1. INTRODUCTION

The study of gamma-ray bursts (GRBs) has been revolutionized due to observations of multiwavelength afterglows in the past few years, but the nature of their progenitors remains unknown (for a review see Mészáros 2000). Two currently popular models of GRB progenitors are the mergers of compact objects (neutron stars or black holes) and the explosions of massive stars. These two cases have distinctive environments for GRBs: compact object mergers occur in the interstellar medium (ISM) and the explosions of massive stars occur in the preburst stellar wind. We here discuss neutrino afterglows from reverse shocks as a result of the interaction of relativistic fireballs with their surrounding wind matter. After comparing with the analytical result of Waxman & Bahcall for the homogeneous ISM case, we find that the differential spectrum of neutrinos with energy from $\sim 3 \times 10^{14}$ to $\sim 3 \times 10^{17}$ eV in the wind case is softer by 1 power of the energy than in the ISM case. Furthermore, the expected flux of upward moving muons produced by neutrino interactions below a detector on the surface of the Earth in the wind case is $\sim 5$ events yr$^{-1}$ km$^{-2}$, which is about 1 order of magnitude larger than in the ISM case. In addition, these properties are independent of whether the fireballs are isotropic or beamed. Therefore, neutrino afterglows, if detected, may provide a way of distinguishing between GRB progenitor models based on the differential spectra of neutrinos and their event rates in a detector.

Subject heading: gamma rays: bursts
The structure of this paper is as follows: In §2 we analyze reverse shocks produced during the interaction of ultrarelativistic fireballs with the surrounding wind matter and discuss the photon emission from these shocks. In §3 we investigate neutrino afterglow emission as a result of photomeson interaction in the reverse shocks, and in §4 we discuss the detectability of such afterglows. In the final section, several conclusions are drawn.

2. SHOCK MODEL AND PHOTON EMISSION

We first assume that a relativistic GRB shell will interact with the surrounding stellar wind via two shocks: a reverse shock and a forward shock. The forward shock runs forward into the wind, while the reverse shock sweeps up the shell material. The recently observed optical flash of GRB 990123 has been argued to come from a reverse shock (Akerlof et al. 1999; Sari & Piran 1999; Mészáros & Rees 1999). We believe that reverse shock emission should be common for all GRBs. The shocked ambient and shell materials are in pressure balance and are separated by a contact discontinuity. We assume that these shocked materials are uniform and move together. Sari & Piran (1995) and Mitra (1998) considered the jump conditions for relativistic shocks and found the common Lorentz factor of the shocked materials measured in the unshocked medium frame,

$$\gamma = \frac{\xi^{1/4} \Gamma^{1/2}}{\sqrt{2}},$$  

(1)

where $\Gamma$ is the Lorentz factor of the shocked shell measured in this frame and $\xi \equiv \rho_w/\rho_{sh}$ is the ratio of proper mass densities of the unshocked shell and the shocked ambient medium. The proper mass density of the ambient medium is expressed as

$$\rho_{\text{w}} = \frac{M_w}{4\pi R^2 V_w} \equiv AR^{-2},$$  

(2)

where $M_w$ and $V_w$ are the mass-loss rate and wind velocity of the progenitor star, $A = M_w/(4\pi V_w) = 10^{-5}$ M$_\odot$ yr$^{-1}/(4\pi \times 10^3$ km s$^{-1})$, $A_s = 5 \times 10^{33}$ g cm$^{-1}$ A$_\odot$, and $R$ is the radius of the shell in units of centimeters (Chevalier & Li 1999, 2000). The proper mass density of the unshocked shell is given by

$$\rho_{sh} = \frac{E_0}{4\pi R^2 \Gamma^2 \xi^{1/2} \Delta},$$  

(3)

