PHYSICAL REVIEW D **90,** 053011 (2014)

Cosmic neutrino background absorption line in the neutrino spectrum at IceCube

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(Received 19 July 2014; published 24 September 2014)

The IceCube experiment has recently reported a high energy neutrino spectrum between the TeV and PeV scales. The observed neutrino flux can be as a whole well fitted by a simple power law of the neutrino energy E_ν, $E_\nu^{-\gamma_\nu}$ ($\gamma_\nu \simeq 2$). As a notable feature of the spectrum, however, it has a gap between 500 TeV and 1 PeV. Although the existence of the gap in the neutrino spectrum is not statistically significant at this point, it is very enticing to ask whether it might hint at some physics beyond the Standard Model. In this paper, we investigate a possibility that the gap can be interpreted as an absorption line in the power-law spectrum by the cosmic neutrino background through a new resonance in the MeV range. We also show that the absorption line has rich information about not only the MeV scale new particle but also the neutrino masses as well as the distances to the astrophysical sources of the high energy neutrinos. Viable models to achieve this possibility are also discussed.

DOI: 10.1103/PhysRevD.90.053011 PACS numbers: 13.15.+g, 14.60.Pq, 14.60.St

I. INTRODUCTION

The IceCube experiment has recently reported high energy neutrinos considered to be coming from an extraterrestrial source since those observed events are significantly large compared to the atmospheric neutrino background [1,2]. Such high energetic neutrinos are expected to come from, for example, the photopion production such as $\gamma p \to \Delta \to \pi^+ X$ followed by the pion decay, $\pi^+ \to \nu_\mu(\mu^+ \to \nu_e \bar{\nu}_\mu e^+)$, which produces the neutrinos of the flavor composition with $\nu_e{:}\nu_\mu{:}\nu_\tau = 1{:}2{:}0$ while it becomes $1{:}1{:}1$ after traveling from some extraterrestrial source.[1] The cosmogenic neutrino flux, however, peaks at around $O(1)$ EeV for γ being the cosmic microwave background (CMB), and it is difficult to explain the observed neutrino flux in the sub-PeV region [5,6]. As other possibilities, there are many candidates to explain the events around the sub-PeV to the PeV region by the high energetic cosmic-ray sources inside our galaxy such as the supernova remnants (SNR) [7,8] and the pulsar wind nebulae (PWN) [9] as well as the extragalactic sources such as the gamma ray bursts (GRB) [10,11], the active galactic nuclei [12], and the star forming galaxies [13] (see also Refs. [14–18] and references therein). More ambitious explanations by physics beyond the standard model (SM) such as decaying dark matter or new interactions of neutrino have also been discussed [19–24].

As a current status of the observed neutrino flux, on the other hand, it is as a whole well fitted by a simple power law $E_\nu^{-\gamma_\nu}$ ($\gamma_\nu \simeq 2$), in the sub-PeV to the PeV range, where E_ν is the observed neutrino energy. This power spectrum is vaguely supported by the source spectrum of the cosmic ray proton accelerated by the first order Fermi acceleration mechanism. As a notable feature of the spectrum, however, it has a gap between 500 TeV and 1 PeV. Although the existence of the gap in the observed neutrino spectrum is not statistically significant at this point (see, e.g., [25]), it is very enticing to ask whether it might hint at some physics beyond the SM.

In this paper, we investigate a possibility that the gap in the power-law spectrum can be interpreted as an absorption line by the cosmic neutrino background (CνB) through a new resonance with a mass in the MeV range. We also show that the neutrino absorption line has rich information about not only the MeV scale new particle but also the neutrino masses as well as the distances to the astrophysical sources of the neutrinos. Viable models to achieve this possibility are also discussed.

II. NEW PARTICLE AND RESONANT ABSORPTION

Let us discuss whether it is possible to interpret the null event regions at the sub-PeV neutrinos as the CνB absorption line in the single power law spectrum of $E_\nu^{-\gamma_\nu}$ with $\gamma_\nu = 2$. In the SM, there are no appropriate interactions that show an absorption line at the sub-PeV region. As we will see shortly, however, such an absorption line interpretation becomes possible by introducing a new

[1]The flavor oscillation of the neutrino being the energy E_ν would take place after traveling the distance $L \sim 2E_\nu/\Delta m_{ij}^2$ where the mass difference is defined by $\Delta m_{ij}^2 \equiv m_{\nu_i}^2 - m_{\nu_j}^2$ for the mass eigenstate of neutrinos ν_i, and if we take $\Delta m_{21}^2 \sim 10^{-3}$ eV^2 and $E_\nu \sim 10^6$ GeV, the distance becomes $L \sim 10^{-10}$ Mpc, which is small enough even if some astrophysical neutrino source locates within the intergalactic scale [3,4].

resonance appearing in the s-channel neutrino-(anti)neutrino scattering.

The CνB is a remnant of the primordial plasma reheated after the inflation, and the temperature of the CνB is predicted to be $T_\nu \simeq 1.96$ K $\simeq 1.69 \times 10^{-4}$ eV. From this temperature, the neutrino number density is given by $n_\nu \simeq 56$ cm^{-3} for each flavor. When the high energy neutrinos accelerated by some astrophysical source collide with the CνB of the masses larger than T_ν, the situation is almost the same as the collision with a fixed target in the laboratory frame. In this case, the center-of-mass energy is given by $\sqrt{2m_\nu E_\nu}$, where m_ν denotes the mass of the target neutrino in the CνB. Thus, if the mass of a new particle, M_s, appearing in the s-channel neutrino collisions is around $M_s \simeq \sqrt{2m_\nu E_\nu}$, the injected neutrinos of E_ν are resonantly scattered by the CνB, which leads to the "absorption line" in the neutrino spectrum. For example, $E_\nu \sim 1$ PeV neutrino absorption predicts a new particle in the mass around $M_s \sim 10$ MeV if we take the neutrino mass $m_\nu = 0.1$ eV.

