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Are Pulsars Externally Triggered?

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Abstract. We show that inverse Compton losses of ultrarelativistic electrons travelling upstream through the outflowing fields of a radio pulsar prevent the entry of *all* cosmic ray electrons into most pulsars' magnetospheres. We also argue that gamma ray triggering is probably not relevant, and conclude that the external triggering of electron-positron cascades proposed by Ruderman & Sutherland and by Radhakrishnan is not likely to be a major contributor to the plasma dynamics of pulsars.

Key words: pulsars, magnetospheres

1. Introduction

Recently, Radhakrishnan (1982, hereafter called R) has proposed that the paucity of pulsars at short rotation period P and small rotation period derivative $\dot{P}$ can be explained by postulating an emission threshold dependent on the influence of an external agent on pulsars' magnetospheric structure, analogous to the γ-ray threshold of Shukre & Radhakrishnan (1982). These authors showed that if pulsar emission is due to time-dependent electron-positron cascades above positively charged polar caps, and these cascades are initiated by external gamma rays from the galactic background (as in the model of Ruderman & Sutherland 1975, hereafter called RS) then there exists a narrow range in $\dot{P}$ over which the cascades can occur, much narrower than the observed range of $\dot{P}$ in radio-emitting pulsars. We show that if the external agent postulated by R is cosmic ray electrons, external triggering cannot occur for the vast majority of pulsars, and argue that external triggering is unlikely to be relevant for any choice of triggering mechanism.

In this paper we recalculate the cross section for scattering of a charged particle by a strong spherical vacuum electromagnetic wave (Arons, Kulsrud & Ostriker 1975; Radhakrishnan 1981) that would be produced by a rotating magnetic dipole. We later argue that the result is not changed by the effects of the outflowing MHD wind expected in pulsars. In either case a minimum particle energy, $\gamma_{\min} m_0 c^2$, exists below

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which a charged particle is unable to penetrate interior to the light cylinder, where its subsequent motion might be precipitation onto the stellar surface. On the other hand, a maximum particle energy exists, $\gamma_{\max} m_0 c^2$, above which an electron suffers radiative energy losses in a time short compared to the time to reach the pulsar. Both $\gamma_{\min}$ and $\gamma_{\max}$ depend on P and the magnetic dipole moment, μ, and we show that for most pulsars $\gamma_{\min}$ exceeds $\gamma_{\max}$. Then no electrons can reach the magnetosphere and trigger a cascade.

2. Scattering and radiation losses of charged particles in the wave zone of a pulsar

The motion of charged particles through an inhomogeneous electromagnetic wave is well known in the non-relativistic case (*e.g.* Schmidt 1979) and has been studied under the relevant circumstances for the ultrarelativistic case by Gunn & Ostriker (1971), Arons (1972) and Blandford (1972). The essential result is that in the absence of radiative losses, a very high energy particle undergoes small amplitude oscillations about an oscillation centre, whose momentum follows an orbit described by

$$\frac{d\mathbf{p}_{oc}}{dt} = -\frac{q^2}{m\omega^2}\nabla\frac{\langle E^2\rangle}{4\pi}. \tag{1}$$

Here q is the particle's charge, ω is the circular frequency of the wave, $\langle E^2/4\pi\rangle$ is the average of the electromagnetic energy density over one cycle of wave phase, and m is the relativistic mass averaged over one cycle of phase, as measured in the frame moving with the oscillation centre. For example, if the wave is circularly polarized, $m = m_0 (1+\nu_0^2/2)^{1/2}$ where m_0 is the rest mass and $\nu_0 = qE_0/m_0 c\omega$ = ratio of formal gyration frequencies to wave frequency, E_0 = peak amplitude (Arons 1972; Brown & Kibble 1964). For other polarizations, $m = m_0 \nu_0 \times$ constants of order unity, when $\nu_0 \gg 1$ (Blandford 1972). The total momentum is $\mathbf{p} = \mathbf{p}_{oc} + \delta\mathbf{p}$, where $\delta\mathbf{p}$ = oscillatory momentum imposed by the wave, with $|\delta\mathbf{p}| \sim \nu_0 m_0 c$. If $\nu_0 \gg 1$, the oscillations are themselves relativistic, although $|\mathbf{p}_{oc}| \gg m_0 c$ is then also required.

