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FOCUSING OF ALFVÉNIC WAVE POWER IN THE
CONTEXT OF GAMMA-RAY BURST EMISSIVITY

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Abstract

Highly dynamic magnetospheric perturbations in neutron star environments can naturally account for the features observed in Gamma-ray Burst spectra. The source distribution, however, appears to be extragalactic. Although noncatastrophic isotropic emission mechanisms may be ruled out on energetic and timing arguments, MHD processes can produce strongly anisotropic $\gamma$-rays with an observable flux out to distances of $\sim 1 - 2$ Gpc. Here we show that sheared Alfvén waves propagating along open magnetospheric field lines at the poles of magnetized neutron stars transfer their energy dissipationally to the current sustaining the field misalignment and thereby focus their power into a spatial region $\sim 1000$ times smaller than that of the crustal disturbance. This produces a strong (observable) flux enhancement along certain directions. We apply this model to a source population of "turned-off" pulsars that have nonetheless retained their strong magnetic fields and have achieved alignment at a period of $\gtrsim 5$ seconds.

*Subject headings*: acceleration of particles – cosmology: miscellaneous – galaxies: evolution – gamma rays: bursts – MHD – pulsars: general
1. Introduction

The lack of a precise determination of a distance scale to Gamma-ray burst (GRB) sources has greatly hindered our theoretical understanding of these objects. Much of what we know about these bursts is based on inferences drawn from clues provided by their spectra, including: (1) Most events exhibit rapid variability, apparently on time scales shorter than the best available instrument temporal resolution, with one burst exhibiting structure on a 200 $\mu$s timescale (Bhat et al. 1992). This variability indicates that the sources must be very compact, characterized by a length scale no larger than $\sim 50-100$ km. (2) Bursts can last anywhere from tens of milliseconds to as long as 900 seconds, and usually have a complex temporal structure. This complicated time dependence would seem to favor mechanisms that invoke highly dynamic perturbations in otherwise stable environments. (3) Typical GRBs emit a substantial fraction of their power at energies in excess of $\sim 1$ MeV, which suggests that nonthermal processes are responsible for the emission of the $\gamma$-rays. In this regard, the inverse-Compton scattering of soft photons by relativistic particles has been very successful in reproducing the observed spectra (e.g., Pozdnyakov, Sobol’ & Sunyaev 1977; Canfield, Howard & Liang 1987; Melia & Fatuzzo 1989; Ho & Epstein 1989; Melia 1990a,b).

It is possible to account for these observations by invoking a model in which the bursts originate within the magnetosphere of strongly magnetized neutron stars. In addition to being very compact, these environments are subject to magnetic fluctuations on sub-millisecond timescales, and the constituent particles can be energized nonthermally via the induced electrostatic forces.

As is well known, however, the neutron-star paradigm must be reconciled with the uniform, yet spatially truncated GRB distribution observed by the BATSE experiment on CGRO (Meegan et al. 1992). These observations seem to rule out nearby (i.e. Galactic) single population models, and have therefore led to renewed speculation that GRBs originate at cosmological redshifts. But a naive estimate of the burst energy required for such distant sources yields a value that is significantly larger than that which a neutron
star could reasonably produce, unless the event was catastrophic (e.g., the coalescence of a neutron-star binary, Narayan et al. 1992), which does not seem to be borne out by the time history of typical bursts.

