HIGH-ENERGY NEUTRINO SIGNALS FROM THE EPOCH OF REIONIZATION

F. Iocco, K. Murase, S. Nagataki, and P. D. Serpico

Received 2007 July 9; accepted 2007 November 12

ABSTRACT

In this paper we generate a new estimate of the high-energy neutrinos expected from GRBs associated with the first generation of stars in light of new models and constraints on the epoch of reionization and a more detailed evaluation of the neutrino emission yields. We also compare the diffuse high-energy neutrino background from Population III stars with the one from “ordinary stars” (Population II), as estimated consistently within the same cosmological and astrophysical assumptions. In disagreement with previous literature, we find that high-energy neutrinos from Population III stars will not be observable at current or near-future neutrino telescopes, failing below both the sensitivity of a $\text{km}^3$ telescope and the atmospheric neutrino background, also under the most optimistic predictions for the GRB rate. This rules them out as a viable diagnostic tool for these still-elusive metal-free stars.

Subject headings: cosmology: theory — gamma rays: bursts — neutrinos — stars: early-type

Online material: color figures

1. INTRODUCTION

The first generation of stars (Population III, or Pop III stars), born after the collapse of the very first structures in the universe, has puzzled the astrophysical community for a long time. The vanishing metallicity of the universe at the epoch of their formation (Iocco et al. 2007) is supposed to give them peculiar properties, all arising from their very high characteristic mass, due to the peculiar cooling of the cloud (Abel et al. 2002). In fact, stars of $M \sim 10^2 {\text{M}_\odot}$ have very short lifetimes, and either explode as pair-instability supernovae (PISNe) or directly collapse to black holes (Heger et al. 2003). They also contribute to the metallicity enrichment of the early universe, shaping the transition to the second generation of stars. Moreover, the very high temperature of these objects makes them efficient engines for the production of Lyman-Werner ultraviolet photons, thus initiating the cosmic reionization process. Unfortunately, all of the traces Pop III stars leave behind are challenging to observe, and so far no unambiguous detection has been recognized, although the detection of anisotropies in the infrared background consistent with the existence of Pop III has been claimed (Kashlinsky et al. 2005).

Despite the difficulty of detecting neutrinos, they may be potentially interesting messengers of the high-redshift universe, since, unlike gamma rays, neutrinos travel unimpeded over cosmological distances. Also, the peculiar initial mass function (IMF) of Pop III stars should increase their supernova (SN) rate with respect to later stellar populations. The peculiarity of Pop III stars may well produce high-energy neutrinos above the expectations for Pop II stars, thus making them an intriguing target for neutrino telescopes. A fairly general expectation is that Pop III stars may emit large amounts of neutrinos, produced both thermally at the time of the collapse (Iocco et al. 2005) and nonthermally during a gamma-ray burst (GRB) phase associated with the explosion (Schneider et al. 2002). GRBs are one of the candidate hadronic accelerators, and any proton accelerator is a potential emitter of high-energy neutrinos, produced via hadronic ($p\overline{p}$) or photo-hadronic ($p\gamma$) reactions in the surrounding medium. Although the details of this whole picture are far from being established, and despite the large uncertainties in the models, the idea of using neutrinos to probe Pop III stars seems promising, and surely deserves further study.

In this paper we perform a new estimate of the high-energy neutrinos expected from GRBs associated with the first generation of stars, in light of new models and constraints on the epoch of reionization (EoR) and a more detailed evaluation of the neutrino emission yields. The goal of this paper is twofold: (1) to compare the diffuse high-energy neutrino background from Pop III with the one from “ordinary” (Pop II) stars that contribute with Pop III to the reionization at $7 \lesssim z \lesssim 12$; and (2) to discuss the chances of detection in light of the performances expected for current or future neutrino telescopes as a function of the astrophysical input. Indeed, both conditions need to be fulfilled to establish whether there is any realistic chance of discerning the high-energy neutrino emission from Pop III stars. We anticipate that we will find that high-energy neutrinos from Pop III stars will not be observable at current or near future neutrino telescopes, falling below both IceCube sensitivity and atmospheric neutrino background under the most optimistic assumptions for the GRB rate. The disagreement with previous literature is mostly due to unrealistically high Pop III GRB rates previously considered, followed, e.g., from values of the ionization efficiency nowadays considered too extreme based on self-consistent reionization scenarios.

The plan of this paper is as follows. After introducing the basic formalism in § 2, we devote § 3 to describing the models used to perform theoretical estimates of the neutrino fluxes, and here we also briefly discuss what a semiempirical estimate of the GRB rate at high redshift can tell us (§ 3.3). We shall pay particular attention to comparing estimates derived consistently within the same cosmological assumptions. In § 4 we present our results and compare with the chances of detection, and finally in § 5 we conclude. In the Appendix we discuss in some detail the GRB model used to compute the neutrino yields.

