The Astrophysical Journal, 609:L17–L20, 2004 July 1


# HIGH-ENERGY EMISSION FROM THE DOUBLE PULSAR SYSTEM J0737−3039

Jonathan Granot and Peter Mészáros[1]
Institute for Advanced Study, Einstein Drive, Princeton, NJ 08540; granot@ias.edu, pim@ias.edu
*Received 2004 March 11; accepted 2004 May 19; published 2004 May 25*

## ABSTRACT

We discuss the effects of particle acceleration at the bow shocks expected in the binary pulsar system J0737−3039, because of the wind from pulsar A interacting with both the interstellar medium (ISM) and the magnetosphere of pulsar B. In this model, we find that the likeliest source for the X-rays observed by *Chandra* is the emission from the shocked wind of pulsar A as it interacts with the ISM. In this case, for favorable model parameter values, better statistics might help *Chandra* marginally resolve the source. A consequence of the model is a power-law high-energy spectrum extending up to $\lesssim$60 keV, at a level of $\sim 2 \times 10^{-13}$ ergs cm$^{-2}$ s$^{-1}$.

*Subject headings:* gamma rays: observations — pulsars: general — pulsars: individual (PSR 0737−3039A, PSR 0737−3039B) — radiation mechanisms: nonthermal — X-rays: binaries

## 1. INTRODUCTION

The double radio pulsar system J0737−3039 (Lyne et al. 2004; Kaspi et al. 2004) is of great interest as a remarkable laboratory for probing strong field gravity and magnetospheric interactions. It has also been detected in a 10 ks *Chandra* observation (McLaughlin et al. 2004), with an X-ray luminosity of $L_X \approx 2 \times 10^{30}(d/0.5\ \mathrm{kpc})^2$ ergs s$^{-1}$ in the 0.2–10 keV range (where $d$ is the distance to the source) and a reported X-ray photon number index of $\Gamma = 2.9 \pm 0.4$. The spin-down luminosity of pulsar A, which is expected to be channeled mainly into its relativistic wind, is $\dot{E}_A \simeq L_A \simeq 6 \times 10^{33}$ ergs s$^{-1}$ (Lyne et al. 2004; Kaspi et al. 2004). Since $L_A \sim 3 \times 10^3 L_X$, only a small fraction of $L_A$ is required in order to produce the observed X-ray emission. Since only 77 ± 9 X-ray photons were detected, the determination of the spectral slope is difficult and might be consistent with a flat $\nu F_\nu$ ($\Gamma \sim 2$), as expected from shock acceleration. Here, we explore whether particle acceleration in the bow shocks of the pulsar A relativistic wind can explain the properties of the X-ray emission. The bow shock on the the magnetosphere of pulsar B involves only a small fraction of the pulsar A wind because of the small solid angle that it extends as seen from pulsar A. Therefore, it must have a very high radiative efficiency in order to explain the observed X-ray luminosity. On the other hand, the bow shock on the interstellar medium (ISM) involves most of the pulsar A wind and thus allows for a significantly smaller and more realistic radiative efficiency. We evaluate the expected high-energy emission from this shock model, which also predicts emission up to tens of keV.

## 2. EMISSION FROM THE BOW SHOCK ON THE ISM

At a sufficiently large distance from the double pulsar system, a bow shock forms because of the interaction of the wind from pulsar A with the ISM.[2] This situation is similar to that for a millisecond pulsar with a close low-mass binary companion (Arons & Tavani 1993), as far as the interaction between the pulsar wind and the ISM is concerned. The relative velocity of the center of mass of the binary pulsar with respect to the ISM is 140.9 ± 6.2 km s$^{-1}$ on the plane of the sky (Ransom et al. 2004). A velocity component along our line of sight could lead to a larger total velocity, $v_{\rm ext} = 200 v_{200}$ km s$^{-1}$ with $v_{200} \gtrsim 1$. The head of the bow shock is at a distance $R$ from pulsar A where the kinetic pressure of the wind balances the ram pressure of the ambient medium, $\rho_{\rm ext} v_{\rm ext}^2$,

$$R = \sqrt{\frac{L_A}{4\pi\rho_{\rm ext} v_{\rm ext}^2 c}} = 4.9 \times 10^{15} n_0^{-1/2} v_{200}^{-1}\ \mathrm{cm}, \qquad (1)$$

where $\rho_{\rm ext} = n_{\rm ext} m_p$ and $n_{\rm ext} = n_0$ cm$^{-3}$ are the ambient mass density and number density, respectively.

