

NEUTRINO OSCILLATIONS AND GAMMA-RAY BURSTS

W. KLUŻNIAK

Copernicus Astronomical Center, ul. Bartycka 18, PL-00-716, Warszawa, Poland; and University of Wisconsin at Madison, Sterling Hall, 1150 University Avenue, Madison, WI 53706; wlodek@astro.physics.wisc.edu

Received 1998 July 22; accepted 1998 August 31; published 1998 October 12

ABSTRACT

If the ordinary neutrinos oscillate into a sterile flavor in a manner consistent with the Super-Kamiokande data on the zenith-angle dependence of atmospheric μ -neutrino flux, an energy sufficient to power a typical cosmic gamma-ray burst (GRB) ($\sim 10^{52}$ ergs) can be carried away from the source by sterile neutrinos and deposited in a region relatively free of baryons. Hence, ultrarelativistic bulk motion (required by the theory and observations of GRBs and their afterglows) can be achieved in the vicinity of plausible sources of GRBs. Oscillations between sterile and ordinary neutrinos might, therefore, provide a solution to the “baryon-loading problem” in the theory of GRBs.

Subject headings: gamma rays: bursts — elementary particles — radiation mechanisms: nonthermal

1. ULTRARELATIVISTIC BULK MOTION AND THE BARYON-LOADING PROBLEM

Available evidence clearly shows that gamma-ray bursts (GRBs) originate in distant reaches of the universe (Fishman & Meegan 1995), with redshifts $z > 0.8$ and $z = 3.4$ reported for two particular sources (Djorgovski et al. 1997; Kulkarni et al. 1998).

Ultrarelativistic bulk motion of matter in the sources is required to account for the observed fluences ($10^{-5 \pm 2}$ ergs cm^{-2}) in cosmic GRBs originating at a redshift $z \sim 1$ (Mészáros & Rees 1993). Otherwise, if the source were sufficiently large to be optically thin to electron pair creation, the typical duration of a GRB would be days rather than the observed seconds or minutes, while a source only light minutes (or less) in size would be opaque to its own gamma-ray radiation. Considerations of such bulk motion led to predictions (Paczynski & Rhoads 1993; Katz 1994; Mészáros & Rees 1997; Vietri 1997) of radio, optical, and X-ray afterglows, which were subsequently observed (Galama et al. 1997; van Paradijs et al. 1997; Frail et al. 1997). Most sources capable of impulsively releasing the 10^{52} ergs or more of energy required to power a GRB, however, contain so much matter around them that if the energy released were used to accelerate even a very small fraction ($\sim 10^{-3}$) of the baryons present, only a nonrelativistic wind would result. This is known as the baryon-loading problem (see, e.g., Piran 1997).

It has been hoped that the geometry of the sources is such that at least some of the energy released is channeled along directions relatively free of baryons, so that relativistic bulk motion and the ensuing beaming of radiation may occur along certain lines of sight. So far, this has not yet been fully demonstrated for any theoretical source of GRBs—although calculations of binary “mergers” of a neutron star with another neutron star or a black hole suggest that this may be the case (Rasio & Shapiro 1992; Kluźniak & Lee 1998; Wilson, Salmonson, & Mathews 1998); for other models, even approximate calculations of the spatial distribution of baryons are not available (Woosley 1993; Paczyński 1998).

The diversity of the observed GRBs is such that it is not at all clear that they all arise in sources of identical, or even similar, geometry. One is tempted to search for a solution to the baryon-loading problem that does not invoke a very special geometry. The transport of energy by noninteracting particles

would provide a natural solution. Sterile neutrinos, if they exist, would be such particles.