2. BASIC FORMALISM

In this section we establish the formalism we will use throughout the paper to estimate the diffuse flux of neutrinos emitted...
by GRBs. The integrated signal observed today at energy $E_\nu$ is
\[ E_\nu^2 \Phi_\nu [\text{GeV cm}^{-2} \text{s}^{-1} \text{sr}^{-1}] =\]
\[ E_\nu^2 \frac{dF_\nu}{dE_\nu} = E_\nu \int dz \frac{d\dot{N}_{\text{GRB}}(z)}{dz} \frac{dN_{\nu}^{\text{iso}}}{dE_\nu} (E_\nu, z), \]
(1)
where $\dot{N}_{\text{GRB}}/dz d\Omega$ is the differential rate of GRBs that beam toward us per unit solid angle, and $dN_{\nu}^{\text{iso}}/dE_\nu d\Omega$ is the average flux emitted by a single source at energy $E_\nu(1+z)$. The function $d\dot{N}_{\text{GRB}}/dz d\Omega$ can be written as
\[ \frac{d\dot{N}_{\text{GRB}}}{dz d\Omega}(z) = \frac{1}{4\pi} \frac{\rho_{\text{GRB}}(z) dV}{(1+z) \frac{dV}{dz}} = b \frac{G(z) dV}{4\pi (1+z) \frac{dV}{dz}}, \]
(2)
where $\rho_{\text{GRB}}(z)$ is the GRB rate in a comoving volume and $dV$ is the comoving volume element, such that
\[ \frac{dV}{dz} = 4\pi \frac{c}{H(z)} r^2(z). \]
(3)
We have introduced the comoving distance $r(z)$, defined as
\[ r(z) = \int_0^z \frac{c}{H(w)} dw. \]
(4)
The Hubble function $H(z)$ can be written in terms of the fractions of the critical energy density in matter $\Omega_M$ and cosmological constant $\Omega_\Lambda$ as
\[ H(z) = H_0 \sqrt{\Omega_M (1+z)^3 + \Omega_k (1+z)^2 + \Omega_\Lambda}, \]
(5)
$H_0$ being the inverse Hubble distance today, $H_0^{-1} \approx 4$ Gpc, and $\Omega_k = 1 - \Omega_M - \Omega_\Lambda$. We also represent by $b = (1 - \cos \theta_{\text{jet}})$ the averaged beaming factor of a jet of opening angle $\theta_{\text{jet}}$, so that $G(z)$ is the beaming-corrected overall GRB rate. The quantity $b$ also fixes the ratio between the physical energy $E_{\text{jet}}$ and the energy released by the GRB if that were isotropic, $b/4\pi = E_{\text{jet}}/E_{\nu}^{\text{iso}}$. If we introduce the function $J_\nu$ of the energy at the source $E_\nu^{\text{iso}}$,
\[ J_\nu[E_\nu^{\text{iso}}] \equiv E_\nu^{\text{iso}} \frac{dN_{\nu}^{\text{iso}}}{dE_\nu} (E_\nu^{\text{iso}}), \]
(6)
the average energy flux, $E_\nu^2 (dN_{\nu}^{\text{iso}}/dE_\nu d\Omega)$ (units GeV cm$^{-2}$) observed at a given energy $E_\nu$ and emitted by a source at redshift $z$ can be expressed as
\[ E_\nu^2 \frac{dN_{\nu}^{\text{iso}}}{dE_\nu d\Omega}(E_\nu^{\text{iso}}) \approx \frac{1}{1+z} \frac{J_\nu[E_\nu(1+z)]}{4\pi r^2(z)}. \]
(7)
In the above expression, $E_\nu(1+z)$ is the emission energy, $J_\nu[E_\nu(1+z)]$ takes into account the redshifted energy spectrum, and possible boost factors are already included in $J_\nu$.

Putting the two pieces together, we finally get
\[ E_\nu^2 \Phi_\nu = \frac{cb}{4\pi} \int dz \frac{J_\nu[E_\nu(1+z)] G(z)}{(1+z)^2 H(z)}. \]
(8)
It remains to estimate $G(z)$ and $J_\nu(E)$, a problem which we shall address in the following section.

### 3. THEORETICAL ESTIMATES OF THE NEUTRINO FLUX

Let us now estimate the GRB rate at high redshifts from current theoretical models of the stellar populations at high redshift, under the assumption that the GRB rate tracks the star formation rate (SFR). In addition to the SFR for the different stellar populations, it is clearly necessary to know their IMFs in order to evaluate the fraction of stars that are likely to end their lives as collapsars, believed to give rise to GRBs. In this section we present the neutrino fluxes expected from the two different populations of stars under different EoR models available in literature. The three models of reionization we consider are summarized in Table 1: the “fiducial” model presented in Choudhury & Ferrara (2005) Figure 1 (p. 586, hereafter CF05a); the one presented in Figure 4 of the same paper (p. 590, hereafter CF05b); and the model in Choudhury & Ferrara (2006, hereafter CF06), which takes into account the new cosmological data on reionization from the third-year data release of the WMAP team (Spergel et al. 2007). These models are used below to derive the star formation rates needed for our estimates. All the CF models share the same physical assumptions: three radiation sources are taken into account during EoR: quasi stellar objects (QSOs), and Population II and Population III stars, which contribute with different characteristics to the radiative and chemical enrichment of the IGM, and to the consequent feedback. Among the many parameters of the models, a key role is played by the fraction of escaping photons per halo (directly related to the ionizing efficiency), which must be deduced from knowledge of the source and halo characteristics of Pop II and Pop III stars; the luminosity function of QSOs is calculated according to standard assumptions. Both the CF05 models assume that all Pop III stars have a mass of 300 $M_\odot$, while Pop II follow a Salpeter IMF, $S(m)$. It is also worth mentioning that the transition between Pop III and Pop II stars, regulated by the chemical feedback, is considered to be instantaneous. The only difference between CF05a and CF05b is the ionizing efficiency of Pop III, which in the latter model is assumed to be 2/7 of the former. This drastically changes the normalization and shape of the Pop III SFR. With respect to the CF05 models, in CF06 the stellar formation in minihalos is suppressed, and a self-consistent calculation is made of the chemical enrichment process, which takes longer and results in a noninstantaneous transition between Pop II and Pop III. For our purposes, however, the crucial difference is that CF06 assumes a Salpeter IMF for Pop III stars as well. The authors implement this choice in order to match the NICMOS high-redshift source counts. We address the reader to the original papers for more details. Also note that none of the models provides a low-redshift $z \lesssim 3$ extrapolation of the SFR. Therefore, for the neutrino fluxes from Pop II shown for comparison in the following, at low-redshift we have used the GRB rate given in equation (12c) of Murase & Nagataki (2006a) normalized to match the observed GRB rate at $z = 0$.
3.1. Rate of GRBs from Pop II

