High Energy Neutrinos from Cosmological Gamma-Ray Burst Fireballs

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Observations suggest that γ-ray bursts (GRBs) are produced by the dissipation of the kinetic energy of a relativistic fireball. We show that a large fraction, $\geq 10\%$, of the fireball energy is expected to be converted by photomeson production to a burst of $\sim 10^{14}$ eV neutrinos. A km^2 neutrino detector would observe at least several tens of events per year correlated with GRBs, and test for neutrino properties (e.g., flavor oscillations, for which upward moving τ's would be a unique signature, and coupling to gravity) with an accuracy many orders of magnitude better than is currently possible. [S0031-9007(97)02796-8]

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Recent observations of γ-ray bursts (GRBs) suggest that they originate from cosmological sources [1] (see, however, [2]). General phenomenological considerations indicate that the bursts are produced by the dissipation of the kinetic energy of a relativistic expanding fireball (see [3] for reviews). The physical conditions in the dissipation region imply [4] that protons may be Fermi accelerated in this region to energies $>10^{20}$ eV. Furthermore, the spectrum and flux of ultrahigh energy cosmic rays (above 10^{19} eV) are consistent with those expected from Fermi acceleration of protons in cosmological GRBs [5]. We show in this Letter that a natural consequence of the dissipative fireball model of GRBs is the conversion of a significant fraction of the fireball energy to an accompanying burst of $\sim 10^{14}$ eV neutrinos, created by photomeson production of pions in interactions between the fireball γ rays and accelerated protons. The neutrino burst is produced by interaction with protons with energies much lower than $\sim 10^{20}$ eV, the maximum acceleration energy. As shown below, 10^{15} eV protons interact with the ~ 1 MeV photons carrying the bulk of γ-ray energy to produce $\sim 10^{14}$ eV neutrinos.

The rapid variability time, ~ 1 ms, observed in some GRBs implies that the sources are compact, with a linear scale $r_0 \sim 10^7$ cm. The high luminosity required for cosmological bursts, $\sim 10^{51}$ ergs s^{-1}, then results in an optically thick (to pair creation) plasma, which expands and accelerates to relativistic velocities [6]. The hardness of the observed photon spectra, which extends to ≥ 100 MeV, implies that the γ-ray emitting region must be moving with a Lorentz factor Γ of order 100 [7], and constitutes independent evidence for ultrarelativistic outflow. The high energy density in the source would result in complete thermalization and a blackbody spectrum [6], in contrast with observations. To overcome this problem, Rees and Mészáros suggested [8] that γ-ray emission results from the dissipation at large radius of the kinetic energy of the relativistic ejecta. Such dissipation is expected to occur due to a collision with the interstellar medium [8], or due to internal collisions within the ejecta [9,10].

Paczyński and Xu suggested [9] that γ rays are emitted by the decay of neutral pions, which are produced in pp collisions once the kinetic energy is dissipated through internal collisions. In this case, an accompanying burst of ~ 30 GeV neutrinos is expected due to the decay of charged pions. **[***Note added.*—S. Pakvasa [*Beyond the Standard Model IV*, edited by J. F. Gunion *et al.* (World Scientific, River Edge, 1995)] and A. K. Mann (to be published) suggested production of higher energy neutrinos without specifying a production mechanism.**]** It is not clear, however, whether this model can account for the observed spectra. A more conservative mechanism for γ-ray production is the emission of synchrotron radiation (possibly followed by inverse-Compton scattering) by relativistic electrons accelerated in the dissipation shocks [8,10]. In this case, the proton density in the wind is too low to allow significant conversion of energy to neutrinos through pp collisions.

In the region where electrons are accelerated, protons are also expected to be shock accelerated. This is similar to what is thought to occur in supernovae remnant shocks, where synchrotron radiation of accelerated electrons is the likely source of nonthermal x rays (recent observations give evidence for acceleration of electrons in the remnant of SN1006 to 10^{14} eV [11]), and where shock acceleration of protons is believed to produce cosmic rays with energy extending to $\sim 10^{15}$ eV (see, e.g., [12] for review). The spectrum of ultrahigh energy cosmic rays (above 10^{19} eV) is consistent [5] with that expected from Fermi acceleration of protons in cosmological GRBs, and the flux is consistent with this scenario [4] provided the efficiency with which kinetic energy is converted to accelerated protons is comparable to the efficiency with which energy is converted to accelerated electrons (and hence to γ rays). We derive below the expected spectrum and flux of high energy neutrinos, produced by photomeson interactions between the wind γ rays and shock-accelerated protons, and discuss the implications for high energy neutrino astronomy.