2. SUPER-KAMIOKANDE RESULT

The Super-Kamiokande collaboration has found strong evidence for a zenith angle dependence of the flux ratio of atmospheric muon and electron neutrinos, ν_μ/ν_e , in a 535 days (33.0 kiloton yr) observation of Cerenkov photons in a water detector (Fukuda et al. 1998). Oscillation between the muon and the electron neutrino is not favored by the data. The data is well fitted by a two-neutrino model of vacuum oscillations of the μ -neutrino “ $\nu_\mu \longleftrightarrow \nu_x$, where ν_x may be ν_τ or a new, noninteracting ‘sterile’ neutrino” (Fukuda et al. 1998). The probability for a flavor change in travel over a distance L is given by

$$P_{a \rightarrow b} = \sin^2 2\theta \sin^2 \left[\frac{1.27 \Delta m^2 (\text{eV})^{-2} L (\text{km})^{-1}}{E (\text{GeV})^{-1}} \right],$$

with the best-fit values inside the physical region ($0 \leq \sin^2 2\theta \leq 1$) given as $\sin^2 2\theta = 1.0$, and $\Delta m^2 = 2.2 \times 10^{-3} (\text{eV})^2$. At the 90% confidence level, $5 \times 10^{-4} < \Delta m^2 < 6 \times 10^{-3} (\text{eV})^2$, and $\sin^2 2\theta > 0.82$ (Fukuda et al. 1998).

If the oscillation is between ν_τ and ν_μ , no dramatic changes to GRB models would be called for since both neutrinos interact with matter through neutral currents. In the remainder of this Letter, I will assume that the detected oscillation is into a sterile flavor that does not interact with matter. Some features of the data suggest that this may indeed be the case. Note that a μ -neutrino of energy E would change flavor, with probability one, from μ to sterile every time it traverses a distance

$$L_0 = \frac{\pi(E/\text{GeV})}{2.54 \Delta m^2 / (\text{eV})^2} \text{ km}, \quad (1a)$$

and the sterile neutrino would revert back to the μ flavor after traveling the same distance again. For an $\sim \text{GeV}$ neutrino, $L_0 \sim 10^3 \text{ km}$, but for $E \sim \text{MeV}$ (which may be the more likely

value in GRB sources), the distance is on the order of 1 km:

$$0.2 \left(\frac{E}{\text{MeV}} \right) < \left(\frac{L_0}{\text{km}} \right) < 2.5 \left(\frac{E}{\text{MeV}} \right), \quad (1b)$$

where the limits correspond to the 90% confidence-level limits on Δm^2 . This is a length scale of some interest in studies of neutron stars (radius $R_* \sim 10$ km) and solar-mass black holes (Schwarzschild radius $r_s = 2.9 \text{ km} \times M/M_\odot$), and, by extension, also in theoretical discussions of GRB models.

3. HYPERNOVA-TYPE MODELS OF GRBs

There is a class of models in which GRBs are ultimately powered by neutrino emission in the gravitational collapse of Earth-sized (but massive) objects. These could be the cores ($M \gg M_\odot$) of very massive stars (Woosley 1993; Paczyński 1998), Type Ic supernovae (possibly as in SN 1988bw; Iwamoto et al. 1998; Wang & Wheeler 1998), or an unstable supermassive ($10^5 M_\odot$) quasi star (Fuller & Shi 1998), etc. While a sufficiently high energy input (up to $\sim 10^{54}$ ergs) in neutrinos or in Poynting flux can be obtained (in these models) in a natural way, it has never been demonstrated how to avoid the baryon-loading problem in these cases.

One may expect a significant fraction of a solar mass in baryonic matter distributed in a “mantle” of at least a 10^4 km extension in every direction in these models. A fraction (up to a few percent) of the released energy that is sufficiently high to power the observed GRBs can be deposited above the mantle if a comparable energy is converted into GeV or TeV neutrinos. By analogy, with active galactic nuclei, or Galactic “mini-quasars” (thought to be stellar mass binary systems), one would imagine the formation of jets in the system and then shock acceleration of protons, which could photoproduce the pions decaying into neutrinos.

For ordinary quasars, μ -neutrino luminosities as high as 20% of the quasar luminosity have been predicted (Stecker 1993). If a comparable fraction of the luminosity of the GRB central engine (cloaked in the mantle) were converted to ν_μ 's, many of these could penetrate the mantle after oscillating into the sterile flavor, oscillate back into μ -neutrinos above the mantle (in agreement with eqs. [1a] and [1b]), and then deposit their energy by electron pair creation and scattering on electrons in a relatively baryon-free region, thus creating the type of ultrarelativistic outflow discussed in GRB scenarios.