If we denote by \( R_{II} \) the SFR of Pop II stars (taken from the models reported in Table 1), an estimate for the number density rate of GRBs from Pop II stars, \( G_{II} \), can be obtained as

\[
G_{II}(z) = \sigma_{II} \gamma_{II} \frac{R_{II}(z)}{M_{II}},
\]

where \( M_{II} \) is the average mass of a Pop II star, \( \sigma_{II} \) is the fraction of Pop II stars that are likely to form core-collapse SNe, and \( \gamma_{II} \) is the fraction of core-collapse SNe that are likely to form a GRB. The average mass of a Pop II star is obtained via an integral over their IMF, \( \mathcal{I}_{II} \), as

\[
M_{II} = \int \mathcal{I}_{II}(M) dM \int \mathcal{I}_{II}(M) dM,
\]

assuming (consistent with all the EoR models considered) that \( \mathcal{I}_{II}(M) \) is a Salpeter-like mass function \( S(M) \), which can be written as (Kroupa 2001)

\[
S(m) \propto m^{-\alpha} \left\{ \begin{array}{ll}
\alpha_0 = 0.3, & 0.01 < m < 0.08, \\
\alpha_1 = 1.3, & 0.08 < m < 0.50, \\
\alpha_2 = 2.3, & 0.50 < m,
\end{array} \right.
\]

where \( m \) is the mass in solar units, \( M_{\odot} \). This is also used to estimate \( \sigma_{II} \); assuming that all Pop II with \( m > 10 \) will end their lives as SNe,

\[
\sigma_{II} \approx \frac{\int_{10}^{125} S(m) dm}{\int_{10}^{125} S(m) dm}.
\]

We allow for a \( z \)-dependence in \( \gamma_{II} \) to take into account the GRB-metallicity anticorrelation and the fact that the metallicity evolves with the age of the universe. Following Yuksel & Kistler (2007) and Yoon et al. (2006),

\[
\gamma_{II}(z) = \gamma_{II}(0)(1+z)^{4} = \frac{(1+z)^{4}}{1250},
\]

where we fixed \( \gamma_{II}(0) \) according to Yoon et al. (2006), and \( \gamma_{II}(z) \) is considered constant for \( z > 4.5 \), implying no metallicity evolution above this redshift (however, this choice affects the final results very little, most of the contributions to the GRB neutrino background coming from bursts that occur at \( z \leq 7 \)).

3.2. Rate of GRBs from Pop III

Formally, the number density rate of GRBs from Pop III stars, \( G_{III} \), can be written identically to equation (9), namely

\[
G_{III}(z) = \sigma_{III} \gamma_{III} \frac{R_{III}(z)}{M_{III}},
\]

where, however, the various coefficients and functions involved now assume different values. The \( R_{III}(z) \) are taken from the models reported in Table 1. Concerning the IMF of Pop III stars, \( \mathcal{I}_{III} \), there are currently two hypotheses: (1) a bimodal IMF, with a first peak at a few \( M_{\odot} \) and the second one at \( \approx 100 M_{\odot} \) (Nakamura & Umemura 2001); or (2) a peaked IMF around the range of a few hundred \( M_{\odot} \), as presented in Abel et al. (2002). Recently it has been pointed out (O’Shea et al. 2005) that primordial stars forming in a halo neighboring the formation site of a Pop III star would have still zero metallicity but smaller accretion rates onto the central dense region, resulting in a final stellar mass of 10–100 \( M_{\odot} \). This “intermediate” generation would have contributed equally to the reionization, and we include this hypothesis by studying the contribution of \( M \approx 60 M_{\odot} \) in our neutrino models. Whatever the details are, hypothesis (2) is now widely accepted in the community, or in any case it is the \( \mathcal{I}_{III} \) most frequently implemented in EoR models. We shall consider the following cases:

A. \( \mathcal{I}_{III} \propto \delta(M_{III} - 300 M_{\odot}) \),

i.e., a toy model of a single, high-mass mode for the IMF. B. \( \mathcal{I}_{III} \propto \delta(M_{III} - 60 M_{\odot}) \),

i.e., a toy model of a single-mass IMF, with the average mass of metal-free stars forming from the collapse of halos triggered by the explosion of the very first stars. C. \( \mathcal{I}_{III} \propto S(M) \),

as assumed, e.g., in Choudhury & Ferrara (2006).

Note that in both case A and case B one has \( \gamma_{III} \approx 1 \), while in case C one obtains an expression analogous to equation (12).

Concerning \( \gamma_{III} \), there are no firm estimates. For the sake of clarity and to simplify the comparison with previous results of Schneider et al. (2002), we assume the same value,

\[
\gamma_{III} \approx 1.
\]

This corresponds to an upper limit to the Pop III GRBs; a smaller GRB efficiency would result in a down-scaling of the results presented below.

3.3. An Empirical Estimate of the GRB Rate

In this framework, we think it is appropriate to briefly comment on a different, empirical approach, adopted by Yonetoku et al. (2004) and Fenimore & Ramirez-Ruiz (2000) to estimate the GRB rate, which they derive from observations and extrapolate to high redshifts. This approach assumes the \( E_{p} \)-luminosity relation (Amati et al. 2002) to infer the redshift of a GRB whose host is at unknown redshift. So, in order to extrapolate the \( E_{p} \)-luminosity relation to high redshifts, one implicitly assumes that although GRBs at high redshift might belong to different stellar populations with different SFRs and IMFs, they constitute “standard candles” carrying no memory of the population of the progenitors they belong to.

