Dark matter, neutrino mass, cutoff for cosmic-ray neutrino, and the Higgs boson invisible decay from a neutrino portal interaction

Wen Yin(殷文)[1,2;1]

[1] Institute of High Energy Physics, Chinese Academy of Sciences, Beijing 100049, China
[2] Department of Physics, KAIST, Daejeon 34141, Korea

Abstract: We study an effective theory beyond the standard model (SM) where either of the two additional gauge singlets, a Majorana fermion and a real scalar, constitutes all or some fraction of dark matter. In particular, we focus on the masses of the two singlets in the range of $O(10)$ MeV-$O(10)$ GeV with a neutrino portal interaction, which plays an important role not only in particle physics but also in cosmology and astronomy. We point out that the thermal dark matter abundance can be explained by (co-)annihilation, where the dark matter with a mass greater than 2 GeV can be tested in future lepton colliders, CEPC, ILC, FCC-ee and CLIC, in the light of the Higgs boson invisible decay. When the gauge singlets are lighter than $O(100)$ MeV, the interaction can affect the neutrino propagation in the universe due to its annihilation with cosmic background neutrino into the gauge singlets. Although in this case it can not be the dominant dark matter, the singlets are produced by the invisible decay of the Higgs boson at such a rate which is fully within reach of future lepton colliders. In particular, a high energy cutoff of cosmic-ray neutrino, which may account for the non-detection of Greisen-Zatsepin-Kuzmin (GZK) neutrino or the non-observation of the Glashow resonance, can be set. Interestingly, given the cutoff and the mass (range) of WIMPs, a neutrino mass can be "measured" kinematically.

Keywords: neutrino portal, dark matter, neutrino mass, lepton collider, cosmic-ray, effective theory, Higgs boson

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1 Introduction

Weakly Interacting Massive Particles (WIMPs) are promising candidates for dark matter [1-4]. However, WIMPs with mass 6 GeV-$O(10^2)$ TeV have been severely constrained by the XENON, LUX and PandaX experiments [5-12]. This situation gives the motivation to investigate WIMPs lighter than a GeV. Such a WIMP should be a singlet of the Standard Model (SM) gauge group to avoid the LEP constraints [13]. If a gauge singlet dark matter is stabilized by a hidden symmetry, its possible interaction with the SM particles is represented by a portal coupling $O_{SM}O_{DM}$, where O_{SM} (O_{DM}) is an SM gauge singlet operator composed only of SM fields (only of hidden fields including WIMP).

It is interesting to study the neutrino portal interaction, i.e. $O_{SM} = \phi_H \cdot L$, where L is a Weyl spinor for a left-handed lepton, and we will take Weyl representation hereafter; ϕ_H is the Higgs doublet field; the dot denotes the contraction of the $SU(2)$ gauge indices, while the Lorentz indices are omitted. This is because this interaction can be not only a window from the SM to a dark sector but also affects neutrino and the Higgs boson physics.

Neutrino portal dark matter has been studied in several contexts: asymmetric dark matter [14, 15], decaying dark matter [16], and WIMP dark matter [17-20]. The first part of this paper can be classified as the last one. In particular, we will focus on the dark matter mass range between ~ 10 MeV and ~ 10 GeV, which differs from the previous studies where the mass is greater than several GeV. More concretely, we take an effective field theory approach based on the strategy of simplicity, and focus on the simplest neutrino portal operator of dimension five,

$$\frac{\phi_H \cdot L \psi \phi}{M}, \tag{1}$$

where ψ (ϕ) is a Majorana fermion (a real scalar) carrying a hidden Z_2 charge, and $\frac{1}{M}$ is a dimensionless coupling. Therefore, the lighter of ψ and ϕ is stable.

We point out that ψ and ϕ are restricted to be nearly degenerate to satisfy the neutrino mass constraints from the observations of cosmic microwave background (CMB) and baryon acoustic oscillations (BAO) [21]. Otherwise, a sizable neutrino mass would be produced radiatively from the neutrino portal interaction. This allows the Higgs boson to decay into ψ, ϕ and neutrino kinematically. Such an invisible decay rate could be measured in several future lepton colliders, such as the Circular Electron Positron Collider (CEPC), International Linear Collider (ILC), FCC-ee, and Compact Linear Collider (CLIC) [22-26], and thus could be a probe of the dark matter or neutrino physics.

We show that the observed thermal dark matter abundance can be explained by (co-)annihilation of WIMPs through the neutrino portal interaction. Furthermore, this dark matter, if heavier than around 2 GeV, can be tested in future lepton colliders.

