Title

Can a Large Neutron Excess Help Solve the Baryon Loading Problem in Gamma-Ray Burst Fireballs?

Journal

Physical Review Letters, 85(13)

ISSN

0031-9007

Authors

Fuller, George M
Pruet, Jason
Abazajian, Kevork

Publication Date

2000-09-25

DOI

10.1103/physrevlett.85.2673

Peer reviewed

Can a Large Neutron Excess Help Solve the Baryon Loading Problem in Gamma-Ray Burst Fireballs?

George M. Fuller, Jason Pruet, and Kevork Abazajian
Department of Physics, University of California, San Diego, La Jolla, California 92093-0319
(Received 24 April 2000)

We point out that the baryon loading problem in gamma-ray burst (GRB) models can be ameliorated if a significant fraction of the baryons which inertially confine the fireball is converted to neutrons. A high neutron fraction can result in a reduced transfer of energy from relativistic light particles in the fireball to baryons. The energy needed to produce the required relativistic flow in the GRB is consequently reduced, in some cases by orders of magnitude. A high neutron-to-proton ratio has been calculated in neutron star-merger fireball environments. Significant neutron excess also could occur near compact objects with high neutrino fluxes.

PACS numbers: 98.70.Rz, 14.60.Pq

In this Letter we show how the baryon loading problem can be alleviated in certain gamma-ray burst (GRB) models when significant numbers of baryons are converted to neutrons. Interestingly, many of the proposed GRB "central engines" involve compact objects which are themselves highly neutronized, or which are accompanied by intense neutrino fluxes. Weak interactions induced by these neutrino fluxes can result in significant proton-to-neutron conversion, especially if resonant neutrino flavor transformation takes place [1–3].

Inferences of the energetics and spectral observations of GRBs imply (i) total energies in gamma rays approaching 10^{53} ergs for the most energetic events (in the absence of beaming), and (ii) large Lorentz factors of the progenitor fireball ($\gamma \sim 10^3$) (for a recent review, see Ref. [4]). Excessive baryon pollution of the fireball precludes attainment of these features for many GRB models. This is a consequence of the conversion of radiation energy in the electron/positron/photon fireball to kinetic energy in baryons [5,6]. However, the relatively small cross sections characterizing the interactions of neutrons with the electron/positron/photon plasma may afford a solution to this problem.

This can be seen by considering the fictitious limit of completely noninteracting neutrons. Imagine that protons inertially tether an electron/positron/photon fireball via photon Thomson drag on $e^{\pm}$, which in turn influences protons through Coulomb interactions. If these protons were suddenly converted to noninteracting "neutrons," then the fireball would expand relativistically, leaving behind the baryonic component. Real neutrons can approximate this limit as they interact with the electron/positron/photon plasma only via the neutron magnetic dipole moment. These cross sections are small compared to the Thomson cross section σ_{T}: neutron-electron (positron) scattering has $\sigma_{ne} \sim 10^{-7}\sigma_{\mathrm{T}}$ [7]; neutron-photon scattering has $\sigma_{n\gamma} \sim 10^{-12}\sigma_{\mathrm{T}}$ [8].

However, the real limit on the efficacy of this mechanism is the strong interaction neutron-proton scattering which will dominate the energy transfer process when conversion of neutrons to protons is incomplete. Therefore, the degree to which the baryon loading burden can be lifted in our proposed mechanism will depend on the neutron excess in the fireball environment. Here we will measure the neutron content of the plasma in terms of the electron fraction Y_e, the net number of electrons ($n_{e^-} - n_{e^+}$) per baryon, or in terms of the neutron-to-proton ratio $Y_e = 1/(n/p + 1)$.

We note that although previous studies have invoked neutrino oscillations to attempt a baryon loading problem solution [9,10], none has exploited the Y_e-changing aspect of the weak interaction.