The most natural hint of a transition from Population III to Population II would then be a break in the power-law dependence of the GRB formation history on the redshift: one should expect in general that the conditions for the GRB appearance will occur at a different rate in stellar populations with different characteristics. This is manifestly not the case for the GRB rate derived empirically from observations in Yonetoku et al. (2004) and Fenimore & Ramirez-Ruiz (2000); no discontinuity appears at \( 6 < z < 12 \), where a Pop III-Pop II transition would suggest that it should be seen.

This general result might be interpreted as (weak) empirical evidence for the dominance of Pop II progenitors in the overall GRB population at any redshift. However, we anticipate that the empirical GRB rate obtained in this section is roughly in agreement with the one we obtain from the CF05 model, as described in the previous sections.
3.4. Neutrino Yields from GRBs

High-energy neutrinos from GRBs are predicted in various scenarios. The one most frequently discussed is neutrino emission in the internal shocks that produce prompt emission (Waxman & Bahcall 1997). In this scenario, accelerated protons interact with gamma rays via photomeson production, and produce pions and kaons, which decay into neutrinos. The observed gamma-ray spectra are usually represented by a broken power law. As for proton spectra, a first-order Fermi-type spectrum is frequently assumed. This possibility has been studied by many authors. Murase & Nagataki (2006a) did such calculations using Geant4 with experimental data. Murase & Nagataki (2006b) and Murase (2007) studied other possibilities than neutrino emission in the prompt emission phase. They also took into account various cooling processes. One parameter set of their results for the prompt emission (which is used in Achterberg et al. 2007 and Murase 2007) will be shown for comparison in this paper, and denoted “Prompt” in the following discussion.

Another model was first suggested by Meszaros & Waxman (2001) and later extended in Razzaque et al. (2003). The GRB progenitor is usually taken to be a massive star with an He core and H envelope. The central engine of GRBs is still a major open problem in astrophysics; the leading model for (long-duration) GRBs is a huge stellar collapse, leading to the formation of a highly rotating black hole with an accretion disk (collapsar model). The jet, which produces successful or choked GRBs, propagates in stars, intersecting with the stellar envelope before it produces prompt emission. In addition, the observed submillisecond variability allows us to expect that internal shocks can occur at sufficiently small radii, where the Thomson optical depth exceeds unity. Protons can be accelerated in such internal shocks, and they can interact with photons from electrons that are accelerated in internal shocks and/or termination shocks. Photons from such inner radii cannot escape, due to large optical thickness, so they are hidden sources as far as gamma rays are concerned. Only neutrinos would be useful to probe physical processes in this region. Since here we want to estimate the high-energy neutrinos from Pop III GRBs, consistent with the empirical suggestion from § 3.3, we shall only evaluate neutrino emission from such inner radii. Indeed, this emission is also present for choked GRBs, and we stress that such “hidden” models are the ones that best match our choice \( \gamma_{\text{III}} \approx 1 \). In fact, models such as the Prompt model would predict a lower ratio of GRBs, thus creating friction with our choice. Throughout the paper, we take into account neutrino oscillations in vacuum by adopting mixing angles \( \theta_{12} = 0.59, \theta_{23} = \pi/4 \), and \( \theta_{13} = 0 \). The yields shown in the rest of this paper refer to the muon neutrino flavor, which is the best channel for detecting a low-statistics signal at neutrino telescopes, given the long tracks of muons produced in charged-current interactions. For diffuse signals competing with the atmospheric background, for which the separation must be done on the basis of the different energy spectrum, the electron neutrino showers may be an interesting observable (Beacom & Candia 2004). However, showers are only detected if contained in the instrumented volume, which lowers the effective collecting area of the telescope with respect to the muon track case. We shall see that our event rate predictions for muon tracks fall below the sensitivity of a \( \text{km}^3 \) instrument, which makes exploring alternative channels rather hopeless.

Since our knowledge of the GRB engine and Pop III properties is still limited, it is unclear how efficient Pop III stars are as GRB progenitors. But rotation and chemical mixing seem to be a key ingredient toward successful collapsar models (Maeder & Meynet 2000; Heger et al. 2000), and recent simulations by Yoon et al. (2006) suggest that GRBs can be born from Pop III stars, a hypothesis we will assume throughout. We expect that more massive star collapses lead to larger released energies (Heger et al. 2003). Observationally, to ease the energy requirements, it is usually considered that a globally asymmetric, relativistic jet is launched from GRB progenitors. In fact, this picture is suggested by some observations, e.g., Frail et al. (2001). The GHIRLANDA relation (Ghirlanda et al. 2004) implies that bursts can have various jet energies, and some GRBs may have \( E_{\text{jet}} \sim 10^{52} \) erg, which is larger than the frequently used value \( E_{\text{jet}} \sim 1.2 \times 10^{51} \) erg. The isotropic energy is highly uncertain, too, and it may be as high as \( E_{\text{iso}} \sim 10^{54} \) erg; we adopt this value also, to allow direct comparison with previous results in Schneider et al. (2002).

Because of these high uncertainties, we explore a wide range of parameter values for \( E_{\text{jet}} \) and \( E_{\text{iso}} \) as reported in Table 2. The other parameters varied are \( \eta \), a factor determining the maximum acceleration energy of protons, and \( r_{\text{ifs}} \), the radius at which the internal shock occurs, which can be estimated by \( r_{\text{ifs}} \approx 2 \Gamma_{\text{ifs}} c \delta t \), where \( \Gamma_{\text{ifs}} \) is the Lorentz factor and \( \delta t \) is the typical observed time variability of the jet. Further details of the model and its parameters are given in the Appendix and references therein.

