 2 – 4 higher than this, consistent with the plasma having a radiating surface characteristic of a modified black body [although the continuum absorption opacity isn't bremsstrahlung, but is more likely to be some combination of Comptonized cyclotron emission (AKL) and a resonant form of double Compton scattering (Kirk and Melrose 1986)].

Because the radiation pressure is strong enough to decelerate the flow, and because the flow is necessarily inhomogeneous across as well as along B, one expects a wide range of physical conditions to contribute to the emergent spectra. In particular, in attempting to model the already measured phase resolved spectra of the accretion powered pulsars, one needs to have some means of dynamically determining the range of photon temperature and chemical potential (which measures the depletion of the radiation field below the Planck level) in order to have a hope of using the observations to probe the flow and magnetic field structure. In addition, as we will see, the flows almost certainly impose short time scale ($\sim$ 10–1000 msec) fluctuations on the emergent intensity, even if the accretion flow is free of "raindrops" (Arons and Lea 1976, 1980, Michel 1977, Elsner and Lamb 1977, 1984), "blobs" (Lamb et al 1985) and other semi–quantitative constructs of theo–rists who have studied the entry of plasma into the magnetosphere.

2.4. Theoretical Models

The inhomogeneities in space and time built into the basic model for polar cap accretion can't be inferred by direct inversion of the data, so far as I know, so a more useful method appears to be the construction of **dynamical** models which allow one to determine the flow structure, within the context of an assumed geometry and profile of mass flux of the flow along **B**. In a completely *a priori* theory, these would be known from the theory of the magnetosphere, and **that** would be known, once one specified the evolutionary state of the binary. Needless to say, one doesn't try to be this general, and instead imposes simplified models of the mass flux profile of freely falling plasma. Various authors have made analytical and numerical models of the polar flow when radiation pressure supports the plasma in the direction along **B** (Davidson 1973, Basko and Sunyaev 1976, Wang and Frank 1981, Braun and Yahel 1984, Kirk 1985). Basko and Sunyaev and Braun and Yahel assumed the scale of transverse gradients is always equal to the width of the flow at each altitude. This allowed them to create an approximate one–dimensional theory by replacing all the differential operators in the directions across **B** by the inverse of their assumed length scale. At high luminosity and high optical depth, this one dimensionalization is misleading, since the actual transverse length scale decreases, in order to allow radiation diffusion across **B** to carry the dynamically determined radiation flux. Davidson, Wang and Frank and Kirk all carried out steady state, two dimensional axisymmetric calculations, in each case dropping all the details of the radiation emission processes and thus leaving out the information needed to determine the photon number density, required if one is to determine the modifications to local black body emission imposed by the scattering character of the atmosphere. In the rest of my discussion, I will concentrate on the most ambitious attempt made so far to model the dynamics of polar flows, being carried out by Richard Klein and myself. For more details, see the poster papers at this meeting (Klein and Arons 1987a, Arons and Klein 1987a).

With suitable restrictions it turns out to be possible to make steady flow, multi–dimensional analytic models (Kirk 1985; Arons and Klein 1987a,b). I will outline the numerical models that we have done so far, since these have the greatest realism. I emphasize that this work is preliminary; we are still in the throes of putting together the last pieces of the physics and numerical methods, but some interesting results are already clear. The program is fully time dependent, and assumes axisymmetry of the flow with respect to the magnetic axis. We include magnetically modified bremsstrahlung emission and Coulomb collisional excitation/deexcitation of quantized cyclotron gyration as the photon emission and absorption processes. The cyclotron line emission is treated using a non–LTE, two Landau level "atom" model with complete redistribution for resonant scattering. Because the optical depths are large in the line **and** in the neighboring continuum in the high optical depth models we consider, the problem of cyclotron line transfer is solved "on the spot", with the result that cyclotron line radiation is primarily a source for the continuum[3]. The continuum photons undergo fully saturated Comptonization, using a rate for energy exchange with electrons that takes account of the magnetic modifications of scattering. The consequence is that in most of the flow, the radiation spectrum can be safely approximated as having a Bose–Einstein distribution with photon temperature and chemical potential varying in space and time. The emergent spectrum at each point is characterized by these two numbers.

In order to find them, we include a new theory of the coupled diffusion and advection of photon energy density **and** number density, using appropriate, flux limited "Rosseland" averages of the angle and frequency dependent anisotropic, coherent scattering opacity in the continuum. The electrons and ions are allowed to have separate temperatures, described by the energy equations for the plasma. The plasma temperatures are coupled by Coulomb energy exchange. Mass and momentum conservation are described by the continuity and one–fluid momentum equations for flow strictly along $\mathbf{B}$. In the momentum conservation equation, we include gravitational acceleration along $\mathbf{B}$, radiation force, thermal pressure gradients and the buoyancy force. The assumed geometry is flow along a dipole magnetic field, and in all of our calculations done to date, the mass flux has been assumed to be independent of distance from the magnetic axis, out to the boundary of the flow where the density is assumed to drop abruptly to zero. At the top of the calculational grid (for now, taken to be 20 km above the magnetic pole of a star whose radius is 10 km), we assume the plasma is in free fall.

3. The high optical depth and the associated saturated Comptonization, appropriate for sources with high luminosity, is the essential difference between our work and the pure transfer models of Meszaros and Nagel (1986), for example.

We emphasize that this mass profile is chosen to keep the problem as simple as possible in the first calculations; more complex flows, such as hollow cones, will be done in the near future. We also assume no heat flows into the star, as is plausible since the photon opacity of the crust immediately below the piled up mound of plasma is much greater than in the accreted plasma itself, and we assume photons escaping from the sides of the column go off into black sky.

I will show a few of the results from two calculations. In both, the flow was initialized as cold plasma in free fall all the way down to the stellar surface. The first example was run with an accretion luminosity onto the cap of 5 x 10^{37} erg/s. The magnetic colatitude of the footprint of the field lines at the edge of the polar cap is θ_c = 0.4 radians; for self–congratulatory reasons, we used the model of Arons and Lea (1980) for the polar cap size in these preliminary calculations. There is no implication intended that this is the correct way to scale the geometry for all accreting neutron stars, although there is some weak evidence that it is relevant to lower luminosity (L < 10^{37} erg/s) accretion (White *et al* 1983). The neutron star radius is 10 km, and the surface magnetic field strength was chosen to be 3 x 10^{12} Gauss (magnetic moment μ = 1.5 x 10^{30} cgs), yielding an energy for the cyclotron line at the surface of 35 keV.

