Research in Astron. Astrophys. **2010** Vol. **10** No. **10**, 943–949
http://www.raa-journal.org *http://www.iop.org/journals/raa*

Research in Astronomy and Astrophysics

LETTERS

Effect of resonant neutrino oscillation on TeV neutrino flavor ratio from choked GRBs

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Received 2010 June 10; accepted 2010 July 19

Abstract In the collapsar scenario of the long duration Gamma-Ray Bursts (GRBs), multi-TeV neutrino emission is predicted as the jet makes its way through the stellar envelope. Such a neutrino signal is also expected for more general "failed" GRBs in which a putative jet is "choked" by a heavy envelope. If the $\nu_{\rm e} \rightarrow \nu_\mu$ neutrino oscillation parameters are in the atmospheric neutrino oscillation range, we show that the resonant oscillation of $\nu_{\rm e} \leftrightarrow \nu_{\mu,\tau}$ can take place within the inner high density region of the choked jet progenitor with a heavy envelope, altering the ν flavor ratio on its surface to $\Phi^s_{\nu_{\rm e}}$:$\Phi^s_{\nu_\mu}$:$\Phi^s_{\nu_\tau}$ =5:11:2. Considering vacuum oscillations of these neutrinos on their way to Earth, the final flavor ratio detected on Earth is further modified to either 1:1.095:1.095 for the large mixing angle solution to the solar neutrino data, or 1:1.3:1.3 for maximal mixing among the muon and tau neutrinos in the vacuum.

Key words: gamma rays: bursts

1 INTRODUCTION

Long duration GRBs (LGRBs) are believed to be associated with deaths of massive stars (Woosley 1993; Paczyński 1998). The evidence in support of such an origin includes associations of several LGRBs with Type Ic supernovae and the prevalence of star forming dwarf host galaxies associated with LGRBs (Woosley & Bloom 2006; Fruchter et al. 2006; Zhang et al. 2009). Observationally, only a small fraction ($\leq 10^{-3}$) of core collapse SNe are associated with GRBs (Berger et al. 2003). They correspond to those jets that break through the stellar envelope and reach a highly relativistic speed (Lorentz factor $\Gamma \geq 100$). Internal shocks are formed in the optically thin regions, and gamma-rays are produced by synchrotron radiation and/or inverse-Compton scattering of Fermi accelerated electrons in these shocks.

On the other hand, it is feasible to envision that a much larger fraction of core collapses may also launch a mildly relativistic jet from the central engine, but the jet never makes its way out of the envelope, due to either a smaller energy budget or a more extended, massive stellar envelope than in a GRB progenitor. In any case, both the successful jets and these "choked" ones can accelerate protons to energies $\geq 10^5$ GeV from the internal shocks well inside the stellar envelope. The interaction between these protons and the ~ 1 keV thermal X-ray photons emitted by the hot cocoon surrounding the jet would generate multi-TeV neutrinos through photopion production (Mészáros & Waxman 2001). For an individual GRB at redshift $z \sim 1$, the predicted upward directed muon event number is 0.1–10 in a km^3 detector (Mészáros & Waxman 2001). There can also be neutrino production due to pp and pn collisions involving relativistic protons from the buried jet and the thermal nucleons from the jet and its surroundings, which can produce more abundant neutrinos for the presupernova stars with a heavy envelope (Razzaque et al. 2004a; Razzaque et al. 2004b). The detection of low luminosity (LL) GRBs, such as GRB 980425 and GRB 060218 (Galama et al. 1998; Campana et al. 2006), suggests that the event rate of gamma-ray dim core collapses is much higher than those associated with high luminosity GRBs (Liang et al. 2007; Murase et al. 2006; Gupta & Zhang 2007). It is conceivable that the gamma-ray "dark" choked GRBs are even more abundant, and would contribute more to the high energy neutrino background.