The radial density profiles \( \delta(r) \) for the Pop III precursor mass models, \( M = 200 M_\odot \) and \( M = 60 M_\odot \), have been calculated in Heger & Woosley (2002) and Rockefeller et al. (2006), respectively. We would like to comment on the apparent contradiction between the \( \delta_{\text{III}} \times \delta(M_{\text{III}} - 300 M_\odot) \) discussed in the previous section and the \( M = 200 M_\odot \) model used for the calculation for the neutrino flux. Physically, stars with masses 140 \( M_\odot \) to 260 \( M_\odot \) are known to directly explode as PISNe without forming a black hole, and thus are inefficient GRB progenitors. While important for the GRB rate, however, we expect that the differences in the density profile for the two star masses are much less pronounced in the neutrino yields, and certainly within the uncertainties and assumptions made in such an uneven ground.

In Figure 1 we show the neutrino emission spectra at the redshift of the source \( J_{\nu,\text{III}}(E) \) from the Pop III GRB models considered in Table 2, also compared with the spectrum \( J_{\nu,\text{II}}(E) \) for the Pop II Prompt model.

In the scenario where internal shocks and termination shocks occur at sufficiently small radii inside the progenitor star, almost all accelerated protons are depleted by hadronic and photomeson reactions. This makes the flux features very weakly dependent on \( r_{\text{ifs}} \), provided that the shock occurs where the magnetic field is “weak” in the sense we explain below. The mesons and muons produced suffer from cooling processes: the strong magnetic field and copious photon field make pions and muons cooler before they decay; hence, the neutrino flux will be suppressed due to cooling processes. So, for the energy range between the lower break
energy (obtained in our case by imposing $f_{p,\text{jet}}^{\text{spr}}/\Gamma_{\text{jet}} \sim 1$, where $f_{p,\text{jet}}^{\text{spr}}$ is the fraction of depleted protons) and the higher break energy (which is determined by equating the pion lifetime and its cooling time), the flux is not very sensitive to $r_{\text{is}}$, nor to $\eta_{e}$; in Figure 2, the relative flux level basically reflects the differences among isotropic energies, $E_{\text{iso}}$, of the models. However, the strong field makes mesons and muons lose their energies, so that such suppression becomes important; for example, the difference in the flux level between model E and model F is reduced, thus showing the effects of the field. Furthermore, the strong field can force the higher break energy to be shifted to lower energies.

This effect also leads to a double-peak structure, as shown in Figure 1 for models A–D. The lower peak corresponds to the contribution from the termination shock. Here, in fact, nonthermal protons accelerated in the internal shocks interact with thermal photons produced by electrons accelerated in the termination shock. On the other hand, the higher peak is due to protons from the internal shocks themselves. If we assume similar equipartition parameters, the termination shock can lead to a stronger field than that of internal shocks. Therefore, neutrinos from this region have lower energies due to pion and muon cooling, and this explains the double-peak structure. However, the prominence of this feature is probably due to the simplicity of the model considered: in realistic situations, the emission regions would not be approximated by the simplest two-zone model (internal shock region and termination shock region), and the gradient of the field strength should be taken into account. Then, the double-peak structure would be smoother, although a detailed treatment of this issue is beyond the scope of this paper.

4. RESULTS AND DISCUSSION

In this section we present our results for the neutrino fluxes expected at the Earth. In Figure 2 we show the contribution of Pop III stars calculated as described in the previous sections, assuming CF05a as a fiducial EoR model. In general, these fluxes present a huge variability due to the large uncertainties in the production parameters, yet they are all much lower than the most optimistic results found in Schneider et al. (2002).

In Figure 3 we compare the high-energy neutrino flux at Earth expected from Pop II stars with that expected from Pop III for different but mutually consistent EoR models, as denoted by the same line style. When accounting for the assumed baryon loading factor of $\xi_{\text{acc}} = 10$ (see Appendix for the role of this parameter), the estimated contribution from Pop II stars agrees with typical results found in the literature, and the dependence on the EoR model used is marginal, reflecting the fact that at low redshifts (where most of the signal comes from) all models must agree with the available data. On the other hand, the situation for Pop III is different. The solid and dashed lines assume the neutrino yields of model A, which as shown in Figure 2 maximizes the neutrino production. The dotted line follows instead from the CF06 model, and would be smoother, although a detailed treatment of this issue is beyond the scope of this paper.

---

$^6$ The slight discrepancy with previous results in the Pop II neutrino flux at Earth is due to the higher GRB rate we use for $z > 3$, arising from the reionization models used here.
where we have implemented a Prompt spectrum for \( J_{\text{III}} \), coherently with the assumption of a Salpeter IMF. We have treated this case as the Pop II one, thus obtaining the fraction of collapsars \( \gamma_{\text{III}} \approx 10^{-2} \), consistent with no metallicity evolution beyond \( z = 4.5 \). This illustrates the high-energy neutrino flux for the case of a hypothetical low-mass Pop III generation, whereas the CF06 SFR coupled with the A model for \( J_{\text{III}} \) gives results almost indistinguishable from CF05b, and we therefore have not plotted it. The solid and dashed lines show that the SFR has only a minor effect on the neutrino flux at Earth; on the other hand, a comparison with the dotted line shows that the IMF plays a fundamental role in shaping the flux at Earth, affecting the fraction of stars that will give rise to a collapsar and the magnitude of their explosion energy.

A fairly robust outcome of our study is, however, that neutrinos from Pop III GRBs will not be detectable with current or near future experiments. In fact, the Pop III flux shown in Figure 3 falls below both the current AMANDA-II bound (4 year data; see Halzen 2006) and future prospects for 5 years of IceCube exposure. Even worse, this contribution is buried beneath the atmospheric neutrino background. With reasonable values of the non-thermal baryon loading factor, an extreme contribution from Pop III GRBs—including possible choked bursts—falls in fact below IceCube’s 5 year sensitivity, and is overwhelmed, at low energies, by the atmospheric neutrino background. This implies that although Pop III stars may contribute to the high-energy diffuse neutrino background and that their spectrum would be indeed sensitive to their IMF and SFR, neutrinos cannot be used as a diagnostic tool to check properties of either the population of stars during the epoch of reionization or the GRB internal shock properties, thus confirming the elusive nature of the earliest generation of stars.

