ON THE ENERGY OF NEUTRINOS FROM GAMMA-RAY BURSTS

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ABSTRACT

Ultrahigh-energy protons that are accelerated at the shocks, causing gamma-ray bursts, photoproduce pions and then neutrinos in situ. I consider here the sources of losses in this process, namely, adiabatic and synchrotron losses by both pions and muons. When the shocks under consideration are external, i.e., when they are between the ejecta and the surrounding interstellar medium, I show that neutrinos produced by pion decay are unaffected by losses; those produced by muon decay, in the strongly beamed emission required by afterglow observations of GRB 971214, are limited in energy, but still exceed 10^{19} eV. In particular, this means that ultrahigh-energy neutrinos will be produced through afterglows.

Subject headings: acceleration of particles — cosmology: miscellaneous — elementary particles — gamma rays: bursts — relativity — shock waves

1. INTRODUCTION

The discovery of the afterglows of gamma-ray bursts (GRBs; Costa et al. 1997; van Paradijs et al. 1997; Frail et al. 1997) and the disappearance of flares in their radio flux (Goodman 1997) have provided strong evidence in favor of the fireball model (Rees & Mészáros 1992), whereby mass is ejected from an as-yet-unknown source with large Lorentz factors, $\gamma \approx 100$. Several authors have pointed out that in the extremely energetic events giving rise to these hyperrelativistic flows, high-energy neutrinos are likely to be produced. A first likely source is the merger of the two neutron stars, which may (Narayan, Paczyński, & Piran 1992) give rise to the bursts themselves, as originally suggested by Eichler et al. (1989). Although insufficient to power the burst itself (Janka & Ruffert 1996), this may yet provide a copious source of low-energy neutrinos, ≈ 1 MeV. Another source of neutrinos comes from p-p collisions inside the shocks, which give rise to the observed burst proper (Paczyński & Xu 1994); however, it seems unlikely that these collisions may give rise to the burst itself, although they may originate a flux of higher energy neutrinos, ≈ 30 GeV, which is, however, currently unobservable (Ostrowski & Zdziarski 1995).

A third, distinct source of high-energy neutrinos exists. It has recently been pointed out that GRBs may be responsible for the acceleration of the highest energy cosmic rays observed so far (Vietri 1995; Waxman 1995). It has been suggested (Waxman & Bahcall 1997) that these ultrahigh-energy cosmic rays (UHECRs) may produce neutrinos through pion photoproduction, i.e., through the reaction

$$p + \gamma \rightarrow n + \pi^+ ,$$
$$\pi^+ \rightarrow \mu^+ + \nu_\mu , \qquad (1)$$
$$\mu^+ \rightarrow e^+ + \nu_e + \bar{\nu}_\mu ,$$

inside the burst itself, i.e., against the γ-ray photons of the burst. While the probability of this happening is not large (Vietri 1995), Waxman & Bahcall (1997) pointed out that significant amounts of neutrinos with energy $\approx 10^{14}$ eV ought to be produced and might be revealed by detectors of the AMANDA class. However, this detection is made difficult by the fact that the Earth becomes opaque for this neutrino energy, so that the technique employing upward-moving muons cannot be used. But, through the same mechanism and simultaneously with these lower energy brethren, neutrinos in the energy range $\approx 10^{19}$ eV (ultrahigh-energy neutrinos, hereafter UHENs) ought to be produced; in Vietri (1998, hereafter Paper I) I showed that a significant fraction ($f_\pi \approx 0.01$) of all energy initially released in UHECRs ought to be lost to neutrinos emitted during the burst proper or during the first week of the burst afterglow. The interesting point of this computation is that it appears that the next generation of satellite-borne detectors of large showers, such as AIRWATCH (Linsley 1997) may cover an area large enough to allow detection of UHENs within the first year or so of operation.