Several other neutrino mechanisms have been discussed in the literature. The internal shocks that power the prompt gamma-ray emission can produce ~ 100 TeV neutrinos (Waxman & Bahcall 1997; Waxman & Bahcall 1999; Wang & Dai 2009; Murase 2009), which should lag behind the TeV neutrinos. Inelastic collisions between decoupled protons and neutrons during the acceleration of the fireball can power a multi-GeV neutrino signal (Bahcall & Mészáros 2000; Mészáros & Rees 2000), but the predicted flux level is below the atmospheric neutrino background, and hence, difficult to detect. The neutrinos produced in the early afterglow phase have a very high energy ($\sim$EeV) (Waxman & Bahcall 2000; Dai & Lu 2001) and are not optimized for detection with current neutrino detectors. The multi-TeV neutrino signal discussed here could be above the atmospheric background, and may be detectable with km^3 detectors for some nearby sources.

The main source of high energy neutrinos is the decay of charged pions, which leads to the neutrino flux ratio at the production site $\Phi^0_{\nu_e}$:$\Phi^0_{\nu_\mu}$:$\Phi^0_{\nu_\tau}$ =1:2:0 ($\Phi^0_{\nu_\alpha}$ corresponds to the sum of neutrino and anti-neutrino fluxes for the flavor α at the source). The vacuum oscillations of these neutrinos on their way to Earth would make the observed ratio 1:1:1. This applies to low energy neutrinos including the TeV neutrinos discussed in this paper. For high energy neutrinos above $\sim$1 PeV, muon energy is degraded before decaying to low energy neutrinos so that high energy neutrinos will be absent. The neutrino flux ratio at the source is modified to 0:1:0, which is further modified to 1:1.8:1.8 at Earth after vacuum oscillations are taken into account (Kashti & Waxman 2005).

Another possibility of modifying the neutrino flavor ratio is the resonant conversion of neutrinos from one flavor to another due to the medium effect. Such an effect is known to be important for solar neutrinos (Mikheyev & Smirnov 1985), and has been discussed in early universe hot plasma (Enqvist et al. 1991), a supernova medium (Sahu & Bannur 2000), and in a GRB fireball (Sahu & Dolivo 2005; Sahu et al. 2009a; Sahu et al. 2009b) and jet (Mena et al. 2007). Here we show that for choked GRBs, the multi-TeV neutrinos discussed by Mészáros & Waxman in 2001 could undergo resonance oscillations in the high density core (typically He core) of the presupernova star, if the neutrino oscillation parameters are in the atmospheric neutrino oscillation range. This would alter the neutrino flavor ratio escaping from the stellar envelope, and hence, the eventual detected flavor ratio on Earth.

2 NEUTRINO OSCILLATION IN THE STELLAR ENVELOPE

As a mildly relativistic jet makes its way through the stellar envelope, internal shocks can develop and can accelerate protons to energy $\sim 10^5$ GeV. These protons would interact with the $\sim$keV thermal X-ray photons to produce ~ 5 TeV neutrinos via the process $p + \gamma \to \Delta^+ \to n + \pi^+ \to n + \mu^+ + \nu_\mu \to n + e^+ + \nu_\mu + \nu_e + \bar{\nu}_\mu$.

Depending on the initial mass and metalicity, the presupernova star can have different compositions with different radii. The LGRB progenitors (Type Ic SNe) have lost the H and most of the He envelopes before explosion. They are too small to have an interesting neutrino oscillation signature. The choked jet progenitors, on the other hand, can retain the He envelope (Type Ib SNe) and even the H envelope (Type II SNe). These presupernova stars are favorable for TeV neutrino production and neutrino oscillation. It is believed that the putative jet is launched along the rotation axis where the centrifugal support is the least, and is powered by either $\nu\bar{\nu}$ annihilation or through some electromagnetic processes. Without exploring the details of jet dynamics, here we parametrically analyze the jet. We assume that the jet has developed and that the TeV neutrinos have been produced at a radius $r_{\rm j} \ll R_*$, where R_* is the radius of the star.