5. CONCLUSION

We have studied the high-energy neutrino diffuse background that is to be expected from Pop III stars, under the assumption that they will end their lives as GRBs. We have compared it with the analog emission expected from Pop II stars, using mutually consistent SFRs obtained from epoch of reionization models available in literature. Our estimate of the Pop II GRB rate has been formed using widely accepted assumptions, and leads to estimates of the high-energy neutrino fluxes that are in agreement with previous results (Murase & Nagataki 2006a; Razzaque et al. 2003). On the other hand, we have presented a maximal model for Pop III stars, assuming that all of them will end their lives with a GRB, either choked or not. Even under such optimistic assumptions, a detection of the diffuse high-energy neutrino background expected from Pop III stars appears out of reach. Even worse, the contribution from Pop II GRBs would contaminate that from Pop III GRBs. In addition, we cannot expect neutrino signals to correlate with gamma rays from GRBs, because it is thought to be very rare to see Pop III GRBs with current satellites such as Swift. In addition, the Pop III neutrino signals are expected to be hidden by atmospheric neutrino background. With reasonable values of the non-thermal baryon loading factor, an extreme contribution from Pop III GRBs—including possible choked bursts—falls in fact below IceCube’s 5 year sensitivity, and is overwhelmed, at low energies, by the atmospheric neutrino background. This implies that although Pop III stars may contribute to the high-energy diffuse neutrino background and that their spectrum would be indeed sensitive to their IMF and SFR, neutrinos cannot be used as a diagnostic tool to check properties of either the population of stars during the epoch of reionization or the GRB internal shock properties, thus confirming the elusive nature of the earliest generation of stars.

F. I. thanks A. Heger and C. L. Fryer for kindly providing the stellar models, and A. Ferrara for kind clarifications about EoR models. P. S. acknowledges support by the US Department of Energy and by NASA grant NAG5-10842. This work is in part supported by a Grant-in-Aid for the 21st Century COE “Center for Diversity and Universality in Physics” from the Ministry of Education, Culture, Sports, Science and Technology of Japan. S. N. is partially supported by Grants-in-Aid for Scientific Research from the Ministry of Education, Culture, Sports, Science and Technology of Japan through grants 19104006 and 19740139. K. M. is partially supported by the Japan Society for the Promotion of Science.

APPENDIX

MODELING NEUTRINO EMISSION FROM JETS INSIDE THE GRB PROGENITOR STAR

Here we summarize the basic mechanisms by which neutrinos are produced in jets inside the GRB progenitor star. We assume that a (long) GRB is a huge stellar collapse, leading to the formation of a highly rotating black hole with an accretion disk, during which a globally asymmetric, relativistic jet is launched from the GRB progenitor. GRB observations imply that the
Lorentz factor of jets is very large ($\Gamma_{\text{jet}} \sim 100$) in the gamma-ray-emitting region. This is the final value that is achieved outside the star, and it may be representative of the intrinsic injection Lorentz factor. Of course, we do not know the intrinsic value, and it may in fact be much smaller than the final value. We shall briefly comment on this particular case later. When the jet propogates inside the star, it will produce a bow shock ahead of it. This jet is capped by a termination shock and a reverse shock. The jet termination radius (where the jet is decelerated) is indicated by $r_\text{H}$. If the jet is highly variable and the variability timescale $\delta t$ is small enough, internal shocks can also occur inside the pre-decelerated jet. The internal shock radius is written as $r_s \approx 2 \Gamma_{\text{jet}}^2 c \delta t$. In the usual internal shock scenario of GRBs, the observed gamma rays are attributed to internal shocks that occur in the optically thin region outside the progenitor star. Lower $\Gamma_{\text{jet}}$ and/or $\delta t$ lead to smaller radii below the stellar surface, which will be optically thick. Here, we are interested in such opaque, subsurface internal shocks, and we assume $r_\text{in} \leq r_s < r_\text{r}$, where $r_\text{r}$ is the radius of the progenitor’s stellar surface.

These internal shocks are expected to be collisionless, so that both electrons and protons may be accelerated. If the magnetic field is strong enough, electrons can be accelerated up to high energies and radiate photons. The photon density in the pre-decelerated jet is given by

$$U_{\gamma}^\text{jet} = \frac{F_{\text{sh}}}{4\pi r_\text{in}^2 \Gamma_{\text{jet}}^2 \Delta l}.$$  \hspace{1cm} (A1)

where $\Delta l$ is the width of subshells, for which we use $\Delta l \approx r/\Gamma_{\text{jet}}$, because we assume that the jet acceleration has already ceased. The magnetic energy density in the field $B$ is expressed as

$$U_B^\text{jet} = \frac{B^2}{8\pi} = \epsilon_B^\text{jet} U_{\gamma}^\text{jet}. \hspace{1cm} (A2)$$

Due to the existence of the strong magnetic field, electrons radiate synchrotron photons. These photons will be thermalized, if the jet is optically thick to Thomson scattering. We assume $\epsilon_B^\text{jet} = 0.1$ and expect that the jet is opaque. Therefore, we can approximate the spectral distribution of the radiation of energy density $U_B^\text{jet}$ by a blackbody spectrum, and associate with it a temperature $T_{\gamma}$. The above approximate treatment is accurate enough for our purpose.