where $E_0$ and $\Delta$ are the energy and the width (measured in the unshocked medium frame) of the initial shell. A typical value of $A_s \sim 1$ for Wolf-Rayet stars is found from stellar mass-loss rates and wind velocities (Willis 1991; Chevalier & Li 2000). Since GRBs are believed to come from internal shocks, $\Delta$ is approximately equal to the speed of light times the GRB durations, and thus its typical value should be $\sim 10$ lt-s. The rapid variability of GRBs and their nonthermal spectra require that $\Gamma$ be a few hundreds (Wood & Loeb 1995). From the observed fluences of some GRBs and their measured redshifts, $E_0$ is estimated to be between $10^{52}$ and $10^{54}$ ergs. A recent analysis by Friedman & Waxman (2001) also gives this estimate. Scaling the involved quantities with these typical values, we find

$$\xi = \frac{E_{53}}{A_s \Delta_{10} \Gamma_{300}^2},$$  

(4)

where $E_{53} = E_0/10^{53}$ ergs, $\Gamma_{300} = \Gamma/300$, and $\Delta_{10} = \Delta/10$ is in units of 10 lt-s. From equations (1) and (4), the Lorentz factor of the shocked shell material is rewritten as

$$\gamma = 62 \frac{E_{53}^{1/4}}{A_s^{1/4} \Delta_{10}^{1/4}}.$$

(5)

Following Sari & Piran (1995) and Mitra (1998), we further derive the Lorentz factor of the shocked shell material measured in the unshocked shell rest frame,

$$\gamma' = \frac{\xi^{1/4} \Gamma^{1/2}}{\sqrt{2}} = \frac{1}{2} \gamma = \frac{2.42 A_s^{1/4} \Delta_{10}^{1/4} \Gamma_{300}}{E_{53}},$$  

(6)

which implies that the reverse shock is relativistic. After the reverse shock passes through the shell, the shock front disappears. Instead of maintaining a constant Lorentz factor (e.g., eq. [6]), the shocked materials slow down with time based on the Blandford-McKee (1976) self-similar solution. In the following we discuss photon emission from the reverse shock.

Because of pressure balance across the contact discontinuity, the shocked shell material and the shocked wind material have not only the same bulk Lorentz factor but also the same internal energy density. According to relativistic shock jump conditions, we obtain the internal energy density of the shocked shell material,

$$e' = 4\gamma^2 \rho_w c^2 = 2\rho_{sh} c^2 \xi^{1/2}.$$

(7)

We assume that $\epsilon_e$ and $\epsilon_B$ are the fractions of the internal energy density (in the shocked material rest frame) that are carried by electrons and magnetic fields, respectively. The minimum Lorentz factor of the electrons accelerated behind the reverse shock is approximated by

$$\gamma_{m} \approx (m_p/m_e)\epsilon_e \gamma'$$

(Waxman & Bahcall 2000), viz.,

$$\gamma_{m} \approx 445 \epsilon_e \epsilon_{-1}{\left(\frac{A_s^{1/4} \Delta_{10}^{1/4} \Gamma_{300}}{E_{53}}\right)},$$  

(8)

where $\epsilon_{e,-1} = \epsilon_e/0.1$. Moreover, the magnetic field strength in the shocked shell material is given by

$$B' = (8\pi \epsilon_B \gamma')^{1/2} = \frac{(16\pi \epsilon_B \Gamma_{300}^{21/2})^{1/2}}{t_b},$$  

(9)

where $t_b = R/c$ is the time in the burster’s rest frame. Substituting the relation between this time and the observed time $t = t_b/(1 + z)$, into the above equation and using equations (4) and (5), we further have

$$B' = 3.8 \times 10^{3} \epsilon_{e,-3} \frac{(1 + z)}{2} \left(\frac{\Delta_{10}^{1/4} A_s^{1/4}}{E_{53}^{1/4} t_b}\right) G,$$

(10)

where $\epsilon_{e,-3} = \epsilon_e/10^{-3}$, $z$ is the redshift of the source, and $t_b = t/1$ s. It should be noted that $\epsilon_e \sim 0.1$ (Freedman & Waxman 2001), but $\epsilon_B$ is highly uncertain, and its reasonable value may be taken from $\sim 10^{-2}$ to $\sim 10^{-6}$. Several previous studies of GRB afterglows (e.g., Galama et al. 1999; Dai & Lu 1999, 2000; Wang et al. 2000) give $\epsilon_B \sim 10^{-4} - 10^{-6}$. A recent detailed study of the afterglows of GRBs 980703, 990123, and 990510 by Panaitescu & Kumar (2001) lead to $\epsilon_B \sim 10^{-4} - 10^{-7}$. In addition, Holland et al. (2001) find that $\epsilon_B$ is as small as $10^{-5}$. Therefore, we choose a typical value: $\epsilon_B \sim 10^{-3}$.