Before introducing a new particle, however, let us first examine what is expected on the neutrino spectrum in the SM. There, most of the cosmic-ray neutrinos accelerated by some astrophysical sources are expected to penetrate astrophysical/cosmological distances since they interact with materials very weakly. As the neutrinos are traveling in the distance, the most relevant target material is the CνB since it is as abundant as the CMB while it has larger interaction rates with the neutrino flux than the CMB. In the SM, the neutrinos interact with themselves via the electroweak interactions, where the relevant processes are $\nu_l \bar{\nu}_{l,\mathrm{C}\nu\mathrm{B}} \to f\bar{f}(f = \nu_l, l, q, \ldots)$, $\nu_l \bar{\nu}_{l',\mathrm{C}\nu\mathrm{B}} \to \nu_l \bar{\nu}_{l'}, l\bar{l}' (l \neq l')$, and $\nu_l \nu_{l',\mathrm{C}\nu\mathrm{B}} \to \nu_l \nu_{l'}$. The cross sections of the SM processes are given in, for example, Refs. [20,26]. Since some of them can be enhanced via s-channel Z-boson exchanges at the energy of the Z boson mass, neutrino absorption may occur for the energy of the neutrino flux around $E_\nu = M_Z^2/(2m_\nu) \sim 10^{13}$ GeV. This absorption line is far above the energy range of the recently observed neutrinos, and hence, we cannot attribute the null event regions in the IceCube spectrum to the absorption line in the SM. The occurrence of such an absorption feature by the Z boson is known as the "Weiler mechanism," which has been studied in Refs. [27–31]. Related topics have also been studied in [32–34].

Now, let us introduce a new light particle to make an absorption line at around the sub-PeV range in the neutrino spectrum. The situation is similar to the Z-boson resonance, while the new particle coupling to the neutrinos are predicted to be around the MeV scale in our case as mentioned above. Suppose that the new scalar particle s with a mass M_s couples to the neutrinos by

$$\mathcal{L}_{s-\nu} = gs\bar{\nu}_i \nu_j \tag{1}$$

with coupling g where we assume that the coupling is flavor universal for simplicity. Here, we do not specify whether the neutrino is the Dirac type or the Majorana type. One caution is, however, that if the above interaction is the Yukawa interaction between the left-handed and the right-handed neutrinos of the Dirac neutrino, the right-handed neutrinos are copiously produced in the early universe through this interaction. Such a possibility is severely restricted by the constraints on the effective number of neutrinos, $N_{\mathrm{eff}} = 3.02 \pm 0.27$ from the big-bang nucleosynthesis and the CMB observations [35], which eventually leads to a constraint on the coupling constant,

$$g \lesssim (M_s/M_{\mathrm{PL}})^{1/4}. \tag{2}$$

Here M_{PL} denotes the reduced Planck mass $M_{\mathrm{PL}} \simeq 2.4 \times 10^{18}$ GeV. Since we will use rather sizable coupling constants, we find that the only possible interactions are

$$\mathcal{L}_{s-\nu} = \begin{cases} gs\nu_{Li}\nu_{Lj}, & (\text{Majorana, Dirac}), \\ gs\bar{N}_{Ri}\bar{N}_{Rj}, & (\text{Dirac}), \end{cases} \tag{3}$$

where ν_L and $\bar{N}_R$ denote the left-handed neutrinos and the right-handed neutrinos. Flavor dependence of the coupling as well as the consistency with the electroweak theory will be discussed inthe next section.

The neutrino-(anti)neutrino scattering cross section $\sigma_{\nu\nu}(S)$ is evaluated as shown in the left panel of Fig. 1. The black solid line shows the SM cross section with the resonance at the Z-boson pole. The parameters (g, M_s) for the cross section by the new resonance are as indicated. The highest value of the cross section at $S = M_s^2$ is determined by the decay width of s given by

$$\Gamma_s = N_\nu \frac{g^2}{16\pi} M_s \left[1 - \frac{2m_\nu^2}{M_s^2}\right] \left[1 - \frac{4m_\nu^2}{M_s^2}\right]^{1/2}, \tag{4}$$

where s is assumed to decay into N_ν neutrinos, and consequently the peak of the cross section is $\sigma_{\nu\nu}(S = M_s^2) \simeq 16\pi/(N_\nu^2 M_s^2)$.

The neutrino mean free path (MFP) λ is an important quantity to evaluate how far the neutrino traveling distance is. The MFP is defined by

$$\lambda(E_\nu) = \left[\int \frac{d^3p}{(2\pi)^3} \sigma_{\nu\nu}(E_\nu, p) f_\nu(p)\right]^{-1}, \tag{5}$$

where $f_{\nu_i}(p)$ is the CνB distribution function given by $f_{\nu_i}(p) = [\exp(|\vec{p}|/T_\nu) + 1]^{-1}$. Examples of the MFP are shown in the right panel of Fig. 1 where M_s and m_ν are set to $M_s = 2.5$ MeV and $m_\nu = 3.2 \times 10^{-3}$ eV, respectively. The black, the blue, and the red solid lines, respectively, show the case of $g = 0.001, 0.01$, and 0.1. If the traveling distance of the neutrinos is below the lines, the neutrino flux at a corresponding energy cannot reach to the Earth. In most of the energy region except for the resonance region, the relative magnitude among those lines is determined by