Depending on the energy of the propagating neutrino and the nature of the background, neutrinos can interact with the background particles via charge current (CC) and neutral current (NC) interactions. For a neutrino energy below $E_\nu \simeq M_W^2/2m_{\rm e} \simeq 10^7$ GeV, an electron neutrino can have both CC and NC interactions with normal matter, whereas muon and tau neutrinos can only have NC interactions. The effective potential of NC interactions is the same for all active neutrinos. Since oscillation depends only on the potential difference, for active-active oscillations, the NC contributions cancel out. So only the CC contribution to the neutrino potential, $V = \sqrt{2}G_{\rm F}N_{\rm e}$, is responsible for neutrino oscillations in the medium, where $G_{\rm F}$ is the Fermi coupling constant and $N_{\rm e}$ is the electron number density in the medium. For anti-neutrinos, $N_{\rm e}$ is replaced by $-N_{\rm e}$. Thus for the process $\nu_{\rm e} \leftrightarrow \nu_{\mu,\tau}$, the neutrino potential is $\sqrt{2}G_{\rm F}N_{\rm e}$, while for the process $\nu_\mu \leftrightarrow \nu_\tau$ it vanishes.

Here we consider the simplified picture of two mixed flavor states $\nu_{\rm e}$ and $\nu_\mu(\nu_\tau)$ with the vacuum mixing angle θ and mass square difference Δm^2. In a uniform medium, the evolution of the flavor states is governed by (Kim & Pevsner 1993; Sahu & Bannur 2000; Gonzalez-Garcia & Nir 2003)

$$i\frac{d}{dt}\begin{pmatrix}\nu_{\rm e}\\ \nu_\mu\end{pmatrix} = \begin{pmatrix} V-\Delta\cos 2\theta & \frac{\Delta}{2}\sin 2\theta \\ \frac{\Delta}{2}\sin 2\theta & 0\end{pmatrix}\begin{pmatrix}\nu_{\rm e}\\ \nu_\mu\end{pmatrix}, \tag{1}$$

where $\Delta = \Delta m^2/2E_\nu$, V is the potential difference between $V_{\nu_{\rm e}}$ and V_{ν_μ} (i.e. $V = V_{\nu_{\rm e}} - V_{\nu_\mu}$) and E_ν is the neutrino energy. The transition probability as a function of distance ℓ is given by

$$P_{\nu_{\rm e}\to\nu_\mu(\nu_\tau)}(\ell) = \frac{\Delta^2\sin^2 2\theta}{\omega^2}\sin^2\left(\frac{\omega\ell}{2}\right), \tag{2}$$

with

$$\omega = \left[(V-\Delta\cos 2\theta)^2 + \Delta^2\sin^2 2\theta\right]^{1/2}. \tag{3}$$

Once the neutrinos are produced due to pion decay at a point $r_{\rm j} \ll R_*$, they will propagate away from the star where the medium effect can be substantial. If the density of the medium is such that the condition $\sqrt{2}G_{\rm F}N_{\rm e} = \Delta\cos 2\theta$ is satisfied, then resonant conversion of neutrinos from one flavor to another with maximum amplitude can occur. For anti-neutrinos, the resonance condition can never be satisfied (for the normal neutrino mass hierarchy). Although the oscillation process $\bar{\nu}_{\rm e} \leftrightarrow \bar{\nu}_{\mu,\tau}$ can take place, it will be suppressed.

The critical density for resonance is called the resonance density. For 5 TeV neutrinos, it reads

$$\rho_{\rm R} = (1.32 \text{ g cm}^{-3})\,\frac{\Delta\tilde{m}^2}{E_{\nu,12.7}}\cos 2\theta, \tag{4}$$

where we have $\Delta\tilde{m}^2$ in units of eV^2 and $E_{\nu,12.7}$ in units of $10^{12.7}$ eV. The resonance length is

$$\ell_{\rm R} = \frac{2\pi}{\Delta\sin 2\theta} = 1.24\times 10^9 \text{ cm}\left(\frac{E_{\nu,12.7}}{\Delta\tilde{m}^2}\right)\frac{1}{\sin 2\theta}. \tag{5}$$

Define the stellar radius $r_{\rm R}$ as the radius at which the local density is $\rho_{\rm R}$. The first condition for resonant oscillation is $\ell_{\rm R} < r_{\rm R}$.

