ARE FAST RADIO BURSTS THE BIRTHMARK OF MAGNETARS?

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ABSTRACT

A model of fast radio bursts, which enlists young, short period extragalactic magnetars satisfying \( B/P > 2 \times 10^{16} \text{ G s}^{-1} \) as the source, is proposed. When the parallel component \( E_\parallel \) of the surface electric field (under the scenario of a vacuum magnetosphere) of such pulsars approaches 5% of the critical field \( E_c = m_e c^3/(\epsilon_0 h) \), in strength, the field can readily decay via the Schwinger mechanism into electron–positron pairs, the back reaction of which causes \( E_\parallel \) to oscillate on a characteristic timescale smaller than the development of a spark gap. Thus, under this scenario, the open field line region of the pulsar magnetosphere is controlled by Schwinger pairs, and their large creation and acceleration rates enable the escaping pairs to coherently emit radio waves directly from the polar cap. The majority of the energy is emitted at frequencies \( \lesssim 1 \text{ GHz} \) where the coherent radiation has the highest yield, at a rate large enough to cause the magnetar to lose spin significantly over a timescale \( \sim \times 10^{-3} \text{ s} \), the duration of a fast radio burst. Owing to the circumstellar environment of a young magnetar, however, the \( \lesssim 1 \text{ GHz} \) radiation is likely to be absorbed or reflected by the overlying matter. It is shown that the brightness of the remaining (observable) frequencies of \( \approx 1 \text{ GHz} \) and above are on a par with a typical fast radio burst. Unless some spin-up mechanism is available to recover the original high rotation rate that triggered the Schwinger mechanism, the fast radio burst will not be repeated again in the same magnetar.

Key words: radiation mechanisms: non-thermal – Radio continuum: general – relativistic processes – stars: magnetars

1. INTRODUCTION

Fast radio bursts (FRBs) are a recently discovered and very interesting phenomenon. The initial discovery of one such event (Lorimer et al. 2007) was followed a few years later by four other similar detections (Thornton et al. 2013, see also the catalog of FRBs to date in Petroff et al. 2016). Their essential properties are a spatial extent consistent with point sources, random transients of brightness \( \lesssim 1 \text{ Jy} \) lasting a timescale \( \sim 1 \text{ ms} \), typically without repetition (but see Spitler et al. 2016), and a large line of sight dispersion measure that suggests an extragalactic origin. The difficulty in modeling FRBs is the conflict between time and distance scales. Timescales of order 1 ms are symptomatic of compact objects like pulsars and stellar mass black holes, and these sources are usually too faint to be observable at extragalactic distances.

The discovery of FRBs led us to revisit the subject of compact sources, specifically pulsars, to see if there could be a physical mechanism responsible for their manifestation as extragalactic transients. Now there exists in the literature an estimate of the brightness of a pulsar (in the radio and other wavelengths). In the case of an aligned rotator—the simplified model adopted here—it is

\[
\frac{dE}{dt} = \frac{16\pi^3 B^2 R^6}{P^4 c^3} = 9.24 \times 10^{36}\left(\frac{B}{4 \times 10^{12} \text{ G}}\right)^2 \times \left(\frac{P}{0.1 \text{ s}}\right)^{-4}\left(\frac{R}{10 \text{ km}}\right)^6 \text{ erg s}^{-1}. \tag{1}
\]

Although Equation (1) agrees with Goldreich & Julian (1969) in their idealized baryonic magnetosphere model, the more realistic model of Ruderman & Sutherland (1975) invokes a spark gap to generate electron–positron pairs, which yields the more conservative result for the radio luminosity of

\[
\frac{dE}{dt} = 5.2 \times 10^{32}\left(\frac{B}{10^{12} \text{ G}}\right)^{6/7}\left(\frac{\rho_c}{10^7 \text{ cm}}\right)^{4/7} \times \left(\frac{P}{0.1 \text{ s}}\right)^{-15/7} \text{ erg s}^{-1} \tag{2}
\]