Let us consider synchrotron radiation of the electrons accelerated behind the reverse shock. We first derive two characteristic frequencies of synchrotron photons: the typical frequency $v_{m}$ corresponding to the minimum elec-
tron Lorentz factor and the cooling frequency \( \nu_c \). From equations (5), (8), and (10), we obtain the typical frequency in the observer frame,

\[
\nu_m = \frac{\gamma \varepsilon_l^2}{1 + \frac{z}{2}} \frac{eB'}{2\pi n_m c} = 6.3 \times 10^{16} \varepsilon_l^{-2} \left(1 + \frac{z}{2}\right) \frac{A_\star^1 \Gamma_3^3}{E_5^{1/5} t_s} \text{Hz.} \quad (11)
\]

The cooling frequency corresponds to the Lorentz factor \( \gamma_t \), at which an electron is cooling on the dynamical expansion time. We believe that this Lorentz factor in the reverse shock will increase with time because of the cooling timescale \( \propto B'^{-2} \propto t^2 \). Initially, \( \gamma_t - 1 \leq 1 \). At this stage, a cooling electron may be nonrelativistic, and its kinetic energy \( E_c \approx (1/2) \gamma_n c^2 B'^2 \). Thus, its cooling timescale due to cyclotron radiation, measured in the observer’s frame, can be estimated as \( t_0 = (1 + z) E_c / P_{cy} (\beta) \), where \( P_{cy} (\beta) = (4/3) \pi \gamma \varepsilon_l^7 \gamma_t^2 B'^2 / (8 \pi) \) is the cyclotron radiation power in the local observer’s frame of an electron with velocity of \( \beta c \) and with \( \sigma_T \) being the Thomson scattering cross section. Using equation (10), we easily find

\[
t_0 = 1.0n_m c(\beta - z) \left(1 + \frac{z}{2}\right) \frac{A_\star^1 \Gamma_3^3}{E_5^{1/5}} \text{s}. \quad (12)
\]

Please note that \( t_0 \) is independent of \( \beta \). This implies that at \( t < t_0 \) an electron accelerated to \( \gamma_m \) in the magnetic field \( B' \) will be able to cool to become nonrelativistic, initially through synchrotron radiation and subsequently through cyclotron radiation, on the dynamical expansion time \( t \). However, when \( t > t_0 \), the magnetic field \( (B' \propto t^{-1}) \) will become weaker and the cooling timescale due to cyclotron radiation must be longer than \( t_0 \), so an electron with \( \gamma_m \) cannot cool to a nonrelativistic velocity on time \( t \). In this case, cyclotron radiation is no longer a cooling mechanism but should be replaced by synchrotron radiation. Since for typical parameters \( t_0 \approx 1 \text{s} \), which is much less than the durations of long GRBs from the collapse of massive stars, we will discuss the photon and neutrino emission from the reverse shock at \( t > t_0 \) in the remaining text.

According to Sari, Piran, & Narayan (1998), the cooling Lorentz factor is defined by

\[
\gamma_c = \frac{\gamma_{yc} n_m c(1 + z)}{\sigma_T \gamma T B'^2 / (8 \pi)} = 2.0e^{1/3} \left(1 + \frac{z}{2}\right)^{-1/3} \frac{E_5^{1/3} t_s}{A_\star^1 \Gamma_3^3} . \quad (13)
\]

It is clear that \( \gamma_c > 1 \) for \( t > 1 \text{s} \), implying that the cooling electrons are indeed relativistic. Using this equation, we further derive the cooling frequency in the observer frame,

\[
\nu_c = \frac{\gamma_c^2 eB'}{1 + \frac{z}{2} \pi n_m c} = 1.3 \times 10^{12} \varepsilon_l^{-3/2} \left(1 + \frac{z}{2}\right)^{-2} \frac{E_5^{1/2} t_s}{A_\star^1 \Gamma_3^3} \text{Hz.} \quad (15)
\]

