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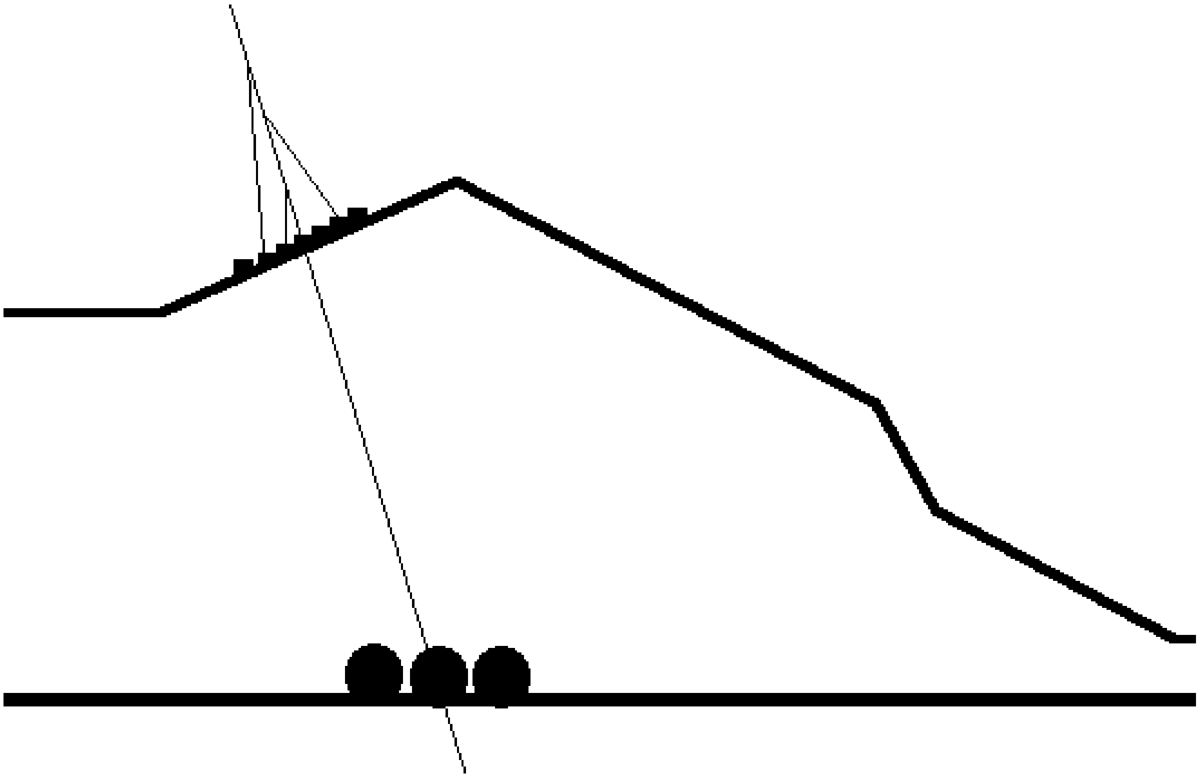
24 Gennaio 2000



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INFN – Laboratori Nazionali del Gran Sasso

Published by ***SIS–Pubblicazioni***
Laboratori Nazionali di Frascati

INFN - Istituto Nazionale di Fisica Nucleare

Laboratori Nazionali del Gran Sasso

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Superconducting Cosmic Strings as Gamma Ray Burst Engines

V.Berezinsky[a,b], B.Hnatyk[a,c], A.Vilenkin[d]

[a] *INFN, Laboratori Nazionali del Gran Sasso, I-67010 Assergy (AQ), Italy*

[b] *Institute for Nuclear Research, Moscow, Russia*

[c] *Institute for Applied Problems in Mechanics and Mathematics, NASU, Naukova 3b, Lviv-53, 290053, Ukraine*

[c] *Physics Department, Tufts University, Medford, MA 02155, USA*

Abstract

Cusps of superconducting strings can serve as GRB engines. A powerful beamed pulse of electromagnetic radiation from a cusp produces a jet of accelerated particles, whose propagation is terminated by the shock responsible for GRB. A single free parameter, the string scale of symmetry breaking $\eta \sim 1 \cdot 10^{14}$ GeV, successfully explains the typical values of GRB duration, fluence, the rate of bursts, as well as the observed ranges of these quantities. The wiggles on the string can drive the short-time structures of GRB. The model predicts reoccurance of GRBs with a period of ~ 20 yrs. The cusp mechanism resolves also some problems of GRBs as sources of Ultra High Energy Cosmic Rays.