where \( \rho_c \) is the curvature radius of the surface magnetic field. In Equation (2), the two parameters that give \( dE/dt \) are the spin period \( P \) and magnetic field \( B \), but there are bounds as well. In particular, to get a large \( dE/dt \) one needs rapidly rotating and highly magnetized neutron stars. Although this is an inviting regime because of the discovery of rotationally driven radio magnetars of the spark gap type (Rea et al. 2012), the lower limit on \( P \) that stems from the need to balance gravity against centrifugal force is \( P \gtrsim 0.5 \text{ ms} \), while magnetic flux conservation during the collapse to a neutron star cannot result in \( B \) values far in excess of \( 10^{15} \text{ G} \), the high field end of the magnetar range.

In addition to these constraints, there also exists a correlation between \( P \) and \( B \), viz. magnetars tend to have large \( P \), of order 1 s or more, e.g., Beloborodov & Thompson (2007); while millisecond pulsars tend to have \( B \lesssim 10^9 \text{ G} \), e.g., Lamb & Yu (2005). As a consequence of all the above, radio pulsars with \( dE/dt \) well in excess of the constraints of Equation (2), i.e., of brightness approaching the Equation (1) limit that may enable their detection from extragalactic distances, will probably have to be powered by a very different physical mechanism (operating either temporarily or steadily) from the spark gap model of Ruderman & Sutherland (1975).

The purpose of this paper is indeed to propose a new mechanism of deceleration and rotational energy loss of a pulsar. If the magnetic field is strong enough to be in the magnetar range, and \( P \rightarrow 0.1 \text{ s} \), the parallel electric field \( E_\parallel \)
imposed by boundary conditions upon an idealized vacuum magnetosphere can partially discharge into positron–electron pairs, even though its magnitude is below the Schwinger critical field of

\[ E_c = \frac{m_e^2 c^3}{\epsilon h} \approx 4.4 \times 10^{13} \text{G}. \]  

(3)

In this regime, one must take account this extra quantum loss effect that can proceed at a higher rate than the classical limit of Equation (1). If an enhancement of \( E \) to such a regime could occur momentarily and briefly, the result would be a sudden and equally short-lived switch of the pulsar emission scheme, from the low mode of Equation (1) to one dominated by the quantum electrodynamics process of vacuum breakdown. Of course, such short period radio magnetars are probably rare; and it is possible their existence can only be inferred from the much larger extragalactic population of pulsars, via the FRBs.

2. SCHWINGER CRITICAL ELECTRIC FIELD

According to the original calculation of Schwinger (1951), a strong electric field is capable of drawing real \( e^+e^- \) pairs out of vacuum. In a pulsar environment where a magnetic field is also present, \( E \) has vacuum components parallel and perpendicular to \( B \). Since the magnitude \( E \) is always below \( B \), it is possible to transform to a frame in which \( \vec{E}[B] \) by boosting at a non-relativistic speed (the drift velocity), i.e., the results obtained in this frame are not very different from the discharge of \( E_0 \) in the laboratory frame. Thus, w.r.t. the \( \vec{E}[B] \) frame, it was shown by Ruffini et al. (2009) that

\[ \frac{dn}{dt} = \frac{\alpha EB}{4\pi^2\hbar} \coth \left( \frac{\pi B}{E} \right) e^{-\epsilon E/E} \text{cm}^{-3} \text{s}^{-1}, \]

(4)

where \( \alpha = e^2/(\hbar c) = 1/137 \) is the fine structure constant. Note that in the limit \( E \ll B, \) which applies throughout this paper, one may ignore the \( \coth(\pi B/E) \) factor.