If the resonance region is wide enough, the transition can be total. We can define a resonance width for which the amplitude of the probability can be $1/2$ instead of unity. In this case, the width can be given as $\Gamma = 2\Delta m^2\sin 2\theta$. This corresponds to a length scale

$$\delta r_{\rm R} = \frac{2\tan 2\theta}{\left|\frac{1}{N_{\rm e}}\frac{dN_{\rm e}}{dr}\right|_{\rm R}}. \tag{6}$$

For $\delta r_{\rm R} > \ell_{\rm R}$, there can be enough time for $\nu_{\rm e}$ to stay in the resonance region and be converted into $\nu_\mu(\nu_\tau)$. This is the second condition for significant resonant oscillation.

In order to evaluate both conditions, one needs to know the matter density profile in the stars (which determines $r_{\rm R}$ and $dN_{\rm e}/dr$). The density profile of a presupernova star is difficult to probe observationally. Numerical models predict a decreasing density with radius. If convective mixing is not important, there is a sharp decrease in density beyond the He core with radius $r_{\rm He} \sim 10^{11}$ cm and local density $\rho_{\rm He} \sim 10^{-3}$ g cm^{-3}. If convective mixing is important, there is no abrupt transition, and the density profile may be roughly described in the analytical form (Razzaque et al. 2004a; Razzaque et al. 2004b; Matzner & McKee 1999; Waxman & Mészáros 2003)

$$\rho(r) = \rho_0 \left(\frac{R_*}{r} - 1\right)^n . \tag{7}$$

The parameters R_* and ρ_0 depend on the type of star. For example, a blue supergiant (BSG) model for SN 1987A gives $R_* = 3 \times 10^{12}$ cm and $\rho_0 = 3 \times 10^{-5}$ g cm^{-3} (Shigeyama & Nomato 1990). In some models, the He core can extend to 10^{12} cm, and the H envelope can extend to 10^{13} cm (Mészáros & Waxman 2001; Razzaque et al. 2004a).

The exponents $n = 3, 3/2$ correspond to the radiative and convective envelope, respectively. In general, it can vary between 2 and 3 for different numerical models. The condition $\delta r_R > \ell_{\rm R}$ can be re-written as a requirement for n (where $r \sim \ell_{\rm R}$ has been adopted which is relevant for resonant oscillation)

$$n < 2\tan 2\theta \left(1 - \frac{\ell_{\rm R}}{R_*}\right). \tag{8}$$

The resonance length $\ell_{\rm R}$ depends on neutrino oscillation parameters and neutrino energy. Apparently, if $\ell_{\rm R} \geq R_*$, the requirement on n ($n < 0$) is unphysical, and no neutrino oscillation is expected inside these stars. On the other hand, if $\ell_{\rm R} \ll R_*$, the constraint $n < 2\tan 2\theta$ may be satisfied in some stars for some oscillation parameters.

In order to evaluate whether the neutrino oscillation conditions are satisfied, one needs to know the neutrino oscillation parameters in matter. Experimentally these are inaccessible. Only oscillation parameters from the solar and atmospheric neutrino experiments are available. For small neutrino mixing angles, the mixing matrix is almost diagonal and each flavor eigenstate nearly overlaps with one of the mass eigenstates. One may then associate $\nu_{\rm e}$ to ν_1, ν_μ to ν_2, and ν_τ to ν_3. While the solar neutrino oscillation parameters are relevant to $\nu_{\rm e} \to \nu_\mu$ oscillations, the atmospheric neutrino oscillation data mostly correspond to $\nu_\mu \to \nu_\tau$ oscillations, and the corresponding neutrino parameters are θ_{23} and Δm^2_{23}, respectively. It is also very possible that the physical properties of neutrinos in a medium could be different from their vacuum values. For example, a neutrino can acquire mass due to its interaction with the background particles even if we consider it to be massless in the vacuum. Similarly, the mixing properties in matter may not follow the vacuum pattern as measured. Nonetheless, since the oscillation parameters of the solar and atmospheric experiments are the best measured, in the following, we test whether these parameters may allow neutrino oscillations to occur in the progenitor stars of choked GRBs. We do not take it for granted that any of these parameters are operating in the oscillation process $\nu_{\rm e} \leftrightarrow \nu_{\mu,\tau}$ in the choked fireball, but just take the only experimentally available parameters to test the conditions for prominent oscillations. Similar analyses have been carried out before to evaluate the possible oscillation effect in GRB fireballs (Sahu et al. 2009a; Sahu et al. 2009b).