A resolution to this apparent conflict was recently proposed by Melia & Fatuzzo (1992, hereafter MF), in which sheared Alfvén waves generated near the polar cap of strongly-magnetized neutron stars produce streams of relativistic particles that are focused by the underlying magnetospheric structure. These energetic charges upscatter the radio-frequency photons (emitted at larger radii) into $\gamma^{-1}$ cones aligned with the underlying magnetic field lines, resulting in an enhanced $\gamma$-ray flux along preferred lines of sight. This anisotropic emission is such that a pulsar glitch releasing $\sim 10^{45}$ ergs of energy could be viewed as a GRB out to a distance of $\gtrsim 1$ Gpc. A key assumption of this scenario is that the Alfvénic power can indeed be emitted anisotropically. We show in this Letter that the required focusing is a natural consequence of the dissipational properties of sheared Alfvén waves whose shear length scales ($s \approx 10$ cm) are much smaller than the size of the region ($\gtrsim 10^4$ cm) encompassing the overall Alfvén wave fluctuation. As such, this work strengthens the case for a non-catastrophic, cosmological origin of GRBs, and supports our view that an improved understanding of the micro-physical processes in neutron-star environments can indirectly, though significantly, influence our study of galactic evolution out to a redshift in excess of $1 - 2$ (Tamblyn & Melia 1993).

2. Sheared Alfvénic Wave Dissipation

The general theory of sheared Alfvén waves (SAWs) has been developed elsewhere (Melia & Fatuzzo 1992; Fatuzzo & Melia 1993). Here we consider the global properties of their dissipation. We idealize the unperturbed polar cap region as a fully ionized, homogeneous plasma threaded by a uniform magnetic field $B_0 = B_0 \hat{z}$. The sheared Alfvén waves may therefore be described by magnetic perturbations of the form

$$B_A = B_a(y) \exp(ikz - i\omega t)\hat{x},$$

(1)

where $B_a(y)$ is an odd function that characterizes the shear geometry. For our purposes
here, we take the shear profile to be

\[
B_a(y) = \begin{cases} 
B_{a0} & S \geq y \geq s \\
B_{a0} g(y) & |y| \leq s \\
-B_{a0} & -S \leq y \leq -s
\end{cases},
\]

(2)

where \( g(y) \) is a continuous function that satisfies the condition \( g(\pm s) = \pm 1 \), and where \( s \ll S \) (\( s \) being the shear lengthscale and \( S \) the lengthscale of the encompassing plane wave regions). For convenience, we define \( \eta \) as the ratio of the nonsheared to sheared surface areas (\( \eta \equiv S/s \) in the present geometry).

It is clear from the form of \( B_a(y) \) and Ampère’s Law that an electric field \( E_{Az} \) must exist inside the sheared region (\( |y| < s \)) until a sufficiently strong current \( J_s \) is produced parallel to the underlying magnetic field \( B_0 \). For these waves, the equilibrium Goldreich-Julian particle density \( n_0 \) is insufficient to support a current large enough to short out \( E_{Az} \). Charges must therefore be copiously stripped off the stellar surface, thereby inducing a charged particle flow to give the required \( J_s \). Since the Alfvén speed is \( v_\alpha \gg c \), the waves travel with a phase velocity \( u_\alpha = \omega/k \approx c \). In order for \( J_s \) and the encompassing magnetic shear to remain in phase, the particle flow must be relativistic, and thus, have an average density

\[
n_s \approx \frac{B_{a0}}{4\pi e s} \approx 1.7 \times 10^{19} \text{cm}^{-3} \left( \frac{B_{a0}}{10^{12} \text{G}} \right) \left( \frac{s}{10 \text{ cm}} \right)^{-1}.
\]

(3)

We note that the stripped particles escape from the system by flowing out along the open magnetospheric field lines.

In standard pulsar theory, radio emission results from the coherent motion of “bunches” of electrons streaming along open field lines with Lorentz factors \( \gamma \sim 10^{4-5} \). As such, strong transient radio emission is expected to be a natural byproduct of sheared Alfvén waves if similar particle energies are reached, and if this emission is produced with front-back symmetry along the local field-line direction, a large fraction of the overall radio flux will naturally be funneled back onto the polar cap. Taking into account the coherent nature of the processes responsible for pulsar emission, we parametrize the flux impinging onto the stellar surface by

\[
F_r = \xi \eta^{-2} (n_s/n_C)^2 \left( \frac{L_C}{\pi R_{pc}^2} \right),
\]

