Decaying neutrinos and implications from the supernova relic neutrino observation

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(Dated: Received 27 May 2003; received in revised form 5 July 2003; accepted 10 July 2003)

Abstract

We propose that supernova relic neutrino (SRN) observation can be used to set constraints on the neutrino decay models. Because of the long distance scale from cosmological supernovae to the Earth, SRN have possibility to provide much stronger limit than the present one obtained from solar neutrino observation. Since the currently available data are only the upper limit on the flux integrated over $E_{\bar{\nu}_e} > 19.3$ MeV, the decay models on which we can set constraints is quite restricted; they must satisfy specific conditions such that the daughter neutrinos are active species, the neutrino mass spectrum is quasi-degenerate, and the neutrino mass hierarchy is normal. Our numerical calculation clearly indicates that the neutrino decay model with $(\tau_2/m, \tau_3/m) < (10^{10}, 10^{10})$ [s/eV], where $\tau_i$ represents the lifetime of mass eigenstates $\bar{\nu}_i$, appears to give the SRN flux that is larger than the current upper limit. However, since the theoretical SRN prediction contains many uncertainties concerning a supernova rate in the universe or simulation of supernova explosions, we cannot conclude that there exists the excluded parameter region of the neutrino lifetime. In the near future, further reduced upper limit is actually expected, and it will provide more severe constraints on the neutrino decay models.

PACS numbers: 13.35.Hb; 14.60.Pq; 98.70.Vc; 95.85.Ry

Keywords: Diffuse background; Supernovae; Neutrino decay; Neutrino oscillation

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I. INTRODUCTION

A number of ground-based experiments, which observed atmospheric [1], solar [2, 3, 4], and reactor neutrinos [5], have revealed the nonzero neutrino masses and flavor mixings, i.e., properties beyond the standard model of the particle physics. Fortunately, our current knowledge of the neutrino mass differences and mixing angles further enables us to consider far more exotic properties of the neutrino, such as a nonzero magnetic moment (see Refs. [6, 7] and references therein) and neutrino decay. In this Letter we show that neutrino decay of a particular type may be ruled out or severely constrained by the current and the future supernova relic neutrino observation at the Super-Kamiokande (SK) detector [8].

We consider two-body neutrino decays such as $\nu_i \rightarrow \nu_j + X$, where $\nu_i$ are neutrino mass eigenstates and $X$ denotes a very light or massless particle, e.g., a Majoron. The strongest limits on this decay modes are obtained from the solar neutrino observation by SK [3] (see also Ref. [10]). It was argued that the limit is obtained primarily by the nondistortion of the solar neutrino spectrum, and the potentially competing distortions caused by oscillations as well as the appearance of active daughter neutrinos were also taken into account. However, owing to the restricted distance scale to the Sun, this limit is very weak, $\tau/m > 10^{-4}$ s/eV, and therefore, the possibility of the other astrophysical neutrino decay via the same modes cannot be eliminated. In fact, detectability of the decay of neutrinos from the high-energy astrophysical sources was discussed [11], and it has been concluded that it should be visible by future km-scale detectors such as IceCube, since the neutrino decay strongly alter the flavor ratios from the standard one, $\phi_{\nu_e} : \phi_{\nu_\mu} : \phi_{\nu_\tau} = 1 : 1 : 1$, expected from oscillations alone.

Our strategy is basically the same as that of the previous studies [9, 11], i.e., the enhancement of the electron neutrino events due to the decay is investigated. As a source of neutrinos, we consider supernovae. Two observational results concerning supernova neutrinos exist; one is the well-known neutrino burst from SN 1987A [12, 13], and the other is the recent upper limit on the flux of supernova relic neutrinos (SRN), which is the accumulation of neutrinos from all the past supernovae, by SK [8].\footnote{From this point on, we consider only $\bar{\nu}_e$ at detection because this kind of flavor is most efficiently detected at SK via $\bar{\nu}_e p \rightarrow e^+ n$ reaction.} Original discussions concerning the SN 1987A signal have been already given, including the effect of the neutrino decay as well
as the pure flavor mixing, in the literatures \[14, 15\] (see also Ref. \[16\]). In this Letter, we use the latter one (SRN) for obtaining implications for the neutrino decay. The advantage of this approach compared to the former ones is that the neutrinos must transit over very long distance scale from the cosmological supernovae to the Earth, and much longer lifetimes would be probed \textit{in principle}. The current SRN upper limit is only a factor three larger than theoretical predictions by Ando et al. \[17, 18\] (hereafter AST), which adopted neutrino oscillations using experimentally inferred parameters. Therefore, if some neutrino decay model predicts the SRN flux which is three times larger than the AST prediction, then it is excluded.