We can see from equations (11) and (15) that for typical parameters the cooling frequency is much lower than the typical frequency, indicating that all radiating electrons cool rapidly down to the cooling Lorentz factor. In other words, the shocked shell material is in the fast cooling regime. It is interesting to note that this conclusion has also been drawn by Chevalier & Li (2000). Therefore, the observed specific luminosity peaks at \( \varepsilon_o \equiv h\nu_c \) rather than \( \varepsilon_m \equiv h\nu_m \), with a peak value approximated by

\[
L_{\nu_c} = (2\pi h)^{-1}(1 + z) \gamma_n P_{cy}' \text{s}^{-1} , \quad (16)
\]

where \( N_c = E_0 c^3 / [(1 + z) \Gamma m_\gamma c^2 \Delta] \) is the number of radiating electrons in the shocked shell region and \( P_{cy}' = m_\gamma c^2 \sigma_T B' / (3e) \) is the power radiated per electron per unit frequency in the shocked shell rest frame.

We turn to derive the synchrotron self-absorption frequency of the reverse shocked matter. In the comoving frame of the shocked matter, the absorption coefficient for \( \nu > \nu_c \) is given by

\[
\alpha_\nu = \frac{\sqrt{3}e^3}{8 \pi n_m c^2} \left(\frac{3e}{2 \pi n_m c^2}\right)^{1/2} (m_\gamma c^2)^{p-1} \kappa L B^{p+2/2}
\]

\[
\times \nu^{-(p+4)/2} \Gamma (3p + 2) 12 ! 12 ! \Gamma \left(\frac{3p + 22}{12}\right) , \quad (17)
\]

where \( \kappa = 4(p - 1)\gamma_{yc} (\beta_{yc}/m_\gamma) \) and \( \lambda = (1/2) \int_0^\infty \left(\sin x\right)^{p+2/2} \sin x dx \) (Rybicki & Lightman 1979, p. 147). In the present case, the electron distribution index \( p = 2 \) and the width of the reverse shock \( R = c \tau / (1 + z) \). Setting \( \nu_{yc} \equiv \nu_c R = 1 \), we can derive the synchrotron self-absorption frequency in the observer’s frame,

\[
\nu_a = \frac{\gamma_{yc}}{1 + z} = 2.2 \times 10^{15} \left(1 + \frac{z}{2}\right)^{1/3} \frac{A_\star^1 \Gamma_3^3}{E_5^{1/2}} \text{Hz}. \quad (18)
\]

Hence, \( \nu_c \ll \nu_a \ll \nu_m \) for typical parameters.

We assume that the electrons behind the reverse shock follow a power-law energy distribution, \( \phi_{\nu_0} = \phi_{\nu_0} / \nu_0 \propto \nu^{-2} \) for \( \nu_c > \nu_m \) (Blanford & Eichler 1987). In this case, the synchrotron radiation spectrum is a broken power law,

\[
L_{\nu_c} = \left(\nu_c \phi_{\nu_0} / \nu_0 \right)^{-1/2} \phi_{\nu_0} \Gamma^{3/2} \text{s}^{-1} , \quad (17)
\]

\[
L_{\nu_c} = \left(\nu_c \phi_{\nu_0} / \nu_0 \right)^{-1/2} \phi_{\nu_0} \Gamma^{3/2} \text{s}^{-1} , \quad (19)
\]

where \( \phi_{\nu_0} = h\nu_c \). The protons behind the reverse shock are expected to be accelerated to the same power-law distribution as the electrons (with the maximum proton energy that will be estimated in the next section).