Shortly before Paper I was accepted, Rachen & Mészáros (1998) presented a detailed analysis of losses limiting neutrinos' energies in the *internal* shock scenario for GRBs, which applies, for instance, to the model of Waxman & Bahcall (1997). They conclude that these losses severely limit both the energies of individual neutrinos and the total energy released by GRBs. Such losses arise because, before decaying, both the photoproduced pions and the muons generated by the pion decay suffer adiabatic and synchrotron losses, thus imparting to neutrinos less than the energy they started out with. Given the short rest lifetimes of pions and muons, one maybe inclined to think these losses negligible, but, because of relativistic time dilation, pions and muons moving with a Lorentz factor typical of UHECRs ($\approx 10^{10}$) may survive more than 10^4 s and cover sizeable distances ($\approx 10^{14}$ cm) in the meantime.

It is the purpose of this paper to carry out the analysis of losses for *external* shocks, both during the burst proper and in the afterglow phase. It will be concluded that in the relatively tame environments generated by external shocks, neutrinos produced by pion decay are unaffected by losses, while neutrinos produced by muon decay are limited in energy by these processes but still manage to exceed 10^{19} eV. The major departure from the discussion in Paper I is due to the sensational discovery (Kulkarni et al. 1998) of the redshift $z = 3.43$ for the burst GRB 971214. This discovery forces us to choose different scaling values for bursts, for the following reason: For cosmological parameters $\Omega = 0$ and $H_0 = 50$ km s^{-1} Mpc^{-1}, the luminosity distance to this redshift is 6.2×10^{28} cm; for $\Omega = 1$ it would be 8.4×10^{28}

cm. The burst as observed by the *BeppoSAX* Gamma-Ray Burst Monitor had a peak luminosity of 6×10^{-7} ergs s^{-1} cm^{-2} and a total duration of ≈ 30 s (Heise et al. 1997), corresponding to a fluence of $\approx 10^{-5}$ ergs cm^{-2}. With the lower, above-stated luminosity distance, this converts to a total energy release of 4×10^{53} ergs, roughly the binding energy of two merging neutron stars. While strictly speaking still not an impossibility, these values seem to imply some beaming, a natural consequence of many GRBs' models. In fact, assuming a beaming factor $\theta \approx 0.1$, the total energy release becomes a more mundane 10^{51} ergs. Similar, although less extreme values have been derived, less cogently, for the burst GRB 970508 (Panaitescu, Mészáros, & Rees 1998). These values will be adopted here.

2. ON THE ACCELERATION OF PROTONS IN FIREBALLS

Before considering losses, I need to establish some properties of the environment within the fireball theory of GRBs. The maximum proton energy in the observer frame was shown to be

$$\epsilon_{\max} = 2 \times 10^{20} \text{ eV } \eta_2^{1/3} E_{51}^{1/3} n_1^{1/6} \theta_{-1}^{-2/3} \xi^{1/2} \qquad (2)$$

(Vietri 1995), where the ejecta shell has, at the moment of impacting with the interstellar medium of density $\rho = n_1 m_p$ g cm^{-3}, a Lorentz factor $\eta = 100 \eta_2$; the shell explosion energy is $E_{51} \times 10^{51}$ ergs, $\theta = 0.1 \theta_{-1}$ the beaming angle; and ξ is the magnetic field energy density in units of the equipartition magnetic field energy density. As mentioned in Paper I, this corrects a small error of Vietri (1995). When seen from the shell, this highest energy proton will have a Lorentz factor γ_p given by

$$\gamma_p = 2 \times 10^9 \eta_2^{-2/3} E_{51}^{1/3} n_1^{1/6} \theta_{-1}^{-2/3} \xi^{1/2} . \qquad (3)$$

The two equations above correspond to equations (29) and (28), respectively, of Vietri (1995). It is perhaps useful to remind the reader that they were obtained under the assumption that the major limiting factor of the protons' energies comes from the finite shell thickness: whenever the proton energy increases so much that g times its Larmor radius r_L exceeds the shell thickness, the proton will cross the whole shell without being deflected backward, and the acceleration cycle stops. Here, the magnetic field is assumed to be a fraction $\xi^{1/2}$ of its equipartition value, and the value $g \approx 40$ is taken from the numerical simulations of Quenby & Lieu (1989).