This Letter is organized as follows. In Section \[II\] we present formulation for the SRN calculation and adopted models. In particular the detailed discussion concerning the supernova model and the neutrino decay model are given in Section \[II A\] and \[II B\] respectively. How the oscillation and decay changes the neutrino spectrum and flux is qualitatively illustrated in Section \[III\] and the numerically calculated results are given in Section \[IV\]. In Section \[V\] we discuss the various possibilities of the neutrino decay, uncertainties of the adopted models, and future prospects.

\section{II. FORMULATION AND MODELS}

The SRN $\bar{\nu}_e$ flux is calculated by

$$\frac{dF_{\bar{\nu}_e}}{dE_{\bar{\nu}_e}} = c \int_0^{z_{\text{max}}} R_{\text{SN}}(z) N_{\bar{\nu}_e}(E'_{\bar{\nu}_e}, z)(1 + z) \frac{dt}{dz} dz,$$

where $E'_{\nu} = (1 + z)E_\nu$, $R_{\text{SN}}(z)$ is a supernova rate per comoving volume at redshift $z$, and $z_{\text{max}}$ is the redshift when the gravitational collapses began (we assume it to be 5). $N_{\nu}(E'_{\nu}, z) = dN_{\nu}(E'_{\nu})/dE'_{\nu}$ is the number of emitted neutrinos per unit energy range by one supernova at redshift $z$. As the supernova rate, we use the most reasonable model to date, which is based on the rest-frame UV observation of star formation history in the universe by the Hubble Space Telescope \[19\], and the model was also used in AST as “SN1.” In this model, the supernova rate exponentially increases with $z$, peaks around $z \sim 1.5$, and exponentially decreases in further high-$z$ region.
A. Supernova model

As original neutrino spectra $N_\nu^0(E, z)$, which is not the same as $N_\nu$, owing to the neutrino oscillation and decay, we adopt the result of a numerical simulation by the Lawrence Livermore group [20], which is the only group that is successful for calculating neutrino luminosities during the entire burst ($\sim 10 \text{ sec}$). The average energies are different between flavors, such as $\langle E_{\nu_e} \rangle \simeq 11 \text{ MeV}$, $\langle E_{\bar{\nu}_e} \rangle \simeq 16 \text{ MeV}$, $\langle E_{\nu_x} \rangle \simeq 22 \text{ MeV}$, where $\nu_x$ represents non-electron neutrinos and antineutrinos. This hierarchy of the average energies is explained as follows. Since $\nu_x$ interact with matter only through the neutral-current reactions in supernova, they are weakly coupled with matter compared to $\nu_e$ and $\bar{\nu}_e$. Thus the neutrino sphere of $\nu_x$ is deeper in the core than that of $\nu_e$ and $\bar{\nu}_e$, which leads to higher temperatures for $\nu_x$. The difference between $\nu_e$ and $\bar{\nu}_e$ comes from the fact that the core is neutron-rich and $\nu_e$ couple with matter more strongly, through $\nu_e n \leftrightarrow e^- p$ reaction.

Our calculations presented in Section IV are strongly sensitive to the adopted supernova model, or in particular, the average energy difference between $\bar{\nu}_e$ and $\nu_x$. Recently, the Livermore calculation is criticized since it lacks the relevant neutrino processes such as neutrino bremsstrahlung and neutrino-nucleon scattering with nucleon recoils, which are considered to make the mean energy difference between flavors less prominent; it has been actually confirmed by the recent simulations (e.g., Ref. [21]). However, we cannot adopt recent models, even though they include all the relevant neutrino microphysics. This is because the SRN calculation definitely requires the time-integrated neutrino spectrum during the entire burst, whereas all the recent supernova simulations terminate at $\sim 0.5 \text{ sec}$ after the core bounce. Since the Livermore group alone is successful to simulate the supernova explosion and calculate the neutrino luminosity during the entire burst, we use their result as a reference model. Again, we should note that our discussions from this point on heavily relies on the adopted supernova model.