After the passage of a shock through the flow in order to adjust to quasi–stationary conditions, a mound, in approximate hydrostatic equilibrium developed with the subsonic plasma supported by radiation pressure; radiation pressure also provides the full deceleration from free fall. A fan profile of emergent radiation is formed after a couple of photon diffusion times. An example is shown in Figure 1, a snapshot taken at 130 μsec into the problem (this is about 50 times the dynamical time). To get this far takes about 2.3 hours on a Cray XMP–48 – this type of calculation is not for the pencil and paper or VAX crowd! The vectors represent the photon flux leaving the point at the foot of each vector, whose direction shows where the flux is going and whose length indicates the flux magnitude. The velocity field at this time is rather dull, showing the same exponential stratification along **B** and Gaussian variation across **B** as is found in the analytic solutions. The rest of the structure – density and pressure, as well as velocity – is also well represented by the steady flow analytic models, if one uses the numerical data to properly choose the magnetic modifications of the opacity.

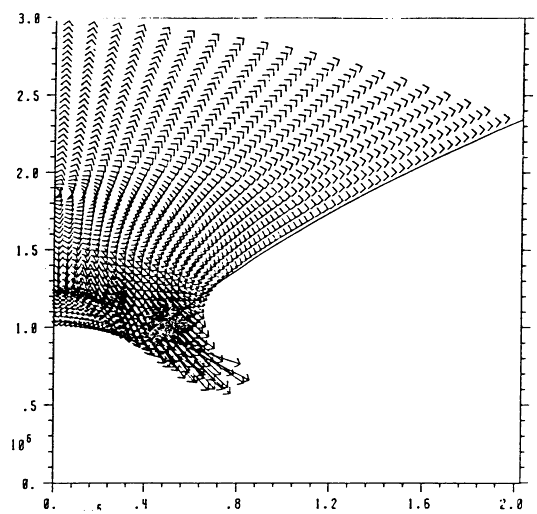

Figure 1: Radiative flux of $10^{37.7}$ erg/s of accretion onto the polar cap. The spatial scales are in units of 10 km.

A bit later, things become more interesting. Along the magnetic axis, the plasma develops a density depression, giving rise to a net buoyancy force.

We believe, although we have not yet conclusively proved, that this is the onset of a form of overstable convection (*e.g.*, Proctor and Weiss 1982). The initial non–linear consequences for the velocity field are shown in Figure 2, a snapshot taken at 600 μsec into the problem. Along the central axis, a lower density, subsonically **rising** "bubble" forms with density lower than the surround–ings; the emergent flux in the fan beam has yet to show a substantial change. The appearance of such structures is not especially surprising, since the polar cap flow is strongly super–Eddington with respect to the **effective** Eddington limit. Many authors have used linear stability arguments (*e.g.*, Hameury *et al.* 1984) to suggest that polar cap flows would form convective motions, and long ago, Hsieh and Spiegel (1976) speculated that the such flows would form "photon bubbles".

The final fate of this structure is still unknown; the lower densities have driven the Courant time down to the point where further evolution using our explicit treatment of the momentum equation requires prohibitive amounts of computer time (the transfer and energy equations are already treated implicitly). Technical developments are in progress which should allow us to get around this problem. If these succeed, we will be able to evolve models to times as long as 1 second. My suspicion is that the bubble will break up into finer scaled structures because of the weight of the accreted plasma piling up on top – the magnetic field precludes the falling away of plasma towards the bubble's sides that one expects in ordinary convection, or was postulated in Hsieh and Spiegel's photon bubbles. Because radiation diffuses faster through the lower density zones, one expects this "convection" to lead to faster, percolative loss of radiative energy across the columns, which **may** lead to observable, short time scale (10 – 100 μsec) fluctuations in the emergent flux. A prediction of the magnitude of such fluctuations, so far not studied observationally in the rotation powered pulsars, requires running calculations to times long enough to determine the number and scale of the bubbles contributing to the percolation, a task to be done as soon as the ability to follow the flow to long times is in place.

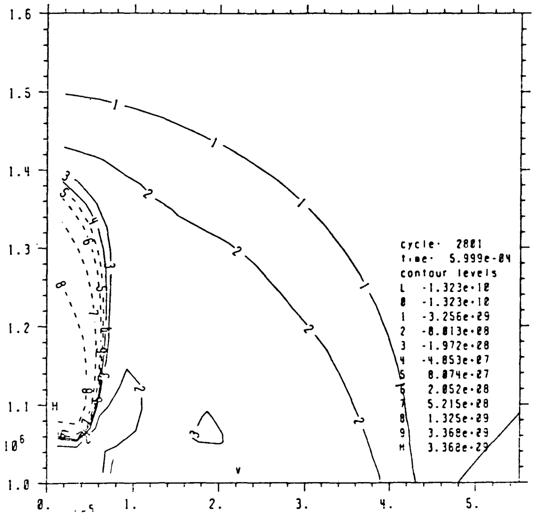

Figure 2: Velocity field of the same model as in Figure 1, but at the later time 600 μsec. The vertical scale is in units of 10 km; the horizontal scale has units of 1 km. Solid curves are contours of down–ward (negative) velocity, in cm/s; dashed curves show of upflow.

The character of these "convective motions clearly changes with vary–ing luminosity. Figure 3 shows the velocity field of a model constructed in the same manner as that of Figure 2, but now for an accretion luminosity onto one cap of 1.5 x 10^{38} erg/s. The snapshot is taken at 0.75 msec into the flow. One sees now three "bubbles" – because of the axisymmetry, these bubbles form toruses in the plasma. As time goes on to 0.9 msec, when we quit trying to follow the flow because of the declining time step, the central bubble had grown to fill in on the magnetic axis, and the satellite bubbles had greatly increased their size.