4. GRB MODELS WITH SOLAR-MASS CENTRAL ENGINES

In another class of models, GRBs are powered by a compact object of approximately $1 M_\odot$. This class includes double neutron star systems in the final stages of binary coalescence, or the interaction of a neutron star with a black hole in a coalescing binary; neutron stars with millisecond rotation periods; neutron stars suffering a phase transition, perhaps to a strange star; and many variations of the above (Latimer & Schramm 1976; Paczyński 1986, 1991; Eichler et al. 1988; Usov 1992; Mochkovitch et al. 1993; Cheng & Dai 1996; Kluźniak & Ruderman 1998; Kluźniak & Lee 1998). Frequently the discussed configuration is that of a solar-mass disk orbiting a stellar-mass black hole. In all these cases, the typical length scale for the distribution of baryons is a few kilometers, and up to few times 10^{53} ergs of energy can be released. In most (but not all) cases, thermal emission of neutrinos is expected, necessarily with energies of a few MeV; neutrino emission in such models bears similarity to that in supernovae. Whether the expected lumi-

nosity of neutrinos is sufficient to power a GRB is a subject of some controversy, but assuming that it is sufficient, the baryon-loading problem is almost always a serious concern.

Clearly, by equations (1a) and (1b), oscillation over a distance comparable to the size of the system is expected, which would be helpful in circumventing the baryon-loading problem since the sterile neutrinos could freely penetrate baryonic matter and even overtake any nonrelativistic wind, and they could deposit energy in a relatively baryon-free region after oscillating back into μ -neutrinos.

5. GEOMETRY OF NEUTRINO ANNIHILATION

Since the ultrarelativistic blast wave would be powered by annihilation of μ -neutrinos into electron-positron pairs, it is necessary to discuss the rate of energy deposition in this process. The overall efficiency of energy conversion into gamma rays has to be rather high in most scenarios, so if sterile (or any) neutrinos are to play a role in the energy transport it helps if the geometry of neutrino annihilation (and simultaneous electron pair creation) is favorable.

In spherical geometry, the luminosity (per unit volume) of neutrino annihilation drops like the radius to the eighth power, $L_{\nu\nu}(r) \propto r^{-8}$. In the discussion above, the neutrino oscillation was taken to occur over a distance comparable to the system size; this would correspond to an efficiency of energy transport by the sterile neutrinos of $\approx 2^{-5} = 0.03$ (this is the ratio of energy deposited in electron-positron pairs at $r > 2r_0$ to that which would be deposited in the absence of baryons and neutrino oscillations at $r > r_0$). However, it is agreed by most authors that angular momentum in the system would lead to significant departures from spherical symmetry. Indeed, disklike or toroidal geometry seems preferred, and it has been pointed out that such geometry is more favorable for annihilation when no baryons are present (Mészáros & Rees 1992). I will show that in disklike geometry, when baryons are present through which sterile neutrinos transport energy away from the source over a distance comparable to the oscillation length, the efficiency is a few percent; i.e., it is comparable to that in the spherical case. In that case, toroidal geometry is much more favorable, and neutrino annihilation could be further enhanced by gravitational focusing.

Neglecting gravitational effects (redshift and trajectory bending), the energy deposition rate per unit volume through neutrino annihilation is (Goodman, Dar, & Nussinov 1987)

$$\dot{q}(\mathbf{r}) = \frac{4KG_F^2\chi(\mathbf{r})}{3A_0^2c} L_\nu L_{\bar{\nu}} \left(\frac{\langle E_\nu^2 \rangle}{\langle E_\nu \rangle} + \frac{\langle E_{\bar{\nu}}^2 \rangle}{\langle E_{\bar{\nu}} \rangle} \right), \quad (2)$$

where L_ν/A_0 ($L_{\bar{\nu}}/A_0$) are the surface emissivities of neutrinos (antineutrinos), $K = (1 - 4 \sin^2 \theta_w + 8 \sin^4 \theta_w)/(6\pi) = 0.027$, $G_F^2 = 5.3 \times 10^{-44} \text{ cm}^2 \text{ MeV}^{-2}$, and c is the speed of light. The expressions in angle brackets are moments of (anti)neutrino energy. Here χ is a position-dependent geometrical factor.

