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PHYSICAL IMPLICATIONS OF THE ECLIPSING BINARY PULSAR

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ABSTRACT

Physical properties of the P1957+20 eclipsing binary pulsar are discussed. We argue that (1) the pulsar-heated companion should appear optically variable with color temperature $\sim 10^3$–10^4 K at peak light, (2) such illumination requires that a significant fraction of the pulsar emission escapes in the form of high-energy particles and radiation, (3) the pulsar radio luminosity may be sufficient to deflect the wind from the companion, and (4) $P \approx 1.5$ ms may be the minimum period for pulsars spun up by accretion.

Subject headings: pulsars — stars: eclipsing binaries — stars: winds

I. INTRODUCTION

Recently, Fruchter, Stinebring, and Taylor (1988) announced the discovery of a new millisecond pulsar, P1957+20, in an eclipsing binary. The observed radio eclipses apparently result from free-free absorption in a stellar wind boiled off the pulsar's low-mass companion by bombarding high-energy particles and radiation from the pulsar. A small dispersion measure fluctuation, $\delta DM \approx 0.017\ \mathrm{cm}^{-3}$ pc, accompanies emersion from the radio eclipse (Fruchter *et al.* 1988).

In this *Letter*, we explore some simple consequences of this general picture. In § II, we derive physical properties of the wind, and by extrapolating to the companion's surface estimate the effective temperature and optical luminosity that should accompany excitation of the wind. In § III, we turn the problem around and discuss very crudely the energetic requirements for generating the wind and implications for pulsar physics. We summarize our results in § IV and also comment on the remarkable concordance of spin periods of P1957+20 ($P_p = 1.61$ ms) and P1937+214 ($P_p = 1.56$ ms), the fastest known pulsars.

Orbital parameters for the binary are given by Fruchter *et al.* (1988). To make numerical estimates, we adopt $\sin i = 1$ (which is reasonable in view of the observed eclipsing) and since the mass function $f = 5.2 \times 10^{-6}\ M_\odot$ is small, assume that the companion mass M_c is much smaller than the pulsar mass M. Then

$$M_c \approx 0.021 M_{1.4}{}^{2/3}\ M_\odot\ , \tag{1}$$

where $M = 1.4 M_{1.4}\ M_\odot$. The semimajor axis of the nearly circular orbit (eccentricity $e < 0.001$) is

$$a \approx 2.48 M_{1.4}{}^{1/3}\ R_\odot\ , \tag{2}$$

where $R_\odot$ is the solar radius. We estimate the radius of the companion using the relation for nearly degenerate $n = 3/2$ polytropes (e.g., Levine *et al.* 1988; Rappaport *et al.* 1987)

$$\begin{aligned} R_c &\approx 0.013 f_c (1 + X_c)^{5/3} (M_c/M_\odot)^{-1/3}\ R_\odot \\ &\approx 0.093 \left(\frac{f_c}{2}\right)(1 + X_c)^{5/3} M_{1.4}{}^{-2/9}\ R_\odot\ , \end{aligned} \tag{3}$$

where X_c is the hydrogen mass fraction of the companion and f_c is a nondegeneracy parameter. Equation (3) is moderately accurate for cold, low-mass stars, with $f_c \sim 2$–3 for pure He composition, and $f_c < 1.4$ for pure H (Levine *et al.* 1988). Note that

$$\frac{M_c/R_c}{M/a} \approx \frac{0.41 M_{1.4}{}^{2/9}}{(f_c/2)(1 + X_c)^{5/3}} \tag{4}$$

so that the escape velocity from the companion surface is a bit smaller than for the binary. As was remarked by Fruchter *et al.* (1988), R_c is somewhat smaller than the Roche lobe radius, $R_{\mathrm{Roche}} \approx 0.28 M_{1.4}{}^{2/9}\ R_\odot$ (Kopal 1959; Paczyński 1971).