3. NEUTRINO EMISSION

For convenience, we hereafter denote the particle energy measured in the shocked shell rest frame with a prime and the particle energy in the observer frame without a prime; e.g., \( \varepsilon_p = \varepsilon_p / (1 + z) \). We now consider neutrino production in the wind case. Assuming \( n_c(\varepsilon_c) d\varepsilon / d\varepsilon' \) to be the photon number density in the shocked shell rest frame and following Waxman & Bahcall (1997), we can write the fractional energy-loss rate of a proton with energy \( \varepsilon_p \) due to pion
production,

\[ t_\pi^{-1}(\varepsilon_p) = -\frac{1}{\varepsilon_p} \frac{d\varepsilon_p}{dt} \]

\[ = \frac{1}{2T_p} \int_0^\infty d\epsilon \sigma_N(\varepsilon) \xi(\varepsilon) \int_{\epsilon/2T_p}^\infty dx x^{-2} n(x), \quad (20) \]

where \( \gamma_p = \epsilon_p/m_p c^2, \sigma_N(\varepsilon) \) is the cross section for pion production for a photon with energy \( \varepsilon \) in the proton rest frame, \( \xi(\varepsilon) \) is the average fraction of energy lost to the pion, \( \epsilon_0 = 0.15 \text{ GeV} \) is the threshold energy, and the photon number density is related to the observed specific luminosity by \( n(x) = L_{\nu_e}(\gamma)/[4\pi R^2 c(1+z)\gamma]x \). Because of the \( \Delta \) resonance, we find that photomeson production is dominated by the interaction with photons in the energy range \( \varepsilon_1 > \varepsilon_m \). Considering the photon spectrum in this energy range, equation (20) leads to

\[ t_\pi^{-1}(\varepsilon_p) = \frac{L_{\nu_e}}{3\pi(1+z)R^2} \left( \frac{\varepsilon_p}{\varepsilon_m} \right)^{1/2} \frac{\sigma_N \xi \Delta \epsilon}{\epsilon_{peak}}, \]

where \( \sigma_N \xi \Delta \epsilon = 5 \times 10^{-28} \text{ cm}^{-2}, \xi_{peak} = 0.2 \) at the resonance \( \epsilon = \epsilon_{peak} = 0.3 \text{ GeV} \), and \( \Delta \epsilon = 0.2 \text{ GeV} \) is the peak width. The fraction of energy loss of a proton with observed energy \( \varepsilon_p \) by pion production, \( f_p(\varepsilon_p) \), is defined by \( t_\pi^{-1} \) times the expansion time of the shocked shell material (\( \sim R/c \)). Thus, we have

\[ f_p(\varepsilon_p) = 2.0 e_{p,c} \left( \frac{1+z}{2} \right)^{-2} \left( \frac{A_4}{E_0} \right)^{1/2} \left( \frac{E_p}{10^{17} \text{ eV}} \right) \]  

(22)

It is interesting to note that \( f_p(\varepsilon_p) \) is independent of \( \varepsilon_p \) and \( \Gamma \). Similarly to the cooling electron Lorentz factor defined by Sari et al. (1998; see eq. [13]), we can define the cooling proton energy \( \varepsilon_{p,c} \) based on \( f_p(\varepsilon_p) = 1 \). According to equation (22), we find \( \varepsilon_{p,c} \approx 5 \times 10^{16} \text{ eV} \) for typical parameters. This implies that the protons with energy \( \geq \varepsilon_{p,c} \) accelerated behind the reverse shock must lose almost all of their energy (viz., significant cooling) due to photomeson interactions, but the protons with energy less than \( \varepsilon_{p,c} \) lose only a fraction (\( \sim f_p \)) of their energy.

We now turn to discuss the neutrino spectrum. The photomeson interactions include (1) production of \( \pi \) mesons, \( p^+ \rightarrow p + \pi^0 \) and \( p^+ \rightarrow n + \pi^+ \), and (2) decay of \( \pi \) mesons, \( \pi^0 \rightarrow 2\gamma \) and \( \pi^\pm \rightarrow \mu^\pm + v_\mu \rightarrow e^\pm + v_\mu + \bar{v}_\mu + v_\mu \). These processes produce neutrinos with energy \( \lesssim 0.05 \varepsilon_p \) (Waxman & Bahcall 1997). Since the protons with energy less than \( \varepsilon_{p,c} \) lose only a fraction (\( \sim f_p \)) of their energy, the differential spectrum of neutrinos below the break energy \( \sim 3 \times 10^{15} \text{ eV} \) is harder than the proton spectrum by 1 power of the energy. But, since the protons with energy \( \geq \varepsilon_{p,c} \) accelerated behind the reverse shock must lose almost all of their energy, the neutrino spectrum above the break traces the proton spectrum. Therefore, if the differential spectrum of accelerated protons is assumed to be a power-law form \( n(\varepsilon) \propto \varepsilon_p^{-2} \), the differential neutrino spectrum is \( n(\varepsilon) \propto \varepsilon^{-1} \) below the break and \( n(\varepsilon) \propto \varepsilon^{-2} \) above the break.