However, at the referee's prompting, I can also show that the acceleration timescale is shorter than the shell light crossing time, which is the timescale on which adiabatic losses set in and also the timescale on which the highest energy protons cross the shell before being deflected backward. By establishing this, I shall in fact demonstrate that the two above equations correctly describe the highest proton energies achievable. The acceleration timescale at relativistic shocks t_r is shorter by a factor $Q \approx 13.5$ than the traditional acceleration timescale at nonrelativistic shocks (Quenby & Lieu 1989); thus

$$t_r = \frac{t_{\text{nr}}}{Q} = \frac{c}{Q(V_1 - V_2)} \int_{p_i}^{p_f} \left(\frac{\lambda_1}{V_1} + \frac{\lambda_2}{V_2} \right) \frac{dp}{p} . \qquad (4)$$

Here, matter on the two sides of the shocks moves with speeds V_1, V_2 with respect to the stationary shock; $\lambda = g r_L$ is the typical deflection length; p_i is the injection momentum; and p_f is the highest achievable momentum in the shock

frame. Since $\lambda = g r_L = g c p / e B$, I find

$$t_r = \frac{c}{Q(V_1 - V_2)} \int_{p_i}^{p_f} \frac{gc}{eB} \left(\frac{1}{V_1} + \frac{1}{V_2} \right) dp$$
$$= \frac{gc^2 p_f}{(V_1 - V_2) QeB} \left(\frac{1}{V_1} + \frac{1}{V_2} \right) . \qquad (5)$$

Focusing now on a relativistic flow, for which $V_1, V_2 \approx c$, and also $V_1 - V_2 = qc \approx c$, I obtain

$$t_r \approx \frac{2gp_f}{QqeB} . \qquad (6)$$

On the other hand, for the highest energy protons, we have set, in deriving equation (2), $\lambda = g r_L = g c p_f / e B = r_{\text{th}}$, with r_{th} being the shell thickness; inserting this into the above equation, we find

$$t_r = \frac{2 r_{\text{th}}}{qQc} \approx 0.1 \frac{r_{\text{th}}}{c} , \qquad (7)$$

which is as long as the shell light crossing time. So the acceleration is prompt enough to propel protons to the highest energies computed in equation (2) within the time-scale r_{th}/c, which is approximately the whole time the proton spends inside the shell before reaching an energy so large that it cannot be turned backward. Furthermore, proton adiabatic losses also become appreciable on the shell expansion timescale (Rachen & Mészáros 1998) r_{th}/c, longer than the acceleration timescale (eq. [7]).

It should be noticed that this result is independent of the specific values of g assumed; nor does it depend on whether the magnetic field has achieved its equipartition value. It is only the highest proton energy that depends on the parameters g and ξ. Nor does it depend explicitly upon the shock Lorentz factor, provided, of course, the shock is relativistic. This is especially important when one considers the application of this same idea to the bursts' afterglows: the maximum proton energy is again provided by the finite shell thickness, which, however, following equation (2), changes as the shell slows down. The reason for the continuing applicability of equation (2) is that, even during the afterglow, the shell thickness is given in the shell frame by $r_{\text{th}} = r/c$, where r is the current shock position exactly as during the burst proper (Mészáros, Laguna, & Rees 1994). This result can be phrased as follows: as the shock decelerates, the acceleration timescale becomes longer, but so does also the shell crossing time. Provided the magnetic field builds up to the same, constant fraction of the equipartition value ξ, the two timescales lengthen by the same amount, thus keeping the result of equation (7) still valid.

The only hidden hypothesis in this argument is that appreciable magnetic fields exist on both sides of the shock to deflect the cosmic rays; in Vietri (1995) I suggested that this occurs when two or more shells collide. I believe the same suggestion is still valid for afterglows; I shall, however, discuss the implications of these multiple collisions in a different paper.

Last, I should remark that the results discussed in this section do depend upon the correctness of the assumption made by Quenby & Lieu (1989) that scattering of high-energy protons at relativistic shocks is isotropic and large pitch angle. Since this provides a satisfactory explanation for the acceleration of cosmic rays at active galactic nuclei (Quenby & Lieu 1989) and is borne out by theoretical

arguments and simulations of interplanetary shocks (Moussas, Quenby, & Valdes-Galicia 1982, 1987), it seems quite a plausible assumption to make.