The Solar Neutrino Oscillation (SNO) salt phase solar neutrino data, combined with the KamLand reactor antineutrino results, constrain the neutrino oscillation parameters to be in the regime $6 \times 10^{-5}\,{\rm eV}^2 < \Delta m^2 < 10^{-4}\,{\rm eV}^2$ and $0.64 < \sin^2 2\theta < 0.96$ (Ahmed et al. 2004; Araki et al. 2005), with the best fit parameters $\Delta m^2 \sim 7.1 \times 10^{-5}\,{\rm eV}^2$ and $\sin^2 2\theta \sim 0.69$ at a 99% confidence level. The best fit values give $\rho_{\rm R,SNO} \simeq 5.2 \times 10^{-5}\,{\rm g\,cm^{-3}} E^{-1}_{\nu,12.7}$ and

$\ell_{\rm R,SNO} \simeq 2.1 \times 10^{13} {\rm cm} E_{\nu,12.7}$. We can see that $\ell_{\rm R}$ is larger than R_* for a typical BSG, suggesting that resonant oscillations would not occur for these neutrino oscillation parameters.

On the other hand, the atmospheric neutrino oscillation parameters reported by the Super Kamiokande (SK) Collaboration are in the range $1.9 \times 10^{-3}\,{\rm eV}^2 < \Delta m^2 < 3.0 \times 10^{-3}\,{\rm eV}^2$ and $0.9 \leq \sin^2 2\theta \leq 1.0$ at the 90% confidence level (Ashie et al. 2004), which corresponds to the oscillations of mostly muon neutrinos to tau neutrinos. If we assume that these parameters apply to $\nu_{\rm e} \rightarrow \nu_{\mu,\tau}$ oscillations in matter, we can get the following constraint. We consider the good fit point $\Delta m^2 \sim 2.5 \times 10^{-3}\,{\rm eV}^2$ and $\sin^2 2\theta \sim 0.9$, and get $\rho_{\rm R,SK} \simeq 1.0 \times 10^{-3}\ {\rm g\,cm^{-3}}\ E_{\nu,12.7}^{-1}$ and $\ell_{\rm R,SK} \simeq 5.2 \times 10^{11}\ {\rm cm}\ E_{\nu,12.7}$. For the nominal BSG model discussed in this paper, the stellar radius at which the density is $\rho_{\rm R,SK}$ is $r_{\rm R} \simeq 7.1 \times 10^{11}$ cm for $n = 3$ and $E_{\nu,12.7} = 1$. We can see that the condition $\ell_{\rm R} < r_{\rm R}$ is satisfied for the typical neutrino energy $E_\nu = 5$ TeV. By using the value of $\tan 2\theta$ from the SK neutrino data and $\ell_{\rm R} = \ell_{\rm R,SK}$, we obtain from Equation (8) $n < 4.96$. Known stellar models have n between 2 and 3. This suggests that the second condition is also fully satisfied. We conclude that resonant oscillations of multi-TeV neutrinos can occur within a nominal BSG progenitor for the neutrino oscillation parameters inferred by the atmospheric neutrino data. A similar analysis suggests that the same conclusion applies to other BSG progenitors or He stars with extended envelopes (with R_* up to 10^{12} cm), but does not apply to typical He stars (with $R_* = 10^{11}$ cm), or other more compact stars. Since GRB observations favor associations with Type Ic SNe (for which the He envelope is mostly stripped off), GRBs are not preferred sources for TeV neutrino oscillations. Instead, we identify choked GRBs, especially those with a heavy envelope, as interesting sources for resonant neutrino oscillation.