II. WIND PROPERTIES AND OPTICAL EMISSION

For simplicity, we ignore the apparent asymmetry between eclipse immersion and emersion evident in the observed excess time delays (Fruchter *et al.* 1988, Fig. 3), except for a brief comment toward the end of the next section. We define the impact parameter b to be the distance of closest approach of the pulsar radio beam relative to the companion at the onset (or end) of eclipse. Then $b \approx \pi a\, \Delta\phi$, where $\Delta\phi \approx 0.08$ is the orbital phase duration of the eclipse. If $n_e(b)$ is the wind electron density at b then

$$\delta DM = \pi f_{DM}\, n_e(b) b\ , \tag{5}$$

where $f_{DM} \lesssim 1$ is a geometrical factor ($f_{DM} = 1$ for an infinite $1/r^2$ wind). The free-free optical depth at the onset (or end) of radio eclipse

$$\tau_{\mathrm{ff}} = \frac{4\pi}{3}\left(\frac{\pi}{6}\right)^{1/2} \frac{e^6 \bar{Z} g(T, \nu, \bar{Z})}{c\nu^2 (m_e kT)^{3/2}}\, n_e{}^2(b) b f_\tau \sim 1 \tag{6a}$$

at frequency ν and temperature T for a wind with mean atomic charge $\bar{Z}$, where the Gaunt factor

$$g(\nu, T, \bar{Z}) \approx 10.6 + \frac{3^{1/2}}{\pi} \ln\left(\frac{T}{\bar{Z}\nu}\right)^{3/2} \tag{6b}$$

(cf. Allen 1973, pp. 102–103) for T in K and ν in Hz, and $f_\tau \lesssim 1$ is another fudge factor ($f_\tau = 1$ for an infinite isothermal $1/r^2$ wind). Combining equations (5) and (6) gives

$$\tau_{\mathrm{ff}} = \frac{4\pi}{3}\left(\frac{\pi}{6}\right)^{1/2} \frac{e^6 \bar{Z} g(T, \nu, \bar{Z})}{c\nu^2 (m_e kT)^{3/2}} \frac{(\delta DM)^2}{b} \frac{f_\tau}{f_{DM}{}^2}\ ; \tag{7a}$$

evaluating at $\nu = 430$ MHz implies

$$\tau_{\mathrm{ff}} \approx 8.0 \times 10^2 F\left(\frac{T}{\bar{Z}^{2/3}}\right)(\delta DM_{0.017})^2 \times M_{1.4}{}^{-1/3}\, \Delta\phi_{0.08}{}^{-1} f_\tau f_{DM}{}^{-2}\ , \tag{7b}$$

where $\delta DM = 0.017\ \delta DM_{0.017}$ cm^{-3} pc, $\Delta\phi = 0.08\ \Delta\phi_{0.08}$ and, $F(\chi) = \chi^{-3/2}(\ln\chi - 0.44)$. Solving for $T/\bar{Z}^{2/3}$ gives

$$\frac{T(\mathrm{K})}{\bar{Z}^{2/3}} \approx 2.6 \times 10^2\ \delta DM_{0.017}{}^{4/3}\tau_{\mathrm{ff}}{}^{-2/3} f_\tau{}^{2/3} f_{DM}{}^{-4/3} \times M_{1.4}{}^{-2/9}\ \Delta\phi_{0.08}{}^{-2/3} \quad (8)$$

for the temperature of the eclipsing gas at a distance $r \approx b$ from the companion.

At such low temperatures, the wind plasma might be expected to be substantially neutral. However, it is easy to show that the wind is far from ionization-recombination equilibrium. As a concrete example, consider a hydrogen plasma for which the recombination rate per proton $\Gamma_{\mathrm{H}} = \alpha_{\mathrm{H}}(t)n_e$, where $\alpha_{\mathrm{H}}(T) = 2.07 \times 10^{-11} T^{-1/2}\phi$ cm^3 s^{-1} with $\phi \approx 13T^{-0.2}$ (cf. Allen 1973, p. 97; our expression for ϕ fits the values tabulated by Allen 1973 to $\lesssim 15\%$ for $2 \leq \log T \leq 5$). At $r \approx b$ we have

$$\Gamma_{\mathrm{H}}(b) \approx 2.1 \times 10^{-6}\ \mathrm{s}^{-1}\ \delta DM_{0.017}{}^{1/15}\tau_{\mathrm{ff}}{}^{7/15} f_\tau{}^{-7/15} f_{DM}{}^{-1/15} \times M_{1.4}{}^{-8/45}\ \Delta\phi_{0.08}{}^{-8/15}\ . \quad (9)$$