3. ON FIREBALLS

Equations (2) and (3) are, however, in an inconvenient form, the reason being that the parameter n_1 is unobservable. It is expedient to substitute for it the burst duration $T = T_2 \times 100$ s, a directly observed quantity that is related to it by the following argument: The shock with the interstellar medium forms at a distance $r_{\rm sh}$ from the explosion site, where

$$r_{\rm sh} = \left(\frac{3E}{4\pi n_1 \, m_p \, c^2 \theta^2 \eta^2}\right)^{1/3} \tag{8}$$

(Mészáros et al. 1993), while the burst duration is of order

$$T = \frac{r_{\rm sh}}{2\eta^2 c}. \tag{9}$$

This is an old argument, originally due to Ruderman (1975) and more recently discussed by Sari & Piran (1997). Eliminating $r_{\rm sh}$ from the above, I find

$$n_1 = \frac{3E}{32\pi m_p \, c^5 \theta^2 \eta^8 T^3} = 8 \text{ cm}^{-3} \, E_{51} \eta_2^{-8} T_2^{-3} \theta_{-1}^{-2}, \tag{10}$$

and

$$r_{\rm sh} = 2\eta^2 cT = 6 \times 10^{14} \text{ cm } T_2 \eta_2^2. \tag{11}$$

The magnetic field behind the shock, in the shell frame, is often assumed (Mészáros et al. 1993) to be some fixed fraction $\xi^{1/2}$ of its equipartition value $B_{\rm eq}$, and it is given by

$$B = (8\pi \xi n_1 \, m_p \, c^2)^{1/2} \eta = 54 G E_{51}^{1/2} \eta_2^{-3} T_2^{-3/2} \theta_{-1}^{-1} \xi^{1/2}, \tag{12}$$

where I have of course used equation (10). It is also convenient to rewrite equation (3), again using equation (10), as

$$\gamma_p = 3 \times 10^9 E_{51}^{1/2} \eta_2^{-2} T_2^{-1/2} \theta_{-1}^{-1} \xi^{1/2} \tag{13}$$

and equation (2) as

$$\epsilon_{\rm max} = 3 \times 10^{20} \text{ eV } E_{51}^{1/2} \eta_2^{-1} T_2^{-1/2} \theta_{-1}^{-1} \xi^{1/2}. \tag{14}$$

4. LOSSES

I have established above that protons suffer no adiabatic losses during their acceleration; in Vietri (1995) I had also shown that synchrotron and photopion losses are negligible. I showed in Paper I that the small photopion production leads, through the decays of equation (1), to the production of a small but significant flux of prompt neutrinos. However, Rachen & Mészáros (1998) have shown that one ought to consider carefully whether synchrotron and adiabatic losses by muons and pions before their decay may significantly limit neutrinos' energies.

The most significant losses will be due to muons, because of their longer lifetime ($\tau_\mu = 2.2 \times 10^{-6}$ s, $\tau_\pi = 2.6 \times 10^{-8}$ s). The maximum muon Lorentz factor γ_μ before significant adiabatic losses set in is given by the following argument: A muon moving with Lorentz factor γ in the shell frame appears to have a lifetime $\gamma\tau_\mu$ because of relativistic time dilation. Within this time, the shell is expanding, thus allowing the comoving magnetic field to decrease on a timescale $\approx r_{\rm th}/\dot{r}_{\rm th}$. The shell typical size $r_{\rm th}$ in the shell frame is r/η (Mészáros et al. 1993), where r is the shock instantaneous

distance from the origin of the explosion; at the moment of the burst, $r = r_{\rm sh}$ (eq. [11]). The shell expansion velocity $\dot{r}_{\rm th} \approx c$. The limiting Lorentz factor γ_μ is found by equating these two timescales:

$$\gamma_\mu = \frac{r_{\rm sh}}{\eta c \tau_\mu} = 9 \times 10^7 T_2 \eta_2. \tag{15}$$