To go beyond the simplistic picture of noninteracting neutrons, we can consider a two-component [(i) neutrons, and (ii) protons/$e^{\pm}$/photons] plasma in the context of a homogeneous fireball with initial radius, temperature, Lorentz factor, and electron fraction, R_0, T_0, γ_0, and Y_{e0}, respectively. Numerical and analytic work has shown the following simple scaling laws for such a configuration [6,11]:

$$\text{for } R < \eta R_0/\gamma_0 \Rightarrow \begin{cases} \gamma = \gamma_0(R/R_0)\,, \\ T = T_0(R_0/R)\,, \end{cases}$$
$$\text{for } R > \eta R_0/\gamma_0 \Rightarrow \gamma = \eta\,. \tag{1}$$

As a matter of convenience we will take the scaling to be such that $\gamma_0 = 1$. In Eq. (1), the ratio of energy in radiation E to total baryon rest mass M is $\eta \equiv E/M$.

One can relate η to the entropy per baryon, s, using the number density of baryons $N = \rho_b/m_p = \rho_{\mathrm{rad}}/m_p\eta$ where ρ_b is the baryon component rest mass energy density and m_p is the proton rest mass. Using this relation, and noting that in terms of the proper entropy density S, the entropy-per-baryon is $s = S/N$, the relation between s, η, and the temperature T_0 is $s \approx 1250\eta\ (1\ \mathrm{MeV}/T_0)$.

For a large enough η, baryon loading is unimportant [5]. In fact, when $\eta \gtrsim 10^5(E/10^{52}\ \mathrm{erg})^{1/3}(10^7\ \mathrm{cm}/R_0)^{2/3}$, the fireball becomes optically thin before transferring its energy to kinetic energy in baryons (here E is the total energy of the fireball).

Written in terms of time t as measured in a frame comoving with the fireball the above relations imply

$$R = R_0 e^{t/\tau_{\rm dyn}}, \quad \gamma = e^{t/\tau_{\rm dyn}}, \quad T = T_0 e^{-t/\tau_{\rm dyn}}, \qquad (2)$$

for $\gamma < \eta$. Here the dynamic time scale is defined to be the initial light crossing time, $\tau_{\rm dyn} \equiv R_0/c$. In fireballs resulting from neutron star mergers, for example, $\tau_{\rm dyn} \sim 2.7 \times 10^{-5}$ s, corresponding to an R_0 of 8 km [12].

A particle comoving with the expanding plasma experiences a 4-acceleration a^μ, with magnitude $\sqrt{a^\mu a_\mu} = d\gamma/dR$. As noted above, the force that drags the neutrons along with the expanding plasma arises principally from n-p collisions. The relative contribution to the total force on the neutrons from collisions with electrons and positrons is roughly $F_{n-e}/F_{n-p} \sim m_e n_e \sigma_{ne}/m_p n_p \sigma_{np} \leq 10^{-10}(s/Y_e)$ and is small for the conditions we consider. The neutron-photon cross section is small enough ($\sigma_{n-\gamma} \sim 10^{-36}[E_\gamma/(1\ \text{MeV})]^2\ \text{cm}^2$ where E_γ is the photon energy in the neutron rest frame [8]) that n-γ interactions are negligible.

The relations in Eqs. (1) imply that an inertial observer with time coordinate t' initially (at $t' = 0$) comoving with the plasma sees the plasma accelerate according to $\gamma v = t'/\tau_{\rm dyn}$. Hereafter we adopt natural units where $c = 1$. If we denote by $\tau_{\rm coll}^{-1}$ the frequency of neutron/proton collisions (per neutron), we expect that the two components of the plasma will achieve a relative velocity given by $v_{\rm rel} \approx 2\tau_{\rm coll}/\tau_{\rm dyn}$, where the factor of 2 arises from the approximate angle independence of the neutron-proton scattering cross section and the near equality of the neutron and proton masses. An equivalent expression is found if one considers the force on the neutrons from collisions with protons [13]. It is clear then that when $\tau_{\rm dyn} \gg \tau_{\rm coll}$ the neutrons are coupled to the rest of the plasma. However, decoupling occurs as these two time scales become comparable. Since the baryon number density in the plasma frame decreases as $e^{-3t/\tau_{\rm dyn}}$, decoupling will occur quickly, i.e., on a time scale shorter than $\tau_{\rm dyn}$.