Protons get accelerated in the shocks as well. It is widely believed that cosmic rays can be accelerated by the first-order Fermi acceleration mechanism. We assume that this mechanism can work efficiently, and adopt the spectral index $\sim 2$. The nonthermal proton energy density can be expressed as $U_{p}^\text{jet} = \xi \text{acc} U_{\gamma}^\text{jet}$, where $\xi \text{acc}$ is the nonthermal baryon loading factor, which for efficient proton acceleration we can express as $\xi \text{acc} \sim 1/\epsilon_e$, where $\epsilon_e$ is the fraction of internal energy carried by the electrons. Because plausible values of this parameter are not yet known, we adopt $\xi \text{acc} = 10$, which corresponds to the assumption that the energy of protons per logarithmic energy bin is comparable to the GRB radiation energy (Waxman & Bahcall 1997; Murase & Nagataki 2006a). Values of $\xi \text{acc}$ that are too large require values of $\epsilon_e$ smaller than usually expected in GRBs (although there is no proof, because we do not know the total explosion energy).

The nonthermal proton spectrum is given by

$$\frac{dN_p^\text{jet}}{d\epsilon_p} = \frac{U_p^\text{jet}}{\ln (\epsilon_p/\epsilon_{\min})} \epsilon_p^{-2}.$$  \hspace{1cm} (A3)

The minimum energy $\epsilon_{\min}$ is set to 10 GeV. Our final results are not so sensitive to the choice of this value. On the other hand, the maximum energy is important for the purpose of knowing neutrino spectra at the highest energies. The maximum energy is determined by the condition

$$\frac{eBc}{\eta \epsilon_{\max}^p} = t_{\text{acc}}^1 \approx \sum_i t_i^{-1},$$ \hspace{1cm} (A4)

where the left-hand side includes the Larmor radius of the proton times the prefactor $\eta[=O(1-10)]$, which depends on the details of acceleration mechanism; $t_{\text{acc}}$ is essentially the acceleration time in the Bohm limit, and the sum on the right-hand side extends over all the energy-loss channel timescales $t_i$.

Accelerated protons can interact with protons themselves or photons, which leads to meson production such as pions and kaons. In this paper, we consider neutrinos produced through photomeson production only. The calculation method is described in Murase & Nagataki (2006a) and Murase (2007). Here, we briefly sketch this process by simple analytic considerations. The photomeson production is a threshold process, with a threshold energy of about 145 MeV in the rest frame of the incident proton. The dominant inelastic channel is $p\gamma \rightarrow \Delta^+$, with cross section $\sigma_{p\gamma} \approx 5 \times 10^{-28}$ cm$^{-2}$ and inelasticity $\kappa_p \approx 0.2$, which is called $\Delta$-resonance. For the cases we consider, the $p\gamma$ optical depth at the $\Delta$-resonance is very large, so that the photomeson production efficiency can also be high. This means that almost all the protons that have sufficiently high energies (above the threshold energy) will be depleted due to photomeson production.

The threshold for $pp$ inelastic interactions is lower, so these processes will also occur. Although this may be an important neutrino source in the TeV energies (in particular for low Lorentz factors $\Gamma_{\text{jet}}$), in this paper we focus on high-energy neutrinos produced by sufficiently high-energy protons above the threshold energy for photomeson production. We take into account the $pp$ process only for estimating the proton energy-loss timescale, treating this process analytically for simplicity.

Next, we consider the interaction between the jet and the progenitor star. As the jet advances through the star, it drives a bow shock ahead of it. The jet is capped by a forward shock, and a reverse shock moves back into the jet, where the relativistic jet is decelerated.
The shocked jet plasma and shocked stellar plasma would advance together with a jet head Lorentz factor \( \Gamma_h \ll \Gamma_{\text{jet}} \). By equating the pressures behind the forward and reverse shocks, one finds the following estimate for the Lorentz factor of the jet head,

\[
\Gamma_h \simeq \Gamma_{\text{jet}}^{1/2} \left( \frac{m_p n_f^{\text{jet}}}{4 \delta} \right)^{1/4},
\]

(A5)

where \( m_p \) is the proton mass and \( \delta \) is the mass density of the environment.

The relative Lorentz factor between the shocked jet plasma and unshocked jet plasma is

\[
\Gamma' \simeq \frac{1}{2} \left( \frac{\Gamma_{\text{jet}}}{\Gamma_h} + \frac{\Gamma_h}{\Gamma_{\text{jet}}} \right).
\]

(A6)

Electrons will be accelerated in these shocks. However, the electrons would give up all their energy on a very short timescale, by synchrotron and inverse-Compton (IC) cooling, converting a large fraction of the shocked plasma internal energy into radiation. These radiated photons in the shocked jet plasma will be thermalized due to the large optical thickness. Hence, the target photon density will be approximately a blackbody radiation, with an overall energy density

\[
U_{\gamma}^h \simeq (4 \Gamma' + 3)(\Gamma' - 1)n_f^{\text{jet}} m_p c^2,
\]

with an associated temperature \( T_h \).

The reverse shock is likely to become radiation dominated, so that the IC cooling by the electrons becomes important and affects the dissipation of jet kinetic energy. If the reverse shock is indeed radiation dominated, the shock thickness would also be of order the mean-free path of thermal photons propagating into the jet. The copious photon field also affects the neutrino spectrum. Charged mesons produced via \( p\gamma \) (and \( pp \)) interactions will suffer IC and synchrotron losses, and neutrino spectra will be suppressed. The magnetic energy density in the shocked jet frame is expressed as

\[
U_B^{h} = \zeta_B \zeta_\gamma T_h.
\]

(A8)

For \( \zeta_B = \mathcal{O}(0.1–1) \), the magnetic field strength is very large, up to \( \sim 10^{2}–10^{8} \) G. Hence, charged particles would suffer from synchrotron loss, and one should take synchrotron losses of protons into account in calculating neutrino spectra, as we do in our computation.