Models of gamma ray bursts (GRBs) face the problem of explaining the tremendous energy released by the central engine [1]. In the case of isotropic emission, the total energy output should be as high as 4×10^{54} ergs. Strongly beamed emission is needed for all known engine models, such as mergers and hypernovae, but such extreme beaming is difficult to arrange. In this paper we show that emission of pulsed electromagnetic radiation from cusps of superconducting cosmic strings naturally solves this problem and explains the observational GRB data using only one engine parameter.

Cosmic strings are linear defects that could be formed at a symmetry breaking phase transition in the early universe [2]. Strings predicted in most grand unified models respond to external electromagnetic fields as thin superconducting wires [3]. As they move through cosmic magnetic fields, such strings develop electric currents. Oscillating loops of superconducting string emit short bursts of highly beamed electromagnetic radiation and high-energy particles [4, 5].

The idea that GRBs could be produced by superconducting strings was first suggested by Babul, Paczynski and Spergel [6] (BPS) and further explored by Paczynski [7]. They assumed that the bursts originate at very high redshifts ($z \sim 100 - 1000$), with the string currents being induced by a primordial magnetic field.

As it stands, the BPS model does not agree with observations. The observed GRB redshifts are in the range $z \leq 3$, and the observed duration of the bursts ($10^{-2}s \leq \tau \leq 10^{3}s$) is significantly longer than that predicted by the model. On the theoretical side, our understanding of cosmic string evolution has considerably evolved since the BPS papers were written.

As in the BPS model we shall use the cusp of a superconducting string as the central engine in GRB. It provides the tremendous engine energy naturally beamed. Other features of our model differ from BPS and allow us to explain the GRB observations.

As it propagates in plasma, low-frequency electromagnetic radiation from the cusp loses its energy accelerating the particles to very large Lorentz factors. Like the initial electromagnetic pulse, these particles are beamed and give rise to a hydrodynamical flow in the surrounding gas, terminated by a shock, as in the standard theory of GRB [10] (for a review see [1]).

The string electric currents are induced by oscillation of string loops in an external magnetic field. Unlike BPS, we do not postulate the existence of a strong primordial magnetic field for which there is little observational evidence. Instead, we use a more realistic model, in which cosmic magnetic fields were generated in young galaxies (*e.g.*, during the bright phase of their evolution [11]) and then dispersed by galactic winds in the intergalactic medium. Our mechanism of GRB operates in galaxies, clusters and superclusters of galaxies, and generally in the sheets and filaments defined by the large-scale structure. Thus, in contrast to the BPS model, GRBs originate at moderate redshifts. Finally, the string symmetry breaking scale η will be the only string parameter used in our calculations. This parameter is sufficient to account for all main GRB observational quantities: the duration τ_{GRB}, the rate of events $\dot{N}_{GRB}$, and the fluence S (or the total energy output $\mathcal{E}$).

An alternative channel of producing GRBs by cusps is through emission and subsequent decay of superheavy particles. This channel is also effective, though model-dependent. It depends on the masses of charge carriers in superconducting string and on their scattering cross-sections. We shall discuss this mechanism in a subsequent publication [12]. Here we consider only electromagnetically driven GRBs.

We begin with a brief review of string properties and evolution. Numerical simulations indicate that strings evolve in a self-similar manner [8, 9, 13]. A horizon-size volume at any time t contains a few long strings stretching across the volume and a large number of small closed loops. The typical distance between long strings and their characteristic curvature radius are both $\sim t$, but, in addition, the strings have small-scale wiggles of wavelength down to

$$l \sim \alpha t, \tag{1}$$

with $\alpha \ll 1$. The typical length of loops being chopped off the long strings is comparable to the scale of the smallest wiggles (1). The loops oscillate and lose their energy, mostly by gravitational radiation. For a loop of invariant length l [14], the oscillation period is $T_l = l/2$ and the lifetime is $\tau_l \sim l/k_g G\mu$. Here, $k_g \sim 50$ is a numerical coefficient, G is the Newton's constant, $\mu \sim \eta^2$ is the mass per unit length of string, and η is the symmetry breaking scale of strings.