When significant decoupling occurs we can neglect the thermal contribution to the collision frequency and write

$$\tau_{\rm coll}^{-1} \approx (9 \times 10^{12}\ \text{s}^{-1})\, \frac{T_{\rm MeV}^3}{s_5}\, v_{\rm rel} Y_e \sigma_{10}\,, \qquad (3)$$

where $s_5 \equiv s/10^5$ and σ_{10} is the neutron-proton cross section in units of 10 fm^2, and $T_{\rm MeV}$ is the temperature in MeV. As the precise energy dependence of σ_{10} is not important here, it suffices to note a few representative values: $\sigma_{10}(v_{\rm rel} = 0.1) \approx 17$, $\sigma_{10}(v_{\rm rel} = 0.3) \approx 2$, and $\sigma_{10}(v_{\rm rel} = 0.6) \approx 0.4$ [14].

The requirement of a non-negligible relative velocity then gives the decoupling time $t = t_{\rm dec}$ as

$$\frac{t_{\rm dec}}{\tau_{\rm dyn}} = 4.6 + (1/3)\ln\!\left(\frac{\sigma_{10}\tau_{-6}T_{0,{\rm MeV}}^3 Y_e}{s_5}\right). \qquad (4)$$

In the above, $\tau_{-6} \equiv \tau_{\rm dyn}/10^{-6}$ s, and $\tau_{\rm coll}$ was evaluated at a terminal velocity of 0.5. The calculated decoupling time is logarithmically sensitive to this choice. In reality the neutrons do not sharply decouple but continue to interact with the plasma over roughly a dynamical time scale. In this sense, the $t_{\rm dec}$ appearing in Eq. (4) is an "effective" decoupling time. An accurate determination of $t_{\rm dec}$ requires solving in detail the neutron and proton transport equations. However, because of the exponential decrease of density with time in the plasma frame, the number 4.6 appearing in Eq. (4) is only uncertain to approximately $\pm(1/3)$.

Once the neutrons decouple they will have an energy $\gamma_{\rm dec}(1 - Y_e)M$. The ratio of kinetic energy in neutrons to the total energy in the fireball is then

$$f_n \approx (1 - Y_e)\,\frac{e^{t_{\rm dec}/\tau_{\rm dyn}}}{\eta} \approx 1.3(1 - Y_e)\left(\frac{Y_e\sigma_{10}\tau_{-6}}{s_5^4}\right)^{1/3}. \qquad (5)$$

(Here "total" energy includes both the thermal $e^\pm/\gamma$ energy and the bulk kinetic energy of baryons.) From this we see that for Y_e less than

$$Y_{e,{\rm crit}} \equiv 0.46\left(\frac{s_5^4}{\sigma_{10}\tau_{-6}}\right) \qquad (6)$$

at the time of decoupling the baryon loading problem is diminished. (Protons and neutrons each move with $\gamma_{\rm dec}$ at the decoupling point; thereafter, protons will possess larger Lorentz factors than do average neutrons.) If the plasma remains optically thick to radiation until the energy in radiation is converted to kinetic energy of the remaining protons, energy conservation gives the final Lorentz factor of the protons

$$\gamma \approx \eta\left(\frac{1 - f_n}{Y_e}\right). \qquad (7)$$

A simple ansatz for the condition that the fireball remains optically thick after decoupling is $\eta \lesssim (Y_e)10^5 \times (E/10^{52}\ \text{erg})^{1/3}(10^7\ \text{cm}/R_0)^{2/3}$. This is obtained by applying the result from [5] and making the replacements $\eta \to (\eta/Y_e)(1 - f_n)$, $E \to (1 - f_n)E$, and $s \to s/Y_e$. Note that even for modest values of f_n, γ can be increased significantly if Y_e is low.