I don't have room to go into the spectroscopy of these models, except to remark that the output spectra are the superposition of the Wien spectra determined by the variable temperature and chemical potential over the surface of the decelerated mound. The beaming pattern is mainly controlled by the trapping of the photons in the optically thick, freely falling plasma above and outside the mound, until at the bottom of the flow, the layer of freely falling plasma becomes optically thin to scattering and photons escape in a narrow ring around the base of the mound, as is shown in Figure 1. One can show that the height of this ring is roughly $\theta_c R_* [v_{ff}(R_*)/c]$, so that the area from which the radiation is beamed is comparable to the area of the polar cap on the stellar surface. The **quality** of the radiation field is determined by the emission from the much larger mound surface, however. So far, the results we have found for the spectrum form a surprisingly good representation of the spectra of some of the rotation powered pulsars – my surprise is a consequence of the simplicity of these calculations' geometric assumptions. One other aspect is that at photon energies well below the cyclotron energy at the surface (35 keV in these models), the emitted radiation is very highly linearly polarized, because the opacity in the extraordinary mode is much smaller than in the ordinary mode, and therefore has an intensity reflecting the Wien intensity of deeper plasma layers. If someone, someday, would finally get the chance to fly an X–ray polarimeter attached to a telescope of decent collecting area, we would be able to observationally probe the geometric structure of the flow, to which the spectroscopy seems to be insensitive.

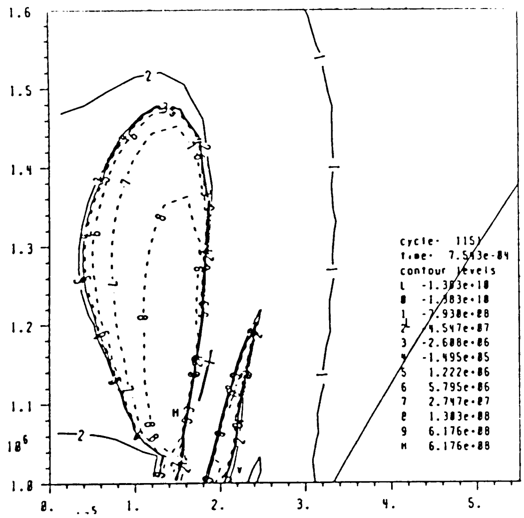

Figure 3: Velocity field of a model with assumptions the same as in Figure 2, except the luminosity is now $10^{38.2}$ erg/s and the polar cap opening angle is 0.33 radians. This snapshot is at 0.75 msec. The meaning and units of the contours are as in Figure 2.

I can't resist leaving this topic without a final speculation, perhaps especially appropriate to this meeting. While we have concentrated on the strongly magnetized, accretion powered pulsars in the galactic disk, the same physics applies, with lower magnetic field strength, to the now popular, more weakly magnetized models for galactic bulge/QPO sources which are likely to be the precursors to the millisecond pulsars. It has not escaped our notice that the likely fluctuation times appearing in the polar flow models are comparable to the fluctuation periods seen in the quasi–periodic oscillators, leading to intriguing thoughts about a model in which the observed fluctuations may not by the direct consequence of orbital phenomena at the magnetopause.

3. CENTRIFUGALLY DRIVEN WINDS FROM NEUTRON STARS

3.1 Standard Ideas

The standard ideas on disk–magnetosphere interaction were invented to explain the average behavior of the rotation period P and dP/dt, as observed during the "American" era of X–Ray astronomy in the 1970's. The basic thoughts were first outlined by Pringle and Rees (1972) in advance of substantial observations of dP/dt, and were applied to the system shown in Figure 4. Here μ is the magnetic moment, Ω_* is the angular velocity, R_m is the magnetopause radius

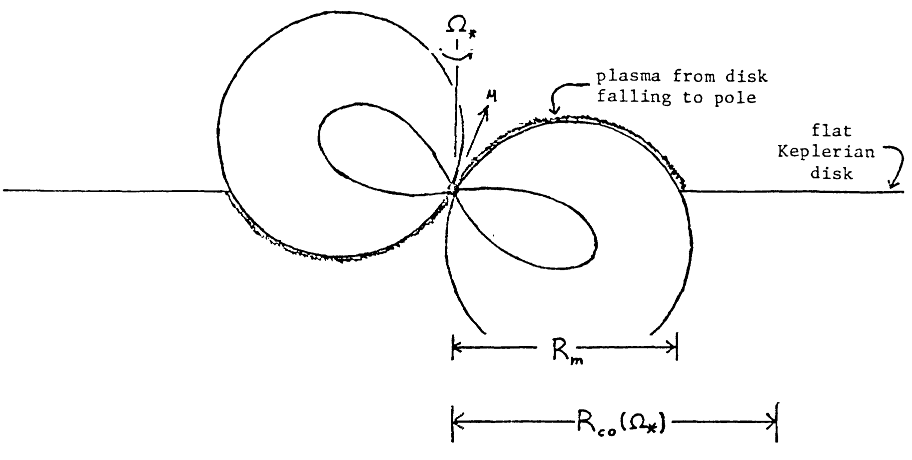

Fig. 4: Cartoon model of a magnetosphere confined by a Keplerian disk, when accretion to the stellar surface is in progress.

in the disk plane, and R_{c*} is the corotation radius, where the Keplerian angular speed $(GM_*/R_{c*}^3)^{1/2}$ equals the stellar angular velocity. This is thought to be relevant to magnetospheric problems since almost all authors have assumed the field lines to have at least one foot on the stellar surface, and to be perfect conductors by virtue of the ideal plasma presumed to completely fill the magnetosphere. If these assumptions are met, the magnetospheric electric field is $\mathbf{E}_{mag} = -(\mathbf{\Omega}_* \times \mathbf{r}) \times \mathbf{B}/c$ and the velocity of plasma across B is $\mathbf{v} = c(\mathbf{E}_{mag} \times \mathbf{B})/B^2 = \mathbf{\Omega}_* \times \mathbf{r}$. Therefore, if $R_{c*}(\mathbf{\Omega}_*) < R_m$, net gravity (= real gravity + centrifugal force) at the magnetopause is outward, and "no" accretion to the stellar surface could occur. Numerically, $R_{c*}(\mathbf{\Omega}_*) = 1500\ (P/1s)^{2/3}(M_*/M_\odot)^{1/3}$ km. The magnetopause radius in the disk plane has been estimated to be at $R_m = 3000[\mu_{30}^4(M\ /M\ )/L_{37}^2]^{1/7}$ km (Pringle and Rees 1972, Lamb, Pethick and Pines 1973, Ghosh and Lamb 1979, Arons et al 1986), a value close to that