The exact value of the parameter α in (1) is not known. We shall assume, following [8], that α is determined by the gravitational backreaction, so that

$$\alpha \sim k_g G\mu. \tag{2}$$

In this case the loops decay within about a Hubble time of their formation, the typical length of loops at time t is $l \sim \alpha t$, and their number density is given by [2]

$$n_l(t) \sim \alpha^{-1} t^{-3}. \tag{3}$$

An electric field E applied along a superconducting string generates an electric current. A superconducting loop of string oscillating in a magnetic field B acts as an *ac* generator and develops an *ac* current of amplitude

$$J_0 \sim e^2 B l. \tag{4}$$

The local value of the current in the loop can be greatly enhanced in near-cusp regions where, for a short period of time, the string reaches a speed very close to the speed of light. Cusps tend to be formed a few times during each oscillation period. Near a cusp, the string gets contracted by a large factor, its rest energy being turned into kinetic energy. The density of charge carriers, and thus the current, are enhanced by the same factor. The contraction factor increases as one approaches the point of the cusp. For a string segment of invariant length $\delta l \ll l$ centered at the cusp, the maximum contraction factor is $\sim l/\delta l$, resulting in a Lorentz factor $\gamma \sim l/\delta l$ and current $J \sim \gamma J_0$. The actual (geometric) length of the segment in its rest frame is $\delta l_r \sim \delta l/\gamma \sim l/\gamma^2$, and this maximum contraction is sustained for a time interval $\delta t_r \sim \delta l_r$ (in the same frame).

The growth of electric current near the cusp due to string contraction is terminated at a critical value J_{max} when the energy of charge carriers becomes comparable to that of the string itself, $(J/e)^2 \sim \mu$. This gives J_{max} and γ_{max} as

$$J_{max} \sim e\eta, \quad \gamma_{max} \sim (e\eta/J_0). \tag{5}$$

Alternatively, the cusp development can be terminated by small-scale wiggles on the string [15]. If the wiggles contribute a fraction $\epsilon \ll 1$ to the total energy of the string, then the maximum Lorentz factor is less than (5), and is given by $\gamma_{max} \sim \epsilon^{-1/2}$. The actual value of γ_{max} is not important for most of the following discussion.

Due to the large current, the cusp produces a powerful pulse of electromagnetic radiation. The total energy of the pulse is given by [4, 5] $\mathcal{E}_{em}^{tot} \sim 2k_{em}J_0J_{max}l$, where $l \sim \alpha t$ is the length of the loop, and the coefficient $k_{em} \sim 10$ is taken from numerical calculations [4]. This radiation is emitted within a very narrow cone of openening angle $\theta_{min} \sim 1/\gamma_{max}$. The angular distribution of radiated energy at larger angles is given by $d\mathcal{E}_{em}/d\Omega \sim 1/\theta^3$ [4]. The radiation emitted between angles θ and 2θ is

$$\mathcal{E}_{em} \sim k_{em}J_0^2 l/\theta, \tag{6}$$

We shall now specify our assumptions about cosmic magnetic fields. We shall assume that magnetic fields were generated in early galaxies by dynamo action at some $z \sim z_B$. A natural possibility is the bright phase of galaxy evolution [11]. The fast evolution of the first generation massive stars is terminated by their fragmentation and by supernova explosions. Large equipartition magnetic fields and powerful galactic winds are expected to be produced in these violent processes [16]. The reduced distance between galaxies at moderately high redshifts facilitates spreading of the magnetic fields throughout the intergalactic space. We shall assume that the magnetic field remains frozen in the plasma, so that

$$B(z) = B_0(1+z)^2, \tag{7}$$

where B_0 is the characteristic field strength at the present time. Eq.(7) applies for $0 \leq z \leq z_B$. For the fraction of the total volume occupied by magnetized plasma we shall use $f_J \sim 0.1$ (f_J also gives the fraction of loops which develop a current). For numerical estimates below we shall use the values $B_0 \sim 10^{-7}\ G$ and $z_B \sim 4$.

We shall now estimate the physical quantities characterizing GRBs powered by superconducting string cusps. The energy output to be observed in a GRB, recalculated for isotropic emission, is for redshift z

$$\mathcal{E}_{iso} \sim \mathcal{E}_{em}/f_B \sim 4k_{em}e^4\alpha^3t_0^3B_0^2(1+z)^{-1/2}\theta^{-3}, \tag{8}$$

where $f_B = \theta^2/4$ is the beaming factor, θ is the angle between the string velocity at the cusp and the direction from the cusp to the observer, and t_0 is the present age of the Universe. The Lorentz factor of the relevant string segment near the cusp is $\gamma \sim 1/\theta$. The duration of the cusp event as seen by a distant observer is [6]

$$\tau_c \sim (1+z)(\alpha t/2)\theta^3 \sim \frac{1}{2}\alpha t_0(1+z)^{-1/2}/\gamma^3. \tag{9}$$

One can expect that the observed duration of GRB is $\tau_{GRB} \sim \tau_c$. This expectation will be justified by the hydrodynamical analysis below.