Neutrino production in dissipative wind models of GRBs.—We consider a compact source producing a wind,

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characterized by an average luminosity $L \sim 10^{51}$ erg s^{-1} and mass loss rate $\dot{M} = L/\eta c^2$. At small radius, the wind bulk Lorentz factor, Γ, grows linearly with radius, until most of the wind energy is converted to kinetic energy and Γ saturates at $\Gamma \sim \eta \sim 100$. Variability of the source on the time scale Δt, resulting in fluctuations in the wind bulk Lorentz factor Γ on a similar time scale, would lead to internal shocks in the ejecta at a radius $r \sim r_d \approx \Gamma^2 c \Delta t$. We assume that internal shocks reconvert a substantial part of the kinetic energy to internal energy, which is then radiated as γ rays by synchrotron and inverse-Compton radiation of shock-accelerated electrons.

The photon distribution in the wind rest frame is isotropic. Denoting by $n(\epsilon_\gamma)d\epsilon_\gamma$ the number density of photons in the energy range ϵ_γ to $\epsilon_\gamma + d\epsilon_\gamma$ in the wind rest frame, the fractional energy loss rate of a proton with energy ϵ_p in the wind rest frame due to pion production is

$$t_\pi^{-1}(\epsilon_p) \equiv -\frac{1}{\epsilon_p}\frac{d\epsilon_p}{dt} = \frac{1}{2\Gamma_p^2}\, c \int_{\epsilon_0}^{\infty} d\epsilon\, \sigma_\pi(\epsilon)\xi(\epsilon)\epsilon \int_{\epsilon/2\Gamma_p}^{\infty} dx\, x^{-2} n(x)\,, \tag{1}$$

where $\Gamma_p = \epsilon_p/m_p c^2$, $\sigma_\pi(\epsilon)$ is the cross section for pion production for a photon with energy ϵ in the proton rest frame, $\xi(\epsilon)$ is the average fraction of energy lost to the pion, and $\epsilon_0 = 0.15$ GeV is the threshold energy. The GRB photon spectrum is well fitted in the Burst and Transient Source Experiment (BATSE) range (30 keV$-$3 MeV) by a combination of two power laws, $n(\epsilon_\gamma) \propto \epsilon_\gamma^{-\beta}$ with different values of β at low and high energy [13]. The break energy (where β changes) in the observer frame is typically $\epsilon_{\gamma b}^{\rm ob.} \sim 1$ MeV, with $\beta \simeq 1$ at energies below the break and $\beta \simeq 2$ above the break. Hereafter we denote quantities measured in the observer frame with the superscript "ob." (e.g., $\epsilon_{\gamma b}^{\rm ob.} = \Gamma \epsilon_{\gamma b}$). The second integral in (1) may be approximated by

$$\int_{\epsilon}^{\infty} dx\, x^{-2} n(x) \simeq \frac{1}{1+\beta}\frac{U_\gamma}{2\epsilon_{\gamma b}^3}\left(\frac{\epsilon}{\epsilon_{\gamma b}}\right)^{-(1+\beta)}, \tag{2}$$

where U_γ is the photon energy density (in the range corresponding to the observed BATSE range) in the wind rest frame, $\beta = 1$ for $\epsilon < \epsilon_{\gamma b}$ and $\beta = 2$ for $\epsilon > \epsilon_{\gamma b}$. The main contribution to the first integral in (1) is from photon energies $\epsilon \sim \epsilon_{\rm peak} = 0.3$ GeV, where the cross section peaks due to the Δ resonance. Approximating the integral by the contribution from the resonance we obtain

$$t_\pi^{-1}(\epsilon_p) \simeq \frac{U_\gamma}{2\epsilon_{\gamma b}}\, c\sigma_{\rm peak}\xi_{\rm peak} \times \frac{\Delta\epsilon}{\epsilon_{\rm peak}} \min(1, 2\Gamma_p \epsilon_{\gamma b}/\epsilon_{\rm peak})\,. \tag{3}$$