Pulsar winds are thought to have a pair plasma composition, perhaps with ions in restricted latitude sectors, and a high asymptotic bulk Lorentz factor (perhaps as high as $\sim 10^4$–$10^6$ in the Crab Nebula and other young pulsar wind nebulae). For simplicity we assume a pure $e^\pm$ pair plasma that holds a fraction $\epsilon_e \approx 1$ of the internal energy behind the shock. We use a fiducial value of $\gamma_w = 10^5 \gamma_{w,5}$ for the wind Lorentz factor just before the shock; however, our main results are rather insensitive to the exact value of $\gamma_w$. We assume $\gamma_w \gg 1$ throughout this work.

The ratio $\sigma$ of Poynting flux to kinetic energy in the wind is believed to be $\sigma \gg 1$ at very small radii, while low values of $\sigma \ll 1$ at large radii are inferred from observations (e.g., $\sigma \sim 3 \times 10^{-3}$ for the Crab; Gallant & Arons 1994; Spitkovsky & Arons 2004). It is hard to estimate the value of $\sigma$ at intermediate radii, which are relevant for our purposes. For the bow shock with the ISM that is at a relatively large radius, one might expect $\sigma \lesssim 1$. The shock jump conditions imply that the fraction $\epsilon_B$ of the internal energy behind the shock in the magnetic field is $\epsilon_B \sim \sigma$. However, amplification of the magnetic field in the shock itself could produce $\epsilon_B \sim 1$ even if $\sigma \ll 1$. Conversely, for $\sigma > 1$ magnetic dissipation behind the shock might decrease the value of $\epsilon_B$ and make it close to unity. Therefore, we assume $\epsilon_B \sim 1$, and to zeroth order we neglect the effect of the magnetic field on the shock jump conditions.

In order to estimate the emission from the shocked wind, we use the values of the hydrodynamical quantities at the head of the bow shock. To first order we neglect the orbital motion of pulsar A. The proper number density in the wind, as a function of the distance $r$ from pulsar A, is $n_w = L_A/4\pi r^2 m_e c^3 \gamma_w^2$. The

[1] Also at the Department of Astronomy and Astrophysics, Department of Physics, and Center for Gravitational Physics at Pennsylvania State University.

[2] The spin-down power of pulsar B is $\sim 3 \times 10^3$ times smaller than that of pulsar A, so that its wind should have a negligible effect on the bow shock with the ISM.

shock jump conditions at the head of the bow shock imply that the shocked pulsar A wind just behind the shock moves away from the shock at $\beta = \frac{1}{3}$ and has a proper energy density $e_{\rm int} = L_{\rm A}/2\pi R^2 c \approx 1.3 \times 10^{-9} n_0 v_{200}^2$ ergs cm$^{-3}$ and a proper number density $n = 2^{3/2}\gamma_w n_w = e_{\rm int}/(m_e c^2 \gamma_w/\sqrt{2}) = 2.3 \times 10^{-8} n_0 v_{200}^2 \gamma_{w,5}^{-1}$ cm$^{-3}$. This implies a magnetic field of $B = 1.8 \times 10^{-4} n_0^{1/2} v_{200} \epsilon_B^{1/2}$ G (in the fluid rest frame). The $e^{\pm}$ pairs are assumed to be accelerated by the shock into a power-law energy distribution $dn/d\gamma_e \propto \gamma_e^{-p}$, with $\gamma_m < \gamma_e < \gamma_{\max}$. Observations of synchrotron emission from electrons accelerated in relativistic collisionless shocks typically imply $p \sim 2$–3. The average random Lorentz factor of the shocked electrons is $\langle\gamma_e\rangle = \gamma_w/\sqrt{2}$, and the minimal Lorentz factor is given by[3]

$$\gamma_m = \frac{p-2}{p-1}\langle\gamma_e\rangle = \frac{g\epsilon_e\gamma_w}{3\sqrt{2}} = 2.4 \times 10^4 g\epsilon_e\gamma_{w,5}, \qquad (2)$$

where $g \equiv 3(p-2)/(p-1)$ equals 1 for $p = 2.5$. The maximal Lorentz factor, from the requirement that the Larmor radius $R_{\rm L} = \gamma_e m_e c^2/eB$ does not exceed the width $\eta R$ of the layer of shocked fluid, is

$$\gamma_{\max,1} = \frac{eB\eta R}{m_e c^2} = 1.2 \times 10^7 \epsilon_B^{1/2}(\eta/0.3). \qquad (3)$$