The orbit of the oscillation centre is found by integrating Equation (1) together with $d\mathbf{r}_{oc}/dt = \mathbf{p}_{oc}/m_0\gamma_{oc}$, $\gamma_{oc} = [1+(p_{oc}/m_0 c)^2]^{1/2}$. Assume $\nu_0(\mathbf{r}) \gg 1$ throughout the wave zone $r > r_L = cP/2\pi = 48000\,P$ km, P = rotation period. Then the force in Equation (1) is derivable from the potential $m_0 c^2 \nu_0(r)/4\pi$. Standard methods (*e.g.* Goldstein 1980) yield the cross-section for scattering of a charged particle by a pulsar (Arons, Kulsrud & Ostriker 1975, Equation 7)

$$\sigma \simeq \pi r_L^2 \left(\frac{\nu_L}{\gamma_\infty}\right)^2 \tag{2}$$

where $\gamma_\infty mc^2$ is the energy of the particle at infinity and $\nu_L \equiv \nu_0(r_L) = (2\pi)^2 e\mu/(m_0 c^4 P^2)$. Here μ = magnetic moment. This relation assumes pulsar spindown is due to vacuum magnetic dipole radiation, as discussed by Ostriker & Gunn (1969). (2) is easily estimated by using energy conservation $\gamma + \nu_0/4\pi$ = constant, plus $\nu_0(\mathbf{r}) \to 0$ as $r \to \infty$. Note, that (2) is an elastic scattering cross-section which does not include the effects of radiation reaction on the particle.

A charged particle is scattered away from the pulsar when $\sigma(\gamma_\infty) \gg \pi r_L^2$. On the other hand, a particle whose energy is so high that $\sigma < \pi r_L^2$ enters the dipolar region interior to r_L without scattering, where it may be captured, just as high energy cosmic rays directly enter the terrestial magnetosphere. Setting $\sigma = \pi r_L^2$ yields the minimum energy at infinity of a particle which can be directly captured

$$\gamma_{\min} = |\nu_L| = 2.6 \times 10^7 \mu_{30} P^{-2} \tag{3}$$

with ν_L evaluated for electrons. Note that for $r \gg r_1$ and $\gamma_\infty \geqslant \gamma_{\min}$, $\gamma(r) \gg \nu_0(r)$. Here $\mu_{30} = \mu/10^{30}$ cgs $\cong (P\dot{P}_{15})^{1/2}$ and $\dot{P}_{15} = \dot{P}/10^{-15}\,\mathrm{s\,s^{-1}}$.

Interstellar electrons with energy $\gamma > \gamma_{\min}$ may be plentiful enough to reach $r = r_L$ at the RS spark rate, if their propagation all the way from infinity to the light cylinder were free of radiation losses (Radhakrishnan 1981). However, a relativistic electron moving through an electromagnetic wave loses the translational energy of motion of its oscillation centre through inverse Compton losses (*e.g.* Blumenthal & Gould 1970; Rybicki & Lightman 1979). In the present circumstances, this scattering takes a special form called Nonlinear Inverse Compton (NIC) radiation (Rees 1971; Arons 1972; Blandford 1972; see also Melrose 1980). The only significant aspect which concerns us here is the total power emitted by an electron with oscillation centre Lorentz factor $\gamma_{oc}(r)$. This is the same as the synchrotron power emitted by an electron of pitch angle θ gyrating in a static B field of strength $B = m_e c\omega\nu_0/e = E_0$, if one replaces the usual $\sin\theta \gg \gamma_{oc}^{-1}$ by $1 - \cos\theta$, where now θ is the angle between the direction of motion of the oscillation centre and the direction of propagation of the strong wave. In the present application, we are concerned with small angular momentum particles heading toward the pulsar with impact parameter $b \lesssim r_L$. Then $\cos\theta \simeq -1$ on the whole inbound orbit and the power radiated is

$$P_{rad} = 4c\sigma_T \frac{E_0^2}{4\pi}\gamma_{oc}^2 \tag{4}$$

while the radiative lifetime of the electron at any $r > r_L$ is

$$t_{rad} = \frac{\gamma_{oc} m_e c^2}{P_{rad}(\gamma_{oc})} = 6 \times 10^6 \frac{P^6}{\mu_{30}^2 \gamma}(r/r_L)^2\ \mathrm{s}. \tag{5}$$