By comparison, for an asymptotic wind speed ($\lambda_w \geq 1$)

$$v_0 = \lambda_w\left(\frac{2GM_c}{R_c}\right)^{1/2} \approx 8.8 \times 10^2\ \mathrm{km\ s^{-1}}\left(\frac{\lambda_w}{3}\right) M_{1.4}{}^{4/9}\left(\frac{f_c}{2}\right)^{-1/2}(1 + X_c)^{-5/6} \quad (10)$$

the "flight time" $t_f(r) = r/v_0$ is

$$t_f(b) \approx 9.4 \times 10^3\ \mathrm{s}\left(\frac{\lambda_w}{3}\right)^{-1} M_{1.4}{}^{-1/9}\left(\frac{f_c}{2}\right)^{1/2}(1 + X_c)^{5/6}\ \Delta\phi_{0.08} \quad (11)$$

at $r \approx b$, so that $\Gamma_{\mathrm{H}}(b)t_f(b) \sim 10^{-2} \ll 1$. Moreover, $\Gamma_{\mathrm{H}}(r)t_f(r) \propto rn_e(r)[T(r)]^{-0.7} \propto r^{-0.07}$ for an adiabatically expanding nonrelativistic ($\gamma = 5/3$) gas flowing radially at constant speed, so that $\Gamma_{\mathrm{H}}(r)t_f(r) \ll 1$ throughout the asymptotic regime of the wind. Of course, ionization-recombination equilibrium is likely to be satisfied in the pulsar-heated stellar atmosphere where the wind speed is still well below v_0 and n_e is correspondingly larger. Thus the ionization state of the wind is largely frozen-in after liftoff. A similar conclusion holds for He-dominated winds (Wasserman 1988).

Because the wind density is low, its total volume emissivity is correspondingly small and adiabatic expansion should be a good first approximation for the radial scaling of the gas temperature. Extrapolating to the surface using equation (8) and $T(r)r^{4/3}$ = constant implies a surface temperature

$$T_s(\mathrm{K}) \approx 3.3 \times 10^3 \bar{Z}^{2/3}\ \delta DM_{0.017}{}^{4/3}\tau_{\mathrm{ff}}{}^{-2/3} f_\tau{}^{2/3} f_{DM}{}^{-4/3} \times M_{1.4}{}^{14/27}\ \Delta\phi_{0.08}{}^{2/3}\left(\frac{f_c}{2}\right)^{-4/3}(1 + X_c)^{-20/9}\ . \quad (12)$$

At this temperature the blackbody luminosity is $L_c = 2\pi R_c{}^2\sigma T_s{}^4$ assuming only half the surface is heated by the pulsar, or

$$\frac{L_c}{L_\odot} \approx 4.6 \times 10^{-4}\bar{Z}^{8/3}\ \delta DM_{0.017}{}^{16/3}\tau_{\mathrm{ff}}{}^{-8/3} f_\tau{}^{8/3} f_{DM}{}^{-16/3} \times M_{1.4}{}^{44/27}\ \Delta\phi_{0.08}{}^{8/3}\left(\frac{f_c}{2}\right)^{-10/3}(1 + X_c)^{-50/9}\ . \quad (13)$$

The heated companion should, therefore, appear visually as a variable star with an effective temperature given by equation (12) and a period equal to the binary period.

Recently, Fruchter *et al.* (1988) have reportedly identified a variable optical counterpart to the binary with peak visual magnitude $m_v \approx 20.35$ (Fruchter 1988) and with minimum light occurring during the pulsar eclipse. Their best estimate of the color temperature is $T_{\mathrm{color}} \approx 5500$ K, which, in view of the numerous uncertain parameters in equation (12), agrees fairly well with our extrapolated value for T_s. Assuming blackbody emission at $T_{\mathrm{color}} \approx 5500$ K and visual extinction $A_v \approx 1.5$ (Fruchter 1988), the observations are consistent with emission by a surface of radius $R \approx 0.2\ (D/1\ \mathrm{kpc})\ R_\odot$ (cf. Allen 1973, p. 197), which is comparable to equation (3). However, Fruchter *et al.* (1988) report no significant variations in T_{color} during the orbit, contrary to what might be expected in our model.

III. PULSAR ILLUMINATION

The solid angle subtended by the companion star at the pulsar is $\Omega_c = R_c{}^2/4a^2$, corresponding to an angular radius

$$\theta_c = (\Omega_c/\pi)^{1/2} \approx 1.1 \times 10^{-2} M_{1.4}{}^{-5/9}\left(\frac{f_c}{2}\right)(1 + X_c)^{5/3} \quad (14)$$