Although the electric fields of interest obey \( E \ll B \) and \( E \ll E_c, \) they are of sufficient strength to accelerate the pairs produced by the discharge, which are initially at rest, in opposite directions. The field is neutralized as a result. In detail, the mechanism induces damped oscillations in the field about zero, and the equation of motion of the field may be written by means of Maxwell’s third and fourth equations as

\[ \vec{E} = -4\pi \vec{J} = -8\pi c e\hbar \frac{2\alpha}{\pi} \frac{e c E B}{\hbar} e^{-\epsilon E/E} \text{G s}^{-2}. \]  

(5)

Note that, as emphasized by Ruffini et al. (2009), Equation (5) describes the back reaction of pairs generated by the decay of an electric field on the field itself. If the pair creation and acceleration rates are high enough to consume the electric field within a cycle, then the oscillation becomes highly damped, and Equation (5) could at best describe only that fraction of the oscillation cycle up to the moment of full dissipation of the field.

\[ 1 \] We can ignore the non-inertial effects of this frame, because the pulsar’s rotation period is much larger than the timescale of pair creation and annihilation.

3. PAIR PRODUCTION BY DIRECT ELECTRIC FIELD DECAY

To obtain the basic parameters of the electric field oscillations, first observe that it is possible to solve Equation (5) analytically to the following point:

\[ t - t_0 = -\frac{\pi}{2} \left( \frac{\hbar}{\alpha e B c} \right)^{1/2} \int_{y_0}^{y} \frac{dy}{y_0 e^{-\gamma y} - y^3 e^{-\gamma y}/y^{1/2}}, \]

(6)

where

\[ y = \frac{E}{E_c} \]

(7)

is the electric field \( E(t) \) normalized to the critical field. The oscillation period may be evaluated by integrating Equation (6) down to the first null point \( y = 0, \) viz.

\[ P_{osc} = \frac{2\pi}{\omega_{osc}} = \frac{2\pi}{\alpha e B c} \int_{y_0}^{\gamma} \frac{dy}{y_0 e^{-\gamma y} - y^3 e^{-\gamma y}/y^{1/2}}, \]

(8)

provided the number of pairs produced per cycle

\[ n = \frac{\alpha E_0 B}{2\pi\hbar} \left( \frac{\hbar}{\alpha e B c} \right)^{1/2} \int_{0}^{y_0} ye^{-\gamma y}dy \]

(9)

is small in the sense indicated above (Equation (5)). Assuming the relativistic equation of motion

\[ \dot{\gamma} = \frac{e E_{c}}{mc}, \]

(10)

for the pairs, it readily follows that they are accelerated to the Lorentz factor

\[ \gamma = \frac{e E_{c}}{mc} \int_{0}^{y_0} \frac{ydt}{y_0 e^{-\gamma y} - y^3 e^{-\gamma y}/y^{1/2}}, \]

(11)

within a quarter cycle of the electric field oscillation. Equations (8)–(11) were represented graphically in Ruffini et al. (2009).

The acceleration of electrons and positrons depicted in Equations (8)–(11) is consistent with that of a large amplitude oscillating electric field. As a simplified and approximate model, consider an electric field that oscillates as \( E(t) = E_0 \cos \omega_0 t \). The equation of motion of an electron, \( dp/dt = eE \), becomes \( \gamma \dot{\gamma} = eE_0 \sin \omega_0 t/(\mu_0 m_0) \). In the ultra-relativistic limit of \( v \approx c \), this means the pair Lorentz factor \( \gamma \) oscillates \( 90^\circ \) out of phase with \( E \), and with the amplitude

\[ \Gamma = \frac{e E_0}{m_\omega_0 c}. \]

(12)

Indeed, if ones sets \( E_0 = \gamma_0 E_c \) and \( \omega_0 = 2\pi/P_{osc} \), one would find that the resulting \( \Gamma \) is in par with Equation (11).