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determined when the infall lacks angular momentum (Davidson and Ostriker 1973, Arons and Lea 1976, Elsner and Lamb 1977, Michel 1977). Here μ_{30} is the magnetic moment in units of 10^{30} cgs, and L_{37} is the accretion luminosity in units of 10^{37} erg/s. Thus, under all circumstances, short period ($P < 1-2$ s) accreting neutron stars with strong fields ($\mu_{30} > 1$) should have trouble accreting mass and angular momentum, as long as the magnetic field lines are rooted in the star and are good conductors. Indeed this limiting period for spin up is one of the ideas behind the probable fact that the shortest period radio pulsars are the consequence of accretion spin up (eg., Damashek et al 1982, Alpar et al 1982).

The simplest, surprisingly successful ideas (Pringle and Rees 1972, Lamb et al 1973, Rappaport and Joss 1977, Mason 1977) of how a pulsar changes its rotational period come from assuming that each bit of plasma enters the magnetosphere from the inner edge of a Keplerian disk, with specific angular momentum $\sim (GM_* R_m)^{1/2}$. This gives a spin up torque $T_+ = \dot{M}_x (GM_* R_m)^{1/2} \propto L^{6/7}\mu^{2/7}$, independent of P, and therefore a spin up rate $\dot{P}/P \propto PL^{6/7}\mu^{2/7}$. This described the spin up data from the 70's for most pulsing X–ray sources in X–ray binaries, if the dispersion in $\mu^{2/7}$ is not too large. Because of the weak dependence on μ, this doesn't prove that these objects have 10^{12} Gauss surface fields, nor is this explanation of spin up unique – the angular momentum accreted from a wind, with no disk formed can give as good an explanation for many of these sources (eg., Arons and Lea 1980, Stella et al 1986). Nevertheless, the simplest idea, of matter being transferred from a disk to a strongly magnetized neutron star through attachment of the plasma to a corotating magnetosphere, seemed to be qualitatively correct.

Unfortunately, in Her X–1, where we are certain (Boynton et al 1980) that mass transfer occurs through a disk, $\dot{P}/P$ is $\sim 1/30$ of the value expected. Most people do not think it possible to obtain this reduction simply through reducing μ by $\sim 2 \times 10^5$, mainly because it is hard to see how to form a well defined polar spot from which the radiation arises without a much stronger surface, **large scale** field. The same argument can be applied if the field is not dipolar but is dominated by some other, higher order multipole. The magnetosphere is smaller (Arons and Lea 1980), with a consequently smaller angular momentum/gram accreted, but the effect is still much too small to explain the reduced spin–up torque observed in this system, while still being consistent with the narrow cyclotron line observed – in particular, the variations of the line shape and centroid energy observed require the presence at the line photosphere of a field of strength $\sim 10^{12.6}$ Gauss, with no magnetic broadening within the instrumental resolution (Voges et al 1982). The rather good fit of simple radiative transfer models of cyclotron line formation in a quasi–homogeneous magnetic field to these data (Meszaros and Nagel 1986) also suggest that the surface field has rather large scale structure. Then one cannot suppose the magnetospheric field to be a very intense ($B > 10^{12}$ Gauss) collection of multipoles, with a very weak field at the magnetopause – the result would be too many different field strengths contributing, which creates an inadmissably broad line. For an alternate view of possible weak field models of Her X–1, see Ruderman (1986).

Therefore, one concludes that some sort of "spin–down" torque operates which is almost in equilibrium with the simple spin–up due to accretion. The first candidate proposed for such a mechanism is the closed, force–free magnetosphere, "resistive" disk theory of Ghosh and Lamb (1979). Since these earlier data and theories, observations carried out primarily on Hakucho and on Tenma have revealed substantially more complex spin–up and spin–down behavior, much of which might be interpretable in terms of spin–down torques competing with the accretion of angular momentum. Our work, as well as that of Anzer and Borner (1980, 1983) and of Ghosh and Lamb (1979), concerns the nature of the interactions between disk and magnetosphere which lead to these torques, as well as models of the torque and the resulting spin behavior of the star, under various assumptions as to the nature of the basic magnetic configuration and approximations to the relevant physics.

3.2 Mass Entry: A Thought Experiment

I now describe some new work I have been doing, in collaboration with C.F. McKee and R. Pudritz. We construct our theory by making thought experiments to answer the question of how a magnetosphere responds to an impinging disk of conducting plasma. We assume the disk has no large scale magnetic field of its own; at present, there is no direct evidence for such fields. In our thought experiments, we first assume ideal MHD describes the interaction, then imagine the effects of dissipation being turned on by the instabilities in the ideal system.

Aly (1980) has found exact analytic solutions to this model problem. The disk is modeled as a perfectly conducting plate with a hole in the center of radius R . The structure of the magnetosphere is outlined in Figure 5, drawn for a rapid rotator so that the corotation radius R_{c*} based on the stellar rotation rate is close to the magnetopause radius – in the jargon of Ghosh and Lamb (1979), the "fastness" $\omega_* = \Omega_*/\Omega_K(R_m) = (R_m/R_{c*})^{3/2}$ is close to unity. A notable feature of Aly's model is the presence of a neutral point at R_n. Field lines which would have closed exterior to R_n are now open, because of the surface currents on the (infinite) disk.