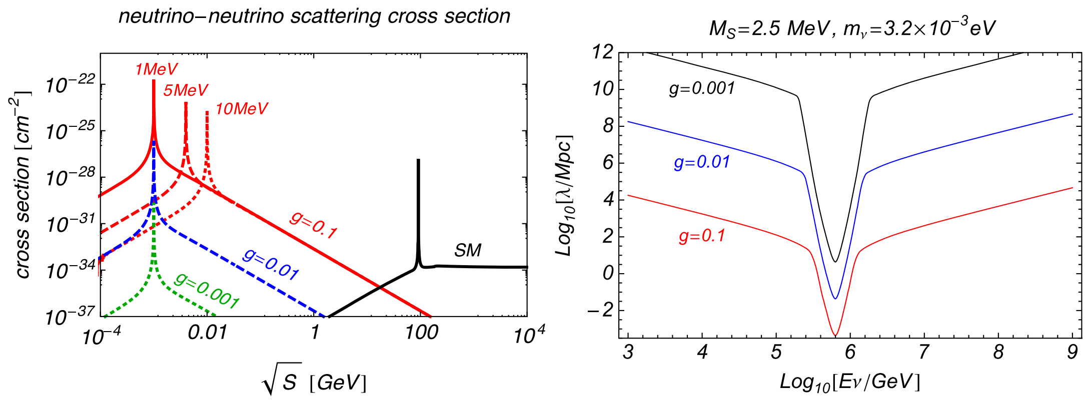

FIG. 1 (color online). Left panel: The neutrino-(anti)neutrino scattering cross sections for the center of mass energy $\sqrt{S}$. The black solid line is the SM. The red solid, the red dashed, the red dotted, the blue dashed, and the green dotted lines depict the contributions from the interaction of Eq. (1) for several parameter samples of the coupling g and the scalar particle mass. Right panel: The neutrino mean free path λ as a function of the energy of the neutrino flux.

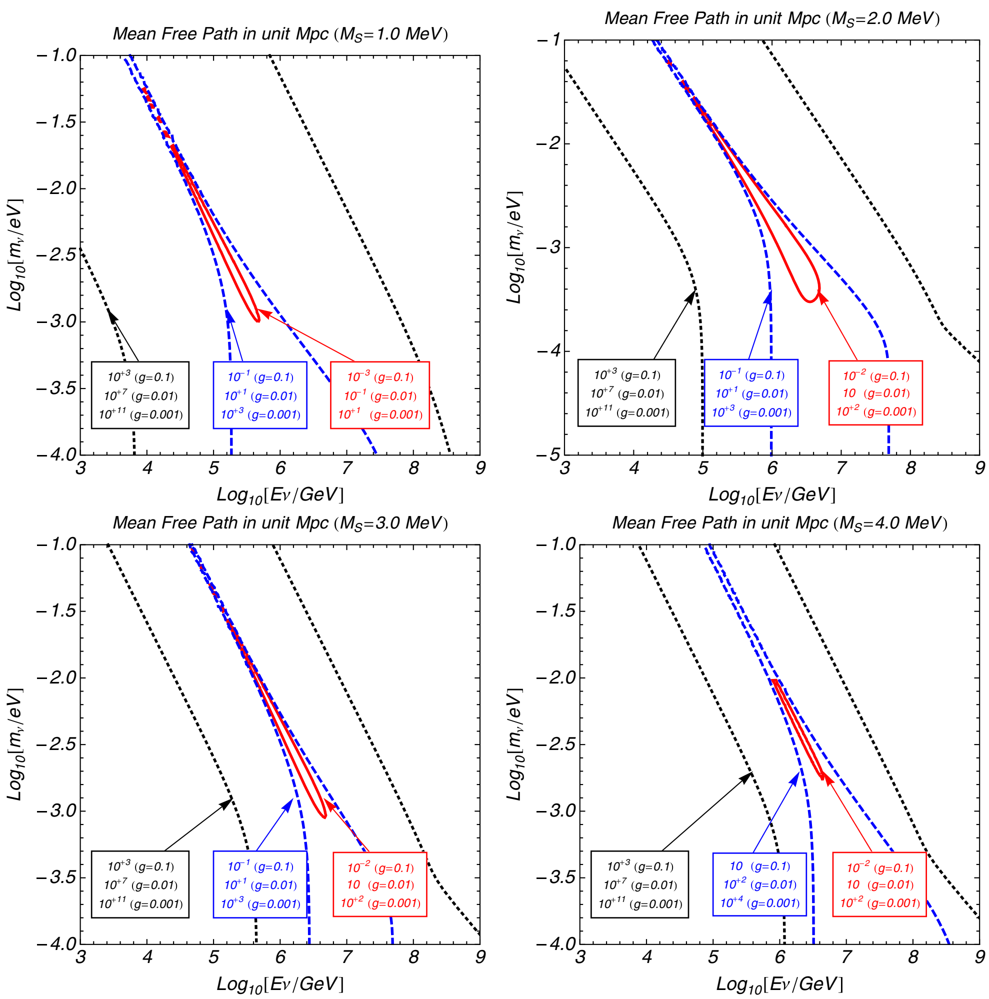

FIG. 2 (color online). The neutrino mean free path for various neutrino masses. The numbers shown in the boxes are the mean free path in the unit Mpc for each coupling. The scalar boson mass is set to be $M_s = 1$ MeV, 2 MeV, 3 MeV, and 4 MeV in the upper left, upper right, bottom left, and bottom right panels, respectively.