B. Neutrino decay model

In this Letter, we consider only the so-called “invisible” decays, i.e., decays into possibly detectable neutrinos plus truly invisible particles, e.g., light scalar or pseudoscalar bosons. The best limit on the lifetime with this mode is obtained from solar neutrino observations
and is \( \tau/m > 10^{-4} \text{ s/eV} \), which is too small to set relevant constraints on discussions below. We do not consider other modes such as radiative two-body decay since they are experimentally constrained to have very long lifetimes (e.g., Ref. [22]).

We note that our approach is powerful only when the decay model satisfies specific conditions such that (i) the daughter neutrinos are active species, (ii) the neutrino mass spectrum is quasi-degenerate \((m_1 \approx m_2 \approx m_3)\), and (iii) the neutrino mass hierarchy is normal \((m_1 < m_3)\), not inverted \((m_1 > m_3)\). This is because at present, only the upper limit of the SRN flux is obtained, and therefore, the decay model which does not give large flux at detection energy range \((E_{\bar{\nu}_e} > 19.3 \text{ MeV})\) cannot be satisfactorily constrained. All these conditions (i)–(iii) must be satisfied to obtain severe constraints on the neutrino lifetime, because (i) if the daughter neutrinos are sterile species, the SRN flux decreases compared with the model without the neutrino decay; (ii) if the neutrino mass spectrum is strongly hierarchical, then the daughter neutrino energy is degraded compared to its parent and the predicted SRN flux at high energy range would not be large; and (iii) in the case of inverted hierarchy, \(\bar{\nu}_e\) most tightly couple to the heaviest mass eigenstates, which decay into lighter states, and it also reduces the SRN flux.

For a while, we assume that the conditions (i)–(iii) are satisfied; all the other possibilities are addressed in detail in Section V A. As discussed in Section V C, we believe that future observational development would enable far more general and model-independent discussions, which are not restricted by the above conditions.

III. NEUTRINO OSCILLATION AND DECAY

Effects of neutrino oscillation and decay are included in \(N_{\bar{\nu}_e}(E, z)\). We address the problem of flavor mixing by the pure supernova matter effect, first for antineutrinos and second for neutrinos, and then we discuss the neutrino decay.

The state of \(\bar{\nu}_e\) produced at deep in the core is coincident with the lightest mass eigenstate \(\bar{\nu}_1\), owing to large matter potential. This state propagates to the supernova surface without being influenced by level crossings between different mass eigenstates (it is said that there are no resonance points). Thus, \(\bar{\nu}_e\) at production becomes \(\bar{\nu}_1\) at the stellar surface \((N_{\bar{\nu}_1} = N_{\bar{\nu}_1}^0)\), and the number of \(\bar{\nu}_e\) there is given by

\[
N_{\bar{\nu}_e}(E, z) \simeq \cos^2 \theta_\odot N_{\bar{\nu}_1}(E, z) + \sin^2 \theta_\odot N_{\bar{\nu}_2}(E, z)
\]
\[ = \cos^2 \theta_\odot N^0_{\nu_e}(E, z) + \sin^2 \theta_\odot N^0_{\nu_x}(E, z), \]

(2)

where $\theta_\odot$ is the mixing angle inferred from the solar neutrino observations ($\cos^2 \theta_\odot \simeq 0.7$), $N^0_\nu$ represents the neutrino spectrum at production. In the above expression [2], we used the fact, $\sin^2 2\theta_{\text{atm}} = 1$, and assumed $\theta_{13} = 0$. These are justified by the atmospheric [1] and reactor neutrino experiments [23].

The situation changes dramatically for neutrino sector. As the case for antineutrinos, $\nu_e$ are produced as mass eigenstates owing to large matter potential, however, the difference is that the produced $\nu_e$ coincide with the heaviest state $\nu_3$. Since in vacuum $\nu_e$ most strongly couples to the lightest state $\nu_1$, there must be two level crossings (or resonance points) between different mass eigenstates during the propagation through supernova envelope; each is labeled by H- and L-resonance, corresponding to whether the density of the resonance point is higher or lower (see, e.g., Ref. [24] for details). It is well-known that at L-resonance the mass eigenstate does not flip (adiabatic resonance) for LMA solution to the solar neutrino problem. However, the adiabaticity of the H-resonance becomes larger than unity when the parameter $\theta_{13}$ is sufficiently large. Instead, we parameterize the flip probability at the H-resonance by $P_H$, i.e., if the resonance is adiabatic (nonadiabatic), $P_H = 0(1)$. Thus, at the stellar surface, the neutrino spectra of mass eigenstates is given by