Now let us get a feel of the oscillation frequency. As a working example, we set the value of the neutron star magnetic field to \( B = 10^{12} \text{G} \) (although our conclusion would remain the same if \( B \) assumes any value within the magnetar range of \( B > E_c \) and the parallel electric field to an initial value \( E_0 = 0.045E_c \), or \( \gamma_0 = 0.045 \) at the polar field line \( \theta = 0 \) corresponding to a rotation period \( P \approx 0.106 \text{s} \) via the aligned
rotator of Goldreich & Julian (1969):

\[
E_{||} = 1.976 \times 10^{12} \left( \frac{P}{0.106 \text{ s}} \right)^{-1} \left( \frac{B}{10^{15} \text{ G}} \right) \times \left( \frac{R}{r} \right)^4 \cos^2 \theta \text{ G.} \tag{13}
\]

A numerical integration of Equation (8) then yields

\[
P_{\text{osc}} = 3.38 \times 10^{-5} \text{ s, at } y_0 = 0.045, \tag{14}
\]

while from Equations (9) and (11) the pair number density and Lorentz factor develop to

\[
n = 1.57 \times 10^{14} \text{ cm}^{-3} \text{ and } \gamma = 5.35 \times 10^{14} \text{ at } y_0 = 0.045, \tag{15}
\]

respectively, within a quarter oscillation cycle of the field. Such a number density is on par with the Goldreich–Julian estimate

\[
n_{\text{proton}} = 7 \times 10^{14} \left( \frac{B}{10^{15} \text{ G}} \right) \left( \frac{P}{0.1 \text{ s}} \right) \text{ cm}^{-3}. \tag{16}
\]

for the protons and electrons pulled out of the stellar surface at comparable speeds (i.e., both speeds being \( \approx c \)), although the subsequent work of Ruderman & Sutherland (1975) presented the difficulties of producing an ionized baryonic magnetosphere and argued for its substitution by pairs, with the number density of primary pair particles at the polar cap surface, i.e., within the “spark gap” region, also commensurate with Equation (16).

When comparing the Ruderman & Sutherland (1975) pairs with the Schwinger pairs proposed here, note that in the former the maximum Lorentz factor is insensitive to \( B \) and \( P \), viz.

\[
\gamma_{\text{spark}} \lesssim 3 \times 10^{6} \left( \frac{B}{10^{12} \text{ G}} \right)^{1/7} \left( \frac{P}{10^{6} \text{ cm}} \right)^{4/7}, \tag{17}
\]

while in the latter the same is \( \propto B/P \) according to Equations (12) and (13), and can become \( \gg 10^{6} \). The difference is because the electric field in the spark gap of the former, which has thickness

\[
h_{\text{spark}} \approx 5 \times 10^{3} \text{ cm,} \tag{18}
\]

develops and discharges on timescales of 1–10 \( \mu \text{s} \) and never reaches the full length and strength of \( E_0 \) as given by Equation (13) in (even) the region immediately outside the star’s surface. The Schwinger pairs, on the other hand, are born directly from the decay of \( E_0 \) and, as a result, the numbers given after Equation (13) indicate they can reach and surpass the number density of Goldreich & Julian (1969) and Ruderman & Sutherland (1975) (viz. Equation (16)) on equally short timescales or shorter. As will be discussed below, this difference between our model and Ruderman & Sutherland (1975) is responsible for not only the magnitude of \( \gamma \), but also the radio luminosity and its emission region.

But for now, the main point is that provided

\[
\frac{B}{P} > 10^{16} \text{ G s}^{-1}, \tag{19}
\]

the dynamics of the neutron star magnetosphere in terms of their back reaction on \( E_0 \) are dictated by Equations (8)–(11) and not Ruderman & Sutherland (1975), and this enables the Schwinger pairs to reach much larger \( \gamma \) values than the spark gap pairs, even if the spark gap voltages are much higher than the value in Equation (17) of Ruderman & Sutherland (1975), such as those depicted in Figure 2(a) of Rea et al. (2012).

4. POWER RADIATED BY THE SCHWINGER PAIRS; RADIO EMISSION MECHANISM

At the field strength of \( y_0 = 0.045 \) considered here, the rest energy density \( 2n_m c^2 = 2.24 \times 10^{18} \text{ erg cm}^{-3} \) of the pairs is much less than the energy density \( E_0^2/(8\pi) = 1.56 \times 10^{23} \text{ erg cm}^{-3} \) of the electric field. So from this viewpoint the field can oscillate for many cycles without noticeable decay, while the pair density \( n \) monotonically increases with time. Yet the pairs do not remain at rest. As we saw in the last section, most of the energy of the pairs is kinetic, because they are accelerated by the electric field.