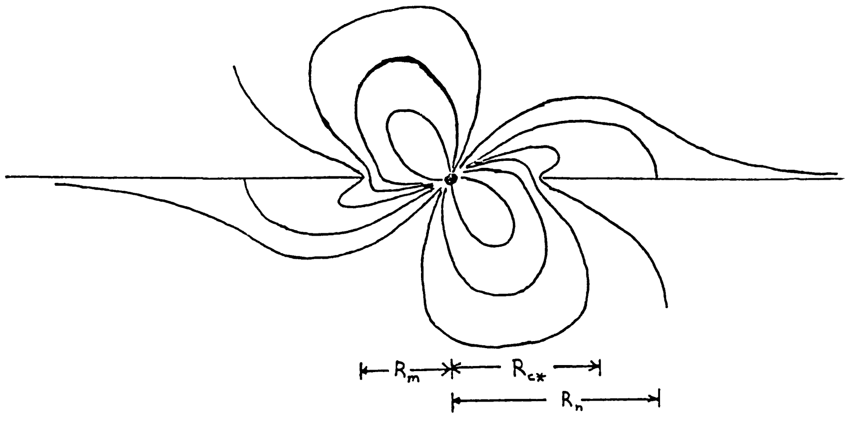

Fig. 5: Field lines of Aly's model. R_m is the magnetopause radius where the disk ends, R_n is the radius of the neutral points, and R_{c*} is the Keplerian corotation radius based on the stellar angular velocity.

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For the rest of my discussion, I assume $R_n > R_m$, as well as assuming R_m to be less than all the other radii which turn out to have significant in forming the wind. Roughly speaking, this implies the angle of obliquity i between the magnetic moment and the rotation axis to be less than $\sim 60°$. Then the topology of the aligned rotator, where R_n is at infinity, is the same as that of the dynamically interesting region near R_m, and we can study the aligned rotator with some confidence that it's physics gives a zeroth order approximation to the oblique problem. When R_n of Aly's model is interior to the radii of interest, the topology is like that assumed in the $i = 90°$ model studied by Anzer and Borner (1980), whose structure is intrinsically 3D. For simplicity, we have studied the aligned case, and have uncovered new physics which we think makes the model relevant to the spin behavior of objects with arbitrary values of the Ghosh–Lamb fastness ω_*. Since the aligned model is relevant to about half the solid angle available to the magnetic moment, we think the aligned case is relevant to the general behavior of accreting magnetized neutron stars.

Interior to R_m, the field is close to the undisturbed dipole, with $B \propto r^{-3}$, while for $r > (2-3)R_m$, the field strength is proportional to $r^{-(2+\nu)}$, with $\nu \simeq 2$. Because Aly models the disk as having no thickness, his solutions have infinite magnetic field at the inner edge. This is artificial, of course, and for a disk of finite half thickness H, his solutions suggest the true value at the inner edge should be $\sim \mu(R_m/H)^{1/2}R_m^{-3}$. We have not attempted to calculate an improved magnetic field model including the finite disk thickness, since we expect that Aly's model provides a reasonable zeroth order estimate of the poloidal field strength, even with dissipation included.

In ideal MHD, the magnetopause separating the disk and the magnetic field lines folded around it is a tangential discontinuity, with no plasma transport across the field lines. However, in general, the magnetospheric field and the disk's plasma move at different speeds, with toroidal relative velocity $v_{rel} = [\Omega_B - \Omega_K(r)]r$. If the field lines corotate with the star, $\Omega_B = \Omega_*$, but we will find that Ω_B can be significantly different than Ω_*. Such relative motion is unstable to the formation of breaking waves (the Kelvin–Helmholtz instability). The mixing induced by such wave breaking is the dominant effect in causing mass to cross the magnetopause boundary, if the disk has no large-scale magnetic field of its own; reconnection between small magnetic loops in the disk and the magnetospheric field acts mainly to trigger the Kelvin–Helmholtz instability by providing a minimum mass density on the magnetospheric side of the interface, if no other mass source is present (Arons et al 1986).

For waves with wavelength along the relative velocity long compared to the thickness of the boundary layer, the linear growth rate is $\gamma \simeq k[\Omega_B - \Omega_K(r)]r(\rho_w/\rho_a)^{1/2}$, where k is the wavenumber, ρ_a is the density in the upper atmosphere of the disk where the boundary layer is formed, and ρ_w is the density in the wind outflow formed within the magnetosphere. From approximate pressure balance, $\rho_a = B^2/8\pi c_s^2(\text{disk}) \sim 10^{-5}$ gm cm^{-3}. Because of X–ray heating, c_s(disk), the sound speed in the disk atmosphere, is $\sim 10^3$ km s^{-1}. We remark that a mass density above 10^{-7} gm cm^{-3} of plasma corotating in the magnetosphere at R_{c*} exerts sufficient stress to burst the magnetosphere, converting closed field to open. Thus, even if only 1% of the disk plasma is scraped off and converted into magnetospheric plasma at radii exceeding R_{c*}, the magnetospheric topology is completely al–

tered. The principle mechanism that does this scrape–off is the breaking of the "cat's eyes" formed by the Kelvin–Helmholtz instability. Hydrodynamic experiments and numerical calculations, and MHD simulations (eg., Miura 1984, 1985) show that after $\eta^{-1} \sim 10$ linear growth times, the waves with wavelength $2\pi/k$ break; the characteristic (vertical) velocity at breaking is therefore the wavelength divided by the time to breaking, or $\eta\gamma/k$. Once the waves break, successive overturns rapidly diffuse the momentum and mass (Arons and Lea 1980), entraining the disk plasma with the magnetosphere. The effective diffusion coefficient for this process is $\sim \eta\gamma/k^2 \sim 0.1\gamma/k^2$; the momentum diffusion coefficient found by Miura is $0.098\gamma/k^2$, once proper account is taken of the distributed shear profile used in his calculations. The instability also generates MHD waves (primarily fast modes); because the Alfven speed on the magnetospheric side of the boundary layer is high compared to the relative speed in the final, self–consistent models, these waves carry off the disk momentum deposited in the magnetosphere with only small amplitude pertubations of the magnetospheric structure – here, we differ from the assertion made by Anzer and Borner, that the magnetic field in and near the boundary layer must be fully turbulent.

The net entry velocity of disk plasma into the magnetosphere is then $\sim \eta\gamma/k = \eta[\Omega_B - \Omega_K(r)]r(\rho_w/\rho_a)^{1/2}$, with the plasma distributed over a layer whose vertical thickness is on the order of the longest wavelengths which can grow and break. For practical purposes, this turns out to be roughly the scale height of the disk atmosphere, since the large inertia of deeper layers in the disk prevents much growth at smaller values of k. The plasma injected onto the field has sufficient rotational inertia to break open the field and be flung away as a centrifugally accelerated wind; at this stage in the thought experiment, we assume this to be so, and find the assumption to be self–consistent at the end. The resulting configuration is shown in the cartoon in Figure 6. Here, the disk lies in the midplane, with the open field lines wrapped around the disk and now disconnected from the star, held in only by the accretion of plasma from the disk's inner edge. Therefore, the angular velocity of the open field lines and of the wind is Ω_B, with $\Omega_* < \Omega_B < \Omega_K(R_m)$. If it turns out that $\Omega_B = \Omega_K(R_m)$, all rotators are fast with respect to angular momentum loss in the wind.