Here, the value of $\eta$ can be estimated by equating the particle injection rate into the hemisphere containing the head of the bow shock ($\theta \leq 90°$), $\dot{N}/2 = L/2\gamma_w m_e c^2$, to the flow of shocked particles behind the shock outside of this hemisphere, $2\pi\eta R^2 nu$, where $n = 2^{3/2}\gamma_w n_w = \dot{N}/\sqrt{2}\pi R^2 c$ and $u = \gamma\beta$ are the proper density and 4-velocity (in the direction perpendicular to the shock) of the shocked wind at $\theta = 90°$. This gives $\eta \approx 1/(2^{3/2}u)$, so that $\eta < 1$ implies $\beta > \frac{1}{3}$. At $\theta = 90°$ we expect $\beta \gtrsim c_s/c \approx 3^{-1/2}$ and $u \gtrsim 2^{-1/2}$ so that $\eta \lesssim \frac{1}{2}$. On the other hand, $\eta < 0.1$ requires $u > 5/\sqrt{2} \approx 3.5$, which begins to be highly supersonic and is therefore not very reasonable. Hence we expect $0.1 \lesssim \eta \lesssim 0.5$ and use a fiducial value of $\eta = 0.3$.

The dominant emission mechanism is synchrotron radiation, and inverse Compton scattering can be neglected. The Lorentz factor of an electron that cools on the dynamical time, $t_{\rm dyn} \sim R/(c/3) = 4.9 \times 10^5 n_0^{-1/2} v_{200}^{-1}$ s, is given by

$$\gamma_c = \frac{6\pi m_e c}{\sigma_{\rm T} B^2 t_{\rm dyn}} = \frac{4.7 \times 10^{10}}{\epsilon_B n_0^{1/2} v_{200}}. \qquad (4)$$

The synchrotron spectral break frequencies corresponding to $\gamma_m$, $\gamma_c$, and $\gamma_{\max}$ are

$$\nu_m = 3.4 \times 10^{11} g^2 \epsilon_B^{1/2}\epsilon_e^2 n_0^{1/2} v_{200}\gamma_{w,5}^2 \text{ Hz}, \qquad (5)$$

$$h\nu_c = 5.5\epsilon_B^{-3/2} n_0^{-1/2} v_{200}^{-1} \text{ GeV}, \qquad (6)$$

$$h\nu_{\max} = 62 n_0^{1/2} v_{200}\epsilon_B^{3/2}(\eta/0.3)^2 \text{ keV}. \qquad (7)$$

We have $\nu F_\nu \propto \nu^{4/3}$ for $\nu < \nu_m$, $\nu F_\nu \propto \nu^{(3-p)/2}$ for $\nu_m < \nu < \min(\nu_c, \nu_{\max})$, and if $\nu_{\max} > \nu_c$ (which is relevant for the next section) we have $\nu F_\nu \propto \nu^{(2-p)/2}$ for $\nu_c < \nu < \nu_{\max}$.

Since $\gamma_c > \gamma_{\max}$, all electrons radiate only a small fraction of their energy. The fraction of energy radiated by an electron is $\sim\min(1, t_{\rm dyn}/t_c) = \min(1, \gamma_e/\gamma_c)$, where $t_c = 6\pi m_e c/\sigma_{\rm T} B^2\gamma_e$ is the synchrotron cooling time. Averaging over the power-law electron energy distribution, we obtain the total fraction $\epsilon_{\rm rad}$ of energy in electrons that is radiated away. For $\gamma_c < \gamma_m$ (fast cooling), $\epsilon_{\rm rad} \approx 1$, since all electrons cool significantly within $t_{\rm dyn}$. For $\gamma_m < \gamma_c < \gamma_{\max}$,

$$\epsilon_{\rm rad} \approx \begin{cases} 1, & p < 2, \\ [1 + \ln(\gamma_{\max}/\gamma_c)]/\ln(\gamma_{\max}/\gamma_m), & p = 2, \\ (3-p)^{-1}(\gamma_m/\gamma_c)^{p-2}, & 2 < p < 3; \end{cases} \qquad (8)$$

for $\gamma_c > \gamma_{\max}$ we have $\epsilon_{\rm rad}(\gamma_c > \gamma_{\max}) \sim (\gamma_{\max}/\gamma_c)\epsilon_{\rm rad}(\gamma_m < \gamma_c < \gamma_{\max})$, or

$$\epsilon_{\rm rad} \approx (\gamma_{\max}/\gamma_c)\begin{cases} (p-2)/(3-p), & p < 2, \\ 1/\ln(\gamma_{\max}/\gamma_m), & p = 2, \\ [(p-2)/(3-p)](\gamma_m/\gamma_{\max})^{p-2}, & 2 < p < 3. \end{cases} \qquad (9)$$