the magnitude of the coupling, for example, the MFP for the case of $g = 0.1$ is four digits smaller than the case of $g = 0.01$ since the cross section is proportional to g^4. As indicated by the peak of the cross section, the bump structure of the MFP reflects the resonance of the singlet scalar. Around the resonance region, the cross section is changing with a strength proportional to g^2, and thus the relative difference among MFPs is two digits magnitude.[2]

Notably, the neutrino masses (of the CνB) are also an important parameter to determine the neutrino MFP. Since the MFP is given by the overlap between the neutrino scattering cross section and the distribution function of the CνB, it is sensitive to the neutrino mass through the center of the mass energy,

$$S \simeq 2E_\nu \left(\sqrt{m_{\nu,\mathrm{C}\nu\mathrm{B}}^2 + p_\nu^2} - p_\nu \cos\theta \right), \tag{6}$$

where $\cos\theta$ denotes the scattering angle and the typical value of p_ν is $O(T_\nu)$. It should be noted that S becomes insensitive to p_ν and is solely determined by E_ν for $m_{\nu,\mathrm{C}\nu\mathrm{B}} \gg T_\nu$, while it takes a wide range for $m_{\nu,\mathrm{C}\nu\mathrm{B}} \ll T_\nu$ due to the p_ν contribution. Therefore the MFP becomes a sharp function of E_ν for $m_{\nu,\mathrm{C}\nu\mathrm{B}} \gg T_\nu$, since $S \simeq M_s$ is achieved only for a particular value of E_ν. On the contrary, the MFP becomes a broad function of E_ν for $m_{\nu,\mathrm{C}\nu\mathrm{B}} \gg T_\nu$ since a wide range of E_ν can achieve $S \simeq M_s$. The neutrino mass dependence of the MFP is shown in Fig. 2, which shows the contours of the MFP for a scalar boson mass of $M_s = 1$ MeV, 2 MeV, 3 MeV, and 4 MeV with $g = 0.1, 0.01$, and 0.001 for each case. The shortest MFP tends to be at a higher E_ν for a larger M_s. In contrast, a heavier neutrino mass makes the shortest MFP be at lower E_ν since the required E_ν to reach $\sqrt{S} = M_s$ becomes small for a heavier m_ν. The figure indicates that the absorption line at around the sub-PeV region can be realized for $M_s \simeq$ 1–3 MeV, $m_\nu = 10^{-(2-2.5)}$ for the neutrino sources at the distance of $O(1)$ Mpc.

Before closing this section, let us comment on the traveling distance of the high energy neutrinos. As mentioned in the Introduction, there are several candidates for the astrophysical source of the high energy neutrinos. One of the promising candidates is the SNRs, which locate typically $O(1-10)$ kpc far from the Earth. The SNRs originated neutrinos are almost the left-handed state even if they are massive since the neutrino has a much higher energy than the neutrino mass. Therefore, the absorption line scenario requires a new interaction involving the left-handed neutrino so that the cosmic-ray neutrino scatters with the CνB, and the interaction should be strong enough to make the MFP shorter than the $O(1\text{–}10)$ kpc scale. Other intriguing sources are the GRBs whose distances are $O(1)$ Gpc from the Earth. In this case, a necessary coupling constant becomes smaller since the required MFP is longer than the case of SNRs.

[2]At an energy near the resonance, the cross section behaves as $\sigma(S = M_s^2 + M_s\Gamma_s) \simeq \sigma(M_s^2) + g^2/M_s^2$.

III. VIABLE MODELS

In the rest of this paper, we discuss viable models that are behind the effective theory considered in Eq. (1). So far, there have been many intriguing models in which neutrinos are interacting with new particles, for example, the Majoron models [36,37], the neutrinophilic Higgs models [38],[3] and the triplet Higgs models [40]. However, straightforward adaptations of those models to our mechanism suffer from cosmological constraints and the constraints from the light meson rare decays since a rather large coupling of the neutrino interaction is required for our purpose.[4]

A. Inverse seesaw model with a neutrinophilic scalar doublet

At first, let us examine a model where a neutrinophilic scalar doublet h_N where, in addition to the usual right-handed neutrino $\bar{N}_R$, we also introduce additional neutrinos N_N that couple not to the Higgs doublet h but only to h_N,

$$\mathcal{L} \supset g h_N l N_N + y h l \bar{N}_R + M \bar{N}_R N_N + m N_N N_N. \tag{7}$$

Here, l denotes the lepton doublet in the SM, g and y denote the dimensionless coupling constants, and m and M are the mass parameters. In this model, we impose charges of the lepton number L and the discrete symmetry Z_2 as shown in Table I. Because of these symmetries, the doublet scalar h_N possesses the neutrinophilic nature. The last three terms induce the tiny neutrino mass, and by assuming $\langle h_N \rangle = 0$, the neutrino mass is given by $m_\nu \simeq y^2 v^2 (m/M^2)$ by the inverse seesaw mechanism [41]. Here, v is the vacuum expectation value (VEV) of the Higgs doublet, and $m \ll yv \ll M$ is assumed. The smallness of the neutrino mass is achieved by assuming that the lepton-number violating mass parameter m is highly suppressed. The neutrinos other than the three active neutrinos have masses of $O(M)$.

Let us emphasize the difference from the conventional model of the neutrinophilic Higgs doublet. In the conventional neutrinophilic model, the neutrino masses are generated by the VEV of h_N, and hence, the neutrinos obtain the Dirac neutrino mass. As discussed in the previous section, however, the Yukawa coupling between the left-handed and the right-handed neutrinos is severely

[3]We use the term "neutrinophilic" coined in Ref. [39].