The total electric energy per unit time spent in creating and accelerating the Schwinger pairs along the open magnetic field lines of the polar cap may be estimated by numerically integrating Equations (8), (9), and (11) over the relevant range of \( r, \theta, \phi \), taking into account the dependence of \( E_0 \) upon \( r \) and \( \theta \), viz. Equation (13), with the integration range for \( \theta \) limited to \( 0 \leq \theta \leq \theta_{\text{max}} \), where \( \theta_{\text{max}} \) is given by

\[
\sin^2 \theta_{\text{max}} = \frac{2\pi R}{P c} = 1.96 \times 10^{-3} \left( \frac{P}{0.106 \text{ s}} \right)^{-1}, \tag{20}
\]

and marks the last open field line that crosses the speed-of-light cylinder to transport radiation to infinity. Moreover, although there is no cutoff height \( r \) for \( E_0 \), the region of sufficient Schwinger breakdown for an appreciable contribution to the total power of the created pairs is found to be limited to

\[
r_{\text{max}} \approx 0.02 R, \tag{21}
\]

or a scale height \( \delta R \approx 2 \times 10^4 \text{ cm} \), a few times larger than the spark gap size of Ruderman & Sutherland (1975), Equation (18). Under the scenario of \( y_0 = 0.045 \) at \( r = R \) and \( \theta = 0 \), we find

\[
\frac{dE_{\text{pairs}}}{dt} = 5.59 \times 10^{41} \text{ erg s}^{-1}, \tag{22}
\]

which is slightly larger than the magnetic dipole radiation rate (Equation (1)) for the corresponding values of \( B = 10^{15} \text{ G} \) and \( P = 0.1 \text{ s} \). As \( y_0 \) increases beyond this point, the pair luminosity rises sharply, as seen in Table 1.

We are ready to discuss the generation of an FRB. In contrast to Ruderman & Sutherland (1975), wherein the radio signals originated downstream at the light cylinder boundary, the current model can accommodate the much simpler scenario of direct in situ emission from the polar caps, where the Schwinger discharge is actually occurring, by coherent curvature radiation. Specifically, the curvature radius for dipole magnetic field lines at the neutron star surface is

\[
\rho_c = \frac{R}{3 \sin \theta} \frac{1 + 3 \cos^2 \theta}{1 + \cos^2 \theta}^{3/2}, \tag{23}
\]

see e.g., Smirnov (1973). For \( y_0 \gtrsim 0.095 \) and \( B = 10^{15} \text{ G} \), the rotation period is \( P \lesssim 0.502 \text{ s} \) from Equation (13), which yields \( \theta_{\text{max}} = 0.0646 \text{ by Equation (20)}, \) hence \( \rho_c \lesssim 4.125 \times 10^7 \text{ cm} \) at \( \theta = \theta_{\text{max}}/2 \). Now the most conservative estimate of the size \( d \) of a region wherein charges are capable of emitting
that if all the charges are positioned within a radiation wavelength of each other to ensure the phase of each amplitude
wavefunctions are essentially the same vector. If the number density of
pairs is \( n \), the total emission from the polar cap becomes

\[
\frac{dE}{d\omega dR} = 4.67 \times 10^{-23} \rho_e^{-2/3} \omega^{3/2} \text{erg s}^{-1} \text{Hz}^{-1},
\]

between \( \omega = 0 \) and \( \omega = 3\gamma^2c/(2\rho_e) \). If the number density of
pairs is \( n \), the total emission from the polar cap becomes

\[
\frac{dE}{dt} = \frac{d}{dt} \int \hat{n} \times (\hat{n} \times \mathbf{v}) e^{i\omega(t-\hat{n} \cdot \mathbf{r})/c} dt,
\]

as the coherence volume.