For our fiducial parameters and $p \approx 2$, we have $\epsilon_{\rm rad} \approx 4.3 \times 10^{-4} n_0^{1/2} v_{200}\epsilon_B^{3/2}(\eta/0.3)$. Most of the radiated energy will be emitted near $\nu_{\max}$ at tens of keV. The fraction $f_{\rm X}$ of the radiated energy in the 0.2–10 keV *Chandra* range is approximately given by the ratio of the average $\nu F_\nu$-value in the *Chandra* range (equal to the $\nu F_\nu$-value at some frequency $\nu_{\rm X}$ within that range) to the peak $\nu F_\nu$-value. In our case, $f_{\rm X} \sim (\nu_{\rm X}/\nu_{\max})^{(3-p)/2} \approx 0.27 n_0^{-1/4} v_{200}^{-1/2} \times \epsilon_B^{-3/4}(\eta/0.3)^{-1}$ (this expression holds for $\nu_{\rm X} < \nu_{\max} < \nu_c$). Therefore, the ratio of the expected X-ray luminosity $L_{\rm X} = f_{\rm X}\epsilon_e\epsilon_{\rm rad} L_{\rm A} \approx 7 \times 10^{29} n_0^{1/4} v_{200}^{1/2}\epsilon_B^{3/4}$ ergs s$^{-1}$ in the *Chandra* range to the observed $L_{\rm X}^{\rm obs} \approx 2 \times 10^{30}(d/0.5 \text{ kpc})^2$ ergs s$^{-1}$ is $\sim 0.35 n_0^{1/4} v_{200}^{1/2}\epsilon_B^{3/4}(d/0.5 \text{ kpc})^{-2}$. This ratio would be unity, e.g., for $n_0 \sim 60$ with $v_{200} \sim 1$ or for $n_0 \sim 10$ and $v_{200} \sim 2.5$. If $d \approx 1$ kpc instead of $\approx$0.5 kpc, then we would need $n_0 \sim 10^3$ and $v_{200} \sim 4$, which are less likely. Conversely, the constraint is easier to satisfy if $d < 0.5$ kpc. According to this interpretation the high-energy emission should peak at tens of keV (near $\nu_{\max}$ that is given in eq. [7]) with a flux of $\sim L_{\rm X}^{\rm obs}/f_{\rm X} 4\pi d^2 \sim 2.5 \times 10^{-13} n_0^{1/4} v_{200}^{1/2}\epsilon_B^{3/4}(\eta/0.3)$ ergs cm$^{-2}$ s$^{-1}$.

Another contribution to the X-ray luminosity might be expected from the shocked ISM in the bow shock. The energy injection rate is[4] $\sim(v_{\rm ext}/c)L_{\rm A} \sim 10^{-3}L_{\rm A}$, which is of the order of the observed X-ray luminosity and perhaps larger by a factor of a few. This could account for the observed X-ray luminosity, provided that $f_{\rm X}\epsilon_e\epsilon_{\rm rad} \gtrsim 0.1$–0.3. Here, the dynamical time is $t_{\rm dyn} \sim R/v_{\rm ext}$ and $h\nu_c = 22\epsilon_B^{-3/2} n_0^{-1/2} v_{200}$ keV, while $\nu_m$ is very low (in fact $\gamma_m \sim 1$), and the expression for $\nu_{\max}$ is the same as in equation (7) with the difference that here $\eta R$ is the width of the shocked ISM layer (instead of the shocked pulsar wind) and that $\epsilon_B$ might be different (probably somewhat smaller) in the shocked ISM. One might expect instabilities near the contact discontinuity between the shocked wind and the shocked ISM, of both the Rayleigh-Taylor and the Kelvin-Helmholtz types, which could bring the magnetic field in the shocked ISM close to equipartition. For $v_{200} \sim 1$ and $n_0 \gtrsim 10$, $\nu F_\nu$ peaks in the *Chandra* range, so that we can have $f_{\rm X} \sim 1$. Since the shock

[3] More generally, this expression should be multiplied by a factor of $(1 + \rho_p/\rho_e)$, which can be as high as $m_p/m_e$ in the limit of a proton-electron plasma. Also, the factor of $(p-2)/(p-1)$ is valid for $p > 2$, while for $p = 2$ it should be replaced by $1/\ln(\gamma_{\max}/\gamma_m)$ so that $g = 3/\ln(\gamma_{\max}/\gamma_m)$.