The above results are summarized in Fig. 1, where we have plotted the smallest entropy s_5 ($\approx$1.25 η_{100}/T_0) for which decoupling occurs as a function of Y_e. For example, if $s_5 = 0.6, T_0 = 2$ MeV (corresponding to $\eta \approx 100$), and $Y_{e0} = 0.02$, then the final Lorentz factor of the plasma after neutron decoupling [Eqs. (5) and (7)] would be $\gamma \approx 1500$, which is 15 times larger than the standard case of $\gamma = \eta$. As another example, consider the Ref. [12] values of $\tau_{-6} = 27$ and $Y_e = 0.1$ and suppose that $T_0 = 10$ MeV and $s_5 = 2.5$ (corresponding to $\eta = 2000$). In this case we find $\gamma = 1.1 \times 10^4$, an increase by a factor of 5.7. Clearly, the importance of this effect depends on how low Y_e can be.

Two conditions must be met in order to achieve a low Y_e at the time of decoupling: (i) Y_e must be low initially and (ii) Y_e must not be unacceptably raised during the evolution of the fireball. We can divide up the discussion of Y_e

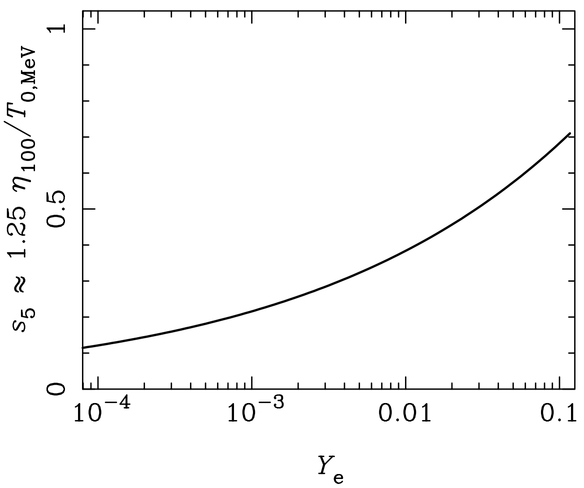

FIG. 1. The smallest entropy per baryon, $s_5 \equiv s/10^5$, for which decoupling occurs as a function of Y_e. We have taken $\tau_{-6}\sigma_{10} = 1$.

in this way because the initial electron fraction depends in detail on the GRB central engine, whereas the later evolution of the fireball is generically given by the relations in Eq. (2). Many proposed GRB central engines involve neutrino heating or are sited in environments subject to intense neutrino fluxes [12,15–22]. General discussions of the relation between neutrino processes and the dynamics of outflow may be found in Refs. [12,23–25]. However, the details of neutron decoupling are insensitive to how Y_e is set and we are not arguing for a specific GRB site.

The processes which have a significant effect on Y_e in the fireball environment are lepton capture/decay involving free nucleons and inelastic $nn \rightarrow np\pi$ scattering (charged pion-nucleon bremsstrahlung),

$$\nu_e + n \rightleftharpoons p + e^-, \tag{8a}$$

$$\bar{\nu}_e + p \rightleftharpoons n + e^+, \tag{8b}$$

$$n \rightarrow p + e^- + \bar{\nu}_e\,, \tag{8c}$$

$$n + n \rightarrow n + p + \pi^-. \tag{8d}$$

In general, Y_e is set by the competition between the above processes [1,26]. For the range of fireball parameters of interest to us, free neutron decay (8c) is unimportant as the fraction of neutrons decaying during the evolution of the fireball is $\sim 10^{-9}\tau_{-6}\ln\eta$. Furthermore, as lepton capture is important only during the early, hot, evolution of the fireball and inelastic nucleon-nucleon scattering occurs only after neutron decoupling, the lepton capture and pion bremsstrahlung processes may be considered separately.