Specific to curvature radiation, the intensity summed over all
directions, \( \int dI/d\omega \) per emitting charge in units of erg s\(^{-1}\) Hz\(^{-1}\),
may be approximated as

\[
\int dI/d\omega = 4.67 \times 10^{-23} \rho_e^{-2/3} \omega^{3/2} \text{erg s}^{-1} \text{Hz}^{-1},
\]

between \( \omega = 0 \) and \( \omega = 3\gamma^2c/(2\rho_e) \). If the number density of
pairs is \( n \), the total emission from the polar cap becomes

\[
\frac{dE}{dt} = \frac{d}{dt} n^2 \mathcal{V}_c \times \frac{\pi \theta_{\text{max}}^2}{2} \delta R
\]

\[
= 2.14 \times 10^{62} \left( \frac{n}{7.60 \times 10^{20} \text{cm}^{-3}} \right)^3 \left( \frac{\mathcal{V}_c}{9.84 \times 10^3 \text{cm}^3} \right)
\]

\[
\times \left( \frac{\theta_{\text{max}}}{2.4 \times 10^{-3}} \right)^3 \left( \frac{R}{10^8 \text{cm}} \right)^2
\]

\[
\times \left( \frac{\delta R}{2 \times 10^4 \text{cm}} \right) \delta R
\]

\[
\times \left( \frac{\rho_e}{4.125 \times 10^9 \text{cm}^3} \right)^{-2/3}
\]

\[
\times (\omega_{\text{min}}^{5/3} - \omega_{\text{max}}^{5/3}) \text{erg s}^{-1}
\]

with the effect of coherence taken into account, where the
parameters \( \theta_{\text{max}} \) and \( \rho_e \) are chosen to suit the \( \gamma_0 = 0.095 \)
scenario for the following reason. When a young (newly born)
magnetar has \( \gamma_0 \) exceeding this value, then, assuming\(^2\) that
\( \omega_{\text{max}} \) corresponds to \( \nu_{\text{max}} = 1.4 \text{ GHz} \) and \( \omega_{\text{min}} \) to
\( \nu_{\text{min}} = 1 \text{ MHz} \), Equation (27) would yield \( dE/dt \approx 10^{51} \)
\( \text{erg s}^{-1} \), which from Table 1 is the same as the energy spent
on producing and accelerating the pairs. One can see that in
such a regime the energy released by the Schwinger discharge
mechanism can all be dissipated into coherent radio emission,
which escapes along the open field lines.

5. CONCLUSION: A MODEL OF FRBS AS GLITCH EVENTS IN FAST-SPINNING YOUNG MAGNETARS

It is now possible to propose an origin of FRB events, as
phenomena ensuing from the birth of a magnetar possessing
too high an initial spin rate, which leads in turn to the creation
of a giant burst of \( e^+ e^- \) pairs via the Schwinger mechanism.

The energy from this burst is ultimately emitted in the radio
wavelengths for reasons given in the previous section, and is
drawn from the star’s rotation. In fact, the energy budget is
large enough to cause the star to undergo a short episode of
decelerating spin-gluitch which is the duration of the FRB.

Quantitatively, here is how the numbers work out. Since the
kinetic energy of the neutron star rotation is

\[
\mathcal{E}_{\text{rot}} = \frac{1}{2} \mathcal{I} \Omega^2 = 6.27 \times 10^{48} \left( \frac{M_{\text{max}}}{M_\odot} \right) \left( \frac{R}{10^6 \text{ cm}} \right)^2
\]

\[
\times \left( \frac{P}{0.0502 \text{ s}} \right)^{-2} \text{erg},
\]

the radio emission, which is dominated by radiation below the
1 GHz range, would significantly slow down the rotation in a
matter of \( 5–6 \times 10^3 \) s, i.e., the timescale of an FRB. Unless some
spin-up mechanism is available to recover the original
high rotation rate that triggered the Schwinger mechanism, the
FRB will not be repeat again in the same magnetar, because the
spin deceleration takes the magnetar outside the regime where
the Schwinger mechanism is powerful enough to deliver the
luminosity of an FRB.