The maximum energy of the resultant neutrinos is estimated as follows. This energy is determined by the maximum energy of the protons accelerated by the reverse shock. The typical Fermi acceleration time is \( t'_a = fR_L/c \), where \( R_L = (1+z)\varepsilon_p/(\gamma e B) \) is the Larmor radius and \( f \) is of order unity (Hillas 1984). The requirement that this acceleration time is equal to the time for energy loss of protons \( (t'_p) \) due to pion production leads to the maximum proton energy,

\[ \varepsilon_{p,max} = 5.6 \times 10^{18} f \frac{\varepsilon_p}{f} \frac{e^{-1/2}}{e_{v,-3}^{1/2}} \frac{E^{3/8}_{3}}{A^{5/8}_{2}} \frac{\Delta_{3/8}^{3/8}}{eV}. \]

(23)

From this equation, we can draw two conclusions: (1) For reasonable parameters, the maximum proton energy is \( \sim 6 \times 10^{18} \text{ eV} \), which is 2 orders of magnitude smaller than the maximum energy of the protons accelerated by the reverse shock in the ISM case (Waxman & Bahcall 2000). The physical conditions in the reverse shock for the ISM case imply that protons can be Fermi accelerated to \( \sim 10^{20} \text{ eV} \) (Waxman 1995; Vietri 1995; Milgrom & Usov 1995; see Waxman 2000a, 2000b for recent reviews). (2) The maximum energy of neutrinos produced in the wind case is \( \sim 3 \times 10^{17} \text{ eV} \).

**4. DETECTABILITY**

We discuss the detectability of the neutrino afterglow emission in the wind case. Since the protons with energy \( \geq 5 \times 10^{16} \text{ eV} \) must lose almost all of their energy due to photomeson interactions, the present day neutrino energy density due to GRBs is approximately given by \( U_\nu = (1/2)(1/2)U_E \), where the first factor 1/2 accounts for the fact that about one half of the proton energy is lost to neutral pions that do not produce neutrinos, the second factor 1/2 accounts for the fact that about one half of the energy in charged pions is transferred to \( \nu_\mu + \bar{\nu}_\mu \) and \( t_H = 10 \text{ Gyr} \) is the Hubble time. Here we assume that \( E = 0.8 \times 10^{44} \text{ ergs Mpc}^{-3} \text{ yr}^{-1} \) is the production rate of GRB energy per unit volume (Waxman & Bahcall 2000). The neutrino flux is thus approximated by

\[ \phi_\nu = \frac{c}{4\pi} \frac{U_\nu}{E_\nu} \approx 4 \times 10^{-15} \left( \frac{E_\nu}{3 \times 10^{15} \text{ eV}} \right)^{-1} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \]

(24)

The resulting high-energy neutrinos may be observed by detecting the Cherenkov light emitted by upward moving muons produced by neutrino interactions below a detector on the surface of the Earth (Gaissier, Halzen, & Stanev 1995; Gandhi et al. 1998). Planned 1 km\(^2\) detectors of high-energy neutrinos include ICECUBE, ANTARES, and NESTOR (Halzen 2000) and NuBE (Roy, Crawford, & Trattner 1999). The probability that a neutrino could produce a high-energy muon in the detector is approximated by \( P_{\nu_{\mu\mu}} \approx 6 \times 10^{-4} (E_\nu/3 \times 10^{15} \text{ eV})^{0.5} \). Using equation (24), we obtain the observed neutrino event rate in a detector,

\[ N_{\text{events}} = 2\pi \phi_\nu P_{\nu_{\mu\mu}} \approx 5 \left( \frac{E_\nu}{3 \times 10^{15} \text{ eV}} \right)^{-0.5} \text{ km}^{-2} \text{ yr}^{-1}. \]