\[ N_{\nu_1}(E, z) = N^0_{\nu_e}(E, z), \]
\[ N_{\nu_2}(E, z) = P_H N^0_{\nu_e}(E, z) + (1 - P_H) N^0_{\nu_x}(E, z), \]
\[ N_{\nu_3}(E, z) = (1 - P_H) N^0_{\nu_e}(E, z) + P_H N^0_{\nu_x}(E, z). \]

(3)  
(4)  
(5)

If we include the neutrino decay, the expected $\bar{\nu}_e$ flux from each supernova changes drastically. Before giving a detailed discussion, we first place some simplifying assumptions. Instead of the lifetime, we define “decay redshift” $z^d_i$ of the mass eigenstate $\nu_i(\bar{\nu}_i)$; if the source redshift $z$ is larger than the decay redshift $z^d_i$, all the neutrinos $\nu_i(\bar{\nu}_i)$ decay, on the other hand if $z < z^d_i$, $\nu_i(\bar{\nu}_i)$ completely survive. We consider the decaying mode $\nu_3(\bar{\nu}_3) \rightarrow \bar{\nu}_1$ and $\nu_2(\bar{\nu}_2) \rightarrow \bar{\nu}_1$, and $z^d_2$ and $z^d_3$ are taken to be two free parameters. The other case that one of them is stable can be realized if we take $z^d > z_{\text{max}}$. With these assumptions and parameterization, the neutrino spectrum which is emitted by the source at redshift $z$ can be obtained. First, we consider the decay mode $\nu_i \rightarrow \bar{\nu}_j + X$, in which the neutrino helicity
The $\bar{\nu}_e$ spectrum is given by

$$
N_{\bar{\nu}_e}(E, z) = \cos^2 \theta \cdot N_{\bar{\nu}_1}(E, z) + \sin^2 \theta \cdot N_{\bar{\nu}_2}(E, z)
$$

$$
= \cos^2 \theta \cdot \left[ N_{\bar{\nu}_1}^0(E, z) + N_{\bar{\nu}_2}^0(E, z) \Theta(z - z_d^d) 
+ N_{\nu_3}(E, z) \Theta(z - z_d^3) \right]
+ \sin^2 \theta \cdot N_{\bar{\nu}_2}^0(E, z) \Theta(z_d^d - z)
$$

$$
= \cos^2 \theta \cdot \left[ N_{\bar{\nu}_e}^0(E, z) + N_{\nu_e}(E, z) \times \left\{ \Theta(z - z_d^d) + \Theta(z - z_d^3) \right\} \right]
+ \sin^2 \theta \cdot N_{\nu_e}^0(E, z) \Theta(z_d^d - z),
$$

(6)

where $\Theta$ is the step function. On the other hand, if the relevant decay mode is $\nu_i(\bar{\nu}_i) \rightarrow \bar{\nu}_j(\nu_j) + X$ (helicity flips), the expected spectrum becomes

$$
N_{\bar{\nu}_e}(E, z) = \cos^2 \theta \cdot N_{\bar{\nu}_1}(E, z) + \sin^2 \theta \cdot N_{\bar{\nu}_2}(E, z)
$$

$$
= \cos^2 \theta \cdot \left[ N_{\bar{\nu}_1}^0(E, z) + N_{\bar{\nu}_2}^0(E, z) \Theta(z - z_d^d) 
+ N_{\nu_3}(E, z) \Theta(z_d^d - z) \right]
+ \sin^2 \theta \cdot N_{\nu_e}^0(E, z) \Theta(z_d^d - z).
$$

(7)

Comparing Eqs. (6) and (7) with Eq. (2), the SRN flux with the neutrino decay is expected to be very different from the case of the pure neutrino oscillation. In particular, for the mode $\bar{\nu}_i \rightarrow \bar{\nu}_j + X$ [Eq. (4)], the SRN spectrum is expected to be hard since it contains a fair amount of $\nu_x$. In the case of $\nu_i \rightarrow \bar{\nu}_j + X$ mode, the $\nu_e$ spectrum is also included, and then the corresponding upper limit is not as strong as that for $\bar{\nu}_i \rightarrow \bar{\nu}_j + X$ mode. In the next section, we only consider the model, whose upper limit is the most severe among all the models considered, i.e., the decay mode $\bar{\nu}_i \rightarrow \bar{\nu}_j + X$.