or an angular diameter $\sim 1°$. Although $\theta_c \ll 1$, the companion may intercept a significant energy flux from the high-energy radiation and particles emitted by the pulsar. To estimate the energy absorbed by the companion we assume that the high-energy emission from the pulsar is beamed into a solid angle $\Omega_B = \pi\theta_B{}^2$, and that the beam axis lies close to the orbital plane, which is consistent with the observed radio eclipsing. (Because the pulsar has presumably been spun up via accretion [cf. Alpar *et al.* 1982] its spin axis is very likely parallel to the binary angular momentum, so the assumed beam [or magnetic] axis is roughly orthogonal to the pulsar spin.) If $\theta_B \gtrsim \theta_c$, the companion intercepts a fraction $(\theta_c/\theta_B)^2$ of all emitted high-energy particles and radiation for a fraction θ_B/π of each spin period. Assuming that a fraction ϵ of the total pulsar luminosity L_p escapes in the form of such high-energy emission, the average rate at which energy is absorbed by the companion is $L \approx \epsilon L_p \theta_c{}^2/\pi\theta_B$. Since $L_p = 4\pi^2 I\dot{P}_p/P_p{}^3$, where I is the pulsar moment of inertia and $\dot{P}_p$ its period derivative

$$L \approx 4\pi\epsilon \frac{I\dot{P}_p}{P_p{}^3}\frac{\theta_c{}^2}{\theta_B} \approx 9.3 \times 10^{-3}\ L_\odot \frac{\epsilon}{\theta_B}\ \dot{P}_{p,19} I_{45} M_{1.4}{}^{-10/9}\left(\frac{f_c}{2}\right)^2 \times (1 + X_c)^{10/3}\ , \quad (15)$$

where $I \equiv 10^{45} I_{45}$ g cm^2 and $\dot{P}_{p,19}$ is the spindown rate normalized to 10^{-19} s s^{-1}, roughly the value for P1937+214 (Ashworth, Lyne, and Smith 1983; Backer, Kulkarni, and Taylor 1983; Cordes and Stinebring 1984).

In steady state, the absorbed energy is presumably converted in part to the kinetic and thermal energy of the wind. If the ionization fraction in the wind is χ then it carries away energy at a rate

$$L_w \approx \frac{2\,\delta DM m_p v_0^3 b}{\chi f_{DM}(1+X_c)}$$
$$\approx 2.6 \times 10^{-6}\, L_\odot\, \chi^{-1}\left(\frac{\lambda_w}{3}\right)^3 \delta DM_{0.017}\, f_{DM}^{-1}$$
$$\times M_{1.4}\, \Delta\phi_{0.08}\left(\frac{f_c}{2}\right)^{-3/2}(1+X_c)^{-3/2}\,, \qquad (16)$$

assuming that the wind is emitted into solid angle 2π and neglecting the relatively small thermal energy in the wind. Comparing equations (15) and (16) we see that $L_w \ll L$ for a fully ionized wind; $L_w \sim L$ only for $\chi \sim 10^{-3}$. Low fractional ionization cannot be ruled out *a priori*, especially if X_c is small. The asymmetry in excess time delays measured at onset and end of radio eclipse (Fruchter *et al.* 1988) may, however, furnish indirect evidence that χ is not small. Radiation pressure due to absorbed radio waves can significantly deflect the eclipsing gas if

$$L_{\rm RADIO} \approx \frac{\Omega_B v_0^2 ac}{\kappa_{\rm RADIO}}\,, \qquad (17)$$

where $L_{\rm RADIO}$ is the radio luminosity of the pulsar and $\kappa_{\rm RADIO}$ is a suitably averaged radio frequency opacity (see also Rasio, Shapiro, and Teukolsky 1988). Since $\kappa_{\rm RADIO}\,\rho(b) b \approx 1$ at the onset (or end) of radio eclipse, equation (17) may be rewritten as

$$L_{\rm RADIO} \approx \frac{2\Omega_B\, \delta DM m_p v_0^2 ac}{\pi \chi f_{DM}(1+X_c)}$$
$$\approx 2.1 \times 10^{29}\ {\rm ergs\ s^{-1}}\ \chi^{-1}\left(\frac{\theta_B}{10^\circ}\right)^2\left(\frac{\lambda_w}{3}\right)^2$$
$$\times M_{1.4}^{11/9}\left(\frac{f_c}{2}\right)^{-1}(1+X_c)^{8/3}\,. \qquad (18)$$

By comparison, we estimate $\log L_{\rm RADIO} \approx 29.7$ for P1937+214, which is much smaller than equation (18) if $\chi \ll 1$ and is very likely of the same order if $\chi \sim 1$.