The condition that neutrino production is not affected, $\gamma_p < \gamma_\mu$, can be expressed as a requirement on the burst parameters: I obtain

$$\eta_2 > 3.1 E_{51}^{1/6} T_2^{-1/2} \theta_{-1}^{-1/3} \xi^{1/6}. \tag{16}$$

For pion adiabatic losses, the only change comes in because of its shorter lifetime:

$$\eta_2 > 0.5 E_{51}^{1/6} T_2^{-1/2} \theta_{-1}^{-1/3} \xi^{1/6}. \tag{17}$$

Synchrotron losses by muons are related to those of protons as follows: Calling $t_s = 1$ yr $(10^{11}/\gamma)(1G/B)^2$ the lifetime against synchrotron losses for a proton of Lorentz factor γ, the largest muon Lorentz factor γ_μ before significant losses set in is given by

$$\gamma t_s \left(\frac{m_\mu}{m_p}\right)^3 = \gamma_\mu^2 \tau_\mu \tag{18}$$

(Rachen & Mészáros 1998). Here the muon mass is $m_\mu \approx 0.1 m_p$. Using equation (12), I find

$$\gamma_\mu = 4 \times 10^{10} \frac{1G}{B} = 7.4 \times 10^8 E_{51}^{-1/2} \eta_2^3 \, T_2^{3/2} \theta_{-1} \, \xi^{-1/2}. \tag{19}$$

Again, the condition that these losses are irrelevant, $\gamma_p < \gamma_\mu$, can be rewritten as a requirement on the burst parameters: I find

$$\eta_2 > 1.3 E_{51}^{1/5} T_2^{-2/5} \theta_{-1}^{-2/5} \xi^{1/5}. \tag{20}$$

For pions the relevant criterion is

$$\eta_2 > 0.4 E_{51}^{1/5} T_2^{-2/5} \theta_{-1}^{-2/5} \xi^{1/5}. \tag{21}$$

It is convenient to remark here that the use of η and T as independent burst parameters allows discussion of afterglow with the same formulae. During afterglows, the quantity T loses meaning, because the afterglow emission is of course continuous. However, let us still define, by analogy with equation (9), $Tr/2\eta^2 c$, where r and η are the instantaneous shock position and Lorentz factor. During the afterglow, for adiabatic expansion, the shell Lorentz factor decreases as $\eta \propto r^{-3/2}$. Thus, the variation of T parametrizes the shell radius. I find that

$$\eta = \eta_i \left(\frac{T}{T_i}\right)^{-3/8} \tag{22}$$

describes the further evolution of the shell. Here η_i and T_i are the initial values of the afterglow at the moment of the formation of the external shock. This equation can also be established for an adiabatic expansion in a constant density environment by imposing the constancy through the afterglow of equation (10). It is also identical to the time evolution law for adiabatic expansion with respect to observer's time, so that the fictitious quantity T can also be identified

with Earth time, a physical quantity. In other words, by considering arbitrary values of η and T, I am considering points that do not really model any known burst, but that correctly parametrize later, afterglow moments of realistic bursts; the realistic initial models are linked to later values of η and T by equation (22).

First I consider neutrinos produced directly by charged pion decay. In Figure 1 I show the constraints (eqs. [17] and [21]) for $E_{51} = \theta_{-1} = \xi = 1$, plotted as solid lines. The dashed lines represent the afterglow evolution tracks (eq. [22]). The figure covers roughly the first three afterglow hours. The dotted lines represent the loci of points of constant highest neutrino energies; when losses are negligible, these are computed by multiplying the highest proton energies (eq. [14]) by 0.05. Neutrinos produced directly by charged pion decay will not be limited by the rather more stringent requirements on muon losses (eq. [16]). This means that, should a burst fail to produce high-energy neutrinos from muon decay, its neutrino flux will only be decreased (roughly) by a factor of 3, because the two neutrinos from the muon decay will be lost at the same energy level. It can easily be seen from Figure 1 that neutrinos of energy as large as 10^{19} eV may be produced.