(25)

This equation shows that a 1 km\(^2\) neutrino detector should detect each year about five neutrinos (with energy of
$\sim 3 \times 10^{15}$ eV) correlated with GRBs. For a GRB, its neutrino emission from the reverse shock in the wind case should be delayed to a few seconds after the main burst. Waxman & Bahcall (2000) have found $f_\nu \sim 0.1$ for neutrino emission from reverse shocks in the ISM case (where the typical energy of neutrinos is $\sim 3 \times 10^{17}$ eV). Using the same expression of $P_{\nu \rightarrow \mu}$, we have rederived their neutrino event rate in a detector and obtained $N_{\nu_{\text{event}}} \sim 0.5$ km$^{-2}$ yr$^{-1}$, which is smaller than our event rate by a factor of $\sim 10$.

5. DISCUSSION AND CONCLUSIONS

Neutrino bursts during the GRB phase were studied in internal shock models by Waxman & Bahcall (1997) and Halzen (1998), who found that the neutrino event rate in a detector (mainly neutrinos with typical energy of a few $10^{14}$ eV) is $\sim 26$ events yr$^{-1}$ km$^{-2}$, which is larger than our event rate by a factor of $\sim 5$. Compared with the analytical result of Waxman & Bahcall (2000), our discussions on neutrino afterglows in the wind case can lead to the following conclusions: (1) The protons with energy $\geq 5 \times 10^{16}$ eV must lose almost all of their energy due to photomeson interactions, and thus the neutrino afterglow emission in the wind case is dominated by neutrinos with energy $\sim 3 \times 10^{15}$ eV. (2) The maximum energy of the protons accelerated behind the reverse shock in the wind case is $\sim 6 \times 10^{18}$ eV, so ultra-high energy cosmic rays cannot be produced in this case. In addition, the maximum neutrino energy is $\sim 3 \times 10^{19}$ eV. (3) The neutrino differential spectrum below $\sim 3 \times 10^{15}$ eV is proportional to $\epsilon^{-1}$, but the spectrum between $\sim 3 \times 10^{15}$ and $\sim 3 \times 10^{17}$ eV steepens by $1 \times 10^{15}$ eV, so ultra-high energy cosmic rays cannot be produced in this case. In addition, the maximum neutrino energy is $\sim 3 \times 10^{19}$ eV. (4) The observed neutrino event rate in the wind case is $\sim 10$ times larger than the one in the ISM case.

If GRB emission is isotropic, the optical afterglow in the wind case must decline more steeply than in the ISM case. This has been suggested as a plausible way of distinguishing between the GRB progenitor models (Chevalier & Li 1999, 2000). It is widely believed that GRBs may come from jets (Kulkarni et al. 1999; Castro-Tirado et al. 1999). As argued by Rhoads (1999) and Sari, Piran, & Halpern (1999), the optical afterglow from a jet is likely to decay more rapidly at late times than at the early stage due to the lateral spreading effect. If this effect is true, however, both ISM and wind cases should show the same emission feature during the lateral spreading phase, and in particular, on a timescale of days the wind density is similar to typical ISM densities, so an interaction with the wind would give results that are not too different from the ISM case (Chevalier & Li 2000; Livio & Waxman 2000). If GRBs are beamed, thus, their optical afterglow emission could not be used to discriminate the massive star progenitor model from the compact binary progenitor model. However, the neutrino afterglow emission discussed here is independent of whether the fireballs are isotropic or highly collimated. Therefore, neutrino afterglows, if detected, may be used to distinguish between GRB progenitor models based on differential spectra of observed neutrinos and their event rates in a detector.

What we want to point out is that the above conclusions are drawn by considering typical values of the parameters entering the fireball shock model. In fact, these parameters may have respective distributions. It is interesting to note that such distributions may lead to an event rate larger than estimated in this paper (Halzen & Hooper 1999).

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