where \( L_C \) and \( n_C \) are the Crab pulsar
luminosity and magnetospheric number density, respectively, and where \( R_{pc} \) is the radius of the open field line polar cap. Assuming a stellar radius of \( R_* = 10^6 \text{cm} \), \( R_{pc} \) can be related to the pulsar period \( P \) via \( R_{pc} = 1.4 \times 10^4 \text{ cm} \left( \frac{P}{1 \text{s}} \right)^{-1/2} \). With \( L_C = 10^{32} \text{ ergs s}^{-1} \) and \( n_C = 10^{13} \text{ cm}^{-3} \), this yields

\[
F_r = 2 \times 10^{30} \text{ ergs cm}^{-2} \text{ s}^{-1} \xi \left( \frac{\eta}{10^3} \right)^{-2} \left( \frac{B_{a0}}{10^{12} \text{ G}} \right)^2 \left( \frac{s}{10 \text{ cm}} \right)^{-2} \left( \frac{P}{5 \text{ s}} \right),
\]

where \( \eta^{-1} \) is the sheared flow “filling factor”. The parameter \( \xi \) encompasses both geometric and emission uncertainties, and as such is poorly known. We note that if \( \xi \) becomes too small (\( \lesssim 0.1 \) for the range of parameters considered here), the wave dissipation lengthscale due to field line annihilation becomes much larger than \( R_* \), and the SAW mechanism becomes inefficient at producing \( \gamma \)-rays (see the discussion after equation [7]). However, since the Crab pulsar is itself very inefficient at converting spin-down energy into radio emission compared to typical pulsars, and since we have made the conservative assumption that the (coherent) radio flux scales as \( \eta^{-2} \) (i.e., the square of the total number of particles), it is reasonable to assume that \( \xi \gg 0.1 \) (see also the discussion after equation [11]).

The presence of \( F_r \) results in a radiative drag on the relativistic current-carrying charges. By analogy with MHD phenomena, the current driving electric field \( (\mathbf{E} = E_A \hat{z}) \) must be generated within the shear at the expense of the magnetic wave energy. However, the sheared waves are distinguished from pure MHD fluctuations for two important reasons. First, the charges which generate \( J_s \) are constrained to always move along the same \( \mathbf{B}_0 \) field lines, so that SAWs cannot easily change their initial structure. Second, the simple concept of Ohm’s law is not valid for the relativistic flow inside the shear. Indeed, once the particles become relativistic, the current quickly decouples from the driving electric field, and since the radiative drag increases rapidly with \( \gamma \) (the particle Lorentz factor), one might expect that a mildly relativistic flow will be favored by the system.

Though \( E_A \hat{z} \) depends on the microphysics of the shear (including all the annihilation processes, such as the tearing mode instability), its value may be estimated with a relatively simple argument under the assumption that the annihilation time scale within the shear
is the shortest of the relevant time scales. The strength of the electric field is limited by
the rate at which the oppositely-directed magnetic fluctuations are driven together by the
large magnetic pressure gradients associated with SAWs. Since the Alfvénic field lines are
strongly coupled to \( B_0 \) via flux freezing with the charged medium, this transfer of Alfvénic
power into the shear is dictated by the diffusion rate within the resistive plasma in the
region \(|y| > s\), where the resistance is provided primarily by \( e^-/\)radio photon scatterings
that occur with a frequency \( \sim n_{ph} \sigma_{MC} c \), in terms of the photon number density \( n_{ph} \) and
the magnetic Compton cross section \( \sigma_{MC} \). With \( \epsilon_0 \) the characteristic radio photon energy,
each event imparts a momentum \( \sim \epsilon_0/c \) to the electron (moving nonrelativistically with a
velocity \( v_e \gg \epsilon_0/m_e c \)), which must therefore interact with \( \sim (m_e v_e c/\epsilon_0)(c/v_e) \) photons in
order to suffer a significant deviation to its path. Thus, the electron deflection frequency is
\[
\nu_e \sim \frac{F_r \sigma_{MC}}{m_e c^2} \approx 4.3 \times 10^{11} \text{s}^{-1} \left( \frac{\eta}{10^3} \right)^{-2} \left( \frac{B_{a0}}{10^{12} \text{G}} \right)^2 \left( \frac{s}{10 \text{ cm}} \right)^{-2} \left( \frac{P}{5 \text{s}} \right),
\]
where \( \sigma_{MC} \approx \sigma_T/4 \) when \( B_{a0} \sim B_0 \sim 10^{12} \text{G} \) (Dermer 1990).