If the neutrinos and antineutrinos are emitted isotropically from the surface of a sphere (the same for both) of radius r_0 , L_ν is the neutrino luminosity of that sphere of area $A_0 = 4\pi r_0^2$, while $\chi(\mathbf{r}) = \chi_s(x)$, where $\chi_s(x) \equiv (1-x)^4(x^2 + 4x + 5)$, and $x = (1 - r_0^2/r^2)^{1/2}$, with r being the spherical radial coordinate (Goodman et al. 1987). This gives the r^{-8} dependence of luminosity mentioned above: $\chi(\mathbf{r}) \approx (5/8)r_0^8/r^8$ for $r^2 \gg r_0^2$. Equation (2) has been derived in the approximation that the neutrinos are sufficiently energetic for threshold effects to be important only over a negligible solid angle (i.e., $\langle E_\nu \rangle \gg 511 \text{ keV}$).

The energy deposition rate at point P on the z -axis above some azimuthally symmetric surface emitting neutrinos isotropically with surface emissivity L_s/A_0 is given by equation (2), with $\chi = \chi_s(x_2) - \chi_s(x_1)$, when the surface is visible from the point P between azimuthal angles $\arccos x_1$ and $\arccos x_2$. Thus, for an infinite disk, $\chi = 5$. There is a slight complication if some of the neutrinos are sterile and do not contribute to the annihilation rate. As an example, let us consider a disk isotropically emitting from its surface sterile neutrinos that oscillate back into the μ -flavor upon traveling a distance L_0 , as given by equations (1a) and (1b). At point P , a height h above the surface, only neutrinos arriving with azimuthal angle θ satisfying $h/[(2n-1)L_0] < \cos \theta < h/[(2n)L_0]$, with $n = 1, 2, 3, \dots$, are μ -neutrinos and hence contribute to \dot{q} . Thus, the energy deposition rate is given by equation (2) with

$$\begin{aligned} \chi &= -\chi_s\left(\frac{h}{L_0}\right) + \chi_s\left(\frac{h}{2L_0}\right) - \chi_s\left(\frac{h}{3L_0}\right) + \dots \\ &= 16\left(\frac{h}{L_0}\right) \ln 2 - \frac{5(\pi h/L_0)^2}{4} + \frac{7(\pi h/L_0)^4}{144} - \frac{31(\pi h/L_0)^6}{6 \times 7!} \end{aligned}$$

for $0 < h \leq L_0$. For example, at height L_0 above the disk, the presence of sterile neutrinos reduces the energy deposition rate by a factor of about 2: $\chi = 2.503$ for $h = L_0$.

To estimate the efficiency of energy transport through a baryonic mantle, let us consider a disk with no baryons outside emitting muon neutrinos and antineutrinos, and compute the energy deposition in electron positron pairs up to a height L_0 above the disk; this is obtained by integrating equation (2) with $\chi = 5$ over this external volume. Now let us consider an identical disk emitting sterile neutrinos and antineutrinos at the same luminosity as before, and let us assume that the disk is cloaked in baryons, which extend to height $L_0/2$ above the disk. Taking into account only those neutrinos that pass through the mantle only in the sterile flavor, let us compute the energy deposition in electron pairs above the baryons, up to the same height as before, and take the ratio of the two results. The answer is 0.026, and I conclude that the efficiency of energy transport by sterile neutrinos in disklike geometry is similar to that in spherical geometry, on the order of a few percent. This is true for the following reason: neutrinos traveling sharply away from the normal to the disk surface oscillate back into the interacting flavor inside the baryonic mantle, while in the spherical case, the mean collision angle of the neutrinos drops sharply with distance.

For a torus cloaked in baryons to a depth less than L_0 , however, the annihilation rate of ν_μ 's oscillating out of the sterile phase above the mantle is comparable to that for nonoscillating μ -neutrinos in the absence of a baryonic mantle; i.e., here the efficiency is close to unity. But it must be noted that in supernova-like emission, the luminosity in muon neutrinos is typically on the order of only 10% of the electron luminosity, so the overall efficiency can be no higher than this in such models.

6. CONCLUSIONS

The Super-Kamiokande collaboration has reported that μ -neutrinos oscillate in a vacuum into another nonelectronic flavor, perhaps into a sterile neutrino. If a significant flux of neutrinos is initially present in GRB sources (as most models would have it), the oscillation of ν_μ into a sterile neutrino would allow the energy to be carried across a region contaminated with baryons and would allow its deposition in a relatively baryon-free region, thus allowing the formation of an ultrarelativistic blast. The efficiency of this deposition (relative to direct annihilation of the neutrinos when neither the oscillation nor baryons are present) is on the order of unity for toroidal geometry of emission, and a few percent for spherical or disklike geometry.