[4] Comparing the energy injection rate per unit area into the wind termination shock, $L_{\rm A}/4\pi R^2$, and into the bow shock going into the ISM, $\frac{1}{2}\rho_{\rm ext} v_{\rm ext}^3$, and balancing the two ram pressures, $L_{\rm A}/4\pi R^2 c$ and $\rho_{\rm ext} v_{\rm ext}^2$, respectively.

going into the ISM is Newtonian, one expects $p \approx 2$, as in supernova remnants (SNRs). For $n_0 \sim 60$ this would imply $\epsilon_{\rm rad} \approx 0.2$. From modeling of collisionless shocks in SNRs, which propagate into a similar ISM with similar shock velocities, a typical value of $\epsilon_e \sim 0.1$ might be adopted. The resulting value of $L_{\rm X} \sim (v_{\rm ext}/c) f_{\rm X} \epsilon_e \epsilon_{\rm rad} L_{\rm A} \sim 2 \times 10^{-5} L_{\rm A} \sim 10^{29} v_{300}(\epsilon_e/0.1) \times (f_{\rm X}\epsilon_{\rm rad}/0.2)(d/0.5\ {\rm kpc})^2$ ergs s$^{-1}$ is only $\sim 0.05 L_{\rm X}^{\rm obs}$. Thus, this emission component may not easily account by itself for the *Chandra* observation (unless $\epsilon_e \sim 1$), although it can contribute to that from the shocked wind of pulsar A.

We note that equation (1) implies that the angular distance between the double pulsar system and the head of the ISM bow shock is $\theta_{\rm bs} = 0.65(d/0.5\ {\rm kpc})^{-1} n_0^{-1/2} v_{200}^{-1}$ arcsec, and the relatively bright emission from the bow shock could extend over an angular scale a few times larger than this value. This angular scale may be resolved with *Chandra*, with longer integration times, even though in the 10 ks *Chandra* detection it was reported as a point source (McLaughlin et al. 2004). If resolved, one might constrain the source angular size to $\lesssim 1''$. However, we note that the observed X-ray emission is best explained from the bow shock with the ISM if $n_0^{-1/2} v_{200}^{-1} \sim 0.35 \epsilon_B^{3/4} (d/0.5\ {\rm kpc})^{-2}$, which in turn implies $\theta_{\rm bs} \approx 0.23 \epsilon_B^{3/4} (d/0.5\ {\rm kpc})^{-3}$ arcsec, which may be difficult to resolve with *Chandra* unless $d \lesssim 0.5$ kpc. On the other hand, this suggests that $d \gtrsim 0.3$ kpc, as otherwise the source should have already been resolved by *Chandra*, despite the poor photon statistics in the current observation.

The emission from the bow shock with the ISM is not expected to show significant modulation at the spin period of pulsar A, $P_{\rm A} = 22.7$ ms, or at the orbital period, $P_{\rm orb} = 2.45$ hr. The former is because $P_{\rm A}$ is ~6–7 orders of magnitude smaller than $R/c$.[5] The latter is because the orbital velocity of pulsar A is $v_{\rm orb} \approx 300\ {\rm km\ s^{-1}} \ll c$, and the distance between pulsars A and B is $R_{\rm AB} = 8.8 \times 10^{10}\ {\rm cm} \ll R$.[6]

### 3. EMISSION FROM THE BOW SHOCK AROUND PULSAR B

Balancing the ram pressure of the pulsar A wind with the magnetic pressure of pulsar B, assuming a predominantly dipole magnetic field and a surface field strength of $B_* = 1.2 \times 10^{12}$ G (Lyne et al. 2004), the distance of the head of the bow shock measured from pulsar B is $R_{\rm bs} \approx 6 \times 10^9$ cm. This is ≈0.07 of the separation between the two pulsars, $R_{\rm AB} = 8.8 \times 10^{10}$ cm (Lyne et al. 2004). Therefore, as seen from pulsar A, the fraction of the total solid angle subtended by the bow shock is $\Omega/4\pi = C\pi(R_{\rm bs}/R_{\rm AB})^2/4\pi \approx 10^{-3} C$, where $C \sim$ a few, its value depending on the exact shape of the bow shock. Thus, producing the X-rays in the shocked wind of pulsar A in the bow shock occurring near pulsar B would require an efficiency of $\sim 0.3/C$, in order to account for the *Chandra* observation.