With a conductivity \( \sigma \approx n_0 e^2/m_e \nu_e \) and a diffusion time scale \( \tau_d \equiv 4\pi \sigma S^2/c^2 \) (the well
known MHD value which is valid as long as \( \tau_d \ll 2\pi/\omega \)), the diffusion velocity \( v_d \approx S/\tau_d \) for the magnetic field lines is given as
\[
v_d = \min \left[ c, 8.7 \times 10^8 \text{cm s}^{-1} \right. \left. \times \xi \left( \frac{\eta}{10^3} \right)^{-2} \left( \frac{B_{a0}}{10^{12} \text{G}} \right)^2 \left( \frac{P}{5 \text{s}} \right)^2 \left( \frac{B_0}{10^{12} \text{G}} \right)^{-1} \left( \frac{s}{10 \text{ cm}} \right)^{-2} \left( \frac{S}{10^4 \text{ cm}} \right)^{-1} \right].
\]
Thus, since \( u_\alpha \approx c \), the waves dissipate over a lengthscale
\[
R_d \equiv \left( \frac{Sc}{v_d} \right) \approx \max \left[ S, 3.5 \times 10^5 \text{cm} \right. \left. \times \xi^{-1} \left( \frac{\eta}{10^3} \right)^2 \left( \frac{B_{a0}}{10^{12} \text{G}} \right)^{-2} \left( \frac{P}{5 \text{s}} \right)^{-2} \left( \frac{B_0}{10^{12} \text{G}} \right) \left( \frac{s}{10 \text{ cm}} \right)^2 \left( \frac{S}{10^4 \text{cm}} \right)^2 \right].
\]
As long as \( R_d \ll R_* \) (i.e., \( \xi \) is sufficiently large), most of the wave energy is channeled
into the shear before the waves break, and we may equate the Alfvénic luminosity (\( \sim B_{a0}^2 c A_w/8\pi \)) generated at the stellar surface with the power (\( \sim E_{Az} J_s A_s R_d \)) dissipated by
the current as it converts magnetospheric energy into upscattered radiation. Here, $A_w$ and $A_s$ are the surface areas corresponding to the wave and shear regions, respectively. This yields an average electric field strength

$$E_{Az} \approx \min \left[ 5 \times 10^{11} \text{S V cm}^{-1} \left( \frac{B_{a0}}{10^{12} \text{G}} \right) \left( \frac{\eta}{10^3} \right) \left( \frac{s}{10 \text{ cm}} \right) \left( \frac{S}{10^4 \text{ cm}} \right)^{-1}, 1.4 \times 10^{10} \text{S V cm}^{-1} \right] \times \xi \left( \frac{\eta}{10^3} \right)^{-1} \left( \frac{B_{a0}}{10^{12} \text{G}} \right)^3 \left( \frac{P}{5 \text{ s}} \right)^2 \left( \frac{B_0}{10^{12} \text{G}} \right)^{-1} \left( \frac{s}{10 \text{ cm}} \right)^{-1} \left( \frac{S}{10^4 \text{ cm}} \right)^{-2} .$$  

(8)