Here σ_T is the Thomson cross section and m_e is the electron rest mass. If an electron of energy $\gamma_\infty > \gamma_{\min}$ is able to reach the magnetosphere interior to r_L, it must not have lost all its energy in NIC radiation at larger radii; if its energy *is* substantially reduced by radiation, the ponderomotive pressure represented by (1) expels the particle before it reaches r_L, even if γ_∞ exceeds $\gamma_{\min}$. By requiring the energy loss time, t_{rad}, to be much greater than the transit time,

$$t_{transit} = r_L c^{-1} = 0.16\,P\ \mathrm{s}$$

we find an upper limit on γ_∞, above which radiative losses will reduce the energy of a cosmic ray until it is less than

$$\gamma_{\max} = 3.9 \times 10^7 P^5 B_{12}^{-2}. \tag{6}$$

Here B is the magnetic field at the stellar surface ($r = 10^6$ cm), $B_{12} = B/10^{12}$ Gauss. We require that $\gamma_{\max} > \gamma_{\min}$ for cosmic ray penetration, implying

$$B_{12} < 1.4\,P^{7/3}. \tag{7}$$

High energy electrons rapidly lose energy radiatively, while lower energy electrons are excluded by the ponderomotive pressure. Therefore, no electrons can survive and reach the magnetosphere if criterion (7) is violated. We have plotted this threshold for cosmic ray penetration in Fig. 1. Clearly, if rapid time dependent discharges at the positively charged polar cap are relevant to pulsar activity, with each discharge triggered by cosmic ray electrons, then only very long period pulsars can function in this manner. Since there is no strong observational distinction between pulsars in the allowed and forbidden region of Fig. 1, we conclude that this mechanism of external triggering is irrelevant.

We have thus far treated the external field of the pulsar as a vacuum strong wave. The externally triggered model of Ruderman & Sutherland, however, implies that the $\sim 10^5$ sparks/per second at the surface amplify and yield an outflowing pair plasma with $\sim 10^{33}$ electrons and positrons per second leaving the star. This is dense plasma in the electrodynamic sense; the outflow from the star is then a relativistic MHD wind, not a vacuum wave (*e.g.* Arons 1981a). However, this makes little difference to our evaluation of the significance of radiation losses. If the wind has the same spirally wound topology of the vacuum magnetic field as would be expected in the starvation scenario outlined by Arons (1981b, 1983a), with alternating segments of oppositely directed toroidal magnetic field spaced with wavelength $\sim r_L$, the motion of an electron with $\gamma_\infty \gtrsim \gamma_{min}$ is still basically a straight line motion virtually the same as in the vacuum wave case, plus a slow deflection of the oscillation centre with B_{wind} replacing E_{wave} in Equation (1), since the Larmor radius in the wind zone magnetic field vastly exceeds the wavelength when $r \gg r_L$. For $r \sim r_L$, even an electron of energy $\gamma \sim \gamma_{min}$ has Larmor radius $< r_L$, and is then easily picked up and swept outwards by the magnetic field lines sweeping outward from the star. The inverse Compton losses here are the same as in the NIC process for

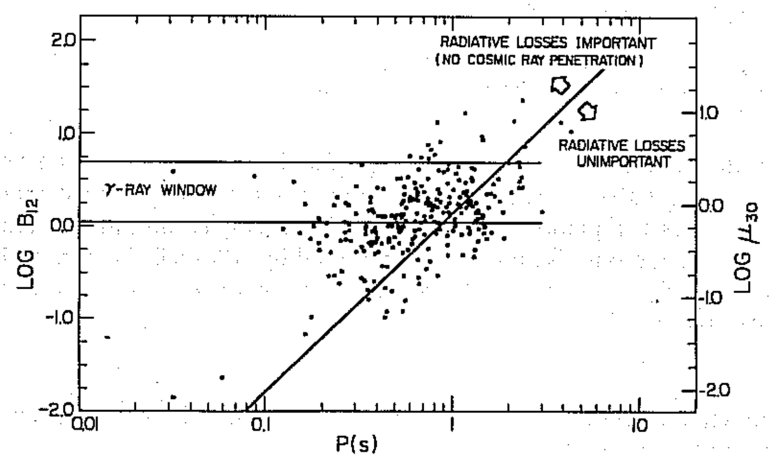

Figure 1. The μ_{30}–P diagram. The cosmic ray penetration threshold (Equation 7) and the γ-ray window of Shukre & Radhakrishnan (1982) are shown. The radiative loss line has no free parameters when plotted as a function of μ_{30} and P. $R = 10^6$ was used to fix the left hand scale. The points themselves have random uncertainty in the vertical direction (due to uncertainty in $I/\sin^2 i$) and possible systematic uncertainty if plasma effects systematically alter the torque from that due to magnetic dipole radiation.