Muon-produced neutrinos will be limited in energy both by synchrotron losses and by adiabatic losses, depending exactly on the model parameter values. The crossover occurs when equation (15) equals equation (19), i.e., for

$$\eta_{2cr} = 0.3 E_{51}^{1/4} T_2^{-1/4} \theta_{-1}^{-1/2} \xi^{1/4} \qquad (23)$$

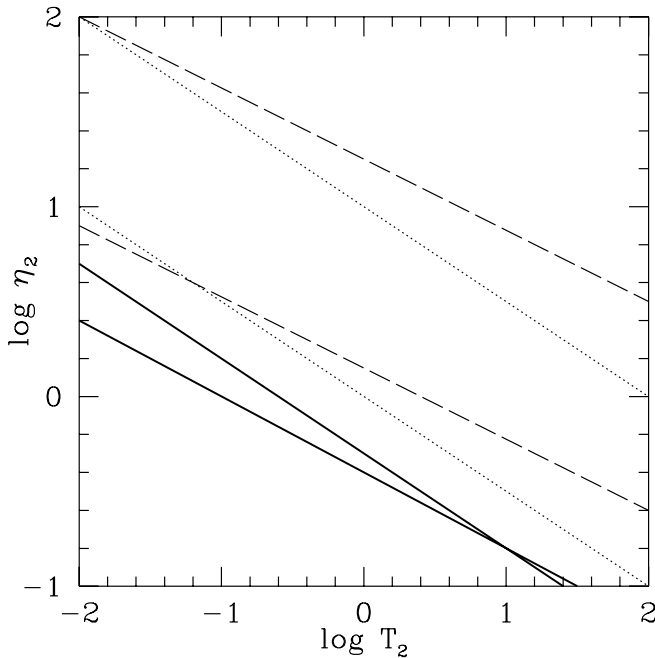with adiabatic losses dominating whenever $\eta_2 > \eta_{2cr}$.

FIG. 1.—Typical energies for neutrinos produced by charged pion decay, as a function of bursts' parameters. Pion adiabatic losses are negligible in the upper region (top solid line, eq. [17]; bottom solid line, eq. [21]). Below this line synchrotron losses become important. Dashed lines show representative tracks for afterglow evolution (eq. [22]); dotted lines show loci of points of constant neutrino energies in the observer's frame. The top dotted line corresponds to $\epsilon_\nu = 1.5 \times 10^{18}$ eV, while the bottom dotted line to $\epsilon_\nu = 1.5 \times 10^{19}$ eV. These values are obtained by multiplying eq. (14) times 0.05, the typical yield for the neutrino energy. The figure covers, roughly, the first 3 afterglow hours.

When synchrotron losses dominate, we see by comparing equation (19) with equation (22) that afterglow evolution occurs along lines of (roughly!) constant γ_μ in the shell frame, so that the maximum achievable energy in the observer frame is given by $100 \times \eta_2$ times equation (19). Multiplying times the typical energy of outgoing neutrinos, ≈ 50 MeV, we find that the neutrino energy is

$$\epsilon_\nu^{sy} = 3.7 \times 10^{18} \text{ eV } E_{51}^{-1/2} \eta_2^4 T_2^{3/2} \theta_{-1} \xi^{-1/2} . \qquad (24)$$

Taking into account that during afterglow the shell Lorentz factor varies according to equation (22), the typical neutrino energy is given by inserting this into the above equation:

$$\epsilon_\nu = 3.7 \times 10^{18} \text{ eV } \eta_{21}^4 T_{2i}^{3/2} \xi^{-1/2} , \qquad (25)$$

where of course η_{2i} and T_{2i} are the initial Lorentz factor and burst duration in units of 100 and 100 s, respectively, and the result is independent of postburst time. Since the dependence on the initial parameter is so steep, it is enough that η_{2i} exceeds unity by a factor of 2 for ϵ_ν to exceed 10^{19} eV. When adiabatic losses dominate, the limiting muon Lorentz factor in the observer frame is given by $\gamma_\mu^{obs} = 9 \times 10^9 T_2 \eta_2^2$. The locus of constant γ_μ^{obs} is thus parallel to the top solid line of Figure 1 (eq. [16]) and to equation (14). This gives a typical neutrino energy of

$$\epsilon_\nu^{ad} = 4.5 \times 10^{17} \text{ eV } T_2 \eta_2^2 . \qquad (26)$$