In general, the decay redshifts depend on the neutrino energy since the lifetime at the laboratory frame $\tau_{\text{lab}}$ relates to that at the neutrino rest frame $\tau$ via a simple relation
However, we believe that it does not make sense in discussing this point strictly, because the estimation of the SRN flux contains many other uncertainties as shown in Section VIB. In order to obtain the lifetime of the mass eigenstates $\bar{\nu}_i$ from the decay redshifts, the typical neutrino energy $E = 10$ MeV is assumed with the following formulation:

$$\tau_i = \frac{m}{E} \int_0^{z_d} \frac{dz}{dz} \frac{dz}{dz} \times \frac{1}{\sqrt{(1 + \Omega_m z)(1 + z)^2 - \Omega_\Lambda (2z + z^2)}},$$

where the Hubble constant $H_0$ is taken to be $70$ km s$^{-1}$ Mpc$^{-1}$, and $\Lambda$-dominated cosmology is assumed ($\Omega_m = 0.3, \Omega_\Lambda = 0.7$). Since the exact value of the neutrino mass $m$ is not known, the relevant quantities is the neutrino lifetime divided by its mass, $\tau/m$.

IV. RESULTS

Figure 1 shows the SRN flux for various parameter sets of decay redshifts ($z_d^2, z_d^3$) as a function of neutrino energy, which is calculated using Eqs. (1) and (6). The solid curve in Fig. 1 shows the SRN flux in the case that ($z_d^2, z_d^3$) = (5.0, 5.0), which represents the same flux as that without the neutrino decay; this SRN flux almost coincides with one obtained by the AST calculation. Then, we included the decay of the heaviest mass eigenstate $\bar{\nu}_3$ by reducing the value of $z_d^3$, with keeping $\bar{\nu}_2$ stable. When $z_d^3 = 1.0$ (dotted curve), the SRN flux at low energy region ($E_{\bar{\nu}_e} \lesssim 35$ MeV) deviates from the pure oscillation model (5.0, 5.0). This is because the neutrinos from supernovae at redshift larger than $z_d^3 = 1.0$ are affected by the $\bar{\nu}_3 \to \bar{\nu}_1$ decay and it results in the increase of $\bar{\nu}_e$. Since the neutrino energies are redshifted by a factor of $(1 + z)^{-1}$ owing to an expansion of the universe, the decay effect can be seen at low energy alone. When the value of $z_d^3$ is reduced to $10^{-2}$, the neutrinos even from the nearby sources are influenced by the $\bar{\nu}_3 \to \bar{\nu}_1$ decay, resulting in the deviation over the entire energy range as shown by the long-dashed curve in Fig. 1. If we add the $\bar{\nu}_2 \to \bar{\nu}_1$ decay, it further enhances the SRN flux.

In Table I we summarize the SRN flux integrated over the energy range of $E_{\bar{\nu}_e} > 19.3$ MeV, for the each decay model. In the second column we placed the lifetime-to-mass ratio which corresponds to each decay redshift, which is obtained using Eq. (8). The corre-
FIG. 1: The SRN flux for various parameter sets of decay redshifts. Each label represents \((z_d^2, z_d^3)\).

TABLE I: The predicted SRN flux for various decay models and the corresponding SK limit (90% C.L.). Integrated energy range is \(E_{\bar{\nu}_e} > 19.3\) MeV. The ratio of the prediction and the limit is shown in the fifth column.