Here, $\sigma_{\rm peak} \simeq 5 \times 10^{-28}$ cm^2 and $\xi_{\rm peak} \simeq 0.2$ are the values of σ and ξ at $\epsilon = \epsilon_{\rm peak}$, and $\Delta\epsilon \simeq 0.2$ GeV is the peak width.

The energy loss of protons due to pion production is small during the acceleration process [4]. Once accelerated, the time available for proton energy loss by pion production is comparable to the wind expansion time as measured in the wind rest frame, $t_d \sim r_d/\Gamma c$. Thus, the fraction of energy lost by protons to pions is $f_\pi \simeq r_d/\Gamma c t_\pi$. The energy density in the BATSE range, U_γ, is related to the luminosity L_γ by $L_\gamma = 4\pi r_d^2 \Gamma^2 c U_\gamma$. Using this relation in (3), f_π is given by

$$f_\pi(\epsilon_p^{\rm ob.}) = 0.20\, \frac{L_{\gamma,51}}{\epsilon_{\gamma b,{\rm MeV}}^{\rm ob.} \Gamma_{300}^4 \Delta t_{\rm ms}} \times \begin{cases} 1, & \text{if } \epsilon_p^{\rm ob.} > \epsilon_{pb}^{\rm ob.}\,; \\ \epsilon_p^{\rm ob.}/\epsilon_{pb}^{\rm ob.}, & \text{otherwise.} \end{cases} \tag{4}$$

Here, $L_\gamma = 10^{51} L_{\gamma,51}$ erg s^{-1}, $\Gamma = 300\Gamma_{300}$, $\Delta t = 10^{-3}\Delta t_{\rm ms}$ s, and the proton break energy is

$$\epsilon_{pb}^{\rm ob.} = 1.3 \times 10^{16} \Gamma_{300}^2 (\epsilon_{\gamma b,{\rm MeV}}^{\rm ob.})^{-1} \text{ eV}\,. \tag{5}$$

Thus, for parameters typical of a GRB producing wind, a significant fraction of the energy of protons accelerated to energies larger than the break energy, $\sim 10^{16}$ eV, would be lost to pion production. Note that since the flow is ultrarelativistic, the results given above are independent of whether the wind is spherically symmetric or jetlike, provided the jet opening angle is $> 1/\Gamma$ (for a jetlike wind, L is the luminosity that would have been produced by the wind if it were spherically symmetric).

Since the constraint on the bulk Lorentz factor, $\Gamma \geq 100$, is derived from the requirement that the wind be optically thin to pair production for photons with observed energy ~ 100 MeV, it is useful to express f_π as a function of the pair production optical depth $\tau_{\gamma\gamma}$. A test photon with energy $\epsilon_{\rm t}$ may produce pairs in interactions with photons with energy exceeding a threshold $\epsilon_{\rm th}$, determined by $\epsilon_{\rm t}\epsilon_{\rm th} = 2(m_e c^2)^2/(1 - \cos\theta)$, where θ is the angle between the photons propagation directions. Pair production interactions involve photons with energy much higher than the break energy, $\epsilon_{\gamma b}$. A test photon with observed energy $\epsilon_{\rm t}^{\rm ob.} = 100$ MeV, for example, has an energy ~ 1 MeV in the wind rest frame, and therefore interacts mainly with photons of similar energy to produce pairs. Assuming that the spectrum of photons produced in the dissipation region extends as $n(\epsilon) \propto \epsilon^{-2}$ from the BATSE range to (observed) energies well above 100 MeV, the mean free path for pair production (in the wind rest frame) for a photon of energy $\epsilon_{\rm t}$ is

$$l_{\gamma\gamma}^{-1}(\epsilon_{\rm t}) = \frac{1}{2}\frac{3}{16}\sigma_T \int d\cos\theta(1 - \cos\theta) \int_{\epsilon_{\rm th}(\epsilon_{\rm t},\theta)}^{\infty} d\epsilon\, \frac{U_\gamma}{2\epsilon^2} = \frac{1}{16}\sigma_T \frac{U_\gamma \epsilon_{\rm t}}{(m_e c^2)^2}\,. \tag{6}$$