However, afterglow evolution (eq. [22]) is shallower than either of these; thus during the afterglow the energy of neutrinos produced becomes higher. This can be checked by computing the highest γ_μ^{obs} for the afterglow evolution, equation (22): I find $\gamma_\mu^{obs} \propto T^{1/4}$. The normalization is such that the highest neutrino energy occurs on the top solid line of Figure 1: inserting equation (16) into equation (14) and considering that typical neutrinos carry away 0.05 of the proton energy, I find

$$\epsilon_\nu = 5 \times 10^{18} \text{ eV } E_{51}^{1/3} \theta_{-1}^{-2/3} \xi^{1/3} . \qquad (27)$$

In Figure 2 I show, as dashed lines, tracks of constant energy for neutrinos produced by muon decay only in the observer frame, together with the previously defined afterglow tracks (eq. [22], dashed lines). The bottom solid line, equation (23), marks the boundary between the region where adiabatic losses (above) and that where synchrotron losses dominate (below). The top solid line marks the region (below the line) where adiabatic muon losses limit neutrinos' energies. At fixed T_2, the neutrino energy is not monotonic with η_2, because, above the region where adiabatic losses operate, each neutrino carries away a fraction (≈ 0.05) of the emitting proton energy, which is a decreasing function of η_2 for fixed n_1 (eq. [14]). In either case, i.e., synchrotron or adiabatic losses, we see that it is still possible to produce high-energy neutrinos, of order 10^{19} eV.

I showed in Paper I that the fraction f_π of the photon luminosity L_γ that goes into neutrinos is

$$f_\pi = 0.01 \eta_2^{-4} \frac{L}{10^{51} \text{ ergs s}^{-1}} \frac{300 \text{ keV}}{\epsilon_\gamma} \frac{1s}{T} , \qquad (28)$$

where $\epsilon_\gamma \approx 300$–$10^3$ keV is the typical burst observed spectral turnover energy. The important parameter ϵ_γ and its time dependence during the afterglow are not currently observed, nor are they in any way predicted by theory. A simple argument allows us to state that if ϵ_γ scales with the
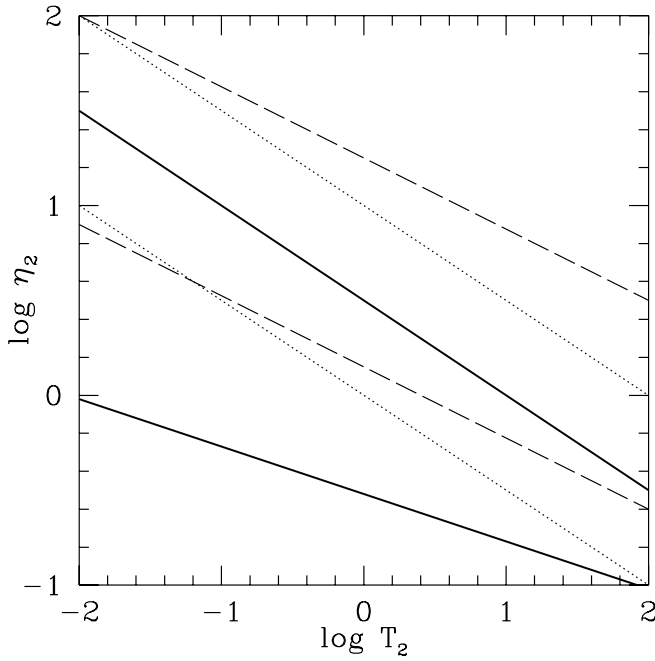
FIG. 2.—Typical energies for neutrinos produced by muon decay as a function of bursts' parameters. The bottom solid line (eq. [23]) separates the region where adiabatic losses prevail (*above the line*) over synchrotron losses. The top solid line (eq. [16]) separates regions where adiabatic losses are important (*below the line*) from that where they are negligible (*above the line*). Dotted lines show loci of constant neutrino energy in the observer's frame. The top dotted line corresponds to $\epsilon_\nu = 1.5 \times 10^{18}$ eV, while the bottom dotted line to $\epsilon_\nu = 4.5 \times 10^{17}$ eV. Dashed lines show representative afterglow tracks (eq. [22]). Since afterglow tracks are shallower than constant neutrino energy tracks, it can be seen that later in the afterglow emitted neutrinos will have larger energies.