| Model \((z_d^2, z_d^3)\) | \((\tau_2/m, \tau_3/m)\) [s/eV] | Predicted flux \(\text{SK limit (90\% C.L.)}\) | Prediction/Limit |
|----------------|----------------|-----------------------------|-----------------|
| \((5.0, 5.0)\) | \((3.9 \times 10^{10}, 3.9 \times 10^{10})\) | 0.43 cm\(^{-2}\) s\(^{-1}\) | < 1.2 cm\(^{-2}\) s\(^{-1}\) | 0.35 |
| \((5.0, 1.0)\) | \((3.9 \times 10^{10}, 2.4 \times 10^{10})\) | 0.55 cm\(^{-2}\) s\(^{-1}\) | < 1.3 cm\(^{-2}\) s\(^{-1}\) | 0.42 |
| \((5.0, 0.2)\) | \((3.9 \times 10^{10}, 7.7 \times 10^9)\) | 0.93 cm\(^{-2}\) s\(^{-1}\) | < 1.2 cm\(^{-2}\) s\(^{-1}\) | 0.75 |
| \((5.0, 10^{-2})\) | \((3.9 \times 10^{10}, 4.4 \times 10^8)\) | 1.0 cm\(^{-2}\) s\(^{-1}\) | < 1.2 cm\(^{-2}\) s\(^{-1}\) | 0.88 |
| \((10^{-2}, 10^{-2})\) | \((4.4 \times 10^8, 4.4 \times 10^8)\) | 1.4 cm\(^{-2}\) s\(^{-1}\) | < 1.2 cm\(^{-2}\) s\(^{-1}\) | 1.2 |
sponding 90% C.L. upper limit given by the SK observation at $E_{\bar{\nu}_e} > 19.3$ MeV and prediction-to-limit ratio are also shown in the fourth and fifth columns, respectively. The predicted flux becomes larger as the decay is included, while the observational upper limit remains unchanged. At first sight, it is expected that the model $(10^{-2}, 10^{-2})$ is already excluded by the current observational data.

![Contour plot](image)

**FIG. 2:** A contour map of a prediction-to-limit ratio of the SRN flux which is projected against the lifetime-to-mass ratios ($\tau_2/m, \tau_3/m$).

Figure 2 shows a contour plot for a ratio of the predicted flux to the observational upper limit, which is projected against the lifetime-to-mass ratios ($\tau_2/m, \tau_3/m$). The area below the solid curve labeled as 1.0 is considered to be an excluded parameter region. We can confirm that our approach is very powerful to obtain the constraint on the neutrino lifetimes, because the best observed lower limit thus far ($\gtrsim 10^{-4}$ s/eV) is much smaller than that shown in Fig. 2 ($\gtrsim 10^{10}$ s/eV), although there are still many uncertainties in this method as discussed in the next section.
V. DISCUSSION

A. Other decaying modes

Our discussions until this point were concerned with one specific decaying mode, $\bar{\nu}_3 \rightarrow \bar{\nu}_1, \bar{\nu}_2 \rightarrow \bar{\nu}_1$, where the daughter neutrinos carry nearly the full energy of their parent. We consider the other possible decaying models; first $\bar{\nu}_3 \rightarrow \bar{\nu}_2, \bar{\nu}_2 \rightarrow \bar{\nu}_1$ is discussed. Since $\bar{\nu}_2$ state contains $\bar{\nu}_e$ state by less fraction than $\bar{\nu}_1$, this decay mode gives smaller SRN flux, resulting in weaker upper limit in general. However, when the lifetime of one mode is much longer (shorter) than the other, i.e., $z_2^d \ll z_3^d$ or $z_2^d \gg z_3^d$, the discussions for our reference decay modes $\bar{\nu}_3 \rightarrow \bar{\nu}_1, \bar{\nu}_2 \rightarrow \bar{\nu}_1$ are basically applicable to this case.

When the daughter neutrino energy is considerably degraded, which is actually the case when the neutrino masses are strongly hierarchical, the observational upper limit for the each model is not as strong as the previous limit shown in Table I. This is because the energetically degraded daughter neutrinos soften the SRN spectrum.

Finally we consider the case that the daughter neutrinos are sterile species which does not interact with matter. If the lifetime of the mode is sufficiently short, then the obtained spectrum from each supernova can be expressed by

$$\mathcal{N}_{\bar{\nu}_e}(E, z) \simeq \cos^2 \theta_{\odot} \mathcal{N}^0_{\bar{\nu}_e}(E, z),$$

(9)

which is smaller than the normal oscillation expression (2). Thus, also in this case, the observational upper limit will be looser. In consequence, all the other possible decaying models give weaker upper limit compared to our reference model.

B. Supernova and supernova rate model uncertainties

Again we restrict our discussion to our standard decay mode, and discuss whether the parameter region $(\tau_2/m, \tau_3/m) \lesssim (10^{10}, 10^{10})$ [s/eV] is really ruled out by the current observational data, as shown in Fig. 2.