[4]For the singlet Majoron model [36], the resultant coupling between the Majoron and the left-handed neutrinos is highly suppressed to achieve light neutrinos. The Majoron model appearing from the triplet Higgs may have a sizable coupling to the left-handed neutrinos, although the model does not work for our purpose as we will comment later.

TABLE I. Charge assignment of the model of Eq. (7). Here we also show the charge assignments of the mass parameters as spurious fields.

	l	N_N	$\bar{N}_R$	h	h_N	m	M
L	$+1$	$+1$	-1	0	-2	-2	0
Z_2	$+$	$+$	$-$	$+$	$-$	0	$-$

restricted. To avoid this problem, we separate the mass generation and the neutrino interaction by evoking the inverse seesaw mechanism. As a result of the inverse seesaw mechanism, the Majorana neutrino masses and the effective coupling between h_N and the left-handed neutrinos are simultaneously generated.

Under the above symmetries, the scalar potential is given by

$$V = -\mu_h^2|h|^2 + \lambda(h^\dagger h)^2 + \mu_N^2|h_N|^2 + \lambda_1(h_N^\dagger h_N)^2 + \lambda_2|h|^2|h_N|^2 - \lambda_3|h^\dagger h_N + \mathrm{H.c.}|^2, \tag{8}$$

where h_N does not acquire a VEV,[5] and parameters $\mu_h^2, \mu_N^2, \lambda, \lambda_1, \lambda_2$, and λ_3 are defined as positive values. We can estimate the scalar mass spectrum of h_N by decomposing into $h_N = (h_N^0 + iA^0, h_N^-)^T$ in which h_N^0, A^0, and h_N^- are neutral CP-even, neutral CP-odd, and charged scalars, respectively. They acquire masses from the third, the fifth, and the last terms; meanwhile, only h_N^0 has an additional mass from the sixth term[6]: $m^2_{A^0,h_N^-} \sim \mu_N^2 + \lambda_2 v^2, m^2_{h_N^0} \sim \mu_N^2 + (\lambda_2 - \lambda_3)v^2$. Therefore, if we take the parameters by $\mu_N^2 \ll v^2$ and $\lambda_2 - \lambda_3 = \mathcal{O}(10^{-6})$ with $\lambda_{2,3} = \mathcal{O}(1)$, a desirable spectrum such as $m_{h_N^0} = \mathcal{O}(1\text{–}10)$ MeV and $m_{A^0,h_N^-} \gtrsim 100$ GeV can be obtained without conflicting with the custodial symmetry.[7]

It should be commented that the above assumption $\mu_N^2 \ll v^2$ is important for two reasons. First, if $\mu_N^2 = O(v^2)$, h_N^0 in the MeV range is achieved by a cancellation between two contributions, μ_N^2 and $(\lambda_2 - \lambda_3)v^2$. In such a case, the vacuum at $h_N^0 = 0$ becomes unstable for a slightly larger field value of $h^0 > v$ due to the negative value of $\lambda_3 - \lambda_2$. To avoid the instability, we need to assume $\mu_N^2 \ll v^2$ so that the lightness of h_N^0 is achieved by the small but positive value of $(\lambda_3 - \lambda_2)$. The second reason for this assumption is the suppression of the invisible decay of the observed Higgs boson into a pair of h_N^0. Under the assumption of $\mu_N^2 \ll v^2$, the value of $(\lambda_3 - \lambda_2)$ is inevitably small. Thus, by remembering that the branching ratio of the mode into a pair of h_N^0 is proportional to $(\lambda_3 - \lambda_2)$, the lightness of the h_N^0 automatically guarantees the small branching ratio to a pair of h_N^0 under the assumption of $\mu_N^2 \ll v^2$.

Once we obtained the above mass splitting in the neutrinophilic Higgs doublet, we obtain the effective theory of h_N^0 and the left-handed neutrinos,

$$\mathcal{L}_{\mathrm{eff}} \simeq \frac{gyv}{M} h_N^0 \nu_L \nu_L, \tag{9}$$

which realizes the model discussed in the previous section by identifying

$$g^{\mathrm{eff}} = \frac{gyv}{M}, \qquad s = h_N^0. \tag{10}$$

By assuming $M = O(1)$ TeV and $g = y = O(1)$, for example, we achieve the effective theory with $g^{\mathrm{eff}} = O(0.1)$.

The experimental limits on the charged Higgs mass are given by using $t \to H^+ b$ for $m_{H^+} < m_t$ and $H^+ \to \tau\nu$ for $m_{H^+} > m_t$ by H^+ production via third generation quarks at the LHC [42]. However, h_N does not couple to quarks in the model, and thus, h_N^- is free from the limit. So only the LEP constrains h_N^- by $e^+e^- \to H^+H^- \to \tau\tau\nu\nu$, and the exclusion limit is $m_{H^+} \gtrsim 100$ GeV by imposing $\mathrm{Br}(H^+ \to \tau\nu) = 1$ [43]. The CP-odd Higgs is still free from any experimental observation since it only couples with the neutrino as long as it is heavier than the Z boson. Lepton flavor violation is also affected by the charged Higgs such as $\mu \to e\gamma$ induced by the effective operator $m_\mu(g^2/\Lambda^2)\bar{\mu}_R\sigma^{\mu\nu}e_L F_{\mu\nu}$ where Λ is a cutoff scale. The experimental limit is given by $\mathrm{Br}(\mu \to e\gamma) \lesssim 10^{-13}$ [44], which reads to $\Lambda \gtrsim \mathcal{O}(100)$ GeV if we take $g = \mathcal{O}(0.1)$ and a loop factor is considered [45].