As before, we assume a typical pulsar spectrum specified as a steep power law with (flux density) index $\mu$ above a break at frequency $\epsilon_0/h \approx 500$ MHz. With $\gamma \gg 1$, a lab frame photon with energy $\epsilon$ will be blue-shifted to $\sim 2\gamma \epsilon$ in the electron rest frame, which is well below the resonant energy $\epsilon_B \equiv (B_0/44.14 \times 10^{12} \text{G}) m_e c^2$, and its angle of propagation relative to the particle direction (and hence $B_0$) is $\sim \gamma^{-1}$. As such, $\epsilon_B/\epsilon \approx 6 \times 10^9 (B_0/10^{12} \text{G}) (\epsilon/h 500 \text{MHz})^{-1} < \gamma^2$, and $\sigma_{MC} \approx 4\sigma_T (\gamma \epsilon/\epsilon_B)^2$ (Melia & Fatuzzo 1989; Dermer 1990). Balancing the accelerating force $eE_{Az}$ by the radiative drag $\gamma^2 F_r \sigma_{MC}/c$, one therefore obtains

$$\gamma_s \approx 10^6 \left( \frac{B_{a0}}{10^{12} \text{G}} \right)^{1/4} \left( \frac{B_0}{10^{12} \text{G}} \right)^{1/4} \left( \frac{P}{5 \text{ s}} \right)^{1/4} \left( \frac{S}{10^4 \text{ cm}} \right)^{-1/4} \left( \frac{\epsilon_0}{h 500 \text{MHz}} \right)^{-1/2} f(\xi) ,$$  

(9)

where $f(\xi) = \min[1, (\xi/\xi_0)^{-1/4}]$ and the parameter $\xi_0$ is defined as the smallest value of $\xi$ for which $v_d = c$ (e.g., $\xi_0 = 35$ for the represented parameter space), and where we have used $\eta = S/s$. This result is consistent with the assumptions discussed above (e.g., that the particle motion is relativistic and sufficiently energetic to produce the required radio luminosity and that $\sigma_{MC}$ have an $\epsilon^2$ dependence). Ultimately, the current $J_s$ must decay in concert with the Alfvénic magnetic field, even though the particles remain relativistic. Evidently, the initially fully charge separated regions must merge together and neutralize. This behavior is expected since the power transferred from the wave to the particles is reduced as the magnetic field decays. As such, an increasing number of charges undergoing collisions will not be energized quickly enough to remain in phase with the wave, and are therefore swept up by the lagging oppositely charged wave region.
The scaling of the magnetic fields $B_0$ and $B_{a0}$ in the above equations was chosen for convenience and not to suggest that $B_{a0} \sim B_0$. We note, however, that the presence of significant reconnection in the nonlinear regime would have the desirable effect of enhancing the annihilation rate within the shear (see the paragraph preceding Eq. 5).

3. The Cosmological Radio Pulsar Model For Gamma-ray Bursts

In applying the above discussion to the cosmological gamma-ray burst model, we must now generalize to a more realistic magnetospheric geometry in which the field lines are more or less radial close to the stellar surface (cf. Eq. 7). Although globally the results of §2 are expected to apply here, we note an important difference between the two geometries. The underlying magnetic field strength $B_0$ now decreases away from the stellar surface as $(R_*/r)^2$, whereas in the absence of wave damping, the Alfvénic field decreases only as $(R_*/r)$, for which the waves eventually lose their linearity and shock. However, since the wave energy diffuses into the shear on a length scale $< R_*$, the Alfvénic field also decreases roughly as $(R_*/r)^2$.

It is evident from §2 that SAWs focus their energy into the internal current flow. Thus, as long as $\gamma^{-1} \ll s/R_*$, the flux is enhanced by a factor $\eta$ in certain directions, for which a source at a distance $D$ will have an observable $\gamma$-ray flux

$$F_\gamma = \eta \frac{B^2_{a0}}{8\pi c} \left(\frac{R_*}{D}\right)^2 = 1.3 \times 10^{-7} \text{ ergs cm}^{-2} \text{ s}^{-1} \left(\frac{\eta}{10^3}\right) \left(\frac{B_{a0}}{10^{12} \text{ G}}\right)^2 \left(\frac{D}{1 \text{ Gpc}}\right)^{-2}. \quad (10)$$

The observability of these bursts at cosmological distances imposes strict (but not unrealistic) conditions on the model parameters, such as the required burst power ($L_{\text{burst}} \sim 10^{44} - 10^{46} \text{ ergs s}^{-1}$), whose magnitude depends on whether the SAWs are generated only near the polar cap (where the field lines are most strongly coupled to toroidal crustal activity), or are generated throughout the entire stellar surface.