shell Lorentz factor η as $\propto \eta^a$ with $a \geq 1$, then f_π remains constant or grows through the afterglow; if f_π remains constant, then the neutrino luminosity $L_\nu = f_\pi L_\gamma \propto t^{-1}$, the last proportionality being observed (Fruchter et al. 1997). Thus, even the bursts that initially suffer important muon losses become, during the afterglow, effective neutrino emitters, and, given that the neutrino luminosity scales only with t^{-1}, the total neutrino fluence will not be much reduced by losses. The reason why I now find that losses are initially important is that when I assume a significant amount of beaming, the maximum proton Lorentz factor (eq. [13]) is greatly increased, while the maximum muon Lorentz factor (eq. [15] for adiabatic losses or eq. [19] for synchrotron losses) is reduced or at most remains constant; it can be seen that in the isotropic case assumed in Paper I the requirement on losses would have been less severe. However, as a partial compensation, for same total energy but significant beaming the maximum proton energy in the observer frame also increases, so that the end result is that,

losses or not, neutrinos with energies exceeding 10^{18} eV are emitted.

5. DISCUSSION

The propagation of UHECRs in intergalactic space is significantly limited by energy degradation via pion photoproduction off the cosmic microwave background radiation (CMBR) photons (Greisen 1966; Zatsepin & Kuz'min 1966). However, it has long been recognized (Wdowczyk, Tkaczyk, & Wolfendale 1972) that observations of high-energy neutrinos, thanks to their negligible cross sections for interaction with matter and photons, will eventually allow us to investigate the sources of cosmic rays way beyond this limit. With the development of AIRWATCH-class experiments (Linsley 1997), capable of monitoring from space atmospheric areas of order 10^6 km^2 in the search for extended air showers and with detection efficiencies close to 1 for neutrinos (L. Scarsi 1997, private communication), it seems the time for extending our search for sources of UHECRs to the whole universe has come.

Two comments are in order. First, as already noted in Paper I, the generic neutrinos (i.e., those produced by cosmic rays that managed to escape unscathed from their sources but that produce pions off CMBR photons) should display dipole and quadrupole moments that, if UHENs come from GRBs, should be totally negligible (Fishman & Meegan 1995). This by itself may help rule out/establish an important alternative idea, that UHECRs have an origin in the Local Supercluster (Stanev et al. 1995). Second, the expected neutrino fluxes computed by Yoshida & Teshima (1993) for cosmologically distributed sources, leading to expected event rates for AIRWATCH-class experiments of 200 yr^{-1}, may be a significant underestimate. In fact, they assumed a distribution of sources of UHECRs arbitrarily limited to a redshift $z_{\mathrm{max}} = 2$. However, the recent, sensational identification of a redshift $z = 3.43$ for the burst GRB 971214 (Kulkarni et al. 1998) makes it clear that significant fluxes of UHENs may be at very large redshifts, thus not violating the constraints on fluxes of UHECRs observed at Earth. Thus the building of AIRWATCH-type experiments with significant neutrino-detection capabilities becomes an even more exciting prospect.

In this paper, I have strengthened the point made in Paper I, that a fraction $f_\pi \approx 0.01$ of all UHENs should correlate (within about a week) with the burst where it was emitted, by showing in some detail that losses do not inhibit the production of high-energy ($\gtrsim 10^{19}$ eV) neutrinos simultaneously with the burst or its afterglow. A search for UHENs ought to be a significant test of the hypothesis that UHECRs are generated in GRBs.

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