In environments where neutrino heating is important the forward reactions (8a) and (8b) can dominate in setting the electron fraction [1]. Integration of the rate equations corresponding to the lepton capture processes gives $n/p \approx \lambda_{\bar{\nu}_e p}/\lambda_{\nu_e n} \approx (L_{\bar{\nu}_e}\langle E_{\bar{\nu}_e}\rangle)/(L_{\nu_e}\langle E_{\nu_e}\rangle)$, where $\lambda_{\bar{\nu}_e p}$ and $\lambda_{\nu_e n}$ are the rates for the reactions in Eqs. (8a) and (8b), $\langle E_{\bar{\nu}_e}\rangle$ and $\langle E_{\nu_e}\rangle$ are the average energies characterizing the energy spectra of the $\bar{\nu}_e$ and ν_e neutrinos, respectively, while $L_{\bar{\nu}_e}$ and L_{ν_e} are the corresponding energy luminosities. Absent neutrino oscillations and flavor/type mixings, any thermal neutrino emission scenario from a compact object, will yield a characteristic average neutrino energy heirarchy for solar mass scale objects: $\langle E_{\nu_\mu}\rangle \approx \langle E_{\bar{\nu}_\mu}\rangle \approx \langle E_{\nu_\tau}\rangle \approx \langle E_{\bar{\nu}_\tau}\rangle > \langle E_{\nu_{\bar{e}}}\rangle > \langle E_{\nu_e}\rangle$. These considerations are consistent with findings in Ref. [12] in which a hard $\bar{\nu}_e$ spectrum from a collapsing neutron star leads to an electron fraction in the fireball of $Y_e \sim 0.1$.

If the ν_e component of the neutrino emission were to disappear or be greatly reduced, then the competition inherent in the above equations would be unbalanced in favor of the reaction $\bar{\nu}_e + p \rightarrow n + e^+$. This, in turn, would result in the wholesale production of neutrons. In fact, several schemes involving matter-enhanced active-sterile neutrino transformation have been proposed as a way of enabling r-process nucleosynthesis in neutrino-heated supernova ejecta: one of these involves matter-enhanced $\nu_e \rightleftharpoons \nu_s$ and $\bar{\nu}_e \rightleftharpoons \bar{\nu}_s$ [2]; the other involves matter-enhanced conversion $\nu_{\mu,\tau} \rightleftharpoons \nu_s$ followed by an active-active matter-enhanced conversion $\nu_{\mu,\tau} \rightleftharpoons \nu_e$ [3]. In either case, the intial ν_e flux can be reduced by more than an order of magnitude and, in turn, this can translate into a substantial decrease in the initial Y_e. (Just how low depends on central engine outflow hydrodynamics and on neutrino background effects [3,26–28].)

If we demand that an initially low Y_e not be raised above $Y_{e,\mathrm{crit}}$, consideration of lepton capture on neutrons allows us to place rough constraints on the fireball and neutrino parameters. We incorporate the uncertainty in the initial fireball evolution by supposing that the relations in Eq. (2) are valid only after the fireball has a Lorentz factor γ_i and temperature T_i. Consideration of positron capture after $\gamma = \gamma_i, T = T_i$ then leads to $T_i < (22\ \mathrm{MeV}) \times (Y_{e,\mathrm{crit}}/\tau_{-6})^{1/5}$. Similarly, consideration of ν_e capture on neutrons leads to $T_{\nu_e} < (40\ \mathrm{MeV})\gamma_i(Y_{e,\mathrm{crit}}/\tau_{-6})^{1/5}$. In deriving this limit above we have taken the ν_e spectrum to be a Fermi-Dirac blackbody with temperature T_{ν_e} and zero chemical potential. This limit could be modified or weakened if ν_e flavor transformation occurs.