Since the particle flow remains optically thin to the radio photons impinging upon the star, the $\epsilon^2$ dependence of the cross-section (see above) implies that the incipient radio spectrum is upscattered to a $\gamma$-ray spectrum with (power) index $\mu + 2 + 1$, and very importantly, that the spectral radio break at $\epsilon_0$ is translated to the corresponding $\gamma$-ray
break at

$$\epsilon_{\text{break}} \sim 2\gamma_s^2 \epsilon_0 \approx 4.4 \text{ MeV} \left( \frac{B_{a0}}{10^{12} \text{ G}} \right)^{1/2} \left( \frac{B_0}{10^{12} \text{ G}} \right)^{1/2} \left( \frac{P}{5 \text{ s}} \right)^{1/2} \left( \frac{S}{10^4 \text{ cm}} \right)^{-1/2} g(\xi) ,$$

(11)

independent of $\epsilon_0$ (we have again used $\eta = S/s$). Here, $g(\xi) = \min[1, (\xi/\xi_0)^{-1/2}]$. This result compares favorably with the observed value of $\epsilon_{\text{break}}$ (which after redshift is taken into account is seen to fall within the range $\sim 100 \text{ keV} - 3 \text{ MeV};$ Schaefer et al. 1992), and suggests that $\xi/\xi_0$ may be as large as 100. A more detailed description of the resulting $\gamma$-ray spectrum is given in MF (see, for example, Figure 2 therein). For completeness, we note that a cylindrical shear would correspond to $\eta \sim (S/s)^2$ and the parameter $S$ in Equation (11) should be replaced by $s$. Such a strictly confined shear region would thus appear to be unlikely, though it cannot be ruled out without a more detailed calculation.

The question of which mechanisms contribute to the generation of SAWs is best addressed by considering the constraints imposed on the model by the different probabilities of observing extragalactic and galactic events. If the source distribution is cosmological, the BATSE data imply a detection rate $R_c \approx 6 \times 10^{-7}$ bursts per year per galaxy, with an actual rate $R_T = R_c / P_c$, where $P_c$ is the probability of detecting the extragalactic burst. Assuming that the number of galactic sources is roughly equal to the average number of sources in all other galaxies, we should therefore expect to see a galactic rate $R_g = R_T P_g$, where $P_g$ is the corresponding probability of detecting a galactic burst in progress. According to Equation (10), a galactic burst would be quite distinguishable, exhibiting fluxes of order $1 \text{ ergs cm}^{-2} \text{ s}^{-1}$, i.e., $\sim 10^7$ times larger than their extragalactic counterparts. Since none of these have ever been detected, we infer that $R_g \ll 0.01$ per year, for which $P_c/P_g \geq 6 \times 10^{-5}$.

What this means in practice is best seen with recourse to a specific scenario. Let us assume that the underlying sources are turned-off pulsars (i.e., neutron stars older than $\sim 10^7$ years), which have nonetheless retained their strong magnetic fields and have achieved alignment at a period $P \sim 5\text{s}$, corresponding to an open field line polar cap radius $R_{pc} \approx 6 \times 10^{3} \text{ cm}$. The stellar rotation implies that a given $\gamma$-ray flux region sweeps in and
out of the observer’s view, but as long as the duration of an individual sweep is longer than the instrument resolution time $\tau_i$, the inferred (average) flux is correctly given by Eq. (10). If, however, the sweep time is shorter than $\tau_i$, the observed (average) flux for the most distant bursts drops below the instrument sensitivity. This means that an extragalactic burst is observable predominantly within a cone centered about the rotation axis, with an opening angle corresponding to a radius $R_{ob} \equiv (s/\pi)(P/\tau_i)$ (two poles). Thus, since the probability that a sheared region occurs within the area enclosed by this “observable” cap is $\sim R_{ob}/S$ (where $R_{ob} < S$), the probability of seeing a burst in progress at cosmological distances must be