The fluence, defined as the total energy per unit area of the detector, is [7]

$$S \sim (1+z)\mathcal{E}_{iso}/4\pi d_L^2(z), \tag{10}$$

where $d_L(z) = 3t_0(1+z)^{1/2}[(1+z)^{1/2} - 1]$ is the luminosity distance and we have assumed for simplicity that the universe is spatially flat and is dominated by non-relativistic matter.

The rate of GRBs originating at cusps in the redshift interval dz and seen at an angle θ in the interval $d\theta$ is given by

$$d\dot{N}_{GRB} \sim f_J \cdot \frac{1}{2}\theta d\theta (1+z)^{-1}\nu(z)dV(z). \tag{11}$$

Here, $\nu(t) \sim n_l(t)/T_l \sim 2\alpha^{-2}t^{-4}$ is the number of cusp events per unit spacetime volume, $T_l \sim \alpha t/2$ is the oscillation period of a loop, $dV = 54\pi t_0^3[(1+z)^{1/2} - 1]^2(1+z)^{-11/2}dz$ is the proper volume between redshifts z and $z + dz$, and we have used the relation $dt_0 = (1+z)dt$.

Since different cusp events originate at different redshifts and are seen at different angles, our model automatically gives a distribution of durations and fluences of GRBs. The angle θ is related to the Lorentz factor of the relevant portion of the string as $\theta \sim 1/\gamma$, and from Eqs.(8),(10) we have

$$\gamma(z;S) \sim \gamma_0 \alpha_{-8}^{-1} S_{-8}^{1/3} B_{-7}^{-2/3} [(\sqrt{1+z} - 1)^2 \sqrt{1+z}]^{1/3}. \tag{12}$$

Here, $\gamma_0 \approx 190$, $\alpha_{-8} = \alpha/10^{-8}$, and the fluence S and the magnetic field B_0 are expressed as $S = S_{-8} \cdot 10^{-8}\ erg/cm^2$ and $B = B_{-7} \cdot 10^{-7}\ G$.

Very large values of $\gamma \sim \gamma_{max}$, which correspond (for a given redshift) to largest fluences, may not be seen at all because the radiation is emitted into a too narrow solid angle and the observed rates of these events are too small.

The minimum value $\gamma(z; S_{min})$ is determined by the smallest fluence that is observed, $S_{min} \sim 2 \cdot 10^{-8}\ erg/cm^2$. Another limit on γ, which dominates at small z follows from the condition of compactness and is given by $\gamma > 100$ [17].

The total rate of GRBs with fluence larger than S is obtained by integrating Eq.(11) over z and over θ from γ_{max}^{-1} to $\gamma^{-1}(z;S)$. For relatively small fluences, $S_{-8} < S_c = 0.3(\gamma_{max}\alpha_{-8}/\gamma_0)^3 B_{-7}^2$, the lower limit of θ-integration can be replaced by 0, and we obtain

$$\begin{aligned} \dot{N}_{GRB}(>S) &\sim \frac{f_J}{2\alpha^2 t_0^4}\int_0^{z_B} dV(z)(1+z)^5\gamma^{-2}(z;S) \\ &\sim 3.3 \cdot 10^2 S_{-8}^{-2/3} B_{-7}^{4/3}\ yr^{-1}. \end{aligned} \tag{13}$$

Remarkably, this rate does not depend on any string parameters and is determined (for a given value of S) almost entirely by the magnetic field B_0. The predicted slope $\dot{N}_{GRB}(> S) \propto S^{-2/3}$ is in a reasonable agreement with the observed one $\dot{N}_{obs}(> S) \propto S^{-0.55}$ at relatively small fluences [18].

For large fluences $S_{-8} > S_c$ integration of Eq.(11) gives $\dot{N}_{GRB}(> S) \propto S^{-3/2}$. Observationally, the transition to this regime occurs at $S_{-8} \sim 10^2 - 10^3$. This can be accounted for if the cusp development is terminated by small-scale wiggles with fractional energy in the wiggles $\epsilon \sim 10^{-6}\alpha_{-8}^2 B_{-7}^{4/3}$. Alternatively, if γ_{max} is determined by the back-reaction of the charge carriers, Eq.(5), then the regime (13) holds for larger S_{-8}, and observed steepening of the distribution at large S can be due to the reduced efficiency of BATSE to detection of short bursts. Indeed, most of the bursts in the distribution (13) originate at $z \sim z_B$, and the ones with large S correspond to small τ_{GRB} (see Eq.(14) below).