Lyutikov (2004) calculated the asymptotic opening angle of the bow shock and finds it to be $\theta \sim 0.11$–$0.13$ rad for the value of $B_*$ from Lyne et al. (2004). This gives $\Omega/4\pi \sim (3$–$4.2) \times 10^{-3}$, which is in agreement with our estimate here and provides an independent cross-calibration for our parameter $C$, namely, $C \sim 2.6$–$3.6$.

As in § 2, the values of the hydrodynamical quantities at the head of the bow shock are used in order to estimate the emission from the shocked wind. To zeroth order, we neglect the orbital motion of the two pulsars and their spins. The bow shock itself is at rest in the lab frame, in our approximation.[7] The expressions for the hydrodynamical quantities are similar to those in § 2, just that here the distance of the head of the bow shock from pulsar A, $R = R_{\rm AB} - R_{\rm bs} \approx 8.2 \times 10^{10}$ cm, is $\sim 10^5$ times smaller. Therefore, we have $e_{\rm int} = 4.7$ ergs cm$^{-3}$, $n = 82 \gamma_{w,5}^{-1}$ cm$^{-3}$, and $B = 11 \epsilon_B^{1/2}$ G. While the dynamical time in this case is much shorter, $t_{\rm dyn} \sim R_{\rm bs}/(c/3) = 0.6$ s, the synchrotron cooling time $t_c \propto R^2$ is smaller by an even larger factor, so that $\gamma_c \lesssim \gamma_{\rm max}$. Here $\gamma_{\rm max}$ is also constrained by radiative losses. This limit may be obtained by equating the cooling time $t_c$ to the acceleration time, $t_{\rm acc} = A(2\pi m_e c \gamma_e / eB)$, where $A \gtrsim 1$, $\gamma_{\rm max,2} = (3e/A\sigma_{\rm T} B)^{1/2} = 1.4 \times 10^7 A^{-1/2} \epsilon_B^{-1/4}$. The limit discussed in § 2 now reads $\gamma_{\rm max,1} = 1.2 \times 10^7 \epsilon_B^{1/2} (\eta/0.3)$, and we have $\gamma_{\rm max} = \min(\gamma_{\rm max,1}, \gamma_{\rm max,2})$.

Since the bow shock around pulsar B is much more compact than the bow shock on the ISM, one might expect inverse Compton scattering of the synchrotron photons to be more important in this case, and therefore we check this. The Compton $y$-parameter is given by $Y \sim \tau_{\rm T} \gamma_m^{p-1} \gamma_c^{3-p}$ for $2 < p < 3$, $Y \sim \tau_{\rm T} \gamma_m \gamma_c [1 + \ln(\gamma_{\rm max}/\gamma_c)]$ for $p = 2$, and $Y \sim \tau_{\rm T} \gamma_m^{p-1} \gamma_c \gamma_{\rm max}^{2-p}$ for $p < 2$ (Panaitescu & Kumar 2000). We expect $p \approx 2$, for which $Y(1+Y) \sim 10^{-2} (\eta/0.3) \epsilon_e/\epsilon_B$. For our values of $\epsilon_e \sim \epsilon_B \sim 1$, this gives $Y \sim 10^{-2}$, so that inverse Compton scattering is not very important and is neglected in our treatment.

We have $\gamma_c = 1.1 \times 10^7 \epsilon_B^{-1}$, and $\gamma_m$ is still given by equation (2). The corresponding synchrotron frequencies are

$$h\nu_m = 86 g^2 \epsilon_B^{1/2} \epsilon_e^2 \gamma_{w,5}^2\ {\rm eV},$$

$$h\nu_c = 17(1+Y)^{-2} \epsilon_B^{-3/2}\ {\rm MeV},$$

$$h\nu_{\rm max} = \min\,[30 A^{-1}(1+Y)^{-1},\ 20 \epsilon_B^{3/2} (\eta/0.3)^2]\ {\rm MeV}.$$

The X-ray luminosity is then $L_{\rm X} = f_{\rm X} \epsilon_e \epsilon_{\rm rad} (\Omega/4\pi) L_{\rm A}$. For $p \approx 2$ typically $\epsilon_{\rm rad} \sim 0.2$, and $f_{\rm X} \sim (\nu_{\rm X}/\nu_c)^{(3-p)/2} \sim 0.02 \epsilon_B^{3/4}$. Altogether, and assuming $\Omega/4\pi \sim 4 \times 10^{-3}$ ($C \sim 4$), we have $L_{\rm X} \sim 10^{29} \epsilon_B^{3/4}$ ergs s$^{-1}$, which is a factor of $\sim 20 \epsilon_B^{-3/4}$ smaller than the observed value.