$$P_c \sim 7.6 \times 10^{-10} \left( \frac{s}{10 \text{ cm}} \right)^3 \left( \frac{P}{5 \text{ s}} \right)^3 \left( \frac{\tau_i}{64 \text{ ms}} \right)^{-3} \left( \frac{S}{10^4 \text{ cm}} \right)^{-1}. \quad (12)$$

The actual event rate would therefore be $R_T \sim 10^3$ per galaxy per year, which in turn implies a stellar repetition time scale of $\sim 10^3 - 10^6$ years if the population is comprised of $\sim 10^6 - 10^9$ objects.

The probability of detecting such a burst in progress within the galaxy is higher since for these events $R_{ob} \sim R_{pc}$ (see above). As such, $P_g$ is simply the ratio of solid angles corresponding to the polar cap region and the entire star (i.e., $4\pi$). For a $P = 5$ s rotator, we therefore infer that $P_g \approx 2 \times 10^{-5}$, so that $P_c/P_g \sim 4 \times 10^{-5}$ for this population, which is consistent with the constraint discussed above. To put this result in another way, we would anticipate that such a population of GRB sources should produce an observable galactic “super” burst roughly once every $\gtrsim 50$ years. In addition, we note that GRBs originating from extragalactic sources are expected to be accompanied by $\sim 0.01 - 1.0$ Jansky radio bursts (see MF). For galactic bursts, the observable radio flux will be roughly 10 orders of magnitude larger and therefore (as is the case for the $\gamma$-ray signal) quite distinguishable. Assuming that the radio emission is produced at $\sim 10R_*$ above the polar cap, the probability of seeing the radio burst will be roughly 10 times greater than of seeing the corresponding $\gamma$-ray burst, suggesting an observable galactic “super” radio burst rate of roughly once every $\gtrsim 5$ years. This result does not conflict with current observations.
since only a small fraction of the sky is monitored by radio telescopes at any given time. A possible link between GRB sources and soft $\gamma$-ray repeater events, which in this picture would be interpreted as bursts viewed outside of the open field line cone, has been discussed elsewhere (Melia & Fatuzzo 1993).

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References

Bhat, P. N. et al. 1992, Nature, 359, 217.
Canfield, E., Howard, W. M. & Liang, E. P. 1987, Ap. J., 323, 565.
Dermer, C. D. 1990, Ap. J., 360, 197
Fatuzzo, M. & Melia, F. 1993, Ap. J., 407, 680.
Ho, C. & Epstein, R. I. 1989, Ap. J., 343, 277.
Meegan, C. A., et al. 1992, Nature, 355, 143
Melia, F. 1990a, Ap. J., 351, 601.
Melia, F. 1990b, Ap. J., 357, 161.
Melia, F. & Fatuzzo, M. 1989, Ap. J., 346, 378.
Melia, F. & Fatuzzo, M. 1992, Ap. J. (Letters), 398, L85 (MF)
Melia, F. & Fatuzzo, M. 1993, Ap. J. (Letters), 408, L9.
Narayan, R., Paczyński, B. & Piran, T. 1992, Ap. J. (Letters), 395, L83.
Pozdnyakov, L. A., Sobol, I. M. & Sunyaev, R. A. 1977, Sov. Astron., 21, 6.
Schaefer, B. E., et al. 1992, Ap. J. (Letters), 393, L51.
Tamblyn, P. & Melia, F. 1993, Ap. J. (Letters), submitted.
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