Moreover, magnetars with the requisite period of \( P < 0.1 \) s
are invariably young, having an upper age limit set by the
magnetic dipole radiation rate (the rightmost column of
Table 1) and total available energy (Equation (28)) of a
few \( 10^6 \) years. Such young magnetars are likely to be

\footnote{The existence of \( \nu_{\text{max}} \) is because if \( \nu_{\text{min}} \) is below \( \approx 1 \text{ MHz} \), the coherence
volume \( \mathcal{V}_c \) would, by Equation (25), span a thickness \( \delta R \) in excess of the
constraint given by Equation (21), to include regions where the field is too
weak to decay into pairs.}

Table 1

| \( \gamma_0 \) | \( n \text{ (cm}^{-3} \) | \( \theta_{\text{max}} \text{ (degrees) \) | \( \rho_e \text{ (cm}^{-3} \) | \( \mathcal{I} \text{ (erg s}^{-1} \) | \( \nu_{\text{max}} \text{ (MHz) \) |
|---|---|---|---|---|---|
| 0.045 | 3.38 \times 10^{-5} | 1.50 \times 10^{12} | 3.94 \times 10^{14} | 4.79 \times 10^{39} | 1.90 \times 10^{41} |
| 0.055 | 5.38 \times 10^{-5} | 1.19 \times 10^{12} | 1.05 \times 10^{12} | 5.59 \times 10^{41} | 3.05 \times 10^{41} |
| 0.065 | 1.84 \times 10^{-12} | 5.37 \times 10^{21} | 5.59 \times 10^{7} | 6.05 \times 10^{34} | 6.81 \times 10^{41} |
| 0.095 | 2.49 \times 10^{-13} | 4.40 \times 10^{22} | 8.51 \times 10^{6} | 1.09 \times 10^{50} | 3.87 \times 10^{42} |

Note. The quantities listed in columns 1, 2, 3, 4, and 6 correspond to (7)–(9), (11), and (1), respectively, while column 5 is derived numerically as described in the text preceding Equation (22). The surface magnetic field and radius of the neutron star are fixed at \( B = 10^{15} \) G and \( R = 10^6 \) cm, respectively, i.e., any value of \( \gamma_0 \) can be reached by adjusting the spin period \( P \) via Equation (13).
embedded in the ionized gas of a supernova remnant, and the Schwinger pairs that constitute the aforementioned jet could snowplow into the supernova ejecta to cause a shell of piled-up matter (Chen et al. 2016); indeed, at least in one case a circumstellar disk was actually observed (Wang et al. 2006). If the $\nu < 1$ GHz frequencies are absorbed or reflected by the shell plasma, as suggested by Kulkarni et al. (2015), the only observable radiation would be in the $\approx 1$ GHz, in which case Equation (27) would deliver $\approx 5 \times 10^{42}$ erg for $\approx 5 \times 10^{-3}$ s duration for the bandwidth of 10% centered at 1.4 GHz. This would explain the total radio luminosity of an FRB (note there is no point in including frequency components too much above 1.4 GHz because, by Equation (27), the luminosity has become too low). It also provides a more substantive basis for the claims of Popov & Postnov (2013), Katz (2014), Kulkarni et al. (2015), Lyubarsky (2014), Pen & Connor (2015), and Katz (2016) of a possible association of FRBs with giant magnetar flares.

In terms of a falsifying test, one possible way is to look for FRBs with duration much longer than 10 ms ($10^{-2}$ s). They would challenge the proposed model, at least in its simplest form as presented here, because the only way for the magnetar to lose a significant amount of rotational kinetic energy on timescales $\gg 10^{-2}$ s is if the ratio $B/P$ as expressed via $y_0$ falls below 0.09, but in that case the 1 GHz luminosity would also have become too low for observability of the radio burst from an extragalactic distance. Another test in the same vein would involve discovering more repeating FRBs of the Spitler et al. (2016) type, as again the simplest version of our model cannot account for such a behavior.

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