The theoretical calculation of the SRN flux contains many uncertainties such as the supernova rate as a function of redshift $z$ as well as the original neutrino spectrum emitted by each supernova. As for the supernova rate, since we have inferred it from the rest-frame UV observation, it may be severely affected by dust extinction, which is quite difficult to
estimate. Thus for example, if we use the supernova rate model which is larger by a factor of 2 (even if this is actually the case, it is not surprising at all), then almost all the region of lifetime-to-mass ratio \((\tau_2/m, \tau_3/m)\) will be excluded as shown in Fig. 2.

The uncertainty concerning the original neutrino spectrum also gives very large model dependence of the SRN flux calculation. Although we adopted the result of the Livermore simulation, as we have noted above it lacks the relevant neutrino processes, which reduce the average energy of \(\nu_x\) to the value close to that of \(\bar{\nu}_e\). If this is the case, the SRN spectrum as a result of the neutrino decay becomes softer compared to that obtained with the Livermore original spectra, which leads to weaker upper limit. In actual, we have calculated the SRN flux for various values of \((z_2^d, z_3^d)\) assuming the original \(\nu_x\) spectrum is the same as that of \(\bar{\nu}_e\), for which we have used the Livermore data. As a result of the calculation for this extreme case, the obtained prediction-to-limit ratio is at most ~ 0.35, even when both of decay redshifts are sufficiently small. In consequence, considering the uncertainties which is included in the supernova rate or the neutrino spectrum models, we cannot conclude that the parameter region \((\tau_2/m, \tau_3/m) \lesssim (10^{10}, 10^{10}) [s/eV]\) is already excluded, although Fig. 2 indicates it.

C. Future prospects

Now, we consider the future possibility to set severer constraint on the neutrino decay models. The largest background against the SRN detection at SK is so-called invisible muon decay products. This event is illustrated as follows. The atmospheric neutrinos produce muons by interaction with the nucleons (both free and bound) in the fiducial volume. If these muons are produced with energies below Čerenkov radiation threshold (kinetic energy less than 53 MeV), then they will not be detected (“invisible muons”), but their decay-produced electrons will be. Since the muon decay signal will mimic the \(\bar{\nu}_ep \rightarrow e^+n\) processes in SK, it is difficult to distinguish SRN from these events. Recent SK limits are obtained by the analysis including this invisible muon background.

In the near future, however, it should be plausible to distinguish the invisible muon signals from the SRN signals; two different methods are currently proposed for the SK detector. One is to use the gamma rays emitted from nuclei which interacted with atmospheric neutrinos \[25\]. If gamma ray events, whose energies are about 5–10 MeV, can be detected before
invisible muon events by muon lifetime, we can subtract them from the candidates of SRN signals. In that case, the upper limit would be much lower (by factor $\sim 3$) when the current data of 1,496 days are reanalyzed. Another proposal is to detect signals of neutrons, which are produced through the $\bar{\nu}_e p \rightarrow e^+ n$ reaction, in the SK detector. Actual candidate for tagging neutrons is gadolinium solved into pure water. Additional neutron signals can be used as delayed coincidence to considerably reduce background events. In addition, future projects such as the Hyper-Kamiokande detector is expected to greatly improve our knowledge of the SRN spectral shape as well as its flux. If a number of data were actually acquired, the most general and model-independent discussions concerning the neutrino decay would be accessible. Finally, we again stress that a great advantage to use the SRN is that the cosmologically long lifetime can be probed in principle.

VI. CONCLUSIONS

We obtained the current lower limit to the neutrino lifetime-to-mass ratio using the SRN observation at SK. Since the available data are only the upper limit on the flux integrated over $E_{\bar{\nu}_e} > 19.3$ MeV, the decay model on which we can set constraints is quite restricted, i.e., the decay models that gives the SRN flux which is comparable to or larger than the corresponding upper limit.

Therefore, our reference decay model in this Letter is the two-body decay $\bar{\nu}_i \rightarrow \bar{\nu}_j + X$, with the daughter neutrinos which are active species and carry nearly full energy of their parent. The neutrino mass hierarchy is assumed to be normal ($m_1 < m_3$). The SRN calculation is also very sensitive to the adopted supernova model, i.e., neutrino spectrum from each supernova explosion; we adopted the result of the numerical simulations of the Lawrence Livermore group. Although their calculation is recently criticized since it lacks several relevant neutrino processes, we adopt it as a reference model. This is because it is only the successful model of the supernova explosion and we definitely need fully time-integrated neutrino spectra.