Comparing L, equation (15) with the blackbody luminosity L_c, equation (13), inferred in § II, we see that

$$\frac{L}{L_c} \approx 20\, \frac{\epsilon}{\theta_b}\, \dot P_{p,19} I_{45} M_{1.4}^{-74/27}\left(\frac{f_c}{2}\right)^{16/3}$$
$$\times (1+X_c)^{80/9}\, \delta DM_{0.017}^{-16/3} \tau_{\rm ff}^{8/3}$$
$$\times f_\tau^{-8/3} f_{DM}^{16/3}\, \Delta\phi_{0.08}^{-8/3}\, \bar Z^{-8/3}\,. \qquad (19)$$

Clearly, this ratio depends quite sensitively on properties of the companion star through f_c and X_c: for solar composition, $X_c \approx 0.7$ and $f_c < 1.4$ so $(f_c/2)^{16/3}(1+X_c)^{80/9} \lesssim 17$, whereas for pure He ($X_c \equiv 0$), $f_c \lesssim 3.3$ so $(f_c/2)^{16/3}(1+X_c)^{80/9} \lesssim 14$ (Levine *et al.* 1988). Nevertheless, it appears that pulsar heating of the companion can only explain the visual emission from the binary if $\epsilon \dot P_{p,19}\, \theta_B^{-1} \sim 10^{-2} - 1$, that is, if a significant fraction of the pulsar luminosity is converted to high-energy particles and radiation before reaching the companion.

IV. SUMMARY AND CONCLUSIONS

Potentially, the eclipsing pulsar P1957+20 can offer unique insights into pulsar physics once the gas dynamics of the binary is understood. The stellar wind responsible for the radio eclipses is undoubtedly excited by intercepted high-energy pulsar emission. In this *Letter*, we have developed one scenario for the properties of the wind. According to this picture, the eclipsing gas is cool but possibly highly ionized, since recombination in the wind is too slow to enforce ionization-recombination equilibrium. Because the wind density and temperature are low, free-free emission from the expanding gas is insignificant (total luminosity $\lesssim 10^{19}$ ergs s^{-1}). Assuming adiabatic expansion, the temperature at the pulsar-heated stellar surface $T_c \sim 10^3$–10^4 K, in reasonable agreement with observations by Fruchter *et al.* (1988), which imply a color temperature $T_{\rm color} \approx 5500$ K.

Intercepted high-energy pulsar emission can only account for this optical brightness if a substantial fraction of the total pulsar luminosity $L \approx 9 \times 10^{35}$ ergs s^{-1} $I_{45} \dot P_{p,19}$ reaches the companion in the form of high-energy particles and radiation. Since the integrated column density of the wind is rather low ($\sim 10^{-7}\chi^{-1}$ g cm^{-2}) such high-energy emission can penetrate to the surface with little attenuation. The absorbed radio emission can, on the other hand, significantly deflect the wind if $\log L_{\rm RADIO} \approx 29$–30 and $\chi \sim 1$.

Finally, we note that the pulse periods of P1937+214 and P1957+20 are remarkably similar: $P_p = 1.56$ ms and 1.61 ms, respectively. There is little doubt that P1957+20 was spun up in an earlier "X-ray binary" phase of its evolution as has been proposed for P1937+214 (Alpar *et al.* 1982). For fixed luminosity accretion, it might be expected that the neutron star approaches "critical fastness" as time proceeds and the neutron star magnetic field decays, with $P_p/P_k(r_A) \to \omega_s^{-1} > 1$, where $P_k(r_A)$ is the Kepler period at the Alfvén radius (Ghosh and Lamb 1978, 1979*a*, *b*). One might then expect P_p to stabilize once r_A approaches the neutron star radius, which happens when the surface field $B_s \sim 10^8$ G for an X-ray luminosity $L_x \sim 10^{37}$ ergs s^{-1} (e.g., Shapiro and Teukolsky 1983). When accretion shuts off, the neutron star would then be left with a spin period $P_p \approx 2\pi\omega_s^{-1} R^{3/2}/GM^{1/2} \approx 1.5(3\omega_s)^{-1} R_6^{3/2} M_{1.4}^{-1/2}$ ms if $R = 10 R_6$ km. Conceivably, then, both P1957+20 and P1937+214 are near the minimum possible spin period for pulsars spun up by accretion. That the required magnetic field strength is $B \sim 3 \times 10^8$ G for $\dot P_p \sim 10^{-19}$ s s^{-1} lends further credence to this hypothesis.

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