It might still be possible to increase $L_{\rm X}$ if somehow $\nu_c$ could be lowered, since this would significantly increase $f_{\rm X}$ and also somewhat increase $\epsilon_{\rm rad}$. This could potentially be achieved if $t_{\rm dyn}$ or the magnetic field experienced by the shocked electrons are increased (as for $p \approx 2$, $f_{\rm X} \propto \nu_c^{-1/2} \propto t_{\rm dyn} B^{3/2}$). This might

[5] This has two effects. First, any variability in the wind with the period $P_{\rm A}$ will be strongly smoothed out by the time it reaches the bow shock. Second, the distance of the bow shock from pulsar A varies, with $\Delta R \sim R$, so that the phase of the pulsar A wind that impinges upon it at any given time changes by $\sim 10^6$–$10^7$ periods. Since the same holds for the observed emission, it significantly averages out a possible modulation with a period of $P_{\rm A}$, even if it exists in the local emission from a given location along the bow shock.

[6] The orbital motion of pulsar A affects the bow shock with the ISM mainly in two ways. First, the distance between pulsar A and the head of the bow shock changes by $\sim \pm R_{\rm AB}/2$, changing the ram pressure by $\Delta p/p \sim 2R_{\rm AB}/R \sim 10^{-5}$ to $10^{-4}$. Second, the wind is highly relativistic and behaves as radiation, so that its intensity in the bow shock rest frame scales as the fourth power of the Doppler factor $\delta \approx 1 + \beta_{\rm orb} \cos\theta$ and will change in the range $(1 \pm \beta_{\rm orb})^4$, resulting in a relative amplitude of $\approx 8\beta_{\rm orb} \approx 0.8\%$. Since $(R/c)/P_{\rm orb} \approx 18 n_0^{-1/2} v_{200}^{-1}$, some additional averaging can occur because of the different phase of this modulation over the different parts of the bow shock, although this effect is not very large for our most promising model, for which $R$ is smaller by a factor of ~8 compared to its fiducial value in eq. (1).

[7] We ignore the slower binary period timescale, which would cause inertial effects, centrifugal, and Coriolis, etc., as well as the possible time variability due to the rotation of the pulsar B magnetosphere.

happen if a reasonable fraction[8] of the shocked wind becomes associated with the closed magnetic field lines of pulsar B for one or more rotational periods of B (where $P_B$ is ~4.6 times larger than the estimate that we used for $t_{dyn}$, i.e., $R_{bs}/(c/3) \approx 0.6$ s). In this case, this material will also pass through regions of higher magnetic field strength. This could be the case if, e.g., interchange instabilities cause mixing of the two fluids across the contact discontinuity.

One might expect a modulation in the emission with the orbital period due to the change in the line of sight with respect to the bow shock (Arons & Tavani 1993). The shocked wind is expected to move away from the head of the bow shock with a mildly relativistic velocity, in a direction roughly parallel to the bow shock (Lyutikov 2004). This might cause a mild relativistic beaming of the radiation emitted by the shocked plasma, resulting in a mild modulation (by a factor of $\lesssim 2$; Arons & Tavani 1993) of the observed emission as a function of the orbital phase. Another possible source of modulation with the orbital period may arise if the luminosity of the pulsar A wind depends on the angle from its rotational axis (Demorest et al. 2004). In this case the wind luminosity in the direction of pulsar B will vary with a period $P_{orb}$. The duration of the *Chandra* observation, $10^4$ s, is close to the orbital period $P_{orb} = 2.45$ hr, and it showed no evidence for variability (McLaughlin et al. 2004). However, the small number of photons ($77 \pm 9$) does not allow to place a strong limit on a possible modulation with the orbital period, which might still have an amplitude of $\lesssim 50\%$.

The rotation of pulsar B, assuming some misalignment of its magnetic pole relative to its spin axis (as expected from the detection of its pulses), would cause a periodic change in $\Omega$ (with a periodicity equal to the spin period $P_B = 2.77$ s), with an amplitude that is typically of the order of unity (Lyutikov 2004). The distance of the bow shock from pulsar A hardly changes, and therefore the values of the thermodynamic quantities in the shocked wind and the resulting values of $f_X$ and $\epsilon_{rad}$ should vary with a smaller amplitude. Thus, the modulation in $L_X$ is expected to largely follow that in $\Omega$ and have a similar amplitude (typically of the order of unity).