3.3 Angular Momentum Loss and a Quantitative Model

The consequence of our thought experiment, therefore, is that for radii larger than $R_c(\Omega_B) < R_{c*}$, where the field lines rotate faster than the underlying disk, the Kelvin–Helmholtz instability feeds plasma from the disk into the overlying magnetosphere at a rate sufficient to blow the field lines open and maintain them in the opened configuration. Indeed, our thought experiment suggests that closed models, such as that of Ghosh and Lamb (1979), are unstable to switching over into the opened state described here. Furthermore, the plasma injected into the poloidal field must be rotationally spun up, which requires the field lines to be swept back into a spiral. Therefore, the wind carries off angular momentum. Since the field lines connect through the boundary layer between the inner edge of the disk and the corotating, closed field region of the Poynting fluxes from this boundary layer region. At the very least, therefore, the wind provides a reduction of the specific angular momentum accreted below the Keplerian value first estimated by Pringle and

Rees (1972). The possibility of producing a true spin–down torque is discussed briefly below.

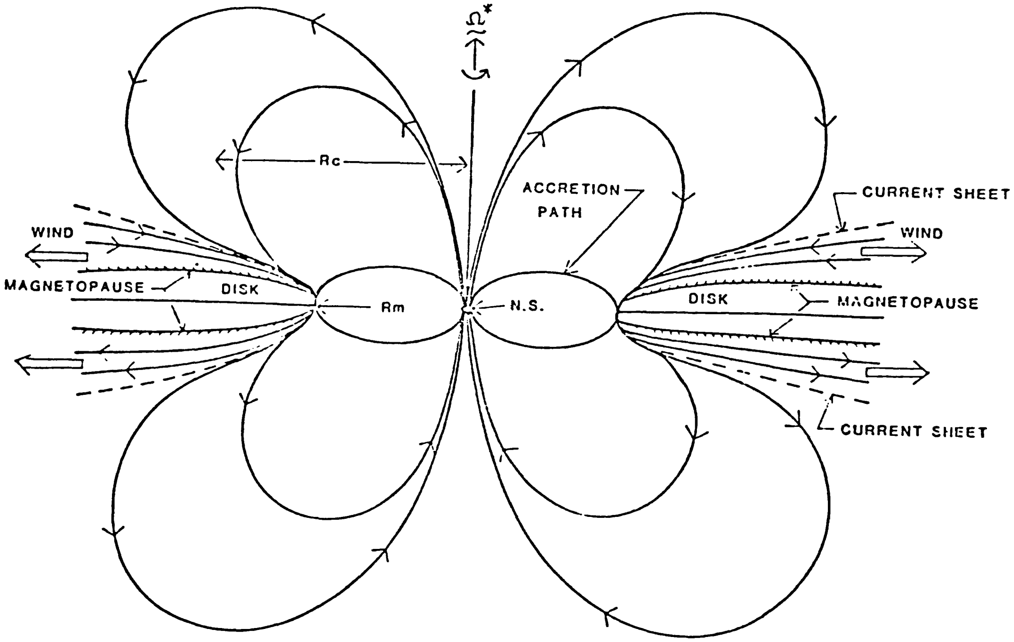

Figure 6: Magnetospheric structure with a wind from a disk.

In our calculations, we assume, and find to be self–consistent, that the region of mass entry is interior to the Alfven radius of the wind R_A, where the poloidal component of the wind velocity equals the poloidal Alfven speed. Then the rate at which angular momentum is lost is given by $T_- = \dot{M}_w \Omega_B [R_A^2 - R_c^2(\Omega\ )]$ (Hartmann and MacGregor 1982). The task of a quantitative model is to find $\dot{M}_w$, Ω_B and R_A, all of which depend on $B[R_c(\Omega_B)]$ and on the temperature of the X–ray heated gas, and also to determine if T_- corresponds to a true spin down torque or is limited to a torque reduction alone; this latter task is closely related to finding Ω_B.

The full details of our calculations are reported elsewhere (Arons *et al* 1986). What we did was to use our simple mixing model as a mass source and a momentum sink for the magnetospheric layers bounding the disk, and integrated the equations for flow along the largely poloidal B to find the asymtototic mass loss rate in the wind. This yields

$$\dot{M}_w \approx \begin{cases} 2 \times 10^{17} \left(\frac{\eta}{0.1}\right)^2 \mu_{30}^2 \left(\frac{10^8 \mathrm{K}}{T_a}\right) P_B^{-10/3} \ \mathrm{g\ sec^{-1}} \\ 10^{18} \left(\frac{\eta}{0.1}\right)^2 \left(\mu/10^{27}\right)\left(\frac{10^8 \mathrm{K}}{T_a}\right) (P_B/10\ \mathrm{msec})^{-10/3} \ \mathrm{g\ sec^{-1}} \end{cases} \qquad (1)$$

these values assume the field strength varies as r^{-3}. We then solve the MHD wind equations as if all the mass is injected with the rate (1) at the corotation radius $R_{co}(P_B)$. This collapse of the mass entry zone to a single injec–

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tion surface works if the Alfven radius is well outside the region of mass entry. The poloidal magnetic field is allowed to vary in proportion to $r^{-(2+\nu)}$. The flow is in approximate corotation with the field lines, at speed $\Omega_B r$, out to the Alfven radius $\approx R_A \approx (\text{flux}^2/\dot{M}_w\Omega_B)^{1/3} \approx 10^{8.2}(0.1/\eta)^{1/3}P_B$ cm ($\nu = 1$). This radius can be guessed by dimensional analysis; all the work of solving the MHD equations yields multiplicative factors nearly equal to unity which depend on ν. The asymptotic wind speed for $\nu = 1$ is $v_\infty = 10^{9.2}\ (0.1/\eta)^{1/2}$ cm s^{-1}. The rate of angular momentum loss is the wind is $T_- = \Omega_B \dot{M}_w(R_A^2 - R_L^2)$. Plugging in yields the ratio of the angular momentum loss rate to the rate stellar angular momentum is gained from accretion

$$\frac{T_-}{T_+} = 1.4\left(\frac{0.1}{\eta}\right)\left(\frac{M_\odot}{M_*}\right)^{23/21} \mu_{30}^{12/7}\left(\frac{10^8 K}{T_a}\right)^{1/2} P_B^{-7/3} L_{37}^{-6/7}. \quad (2)$$

As before, P is the rotation period of the open field lines.