The duration of a GRBs originating at redshift z and having fluence S is readily calculated as

$$\tau_{GRB} \approx 200 \frac{\alpha_{-8}^4 B_{-7}^2}{S_{-8}} (1+z)^{-1} (\sqrt{1+z} - 1)^{-2} \ s \tag{14}$$

Estimated from Eq.(9), $\tau_{GRB}^{max} \sim 10^3 \alpha_{-8} \ s$, while from Eq.(14) using $S_{max} \sim 1 \cdot 10^{-5} \ erg/cm^2$ and $z \sim z_B \sim 4$, one obtains $\tau_{GRB}^{min} \sim 3\alpha_{-8}^4 B_{-7}^2 \ ms$. This range of τ_{GRB} agrees with observations.

The best fit for the GRB rate, fluence, duration and for the predicted ranges of these quantities corresponds to $\alpha_{-8} = 0.4$ ($i.e.$ $\eta = 1.1 \cdot 10^{14} \ GeV$) and $B_{-7} = 2$.

A few remarks are now in order. *(i)* Small-scale wiggles on the string can naturally produce short-time variation in GRBs. *(ii)* The cusps reappear on a loop, producing nearly identical GRBs with a period $T_l \sim \alpha t_0/2 \sim 20 \ yr$. This is a testable prediction of our model. *(iii)* The loop distribution (3) is based on a simplified model of string evolution, which does not take into account the fragmentation of loops in the process of their decay. Therefore, we expect some corrections to the estimates made above.

Let us now turn to the hydrodynamical phenomena in which the gamma radiation of the burst is actually generated. The low-frequency electromagnetic pulse interacting with surrounding gas produces an ultrarelativistic beam of accelerated particles. This is the dominant channel of energy loss by the pulse. The beam of high energy particles pushes the gas with the frozen magnetic field ahead of it, producing an external shock in surrounding plasma and a reverse shock in the beam material, as in the case of "ordinary" fireball (for a review see [17, 1]). The difference is that the beam propagates with a very large Lorentz factor $\gamma_p >> \gamma_c$ where γ_c is the Lorentz factor of the cusp and γ_p is the mean Lorentz factor of accelerated particles. (The precise value of γ_p is not important for this discussion; it will be estimated in a subsequent publication [12].) Another difference is that the beam propagates in a very low-density gas. The beam can be regarded as a narrow shell of relativistic particles of width $\Delta \sim l/2\gamma_c^3$ in the observer's frame.

The gamma radiation of the burst is produced as synchrotron radiation of electrons accelerated by external and reverse shocks. Naively, the duration of synchrotron radiation, i.e. τ_{GRB}, is determined by the thickness of the shell as $\tau_{GRB} \sim \Delta$. This is confirmed by a more detailed analysis, as follows. The reverse shock in our case is ultrarelativistic [19, 17]. The neccessary condition for that, $\rho_B/\rho_g < \gamma_p^2$, is satisfied with a wide margin (here ρ_B is the baryon density in the beam and ρ_g is the density of unperturbed gas). In this case, the shock dynamics and the GRB duration are determined by two hydrodynamical parameters [17]. They are the thickness of the shell Δ and the Sedov length, defined as the distance travelled by the shell when the mass of the snow-ploughed gas becomes comparable to the initial energy of the beam. The latter is given by $l_{Sed} \sim (\mathcal{E}_{iso}/\rho_g)^{1/3}$.

The reverse shock enters the shell and, as it propagates there, it strongly decelerates the shell. The synchrotron radiation occurs mainly in the shocked regions of the shell and of the external plasma. The surface separating these two regions, the contact discontinuity (CD) surface, propagates with the same velocity as the shocked plasma, where the GRB radiation is produced.

The synchrotron radiation ceases when the reverse shock reaches the inner boundary of the shell. This occurs at a distance $R_\Delta \sim l_{Sed}^{3/4}\Delta^{1/4}$ when the Lorentz factor of the CD surface is $\gamma_{CD} \sim (l_{Sed}/\Delta)^{3/8}$. Note that these values do not depend on the Lorentz factor of the beam γ_p and are determined by the cusp Lorentz factor γ_c through l_{Sed} and Δ. The size of the synchrotron emitting region is of the order R_Δ, and the Lorentz factor of this region is equal to γ_{CD}. Thus, the duration of GRB is given by

$$\tau_{GRB} \sim R_\Delta/2\gamma_{CD}^2 \sim l/2\gamma_c^3, \tag{15}$$

i.e. it is equal to the duration of the cusp event given by Eq.(9). The energy that goes into synchrotron radiation is comparable to the energy of the electromagnetic pulse.