A crucial experimental limit is for the coupling among h_N^0 and neutrinos from the rare meson decay rates emitting h_N^0. In particular, the null observations of $\pi/K \to l\nu_{l'}h_N^0$ put stringent constraints on the coupling g [46,47]. Hereafter, we denote $g^{\mathrm{eff}}_{ab}h_N^0\bar{\nu}_a\nu_b (a, b = e, \mu, \tau)$ as the flavor basis, and the coupling is converted into $g^{\mathrm{eff}}_{ij} = (U^\dagger_{\mathrm{PMNS}})_{ia} g^{\mathrm{eff}}_{ab} (U_{\mathrm{PMNS}})_{bj} (i, j = 1, 2, 3)$ in the mass basis using the unitary Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix U_{PMNS},

[5]Suppose that m in Eq. (7) are spurions of explicit breaking of the lepton number, and it has the charge $L = +2$. Therefore, h_N and h mix with each other in a form of $m^* h_N \leftrightarrow h$, whose mixing is of order $m/M \ll 1$ via a one-loop diagram; consequently, the VEV of h_N is negligible, and the contribution to the neutrino masses is also suppressed.

[6]Generically, the term proportional to λ_3 contains two independent terms that are allowed any symmetries other than the custodial symmetry. In our model, to evade the constraints from the electroweak precisions, we fine-tune the potential so that the scalar potential respects the custodial symmetry.

[7]In the triplet Higgs model, the mass splittings in the triplet Higgs multiplet lead to the custodial symmetry breaking. Thus, we cannot obtain a light particle with a mass in the MeV range while keeping other modes such as the charged Higgs bosons in the $O(100)$ GeV range without conflicting with the custodial symmetry. The same problem arises in the Majoron model appearing from the triplet Higgs boson [37].

$$U_{\mathrm{PMNS}} = \begin{bmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{bmatrix} \begin{bmatrix} c_{13} & 0 & s_{13}e^{-i\delta} \\ 0 & 1 & 0 \\ -s_{13}e^{-i\delta} & 0 & c_{13} \end{bmatrix} \times \begin{bmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{bmatrix} \begin{bmatrix} e^{-\alpha_1/2} & 0 & 0 \\ 0 & e^{i\alpha_2/2} & 0 \\ 0 & 0 & 1 \end{bmatrix} \quad (11)$$

Here, $s_{ij} \equiv \sin(\theta_{ij}), c_{ij} \equiv \cos(\theta_{ij})$, and δ and α_i are Dirac and Majorana phases, respectively, and we take $\delta = \alpha_1 = \alpha_2 = 0, s_{12}^2 = 0.31, s_{23}^2 = 0.51$, and $s_{13}^2 = 0.023$ in our analysis for simplicity. In the flavor basis, the constraints from the rare meson decays put a limit on g_{ab}^{eff} [47]

$$\sum_{l=e,\mu,\tau} |g_{el}^{\mathrm{eff}}|^2 < 5.5 \times 10^{-6}, \qquad \sum_{l=e,\mu,\tau} |g_{\mu l}^{\mathrm{eff}}|^2 < 4.5 \times 10^{-5}, \quad \text{and} \quad \sum_{l=e,\mu,\tau} |g_{\tau l}^{\mathrm{eff}}|^2 < 3.2. \quad (12)$$

As we have discussed in the previous section, we need to assume $g^{\mathrm{eff}} = O(0.1–1)$ to obtain a short enough MFP for the neutrino flux from the sources inside our galaxy. Because of the above constraints in Eq. (12), the only allowed coupling of $O(0.1–1)$ is $g_{\tau\tau}^{\mathrm{eff}}$ in the flavor basis. It should be noted though that in the mass basis, $g_{\tau\tau}^{\mathrm{eff}}$ leads to $O(1–0.1)$ couplings between three neutrinos in the mass basis according to the PNMS matrix. If the sources of the neutrinos are the extragalactic ones, on the other hand, the couplings of $g^{\mathrm{eff}} = O(0.01)$ are large enough to achieve the short MFP, which can easily evade the constraints from the rare meson decay.

Now, let us calculate the resultant neutrino spectrum by assuming a single power-law flux at the neutrinos sources. The number of neutrinos reaching to the Earth is approximately estimated by[8]

$$\frac{dN_\nu}{dL}(E_\nu, z) \simeq -\frac{N_\nu(E_\nu, z)}{\lambda(E_\nu)}, \quad (13)$$

where L is the length of the neutrino traveling path defined by

$$L = \frac{c}{H_0}\int dz(\Omega_m(1+z)^3 + \Omega_\Lambda)^{-1/2}\ \mathrm{Mpc}, \quad (14)$$

where z denotes the redshift parameter, $c = 3 \times 10^5$ km/s, $H_0 = 100h$ km/s/Mpc, and Ω_m and Ω_Λ are energy densities of matter and dark energy, respectively. In our analysis, we use $h = 0.67, \Omega_m = 0.32$, and $\Omega_\Lambda = 0.68$ [48].

In Fig. 3, we show some examples of the neutrino spectrum for the extragalactic sources locating at the distance of 1 Gpc. The figure shows that the absorption line can be achieved for $g_{11}^{\mathrm{eff}} = 10^{-3}$, which easily satisfies the constraints from the rare decay in Eq. (12). The left panel of the figure shows the neutrino mass dependence of the absorption line, where we take the neutrino mass as a free parameter and focus on the dominant contribution to the absorption process. From the figure, we find that the neutrino mass about $m_\nu \simeq 5.6 \times 10^{-3}$ eV provides a nice fit to the null regions of the IceCube flux. Interestingly, this mass is close to the square root of the squared mass differences of the first two neutrinos in the normal hierarchy, $\Delta m_{21}^2 \simeq 7.6 \times 10^{-5}$ eV^2 [49]. Thus, the result favors the neutrino mass spectrum in which the first two neutrinos are rather degenerated. The right panel shows the dependence on the resonance mass. The figure shows that the nice fit is achieved for $M_s \simeq 3$ MeV.