the vacuum wave, since the orbits are the same, to lowest order in the inverse square of the wind Lorentz factor. In a wind scenario possibly appropriate to the approximately aligned rotator (Kennel, Fujimura & Pellat 1979), a particle propagates inward along a slightly curved path in a nonsegmented magnetic field until its Larmor radius is $\sim r$, where again it is turned around and expelled by the momentum flux of the wind. $\gamma_\infty = \gamma_{min}$ corresponds to an electron having a Larmor radius $\simeq r_L$ at r_L, and the synchrotron losses from inbound particle in the B field of strength $\propto 1/r$ are the same as for the NIC losses used in our calculations above (except for a factor of 4; since the radially propagating particle is directed across the toroidally wrapped B, $\sin^2\theta \simeq 1$ in the synchrotron case, while $(1-\cos\theta)^2 \simeq 4$ in the NIC case). Then the MHD winds are just as important in excluding cosmic ray electrons from the magnetosphere as NIC losses are in an outflow of vacuum waves.

The pulsars within our small allowed region might be influenced by external triggering in the manner suggested by Radhakrishnan (1981, 1982). We think this is unlikely, since even if a relativistic electron survives to reach $r = r_L$, the strong synchrotron losses present as it tries to spiral down the polar field lines cause it to decay to *non-relativistic* energy while it is still at high altitude ($r_L \gg r \gg$ stellar radius) within the flow of pair plasma. If these few electrons do have any interaction with the particles and fields of the outflow, the overwhelmingly larger momentum flux of the outflow easily sweeps the trickle of cosmic rays back out again (the electric field of the RS model which could drag the trigger electrons down exists only within the polar cap zone at heights $\leqslant 1$ km above the surface, far below the outer magnetosphere at $r \sim r_L$ where the cosmic ray electrons lose all their energy).

γ-ray triggering is not excluded by these considerations, but we think it is also unlikely for the following reasons:

(1) If the binding energy of ions to the stellar crust is large, theory shows that discharges may be maintained without external triggering (Cheng & Ruderman 1980; Jones 1981).

(2) If the binding energy of ions is small or if the rotation and magnetic axes are inclined by an angle less than 90°, such that electrons are accelerated, then the direct emission of γ-rays via curvature radiation from particles extracted from the surface will dominate the formation of electron-positron pairs (Arons & Scharlemann 1979; Arons 1979, 1981c, 1983b; Ruderman 1971).

(3) Observationally there exists long-term stable modulation features observed in the arrival times of the subpulses of some pulsars (*cf.* Manchester & Taylor 1977; Unwin *et al.* 1978). These have modulation timescales greater than seconds and stability timescales up to hours, far in excess of the memory and linkage possible between $\sim 10^5$ sparks per second initiated by γ-rays from a pool of uncorrelated photons such as those from the interstellar medium or a nearby SNR. Since pulsars with marching subpulses exhibit no special behavior in other ways, we expect the pair production mechanism of these objects to be the same as in other pulsars. Thus we conclude that pulsars are unlikely to be triggered by γ-rays.

We should also point out that a straightforward calculation of the radiative loss of cosmic rays with a cosmic flux given by Radhakrishnan 1981 yields a γ-ray flux at the neutron star surface that is approximately 12 orders of magnitude smaller than the background γ-ray flux, for typical pulsar parameters, and so is certainly irrelevant as a triggering mechanism.

Other conceivable triggering agents such as neutrinos, gravitinos, *etc.*, have crosssections and background particle fluxes *many* orders of magnitude smaller than would be required to interact in any electrodynamically interesting way with the magnetic field or stellar crust of the neutron star surface.

3. Conclusions

The main conclusion of this paper is that external cosmic rays cannot penetrate into most pulsar magnetospheres and cannot trigger electron-positron cascades for most pulsars. We have also argued that other external agents are not likely to act as triggering mechanisms, and so are unlikely to account for the maintenance of electron-positron flow and associated coherent radio radiation from pulsars.

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