[8] In the bow shock of the solar wind around the Earth, only ~$10^{-3}$ of the wind particles get captured by the Earth's magnetic field. However, the situation there is different in several respects from our case. For example, the Earth's magnetic field is nearly aligned with its rotational axis, while the solar wind is Newtonian (~400 km $s^{-1}$) with relatively low magnetization and includes protons and electrons in roughly equal numbers. Therefore this fraction might be larger in our case and could possibly be sufficiently large for our purposes, although this is uncertain.

## 4. DISCUSSION

Particle acceleration is expected in the binary pulsar system J0737−3039 from both the bow shock of the pulsar A wind as it interacts with the ISM and the bow shock of the wind of pulsar A interacting with the magnetosphere of pulsar B. The rotational energy loss rate, the systemic velocity, and the orbital separation determine the effective angles subtended by these bow shocks, as well as the synchrotron peak energies in the forward and reverse shock systems and the radiation efficiencies at various frequencies. In this model, the likeliest explanation for the *Chandra* emission (McLaughlin et al. 2004) is the pulsar A wind just behind the bow shock caused by the systemic motion in the ISM. In this case, we predict a power-law spectrum that extends up to $\lesssim 60$ keV.

The eclipse of the pulsar A radio emission near superior conjunction is best explained as synchrotron absorption by the shocked pulsar A wind in the bow shock around pulsar B (Kaspi et al. 2004; Demorest et al. 2004; Lyutikov 2004; Arons et al. 2004). This explanation requires a relatively large number density of $e^{\pm}$ pairs, which in turn requires a relatively low wind Lorentz factor, $\gamma_w \lesssim 100$. However, the X-ray emission from both of the bow shocks is not very sensitive to the exact value of $\gamma_w$, and $\gamma_w \sim 10$–100 would only lower the radiative efficiency $\epsilon_{rad}$ and the X-ray luminosity $L_X$ by a factor of ~2 (for $p \approx 2$) compared to $\gamma_w \sim 10^5$.

An alternative explanation for the X-ray emission is simply emission from pulsar A (McLaughlin et al. 2004; Zhang & Harding 2000).[9] In this case a large part of the X-ray emission is expected to be pulsed with a period $P_A$. In contrast, the emission from the bow shock around pulsar B might be modulated[10] at $P_{orb}$ or $P_B$, while the emission from the bow shock with the ISM is not expected to be modulated but might be angularly resolved by *Chandra*.

We thank Roman Rafikov and Peter Goldreich for useful discussions and the referee for a careful reading and constructive comments. J. G. is supported by the W. M. Keck Foundation and by NSF grant PHY-0070928. P. M. is supported by the Monell Foundation, NASA NAG5-13286, and NSF AST 00-98416.

[9] Emission from pulsar B is unlikely, since $\dot{E}_{rot,B} \sim L_X$, which would require a very high efficiency in producing X-rays.

[10] Although we find that the emission from the bow shock around pulsar B is likely to contribute only a few percent of the total X-ray luminosity from this system, it can still produce an overall modulation of up to several percent, which might still be detectable.

## REFERENCES

Arons, J., Backer, D. C., Spitkovsky, A., & Kaspi, V. M. 2004, in Binary Radio Pulsars, ed. F. A. Rasio & I. H. Stairs (San Francisco: ASP), in press (astro-ph/0404159)
Arons, J., & Tavani, M. 1993, ApJ, 403, 249
Demorest, P., Ramachandran, R., Backer, D. C., Ransom, S. M., Kaspi, V., Arons, J., & Spitkovsky, A. 2004, ApJ, submitted (astro-ph/0402025)
Gallant, Y. A., & Arons, J. 1994, ApJ, 435, 230
Kaspi, V. M., Ransom, S. M., Backer, D. C., Ramachandran, R., Demorest, P., Arons, J., & Spitkovsky, A. 2004, ApJ, submitted (astro-ph/0401614)
Lyne, A. G., et al. 2004, Science, 303, 1153
Lyutikov, M. 2004, MNRAS, submitted (astro-ph/0403076)
McLaughlin, M. A., et al. 2004, ApJ, 605, L41
Panaitescu, A., & Kumar, P. 2000, ApJ, 543, 66
Ransom, S. M., Kaspi, V. M., Ramachandran, R., Demorest, P., Backer, D. C., Pfahl, E. D., Ghigo, F. D., & Kaplan, D. L. 2004, ApJ, submitted (astro-ph/0404149)
Spitkovsky, A., & Arons, J. 2004, ApJ, 603, 669
Zhang, B., & Harding, A. K. 2000, ApJ, 532, 1150