For a full oscillation, $\nu_{\rm e}$ can oscillate to ν_μ and to ν_τ with equal probability but ν_μ can oscillate only to $\nu_{\rm e}$. On average, we can have $1/3$ survival probability for $\nu_{\rm e}$, ν_μ and ν_τ for each $\nu_{\rm e}$ oscillation, but we have 1/2 survival probability for $\nu_{\rm e}$, ν_μ for the $\nu_\mu \leftrightarrow \nu_{\rm e}$ resonant oscillation. The $p\gamma$ process also produces $\bar{\nu}_\mu$ which does not resonantly oscillate. Putting these results together, on the surface of the presupernova star, the survival probability of each flavor (both neutrino and anti-neutrino) is in the ratio of $(\frac{1}{3}+\frac{1}{2})$:$(\frac{1}{3}+\frac{1}{2}+1)$:$\frac{1}{3} =$ 5:11:2. Above r_R (mostly in the H envelope), the density is much lower than the resonance density, and no further reconversion of neutrinos can take place. So the ratio 5:11:2 is the final neutrino flavor ratio escaping from the star. This ratio is notably different from the nominal 1:2:0 ratio for multi-TeV neutrinos without considering the resonant oscillation effect.

3 NEUTRINO OSCILLATION IN VACUUM

Since these choked GRB neutrino sources are typically at large distances, the TeV neutrinos escaping from the star would undergo vacuum oscillations on their way to Earth. The neutrino flux for a particular flavor α on Earth is given by

$$\Phi_{\nu_\alpha} = \sum_\beta P_{\alpha\beta} \Phi^s_{\nu_\beta}, \tag{9}$$

where $\Phi^s_{\nu_\beta}$ signifies the flux of ν_β at the surface of the He envelope after resonant oscillation, and $P_{\alpha\beta}$ corresponds to the oscillation probability from ν_α to ν_β in the vacuum. For the matter effect on resonant oscillation in the stellar envelope, we have applied a two-flavor neutrino analysis. This is because, on one hand, no ν_τ's are generated in the $p\gamma$ process, and on the other hand, the two-flavor neutrino analysis is simpler, as it depends only on one mass square difference. The limited size of the oscillation baseline (the stellar envelope) also makes the three-flavor oscillation effect unimportant. Such a two-flavor neutrino analysis has been applied in most previous resonant oscillation discussions (Sahu & Dolivo 2005; Gonzalez-Garcia & Nir 2003).

When discussing the vacuum oscillation effect along a long base line from the source to Earth, one needs, however, to fully take into account the three-flavor neutrino oscillation effect. This is demanded by the combined analyses of both the solar and the atmospheric neutrino anomalies. For the best fit to the SNO data from the large mixing angle (LMA) solution, one can take the mixing angles $\theta_{12} = 34^\circ \pm 2.5^\circ$, $\theta_{23} = 45^\circ \pm 6^\circ$, $\theta_{13} = 0^\circ \pm 8^\circ$ and the Dirac phase $\delta = 0$ (Kashti & Waxman 2005; Strumia & Vissani 2005). This gives $P_{ee} \simeq 0.57$, $P_{e\mu} = P_{e\tau} \simeq 0.215$ and $P_{\mu\mu} = P_{\mu\tau} = P_{\tau\tau} \simeq 0.393$. Inserting these probabilities in Equation (9) and considering the error of θ_{12}, the flux ratio at Earth is 1:(1.095$\pm$0.012):(1.095 ± 0.012). On the other hand, if we consider the maximal mixing among the ν_μ and ν_τ in a vacuum, then the $\nu_{\rm e}$ oscillation to ν_μ and ν_τ is largely suppressed. One can then have $P_{ee} \simeq 1$, $P_{\mu\mu} = P_{\tau\tau} = P_{\mu\tau} = 1/2$, and all other transition probabilities are negligible. Using these oscillation probabilities, we obtain the flux ratio at Earth to be 1:1.3:1.3.