Here we have used a constant cross section, $3\sigma_T/16$, where σ_T is the Thomson cross section, above the threshold $\epsilon_{\rm th}$. [The cross section drops as $\ln(\epsilon)/\epsilon$ for $\epsilon \gg \epsilon_{\rm th}$; however, since the number density of photons drops rapidly with energy, (6) is a good approximation.]

The optical depth is given by $\tau_{\gamma\gamma} = r_d/\Gamma l_{\gamma\gamma}$. Using (6), we obtain

$$f_\pi(\epsilon_p^{\text{ob.}}) = 0.15\tau_{\gamma\gamma}(\epsilon_{\text{t}}^{\text{ob.}} = 100 \text{ MeV}) \frac{\min(\epsilon_p^{\text{ob.}}, \epsilon_{pb}^{\text{ob.}})}{10^{16} \text{ eV}} . \tag{7}$$

Only a small fraction of bursts show a power law spectrum extending to >100 MeV. Most bursts may therefore have $\tau_{\gamma\gamma}(\epsilon_{\text{t}}^{\text{ob.}} = 100 \text{ MeV})$ larger than unity, leading to higher efficiency of pion production. Furthermore, variability of the source over different time scales would lead to dissipation shocks over a range of radii [10], where the optical depth may become small only at the largest dissipation radii. While photons can escape only from radii where the optical depth is small, neutrinos can escape from almost any depth. Thus, most of the burst energy may actually come out in neutrinos. For $\Gamma = 100$ and $L_\gamma = 10^{51}$ erg s^{-1}, for example, $\tau_{\gamma\gamma}(\epsilon_{\text{t}}^{\text{ob.}} = 100 \text{ MeV}) \sim 1$ for $r_d \sim 10^{14}$ cm, where internal shocks result from variability over a 0.1 s time scale. Variability on shorter time scales results in collisions at smaller radii with larger pion production efficiency. The conversion to high energy pions of a significant fraction, $\geq 20\%$, of the energy of protons accelerated to energy similar to or larger than the break energy ϵ_{pb}, would be avoided only if $\Gamma \gg 100$.

Neutrino spectrum and flux.—Roughly half of the energy lost by protons goes into π^0's and the other half to π^+'s. Neutrinos are produced by the decay of π^+'s, $\pi^+ \rightarrow \mu^+ + \nu_\mu \rightarrow e^+ + \nu_e + \overline{\nu}_\mu + \nu_\mu$ [the large optical depth for high energy γ's from π^0 decay, cf. Eq. (6), would not allow these photons to escape the wind]. The mean pion energy is 20% of the energy of the proton producing the pion. This energy is roughly evenly distributed between the π^+ decay products. Thus, approximately half the energy lost by protons of energy ϵ_p is converted to neutrinos with energy $\sim 0.05\epsilon_p$. Equation (4) then implies that the spectrum of neutrinos above $\epsilon_{\nu b} = 0.05\epsilon_{pb}$ follows the proton spectrum, and is harder (by one power of the energy) at lower energy. The break energy is $\epsilon_{\nu b}^{\text{ob.}} = 5 \times 10^{14}\Gamma_{300}^2(\epsilon_{\gamma b,\text{MeV}}^{\text{ob.}})^{-1}$ eV. For a power law differential spectrum of accelerated protons $n(\epsilon_p) \propto \epsilon_p^{-2}$, as typically expected for Fermi acceleration and which would produce the observed spectrum of ultrahigh energy cosmic rays [5], the differential neutrino spectrum is $n(\epsilon_\nu) \propto \epsilon_\nu^{-\alpha}$ with $\alpha = 1$ below the break and $\alpha = 2$ above the break. The spectrum may be modified above the energy where $\overline{\nu}_\mu$ are produced by the decay of muons with lifetime $\Gamma_\mu t_\mu$ (where Γ_μ is the muon Lorentz factor and $t_\mu = 2 \times 10^{-6}$ s its rest lifetime) comparable to the wind dynamical time $r_d/\Gamma c$, i.e., above $\epsilon_{\nu\mu}^{\text{ob.}} = m_\mu c^2 r_d/3t_\mu c \approx \Gamma^2 m_\mu c^2 \Delta t/3t_\mu = 5 \times 10^{15}\Gamma_{300}^2\Delta t_{\text{ms}}$ eV. Muons producing $\overline{\nu}_\mu$ with higher energy may adiabatically lose a significant part of their energy before decaying (other energy loss processes, e.g., inverse-Compton scattering, can be shown to be less important).