When the neutrino sources are inside our galaxies with the distance of $O(10)$ kpc, on the other hand, we need to have $g_{\tau\tau}^{\mathrm{eff}} = O(1)$. In this case, the meson decay constraints only allow $g_{\tau\tau}^{\mathrm{eff}} = O(1)$ in the flavor basis. The coupling constant in the mass basis is, on the other hand, determined according to the PNMS matrix with $g_{\tau\tau}^{\mathrm{eff}} = O(1)$. As a result, the ratios between the coupling constants are fixed by the PNMS matrix, which leads to a nontrivial relation between the absorption lines made by the three neutrinos. In Fig. 4, we show an example of the neutrino spectrum for $g_{\tau\tau}^{\mathrm{eff}} = 0.5$. The figure shows that the spectrum has not only the broad absorption lines by the first two neutrinos but also a sharp line by the third neutrino. Here, again, the degenerated first two neutrinos are favored. The detailed observation of the neutrino spectrum is required to test the existence of such multiple absorption lines in the neutrino spectrum.[9]

B. Another model

Finally, let us discuss another possibility to induce the CνB absorption line at the sub-PeV scale. As previously mentioned, the high energy neutrinos produced by astrophysical sources are mostly left-handed even if the neutrinos are the Dirac type, and the chirality flip hardly takes place as they travel because its energy is much higher that

[8] To be more accurate, there exist other contributions such as the expansion effect of the universe and secondary neutrino scattering. However, in our case, the flight distance of neutrinos is small enough not to be affected by the expansion of the universe. For the secondary neutrino scattering effect, inelastic scattering is not sufficient since the SM cross section is negligible in the resonance region. In the elastic scattering case, scattered neutrinos settle where their energy is around $E_\nu/2$; however, this contribution is also negligible and does not change our result where Eq. (13) is utilized.

[9] Such multiple absorption lines are also possible for the extragalactic neutrino sources depending on the structure of the Yukawa couplings.

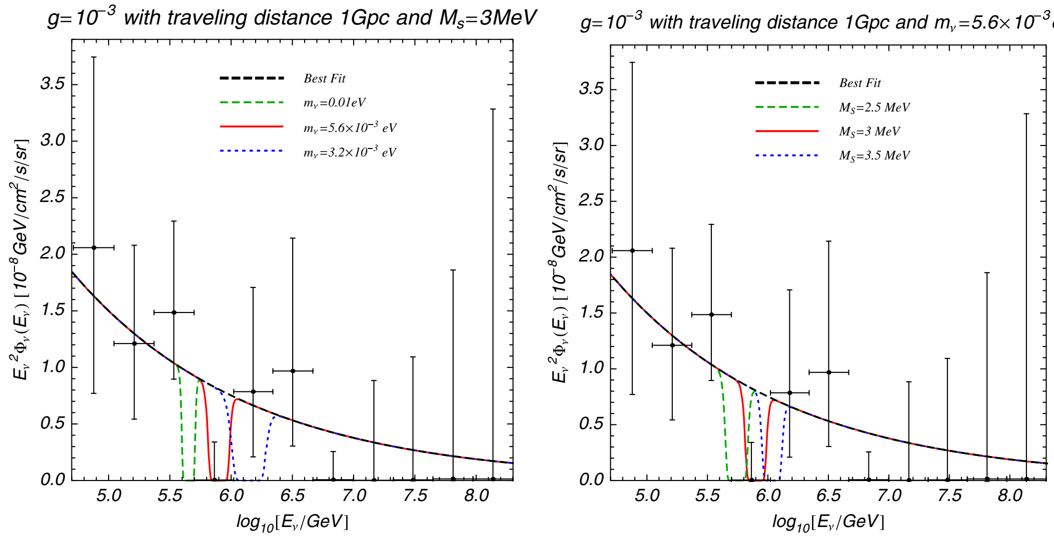

FIG. 3 (color online). Absorption line with the sample parameters assuming the source distance to be $O(1)$ Gpc. Left panel: The neutrino mass dependences. The black dots with error bars are observed data, and the best-fit power law is $E_\nu^2\Phi_\nu(E_\nu) = 1.5 \times 10^{-8}(E_\nu/100\ \mathrm{TeV})^{-0.3}\ \mathrm{GeV/cm^2/s/sr}$ [2]. Right panel: The dependence on the resonance mass. In both figures, we assumed $g \equiv g_{11}^{\mathrm{eff}} = O(10^{-3})$.

the mass. Therefore, it is simple to assume that the resonances appear in the collisions between the left-handed neutrinos.

When the neutrinos pass through the magnetic field, however, the chirality flip is potentially possible since the neutrinos have a finite magnetic momentum, μ_ν. In the magnetic fields B, the Larmor frequency of the neutrino is given by $B\mu_\nu$, and hence, the Dirac neutrinos flip their chirality when the travel time is longer than the Larmor frequency.[10] Thus, once the chirality flip occurs due to a strong magnetic field, the neutrino absorption can be achieved by the resonance appearing in the collisions between the right-handed neutrinos,

$$\mathcal{L} = gs\bar{N}_R\bar{N}_R, \tag{15}$$

in the case of the Dirac neutrino. The required masses of the resonance and the size of the coupling to obtain the visible absorption line are similar to the results in the previous section. It should be noted that the size of the coupling g is hardly constrained by any other experiments including the rare meson decay.