The final step needed is to calculate Ω_B. This requires solving a model of the boundary layer between the disk's inner edge and the co–rotating region of the magnetosphere, which we do by treating the effects of the Kelvin–Helmholtz instability as a mass and momentum diffusion, as well as including the loss of angular momentum from the boundary layer by Poynting fluxes directed along the open field lines which connect to the wind zone. This work is not entirely complete, but present indications are that Ω_B is closer to $\Omega_K(R_m)$ than it is to Ω_*, so that all rotators with obliquity not too large should be interpreted as "fast" rotators. Furthermore, when the accretion rate is sufficiently small, corresponding to accretion luminosities below 10^{37} erg/s, the angular momentum loss into the wind can drive the angular velocity of the boundary layer flow below $\Omega_B R_m$. Since the plasma is brought into co–rotation by viscosity between the co–rotating region and the boundary layer, angular momentum can be extracted from the co–rotating part of the magnetosphere and therefore from the star. When the wind's ability to transport angular momentum is sufficiently large, therefore, the torque T_- **does** correspond to a stellar spindown torque. In its current state of development, the model looks like an interesting candidate to interpret some aspects of the spin behavior of Vela X–1. The quantitative details of these results, and their application to rotational evolution of accreting neutron stars, is described elsewhere (Arons, McKee and Pudritz 1986; Arons and McKee 1986).

3.4. Applications and Conclusions

We have made several preliminary applications of our models to observations, existing or potential, of X–ray pulsars. In the case of Her X–1, the net average spin–up torque is 1/30 of the value expected if no spin–down torques operate. Application of our model to the data yields, assuming $P_B \cong P_*$, a predicted magnetic moment of 1.1 x 10^{30} Gauss–cm^3, with a dipole component of the surface magnetic field of about 2 x 10^{12} Gauss. The result doesn't differ much if the rotation rate of the magnetic field is taken to be the Kepler rate, since in this case R_{c*} is already close to R_m.

A more interesting application of the model is to the interpretation of the broad Fe fluoresence line observed in this and other X–ray pulsars. The fluorescing plasma must have temperature well below 10^7 K, a Thomson optical depth of order unity (eg., Pravdo et al 1979, White et al 1983), and a full width at half–maximum of $\sim 2 \times 10^4$ km/s (Trumper 1986). In fact, our wind model has all these properties, without forcing any of the parameters. The wind plasma is cooled by cyclotron emission to a temperature well below 1 keV. From the mass loss rate (1), the Thomson optical depth is close to unity, and the magnetic moment is inferred by applying the theory to the reduced spin up observed in Her X–1. The predicted velocity full width is then about 15,000 km/s, using the theory without taking proper account of the toroidal fields in the mass entry zone [ie., without worrying about $R_A/R_c(\Omega_B) \sim 1$]. In fact, the theory constructed with proper account of $B_{toroidal}$ gives even better results. Since the basic parameters of our model are already constrained by the dynamical observations (ie., by studying P and dP/dt), the spectroscopic observations provide a test and the possibility of refined modeling of the system.

A third application can be made to the nonconservative evolution of binaries, since the plasma is flung off at 10^4 km s^{-1} and must leave the system. Then it also carries off the orbital angular momentum characteristic of the neutron star, which leads to evolution of the binary on the time scale $\approx 10^8\ (0.1/\eta)^2(10^{27}/\mu)^2(P_B/10\ \text{msec})^{10/3}$ years, a time of possible relevance to the evolution of low mass X–ray binaries.

A final application is to the formation of either enhanced mass loss from the outer regions of the disk, or to the formation of "bubbles" in the interstellar medium. Depending on the shape of the disk's outer region, the wind may or may not collide with the disk. If it misses, the wind goes on and blows a bubble in the interstellar medium, of radius $R_{bubble} \sim 5\ (n_{ISM}/1\ \text{cm}^{-3})^{-1/5}\mu^{2/5}(t/10^5 y)^{3/5}P_B^{-2/3}$ parsecs, rather like the wind from an O star. In this case, a possible test of the theory would be to find bubbles around X–ray sources with B or later optical characteristics, but with bubble properties like those expected for the winds from very early stars (eg., McKee et al 1984). An alternate possibility is that the wind does encounter the disk's outer regions. Then the shock heating might cause the disk plasma to be lifted up above the disk plane, possibly causing the peculiar absorption features seen in some low mass sources whose luminosity is too low to drive the formation of winds from the outer regions by photon heating from the central source.

Our preliminary conclusion is that centrifugally driven winds from the inner edges of disks around magnetized neutron stars are likely to form at a level interesting to the interpretation of dynamical and spectroscopic phenomena observed in these systems.

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References

Alpar, A., Cheng, A.F., Ruderman, M.A., and Shaham, J. 1982, *Nature*, **300**, 728

Aly, J.–J. 1980, *Astron. and Ap.*, **86**, 192

Anzer, U., and Borner, G. 1980, *Astron. and Ap.*, **83**, 133

_______. 1983, *op. cit.*, **122**, 73

Arons, J., Burnard, D.J., Klein, R.I., McKee, C.F., Pudritz, R.E., and Lea, S.M. 1984, in *Proc. Santa Cruz Workshop 'High Energy Transients in Astrophysics'*, S.E. Woosley, ed. (New York: American Institute of Physics), 215

Arons, J., Hernquist, L., and Klein, R.I. 1985, *Bull. Am. Phys. Soc.*, **30**, 1471

Arons, J., and Klein, R.I. 1987a, in *Proc. IAU Symp. No. 125, 'Origin and Evolution of Neutron Stars'*, D.J. Helfand, ed. (Dordrecht: D.Reidel).