If nonthermal protons that are accelerated in internal shocks are not completely depleted, they can enter the shocked jet. Such protons are expected to interact with photons in the shocked jet region, and to produce neutrinos. In principle, protons can be accelerated at the termination shocks as well. However, cooling processes such as IC loss will significantly reduce the maximum proton energy there. Hence, we treat the contribution from protons that are accelerated in internal shocks only. The proton spectrum in the frame of the shocked jet plasma is given by

\[
\frac{dn_p^h}{d\varepsilon_p} \simeq (1 - f_p^{\text{jet}}) \Gamma'^2 \frac{dn_p^{\text{jet}}}{d\varepsilon_p},
\]

(A9)

where \( f_p^{\text{jet}} \) is a fraction of depleted protons in the internal shocks, which can be estimated as the ratio of the dynamical timescale involved to the typical energy-loss time. As protons approach the reverse shock, they can interact with photons via photomeson production, if proton energy is above the threshold for photomeson production. It turns out that in this region as well, the proton density is high enough to ensure that almost all protons that have sufficiently high energies are depleted due to photomeson production. However, neutrino energy from this region is smaller than that from the pre-decelerated jet.

The \( pp \) interaction also occurs in the shocked jet plasma, and a further neutrino source is nonthermal proton interactions with cold protons in the star after they escape the shocked jet plasma. For more detailed predictions, these processes should also be taken into account. Yet, given the high efficiencies of the photohadronic processes for our typical parameters, we expect them to be at most subleading at the high energies we are mostly interested in.

Finally, note that synchrotron, IC, and adiabatic cooling of pions and muons are important in shaping our final result. The treatment of these cooling processes is also similar to that used in Murase & Nagataki (2006b) and Murase (2007). Analytic considerations of the effects of these cooling processes can be found in Razzaque et al. (2003).

REFERENCES

Abel, T., Bryan, G. L., & Norman, M. L. 2002, Science, 295, 93
Achterberg, A., et al. 2007, ApJ, 664, 397
Amati, A., et al. 2002, A&A, 390, 81
Beacom, J. F., & Caudill, J. 2004, J. Cosmol. Astropart. Phys., 0411, 009
Choudhury, T. R., & Ferrara, A. 2005, MNRAS, 361, 577 (CF05)
Evoli, C., Grasso, D., & Maccione, L. 2007, J. Cosmol. Astropart. Phys., 0706, 003
Fenimore, E. E., & Ramirez-Ruiz, E. 2000, preprint (astro-ph/0004176)
Frali, D. A., et al. 2001, ApJ, 562, L55
Ghirlanda, G., Ghisellini, G., & Lazzati, D. 2004, ApJ, 616, 331
Halzen, F. 2006, preprint (astro-ph/0611915)
Heger, A., Fryer, C. L., Woosley, S. E., Langer, N., & Hartmann, D. H. 2003, ApJ, 591, 288
Heger, A., Langer, N. & Woosley, S. E. 2000, ApJ, 528, 368
Heger, A., & Woosley, S. E. 2002, ApJ, 567, 532
Iocco, F. 2001, AIP Conf. Proc. 921, The First GLAST Symposium, ed. S. Ritz, P. Michelson, & C. A. Meegan ( Berlin: Springer), 484
Iocco, F., Mangano, G., Miele, G., Pisanti, O., & Serpico, P. D. 2007, Phys. Rev. D, 75, 087304
Iocco, F., Mangano, G., Miele, G., Raffelt, G. G. & Serpico, P. D. 2005, Astropart. Phys., 23, 303
Kashlinsky, A., Arendt, R. G., Mather, J. C., & Moseley, S. H. 2005, Nature, 438, 45
Kroupa, P. 2001, MNRAS, 322, 231
Maeder, A. & Meynet, G. 2000, A&A, 361, 101
Meszaros, P. & Waxman, E. 2001, Phys. Rev. Lett., 87, 171102

Iocco, F., Mangano, G., Miele, G., Pisanti, O., & Serpico, P. D. 2007, Phys. Rev. D, 75, 087304
Iocco, F., Mangano, G., Miele, G., Raffelt, G. G. & Serpico, P. D. 2005, Astropart. Phys., 23, 303
Kashlinsky, A., Arendt, R. G., Mather, J. C., & Moseley, S. H. 2005, Nature, 438, 45
Kroupa, P. 2001, MNRAS, 322, 231
Maeder, A. & Meynet, G. 2000, A&A, 361, 101
Meszaros, P. & Waxman, E. 2001, Phys. Rev. Lett., 87, 171102
Murase, K. 2007, Phys. Rev. D, 76, 123001
Murase, K., & Nagataki, S. 2006a, Phys. Rev. D, 73, 063002
———. 2006b, Phys. Rev. Lett. 97, 051101 (2006)
Nakamura, F., & Umemura, M. 2001, ApJ, 548, 19
O’Shea, B. W., Abel, T., Whalen, D., & Norman, M. L. 2005, ApJ, 628, L5
Razzaque, S., Mészáros, P., & Waxman, E. 2003, Phys. Rev. D, 68, 083001
Rockefeller, G., Fryer, C. L., & Li, H. 2006, preprint (astro-ph/0608028)

Schneider, R., Guetta, D., & Ferrara, A. 2002, MNRAS, 334, 173
Spergel, D., et al. 2007, ApJS, 170, 377
Waxman, E., & Bahcall, J. 1997, Phys. Rev. Lett., 78, 2292
Yonetoku, D., Murakami, T., Nakamura, T., Yamazaki, R., Inoue, A. K., & Ioka, K. 2004, ApJ, 609, 935
Yoon, S. C., Langer, N., & Norman, C. 2006, A&A, 460, 199
Yuksel, H., & Kistler, M. D. 2007, Phys. Rev. D, 75, 083004