Most of the neutrino energy is carried by neutrinos with energy close to the break energy, $\epsilon_{\nu b}^{\text{ob.}} \sim 10^{14}$ eV. The neutrino flux depends on the relative efficiency with which the wind kinetic energy is converted to accelerated protons, compared to the efficiency with which energy is converted to accelerated electrons (and therefore to γ rays). If cosmological GRBs are the sources of ultrahigh energy cosmic rays, then the efficiency of converting energy to γ rays and to accelerated protons should be similar. The energy production rate required to produce the observed flux of ultrahigh energy cosmic rays, assuming that the sources are cosmologically distributed, is $\sim 4 \times 10^{44}$ erg Mpc^{-3} yr^{-1} over the energy range 10^{19}–10^{21} eV, comparable to the rate of energy production by GRBs as γ rays in the BATSE range [5]. For a proton generation spectrum $n(\epsilon_p) \propto \epsilon_p^{-2}$, similar energy is contained in equal logarithmic energy intervals, implying that the total energy converted to accelerated protons (i.e., over the energy range $\Gamma_p m_p c^2 \sim 10^{12}$–$10^{21}$ eV) is a few times that converted to γ-ray energy in the BATSE range. We therefore assume below that the conversion efficiency is similar for electrons and protons, leading to an energy production rate above the proton energy break, $\sim 10^{16}$ eV, $\dot{E} \sim 4 \times 10^{44}$ erg Mpc^{-3} yr^{-1} ($\dot{E}$ depends only logarithmically on the value of the break energy).

The present day neutrino energy density due to GRBs is approximately given by $U_\nu \approx 0.5 f_\pi(\epsilon_{pb}) t_H \dot{E}$, where $t_H \approx 10^{10}$ yr is the Hubble time. The neutrino flux is therefore approximately given by

$$J_\nu(\epsilon_\nu > \epsilon_{\nu b}) \approx \frac{c}{4\pi}\frac{U_\nu}{\epsilon_{\nu b}} \approx 10^{-13}\frac{f_\pi(\epsilon_{pb})}{0.2}\dot{E}_{44} \times \left(\frac{\epsilon_{\nu b}^{\text{ob.}}}{10^{14} \text{ eV}}\right)^{-1} \text{cm}^{-2}\,\text{s}^{-1}\,\text{sr}^{-1}, \tag{8}$$

where $\dot{E} = 10^{44}\dot{E}_{44}$ erg Mpc^{-3} yr^{-1}. The high energy neutrinos predicted in the dissipative wind model of GRBs may be observed by detecting the Cherenkov light emitted by high energy muons produced by neutrino interactions below a detector on the surface of the Earth (see [14] for a recent review). This technique works only for muons entering the detector from below, "upward moving muons," due to the large background of downward atmospheric muons. The probability $P_{\nu\mu}$ that a neutrino would produce a high energy muon in the detector is approximately given by the ratio of the high energy muon range to the neutrino mean free path. At the high energy we are considering, $P_{\nu\mu} \simeq 10^{-6}(\epsilon_\nu/1 \text{ TeV})$ [14]. Using (8), the expected flux of upward moving muons is

$$J_{\mu\uparrow} \approx 50\frac{f_\pi(\epsilon_{pb})}{0.2} \times \left(\frac{\dot{E}}{4 \times 10^{44} \text{ erg Mpc}^{-3}\,\text{yr}^{-1}}\right) \text{km}^{-2}\,\text{yr}^{-1}. \tag{9}$$

The rate is almost independent of $\epsilon_{\nu b}$, due to the increase of $P_{\nu\mu}$ with energy.