_____. 1987b, *Ap. J.*, submitted

Arons, J., Klein, R.I., and Lea, S.M. 1986, *Ap. J.*, in press

Arons, J., and Lea, S.M. 1976, *Ap.J.*, **207**, 914

_______. 1980, *Ap.J.*, **235**, 1016

Arons, J., and McKee, C.F. 1986, in preparation

Arons, J., McKee, C.F., and Pudritz, R. 1986, in preparation

Basko, M.M., and Sunyaev, R.A. 1976, *MNRAS*, **175**, 395

Boynton, P.E., Crosa, L.M., and Deeter, J.E. 1980, *Ap.J.*, **237**, 169

Braun, A., and Yahel, R.Z. 1984, *Ap. J.*, **278**, 349

Burnard, D.J., Lea, S.M., and Arons, J. 1983, *Ap.J.*, **266**, 175

Damashek, M., Backus, P.R., Taylor, J.H., and Burkhardt, R.K. 1982, *Ap.J.(Lett.)*, **253**, L57

Davidson, K. 1973, *Nature Phys. Sci.*, **246**, 1

Davidson, K., and Ostriker, J.P. 1973, *Ap.J.*, **179**, 585

Elsner, R.F., and Lamb, F.K., 1977, *Ap.J.*, **215**, 897

_____. 1984, *Ap. J.*, **278**, 326

Ghosh, P., and Lamb, F.K. 1979, *Ap.J.*, **232**, 259

Hartmann, L., and MacGregor, K.B. 1982, *Ap.J.*, **259**, 180

Hameury, J.M., Bonnazola, S., Heyvaerts, J., and Lasota, J.P. 1984, *Adv. Space Res.*, No. 10–12, 297

Hsieh, S.–H., and Spiegel, E.A. 1976, *Ap.J.*, **207**, 244

Kirk, J. 1985, *Astron. and Ap.*,**142**, 430

Kirk, J., and Melrose, D.B. 1986, *Astron. and Ap.*, **156**, 257

Klein, R.I., and Arons, J. 1987a, in *Proc. IAU Symp. No. 125, 'Origin and Evolution of Neutron Stars'*, D.J. Helfand, ed. (Dordrecht: D.Reidel).

_____. 1987b, to be submitted to *Ap. J.*

Lamb, F.K. 1984,in *Proc. Santa Cruz Workshop 'High Energy Transients in Astrophysics'*, S.E. Woosley, ed. (New York: American Institute of Physics), 179

Lamb, F.K., Pethick, C.J., and Pines, D. 1973, *Ap.J.*, **184**, 271

Lamb, F.K., Shibazaki, N., Alpar, A. and Shaham, J. 1985, *Nature*,

Mason, K. 1977, *MNRAS*, **178**, 818

McKee, C.F., van Buren, D., and Lazareff, B. 1984, *Ap.J.*, **278**, L115

Meszaros, P., and Nagel, W. 1986, *Ap.J.*, in press
Michel, F.C. 1977, *Ap.J.*, **214**, 261
Miura, A. 1984, *J. Geophys. res.*, **89**, 801
______. 1985, *Geophys. Res. Lett.*, **12**, 635
Pravdo, S.H., *et al.* 1979, *Ap.J.*, **231**, 912
Proctor, M., and Weiss, N. 1982, *Rep. Prog. Phys.*, **45**, 1317
Pringle, J.E., and Rees, M.J. 1972, *Astron. and Ap.*, **21**, 1
Rappaport, S., and Joss, P.C. 1977, *Nature*, **266**, 683
Ruderman, M.A. 1986, *Cargese Lectures (1985)*, to be published
Shapiro, S.L., and Teukolsky, S. 1983, *Black Holes, Neutron Stars and White Dwarfs* (New York: John Wiley)
Stella, L., White, N.E., and Rosner, R. 1986, *Ap.J.*, in press
Trumper, J. 1986, in *Proc. Mediterranean School on Plasma Astrophysics*, N. Kylafis and J. Ventura, eds. (Iraklion: University of Crete), in press
Voges, W., *et al.* 1982, *Ap.J.*, **263**, 803
Wang, Y.–M., and Frank, J. 1981, *Astron. and Ap.*, **93**, 255
White, N.E., Swank, J.H., and Holt, S.S. 1983, *Ap.J.*, **270**, 711

DISCUSSION

A. Burrows: Have you considered energetic particle production in this coupled neutron star/disk context?

J. Arons: Yes, I don't think disk electrodynamic leads to much interesting acceleration in the X-ray source context. Dissipative MHD instability at the bottom of the accretion column is more interesting, but the story is too long for discussion here.

J. Grindlay: Does your work provide any clues as to why cyclotron line features are apparently so difficult to detect in X-ray pulsars with magnetic fields presumably in the 10^{12} - 10^{13} gauss range (e.g. from spin-equilibrium arguments)?

J. Arons: I don't think spin equilibrium or other large scale aspects of the magnetosphere have much to do with line visibility. The dominant effect is likely to be velocity shear, which is enormous in the models I showed. Unless one looks almost across the magnetic field, the variable Doppler shift from the bulk motion smears out the line. If the field is simple, and if one could look at it broadside, then one has a chance to see a narrow feature. Clearly, the phase space for this is small, so finding line features in only a few objects is consistent with the kind of theory I've described.

D. Arnett: What is the ion temperature in the wind?

J. Arons: The collisional equilibration time is short, so it's probably about the same as the electrons, $\sim 10^6$ K.

X.J. Mao: What kind of material is in the disk of your model? According to the Goldreich-Julain model, the density of plasma is getting lower and lower and your disk doesn't go into the magnetosphere of the pulsar.

J. Arons: The disk is made of normal plasma accreted from the normal star in the binary. The Goldreich-Julian charge density occurs in the corotating region of the magnetosphere, and is easily maintained by the accreting plasma which passes through the magnetopause onto closed field lines.

J.H. You: If the temperature in the accreting column is $\gtrsim 10^8$ K, pair production will occur. It is so important for cooling, and radiation from pair-annihiliation may be dominant over bremstrahlung and cyclotron radiation. Don't you consider this effect?

J. Arons: Yes, maybe it can't be neglected and it should be taken into account.