Determining the increase in Y_e due to pion production requires a proper treatment of neutron transport in the plasma. However, an upper limit on the increase is readily obtained by considering the extreme case where (i) the protons are frozen into the accelerating plasma, (ii) nonforward n-p collisions are assumed to result in maximal momentum exchange, (iii) n-n collisions are ignored except as a postprocessing step to determine π production, and (iv) the change in Y_e due to inelastic n-p and inelastic p-p scatterings is ignored. This simple picture gives an upper limit on the increase in Y_e because an exchange of any of the assumptions (i)–(iii) for more realistic ones has the effect of decreasing the velocity dispersion of the neutrons. By a calculation with the above assumptions we obtain the upper limit on the increase in Y_e to be $\Delta Y_e \lesssim 10^{-3}/Y_{e0}$, where Y_{e0} is the initial electron fraction. Figure 2 displays the evolution of the neutron distribution function as

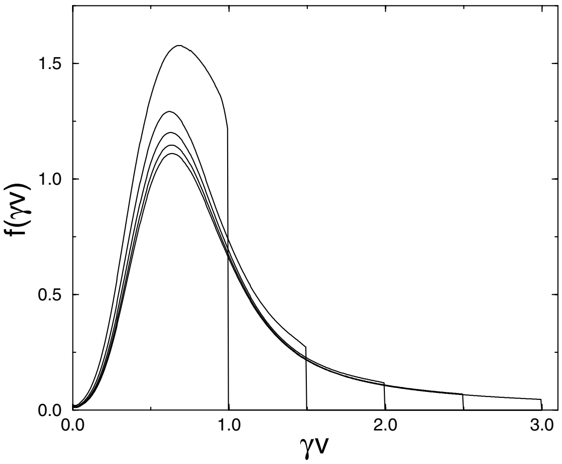

FIG. 2. The neutron distribution function $f(\gamma v)$ [normalized so that $\int f d(\gamma v) = 1$] at several time slices as calculated under the assumptions given in the text. The different time slices correspond to a plasma γv of 1, 1.5, 2, 2.5, and 3, as measured in a frame comoving with the plasma near the decoupling point.

calculated with the above assumptions. The distribution function drops sharply at the instantaneous plasma velocity because there is no mechanism for boosting neutrons to higher velocity.

The ineffectiveness of pion bremsstrahlung in increasing Y_e may be attributed to the fact that the processes which increase Y_e result from a two step process (i.e., an n-p scattering boosts a neutron which then inelastically scatters with another neutron), the decrease of $\sigma(v)v$ with increasing velocity (the product of cross section and relative velocity decreases by a factor of 5 as v increases from 0 to c [14]), and the late onset of this process in n-n scattering. Pion production does not begin until the pion mass threshold is reached at a relative velocity of 0.645 and even at a relative velocity of 0.728 (center of mass energy 2.08 GeV), the inelastic contribution to the cross section is only 8% of the total [14]. For $\Delta Y_e > Y_{e0}$ our perturbative approach to calculating ΔY_e breaks down and the increase in Y_e may cause a recoupling of the proton and neutron flows. Note that our calculation only gives an upper limit on the increase in Y_e due to inelastic n-n scattering. A more careful transport calculation will likely show a smaller increase in Y_e. It is consistent then to discuss decoupling at low Y_e for a wide range of fireball parameters. The final Lorentz factor of the plasma may then be substantially increased for a given energy input and baryon load.

Aside from the fireball energetics versus Y_e issue addressed here, neutron-proton separation recently has been shown to have implications for the electromagnetic and high energy neutrino signatures of GRBs [13,29,30] and may well have implications for afterglow energetics. Future large volume detectors such as AMANDA/ICECUBE will be able to provide high energy neutrino data on GRBs [31]. Perhaps the details of neutron decoupling and the associated electromagnetic/neutrino signature could allow a diagnostic of the weak interaction physics deep in GRB central engine environments.