The rate (9) is comparable to the background expected due to atmospheric neutrinos [14]. However, neutrino bursts should be easily detected above the background, since the neutrinos would be correlated, both in time and angle, with the GRB γ rays. A km^2 neutrino detector should detect each year $\sim$10 to 100 neutrinos correlated with GRBs. Furthermore, nearby bright GRBs, although rare ($\sim$0.1 per year), would produce a burst of several neutrinos in a km^2 detector. A burst at a distance of 100 Mpc producing 0.4×10^{51} erg in $\sim 10^{14}$ eV neutrinos would produce ~ 3 km^{-2} upward muons.

Implications.—Detection of neutrinos from GRBs would corroborate the cosmological fireball scenario for GRB production (acceleration of protons to ultrahigh energy by the processes discussed above is not possible if GRBs are galactic [15]). Neutrinos from GRBs could be used to test the simultaneity of neutrino and photon arrival to an accuracy of $\sim$1 s ($\sim$1 ms for short bursts), checking the assumption of special relativity that photons and neutrinos have the same limiting speed. [The time delay for neutrino of energy 10^{14} eV with mass m_ν traveling 100 Mpc is only $\sim 10^{-11}(m_\nu/10 \text{ eV})^2$ s]. These observations would also test the weak equivalence principle, according to which photons and neutrinos should suffer the same time delay as they pass through a gravitational potential. With 1 s accuracy, a burst at 100 Mpc would reveal a fractional difference in a limiting speed of 10^{-16}, and a fractional difference in a gravitational time delay of order 10^{-6} (considering the Galactic potential alone). Previous applications of these ideas to supernova 1987A (see [16] for review), where simultaneity could be checked only to an accuracy of order several hours, yielded much weaker upper limits: of order 10^{-8} and 10^{-2} for fractional differences in the limiting speed [17] and time delay [18], respectively.

The model discussed above predicts the production of high energy muon and electron neutrinos with a 2:1 ratio. If vacuum neutrino oscillations occur in nature, then neutrinos that get here should be almost equally distributed between flavors for which the mixing is strong. In fact, if the atmospheric neutrino anomaly has the explanation it is usually given, oscillation to ν_τ's with mass $\sim$0.1 eV [19], then one should detect equal numbers of ν_μ's and ν_τ's. Upgoing τ's, rather than μ's, would be a distinctive signature of such oscillations. Since ν_τ's are not expected to be produced in the fireball, looking for upgoing τ's would be an "appearance experiment" (ν_τ's may be produced by photoproduction of charmed mesons; however, the high photon threshold, $\sim$50 GeV, and low cross section, $\sim$1 μb [20], for such reactions imply that the ratio of charmed meson to pion production is $\sim 10^{-4}$). To allow a flavor change, the difference in squared neutrino masses, Δm^2, should exceed a minimum value proportional to the ratio of source distance and neutrino energy [16]. A burst at 100 Mpc producing 10^{14} eV neutrinos can test for $\Delta m^2 \geq 10^{-16}$ eV^2, 5 orders of magnitude more sensitive than solar neutrinos. Note that due to the finite pion lifetime, flavor mixing would be caused by decoherence, rather than by real oscillations, for neutrinos with masses $>$0.1 eV.

The ν_μ's may resonantly oscillate to ν_e's [the Mikheyev-Smirnov-Wolfenstein (MSW) effect] as they escape the fireball, provided the electron number density exceeds the MSW resonance density (e.g., [16]). The fireball electron density is approximately $n_e \sim L/4\pi r_d^2 \Gamma^2 m_p c^3 \sim 10^{10}$–$10^{12}$ cm^{-3}, implying that the MSW effect could be important only for neutrino masses smaller than $\sim 10^{-12}$ eV^2. For such low masses, the spatial scale of the fireball implies that n_e changes too rapidly to allow the use of the adiabatic approximation, and that the influence of the MSW effect would depend on the detailed structure of the fireball.

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