Unfortunately, however, the neutrino magnetic moment predicted in the SM is very small,

$$\mu_\nu \simeq 3 \times 10^{-19}\left(\frac{m_\nu}{1\ \mathrm{eV}}\right)\mu_B, \tag{16}$$

where $\mu_B \equiv e/(2m_e) \simeq 0.6 \times 10^{-13}$ GeV/T is the Bohr magneton. Therefore the necessary distance for the chirality flip is very long,

$$L_{\mathrm{cf}} = \pi/\mu_\nu B \simeq 10\left(\frac{0.1\ \mathrm{eV}}{m_\nu}\right)\left(\frac{\mu G}{B}\right)\ \mathrm{Gpc}. \tag{17}$$

Thus, it is difficult to flip the chirality by the galactic magnetic field $B \simeq 10^{-6}$ G [50,51]. As a result, in order for the chirality flip to take place, we need a new physics that

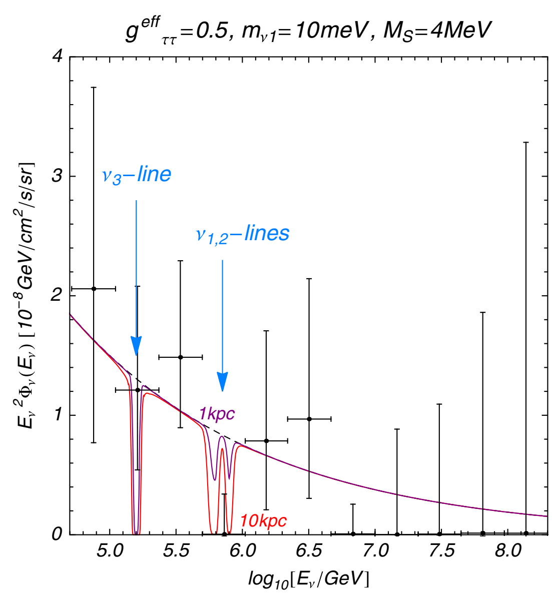

FIG. 4 (color online). Absorption line with the sample parameters assuming the source distance to be $O(10)$ kpc. Here, we have taken $g_{\tau\tau}^{\mathrm{eff}} = 0.5$.

[10]In the rest frame of the injecting neutrino, the travel time is suppressed by a large Lorentz boost factor, where the Larmor frequency is enhanced by the boost magnetic field in the rest frame.

enhances the neutrino magnetic moment significantly (see, e.g., Ref. [52]). For example, if we assume the current experimental upper limit on the neutrino magnetic field, $\mu_\nu < 5.4 \times 10^{-11}\mu_B$ [53], the chirality flip is possible within the traveling distance of $\mathcal{O}(1)$ kpc under the galactic magnetic field $B \simeq 10^{-6}$ G.[11]

IV. SUMMARY AND DISCUSSION

In this paper, we have discussed the possibility whether the null-event region around the sub-PeV scale in the neutrino spectrum observed at the IceCube experiment can be interpreted as an absorption line by the CνB in the power-law spectrum. To achieve such a possibility, we proposed two viable models where the MeV resonance appears in the neutrino-neutrino interactions. For the models with Majorana neutrinos, we found that the resonance is embedded in the neutrinophilic doublet boson that will be tested by future collider experiments. For the models with the Dirac neutrinos, we found that the resonance appearing in the interaction of the right-handed neutrinos is also a possibility, although we need an enhancement of the neutrino magnetic moment to flip the chirality of the neutrinos during the flight to hit the resonance. Such an enhanced neutrino magnetic momentum requires an additional new physics beyond the SM, which will also be tested by future collider experiments.

It should be noted that the shape of the absorption line depends not only on the mass of the new resonance but also on the neutrino masses. Thus, in principle, it is possible to extract the masses of the neutrinos by investigating the absorption lines in the neutrino spectrum, although it requires very high energy resolution. The identification of the astrophysical sources of the high energy neutrinos is also crucial to determine the absorption line, since it depends on the relative magnitude between the MFP and the distance to the neutrino source from the Earth.

Finally let us comment on an implication for cosmology. Nonstandard neutrino interactions can affect the CMB power spectrum and/or the structure formation of the universe since it might change the decoupling temperature of the neutrinos and/or the neutrino free-streaming scale. Interestingly, the recent CMB analysis [54] reported a slight preference for additional neutrino interactions with the magnitude of $g^2/M_s^2 \simeq 1/(10\ \mathrm{MeV})^2$, which is surprisingly close to the ones we are assuming. Since the conclusion has a prior dependence [55], it is premature to say that the existence of the nonstandard neutrino interactions are supported by the CMB observation. However, such cosmological observations are expected to provide significant synergy of the IceCube experiment in future studies.

[11]If the neutrinos are accelerated at the PWN surrounding a very strong magnetic field $B \sim 10^{12}$ G, the required enhancement of the neutrino magnetic moment can be smaller, although we do not pursue this possibility any more in this paper.

ACKNOWLEDGMENTS

We thank T. Fujita for useful discussions. We also thank S. Yoshida for useful discussion on the neutrino absorption in the SM. This work is supported by Grant-in-Aid for Scientific research from the Ministry of Education, Science, Sports, and Culture (MEXT), Japan, Grants No. 24740151 and No. 25105011 (M. I.), and from the Japan Society for the Promotion of Science (JSPS), Grant No. 26287039 (M. I.). This work is also supported by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan.

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