It is a pleasure to acknowledge discussions with N. Dalal, M. Patel, X. Shi, and J. R. Wilson. This research was supported in part by NSF Grant No. PHY98-00980, an IGPP grant, and a NASA GSRP for K. A.

[1] Y.-Z. Qian, G. M. Fuller, G. J. Mathews, R. W. Mayle, J. R. Wilson, and S. E. Woosley, Phys. Rev. Lett. **71**, 1965 (1993).
[2] D. O. Caldwell, G. M. Fuller, and Y.-Z. Qian, Phys. Rev. D **61**, 123005 (2000).
[3] G. C. McLaughlin, J. M. Fetter, A. B. Balantekin, and G. M. Fuller, Phys. Rev. C **59**, 2873 (1999).
[4] T. Piran, Phys. Rep. **314**, 575 (1999).
[5] A. Shemi and T. Piran, Astrophys. J. **365**, L55 (1990).
[6] S. Kobayashi, T. Piran, and R. Sari, Astrophys. J. **513**, 669 (1999).
[7] J. H. Applegate, C. J. Hogan, and R. J. Scherrer, Phys. Rev. D **35**, 1151 (1987).
[8] R. J. Gould, Astrophys. J. **417**, 12 (1993).
[9] W. Kluzniak, Astrophys. J. **508**, L29 (1998).
[10] R. R. Volkas and Y. Y. Y. Wong, Astropart. Phys. **13**, 21 (2000).
[11] T. Piran, A. Shemi, and R. Narayan, Mon. Not. R. Astron. Soc. **263**, 861 (1993).
[12] J. D. Salmonson, J. R. Wilson, and G. J. Mathews, astro-ph/0002312.
[13] E. V. Derishev, V. V. Kocharovsky, and Vl. V. Kocharovsky, Astrophys. J. **521**, 640 (1999).
[14] The Particle Data Group, C. Caso *et al.,* Eur. Phys. J. C **3**, 1 (1998).
[15] B. Paczynski, Astrophys. J. **363**, 218 (1990).
[16] J. Goodman, Astrophys. J. **308**, L47 (1986).
[17] B. Paczynski, Astrophys. J. **308**, L43 (1986).
[18] P. Meszaros and M. Rees, Astrophys. J. **397**, 570 (1992).
[19] S. E. Woosley, Astrophys. J. **405**, 273 (1993).
[20] A. MacFadyen and S. E. Woosley, Astrophys. J. **524**, 262 (1999).
[21] G. E. Brown, C.-H. Lee, R. A. M. J. Wijers, H. K. Lee, G. Israelian, and H. A. Bethe, astro-ph/0003361.
[22] G. M. Fuller and X. Shi, Astrophys. J. **502**, L5 (1998).
[23] R. C. Duncan, S. L. Shapiro, and I. Wasserman, Astrophys. J. **309**, 141 (1986).
[24] Y.-Z. Qian and S. E. Woosley, Astrophys. J. **471**, 331 (1996).
[25] C. Y. Cardall and G. M. Fuller, Astrophys. J. **486**, L111 (1997).
[26] Y.-Z. Qian and G. M. Fuller, Phys. Rev. D **51**, 1479 (1995).
[27] M. Patel and G. M. Fuller, in *Proceedings of the American Physical Society (APS) Meeting of the Division of Particles and Fields (DPF 99), Los Angeles, California* (University of California, Los Angeles, 1999), http://www.dpf99.library.ucla.edu/, hep-ph/0003034.
[28] G. C. McLaughlin, G. M. Fuller, and J. R. Wilson, Astrophys. J. **472**, 440 (1996).
[29] E. Waxman and J. N. Bahcall, Phys. Rev. Lett. **78**, 2292 (1997).
[30] J. N. Bahcall and P. Meszaros, Phys. Rev. Lett. (to be published).
[31] F. Halzen and D. W. Hooper, Astrophys. J. **527**, L93 (1999).