Our calculations with these models shows that the neutrino decay model with $(\tau_2/m, \tau_3/m) \lesssim (10^{10}, 10^{10})$ [s/eV] apparently gives the SRN flux that is larger than the current upper limit (see Fig. 2). Since this value $10^{10}$ s/eV is much larger than that of the limit obtained by the solar neutrino observation $\gtrsim 10^{-4}$ s/eV, our approach is shown to be
very powerful to obtain the implications for the neutrino decay. At present, however, owing
to the large uncertainties such as supernova models, this lower limit $10^{10} \text{s/eV}$ cannot be
accepted without any doubt. Future experiments with the updated detectors and reanalysis
with the refined method is expected to give us greatly improved information on the SRN
flux and spectrum. In that case, the most general model-independent discussions concerning
the neutrino decay would be possible.

Acknowledgments

The author is supported by Grant-in-Aid for JSPS Fellows.

[1] Y. Fukuda, et al., Phys. Rev. Lett. 82 (1999) 2644.
[2] S. Fukuda, et al., Phys. Lett. B 539 (2002) 179.
[3] Q. R. Ahmad, et al., Phys. Rev. Lett. 89 (2002) 011301.
[4] Q. R. Ahmad, et al., Phys. Rev. Lett. 89 (2002) 011302.
[5] K. Eguchi, et al., Phys. Rev. Lett. 90 (2003) 021802.
[6] S. Ando, K. Sato, Phys. Rev. D 67 (2003) 023004.
[7] S. Ando, K. Sato, Phys. Rev. D 68 (2003) 023003.
[8] M. Malek, et al., Phys. Rev. Lett. 90 (2003) 061101.
[9] J. F. Beacom, N. F. Bell, Phys. Rev. D 65 (2002) 113009.
[10] A. Bandyopadhyay, S. Choubey, S. Goswami, Phys. Lett. B 555 (2003) 33;
    A. Bandyopadhyay, S. Choubey, S. Goswami, Phys. Rev. D 63 (2001) 113019;
    S. Choubey, S. Goswami, D. Majumdar, Phys. Lett. B 484 (2000) 73.
[11] J. F. Beacom, N. F. Bell, D. Hooper, S. Pakvasa, T. J. Weiler, Phys. Rev. Lett. 90 (2003)
    181301.
[12] K. Hirata, et al., Phys. Rev. Lett. 58 (1987) 1490.
[13] R. M. Bionta, et al., Phys. Rev. Lett. 58 (1987) 1494.
[14] J. A. Frieman, H. E. Haber, K. Freese, Phys. Lett. B 200 (1988) 115.
[15] R. S. Raghavan, X. G. He, S. Pakvasa, Phys. Rev. D 38 (1988) 1317.
[16] M. Lindner, T. Ohlsson, W. Winter, Nucl. Phys. B 622 (2002) 429;
M. Lindner, T. Ohlsson, W. Winter, Nucl. Phys. B 607 (2001) 326.

[17] S. Ando, K. Sato, T. Totani, Astropart. Phys. 18 (2003) 307.

[18] S. Ando, K. Sato, Phys. Lett. B 559 (2003) 113.

[19] P. Madau, H. C. Ferguson, M. E. Dickinson, M. Giavalisco, C. C. Steidel, A. Fruchter, Mon.
Not. R. Astron. Soc. 283 (1996) 1388.

[20] T. Totani, K. Sato, H. E. Dalhed, J. R. Wilson, Astrophys. J. 496 (1998) 216.

[21] T. A. Thompson, A. Burrows, P. Pinto, astro-ph/0211194.

[22] D. E. Groom, et al., Eur. Phys. J. C 15 (2000) 1.

[23] M. Apollonio, et al., Phys. Lett. B 466 (1999) 415.

[24] A. S. Dighe, A. Y. Smirnov, Phys. Rev. D 62 (2000) 033007.

[25] Y. Suzuki, Talk at: Tokyo-Adelaid Joint Workshop on Quarks, Astrophysics, and Space Science, Tokyo, 2003.

[26] M. Vagins, Talk at: The 4th Workshop on “Neutrino Oscillations and their Origin” (NOON2003), Kanazawa, 2003.