4 DISCUSSION

GRBs have a wide redshift distribution (from $z = 0.0085$ to $z = 8.2$). Observations suggest that the nearby low-luminosity GRBs have a local event rate ~ 100 times higher than that of high-luminosity GRBs (Liang et al. 2007; Murase et al. 2006; Gupta & Zhang 2007). If GRB jets become progressively successful in progressively rarer progenitors, it is conceivable that there could be even more choked GRBs in the nearby universe. Although the gamma-ray luminosity becomes progressively smaller as the envelope becomes progressively heavier, the TeV neutrino luminosity may not decrease, and could even follow an opposite trend. Assuming that the choked GRB progenitor has a local event rate similar to that of LL-GRBs, i.e. $\mathcal{R} \sim 200\ {\rm Gpc}^{-3}\ {\rm yr}^{-1}$, one would expect ~ 14 neutrino bursts (without gamma-ray counterparts) per year over the whole sky at $z < 0.1$. If the TeV neutrino luminosities of these events are similar to those of successful GRBs (Mészáros & Waxman 2001; Razzaque et al. 2004a), then each event would have hundreds of TeV neutrinos detected by a km^3 detector such as IceCube. Such a possibility has already been ruled out by the current upper limits placed on the IceCube observations. This suggests that either there are not many nearby choked GRBs, or that the choked GRBs are not as neutrino-luminous as predicted (Mészáros & Waxman 2001). Going to the conservative extreme, i.e. if the neutrino luminosity is correlated with gamma-ray luminosity, then the detected event rate for these nearby neutrino burst sources would be 0.001–0.01 neutrinos per event in a km^3 detector. In that case, it would be essentially impossible for IceCube to detect individual sources. The real detected neutrino event rate may be between these two extreme values. IceCube or a similar detector would be able to detect these neutrino bursts or to place even more stringent upper limits in the near future.

If the nearby neutrino bursts are bright enough, the deviation of the observed neutrino flavor ratio from 1:1:1 may be tested by IceCube or similar detectors. The flavor ratios can, in principle, be deduced from the relative rates of showers, muon tracks, and the unique tau lepton induced signals (Beacom et al. 2003). The possibility of detecting a tau signal by IceCube is low, especially in the multi-TeV energy range. On the other hand, IceCube can distinguish between shower-like events and the μ-track events, although it is hard to identify $\nu_{\rm e}$ and ν_τ through their electromagnetic and hadronic showers. Nonetheless, assuming a flavor-independent neutrino spectrum and $\nu_\mu - \nu_\tau$ symmetry (as is the case in our two predicted ratios), the $\nu_{\rm e}$ fraction may be extracted from the measured muon to shower ratio (Beacom et al. 2003). The 10% difference in the flavor ratio reduces the $\nu_{\rm e}$ fraction from 1/3 to 0.313 (for flavor ratio 1:1.095:1.095). This corresponds to a slight increase in the muon to shower ratio. With the uncertainty (20%) for the nominal diffuse flux $(E^2_{\nu_\mu} dN_{\nu_\mu}/dE_{\nu_\mu} = 10^{-7}\ {\rm GeV\ cm^{-2}\ s^{-1}}$ for one year) adopted in Beacom et al. (2003), the small change in the muon to shower ratio may not be differentiated. If nearby neutrino bursts are bright enough, the flux would be increased and the uncertainty would be reduced significantly. This would make a better case for detecting the flavor ratio change. For the 1:1.3:1.3 ratio, the $\nu_{\rm e}$ fraction is reduced from 1/3 to 0.28, making the muon to shower ratio as high as $\sim$ 3.5 (as compared to ~ 3 for 1:1:1). The effect may be detectable for the putative bright neutrino burst events discussed above, even if they may be very rare.

Since the parameters ($\ell_{\rm R}$, $\rho_{\rm R}$, $r_{\rm R}$, and $\delta r_{\rm R}$) all depend on E_ν, we expect that the flavor ratio would also depend on neutrino energy. This aspect has been extensively discussed in Razzaque & Smirnov (2010).

Acknowledgements This work is partially supported by DGAPA-UNAM (Mexico) Project No. IN101409 and Conacyt project No. 103520 (SS) and by NASA NNX09AO94G and NSF AST-0908362 (BZ).

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