GRBs have been suggested as possible sources of the observed ultrahigh-energy cosmic rays (UHECR) [20, 21, 22]. This idea encounters two difficulties. *(i)* If GRBs are hosted by galaxies, the sources of UHECR are distributed uniformly in the Universe and such distribution is a classical case of Greisen-Zatsepin-Kuzmin (GZK) cutoff. However, this cutoff is absent in UHE particle spectra. *(ii)* The acceleration by an ultrarelativistic shock is possible only in the one-loop regime (*i.e.* due to a single reflection from the shock) [23]. For a standard GRB with a Lorentz factor $\gamma_{sh} \sim 300$ it results in the maximum energy $E_{max} \sim \gamma_{sh}^2 m_p \sim 10^{14}\ eV$, far too low for UHECR.

Our model resolves both of these difficulties.

If the magnetic field in the Local Supercluster (LS) is considerably stronger than it is outside, then the cusps in LS are more powerful sources of UHECR, and the GZK cutoff is less pronounced.

The maximum Lorentz factor at the cusp reaches $\gamma_{cusp} = 6 \cdot 10^7$, Eq.(5), and the maximum energy produced in one-loop acceleration is $E_{max} \sim 10^{24}\ eV$. Particles with these energies are propagating rectilinearly and generally are not seen, while particles with $E \leq 10^{20}\ eV$ are deflected by magnetic field in the halo and can be observed. The total energy released in such events is $\mathcal{E}_{tot} = 2.1 \cdot 10^{46}\ erg$.

Further details of UHECR production in this model will be given elsewhere [12].

We are grateful to Ken Olum for useful discussions. The work of AV was supported in part by the National Science Foundation. The work of BV and BH was supported in part by the INTAS through grant N0. 1065.

References

[1] T. Piran astro-ph/9907392

[2] A. Vilenkin and E.P.S. Shellard, *Cosmic strings and other topological defects*, Cambridge University Press, Cambridge, 1994.

[3] E. Witten, Nucl. Phys. **B249**, 557 (1985).

[4] A. Vilenkin and T. Vachaspati, Phys. Rev. Lett. **58**, 1041 (1987).

[5] D.N. Spergel, T. Piran and J. Goodman, Nucl. Phys. **B291**, 847 (1987).

[6] A. Babul, B. Paczynski and D.N. Spergel, Ap. J. Lett. **316**, L49 (1987).

[7] B. Paczynski, Ap. J. **335**, 525 (1988).

[8] D. Bennett and F. Bouchet, Phys. Rev. Lett, **60**, 257 (1988).

[9] B. Allen and E.P.S. Shellard, Phys. Rev. Lett. **64**, 119 (1990).

[10] P.Meszaros and M.J.Rees, MNRAS, **258**, 41 (1992).

[11] R.B.Partridge and P.J.Peebles, Ap.J, **147**, 868 (1967).

[12] V. Berezinsky, B. Hnatyk and A. Vilenkin, in preparation.

[13] A. Albrecht and N. Turok, Phys. Rev. Lett. **54**, 1868 (1985).

[14] The actual length of the loop changes as the loop oscillates. The invariant length is defined as $l = E/\mu$, where E is the loop's energy in its center-of-mass frame.

[15] B Carter, Phys. Rev. **D41**, 3869 (1990); A. Vilenkin, Phys. Rev. **D41**, 3038 (1990).

[16] J.Schwartz, J.P.Ostriker and A.Yahil, Ap.J. **202**, 1 (1975).

[17] T. Piran astro-ph/9810256

[18] Catalog BATSE 4B, cited in G.Pugliese et al astro-ph/9912019.

[19] S. Kobayashi, T. Piran and R. Sari, astro-ph/9803217

[20] M.Vietri, Ap.J.**453**, 883 (1995).

[21] E.Waxman, Phys. Rev. Lett. **75**, 386 (1995).

[22] M.Milgrom and V.Usov, Ap.J. Lett. **449**, L37 (1995).

[23] Y.A.Gallant and A.Achterberg, astro-ph/9812316