In models involving a solar-mass compact object (neutron star phase transitions, mergers, etc.; see § 4), the oscillation length for the expected neutrino energy of several MeV well matches the size of the system. For these models, the energy requirements are usually so extreme that toroidal geometry is preferred. In hypernovae models (§ 3), with the characteristic size of the system on the order of thousands of kilometers, GeV neutrino energies are required for this mechanism of energy transport to work; the available power is usually so large that quasi-spherical or disklike geometries with their efficiency of a few percent suffice.

In this Letter, I have assumed that it is ν_μ that oscillates into a sterile flavor. The solar neutrino deficit allows (but does not require) that the electron neutrino oscillates into a sterile flavor. However, the minimum oscillation length required in this case, of about 0.2 AU for \sim MeV neutrinos, would be too large to be useful in current GRB models.

This work was supported in part by the KBN through grant 2P03D01311. I thank the staff of the Aspen Center for Physics for their hospitality, and Mal Ruderman for comments on the manuscript.

REFERENCES

- Cheng, K. S., & Dai, Z. G. 1996, *Phys. Rev. Lett.*, 77, 1210
 Djorgovski, S. G., et al. 1997, *Nature*, 387, 876
 Eichler, D., Livio, M., Piran, T., Schramm, D. N. 1988, *Nature*, 340, 126
 Fishman, G. J., & Meegan, C. A. 1995, *ARA&A*, 33, 415
 Frail, D. A., et al. 1997, *Nature*, 389, 261
 Fukuda, Y., et al. 1998, preprint (hep-ex/9807003)
 Fuller, G. M., & Shi, X. 1998, preprint
 Galama, T. J., et al. 1997, *IAU Circ.* 6655
 Goodman, J., Dar, A., & Nussinov, S. 1987, *ApJ*, 314, L7
 Iwamoto, K., et al. 1998, preprint (astro-ph/9806382)
 Katz, J. I. 1994, *ApJ*, 422, 248
 Kluźniak, W., & Lee, W. H. 1998, *ApJ*, 494, L53
 Kluźniak, W., & Ruderman, M. 1998, *ApJ*, 505, L113
 Kulkarni, S. R., et al. 1998, *Nature*, 393, 35
 Lattimer, J. M., & Schramm, D. N. 1976, *ApJ*, 210, 549
 Mészáros, P., & Rees, M. J. 1992, *MNRAS*, 258, 41P
 ———. 1993, *ApJ*, 405, 278
 ———. 1997, *ApJ*, 482, L28
 Mochkovitch, R., Hernanz, M., Isern, J., & Martin, X. 1993, *Nature*, 361, 236
 Paczyński, B. 1986, *ApJ*, 308, L43
 ———. 1991, *Acta Astron.*, 41, 257
 ———. 1998, *ApJ*, 494, L45
 Paczyński, B., & Rhoads, J. E. 1993, *ApJ*, 418, L5
 Piran, T. 1997, in *Unsolved Problems in Astrophysics*, ed. J. N. Bahcall & J. P. Ostriker (Princeton: Princeton Univ. Press), 343
 Rasio, F., & Shapiro, S. L. 1992, *ApJ*, 401, 226
 Stecker, F. W. 1993, in *Int. Workshop on Neutrino Telescopes*, ed. M. B. Ceolin (Padova: Univ. Padova), 443
 Usov, V. V. 1992, *Nature*, 357, 472
 van Paradijs, J., et al. 1997, *Nature*, 386, 686
 Vietri, M. 1997, *ApJ*, 478, L9
 Wang, L., & Wheeler, J. C. 1998, *ApJ*, 504, L87
 Wilson, J. R., Salmonson, J. D., & Mathews, G. J. 1998, in *AIP Conf. Proc.* 4, *Gamma-Ray Bursts*, ed. C. A. Meegan, R. D. Preece, & T. M. Koshut (New York: AIP), 788
 Woosley, S. E. 1993, *ApJ*, 405, 273