In the second part, we study the neutrino propagation in the universe with the neutrino portal interaction. We show that neutrino propagation is affected only when the invisible decay of the Higgs boson is at such a rate which is fully within the sensitivity reach of future lepton colliders. This possibility is interesting because in the IceCube neutrino observatory [27, 28] the cosmic-ray neutrino event above PeV has not yet been detected, especially of the Greisen-Zatsepin-Kuzmin (GZK) neutrinos [29-31]. Also, the Glashow resonance [32] has not been observed. We point out that the absence of high energy cosmic-ray neutrinos can be explained if annihilation of the neutrino-(anti)neutrino into WIMPs takes place before the neutrino arrives to Earth. Namely, a cutoff can be set from the neutrino portal interaction. Moreover, neutrino mass is constrained kinematically from the mass range of WIMPs with a given cutoff, e.g. for a cutoff of a few PeV, which could explain the non-observation of the Glashow resonance, one of the neutrino masses is within 0.01-0.2 eV. On the other hand, for a cutoff around 10 PeV, which may explain the non-detection of GZK neutrino, one of the neutrino masses is within 0.008-0.1 eV. Therefore, the neutrino mass can be "measured" kinematically through the neutrino portal interaction.

A UV model is built to justify the setup and to study the experimental constraints for heavy particles relevant for generation of the higher dimensional terms. In this model, the neutrino mass can be dominantly obtained from the neutrino portal interaction.

The paper is organized as follows. In Section 2, we explain the model with several constraints and show that ϕ or ψ can explain the thermal dark matter abundance. In Section 3, the propagation of cosmic-ray neutrinos with the neutrino portal interaction is discussed. In Section 4, the UV model is discussed. The last section is devoted to discussion and conclusions.

2   A simple effective theory for WIMP

To simplify the discussion, suppose that the Lagrangian additional to the SM, $\mathcal{L}_{SM}$, has only one generation of neutrino,

$$\delta\mathcal{L} = \psi\bar{\sigma}_\mu\partial^\mu\bar{\psi} + \frac{1}{2}\partial^\mu\phi\partial_\mu\phi - \frac{\phi_H \cdot L\psi\phi}{M} - \frac{M_\psi}{2}\psi\psi$$
$$+ h.c - \frac{m_\phi^2}{2}\phi^2 - V(\phi, \phi_H), \qquad (2)$$

where the total Lagrangian is given by $\mathcal{L} = \mathcal{L}_{SM} + \delta\mathcal{L}$; M_ψ (m_ϕ) is the mass of ψ (ϕ); $V(\phi, \phi_H)$ is the potential of the scalar fields which is supposed to give a vanishing vacuum expectation value (VEV), $\langle\phi\rangle = 0$, and additional mass squared, $\left\langle\frac{\partial^2 V}{\phi^2}\right\rangle = 0$, to ϕ. We will neglect the Higgs portal term, $\lambda_H\phi^2|\phi_H|^2$, in $V(\phi, \phi_H)$, because the scalar mass is lighter than 10 GeV, and λ_H is sufficiently small if $\lambda_H \lesssim \frac{m_\phi^2}{v^2}$, [1] where v=174 GeV is the VEV of the Higgs field. A small portal coupling larger than the order of $\frac{1}{16\pi^2}\left(\frac{\Lambda_{c.o}}{M}\right)^2$ is stable under quantum corrections, where $\Lambda_{c.o}$ is the cutoff scale of the model, which could be smaller than M. The other dimension-five operators, $(\phi_H \cdot L)^2$, $F_Y^{\mu\nu}\psi\bar{\sigma}_\mu\sigma_\nu\psi$, $|\phi_H|^2\psi^2$, and $\phi^2\psi^2$ are suppressed due to approximate lepton number symmetry under which L and $\bar{\psi}$ are 1 while the others are 0. [2] In particular, we suppose that a tree-level $(\phi_H \cdot L)^2$–term induces a neutrino mass that is smaller, or of the same order, as the physical one. Note that the tree-level $(\phi_H \cdot L)^2$ term is not generated in a UV model if all heavy particle masses and interactions preserve the lepton number (see Section 4).

2.1   Constraint from the neutrino mass

For the broken phase of electroweak symmetry, one obtains an interaction $\frac{v}{M}\nu\psi\phi$. It was pointed out that the neutrino mass is generated at the 1-loop level in this broken phase interaction [33-36]:

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1) The neutrino portal models with an efficient Higgs portal interaction are studied in [17-20].

2) The coefficient of these terms are stable under quantum corrections if they are greater than $\frac{1}{16\pi^2}\frac{M_\psi}{M^2}$. The quantum correction for $(\phi_H \cdot L)^2$ will be discussed in the following.

$$m_\nu = \frac{1}{16\pi^2}\frac{\text{v}^2}{M^2}KM_\psi + O\left(\left(\frac{1}{16\pi^2}\right)^2\right)\frac{\text{v}^2}{M^2}M_\psi, \qquad (3)$$

where $K \equiv K\left(\dfrac{m_\phi^2}{M_\psi^2}\right)$ with $K(x) = 1 - \dfrac{x}{x-1}\log(x)$ satisfying $\lim_{x\to 1}K(x) = 1-x$. We have taken the renormalization scale $\mu = M_\psi$ so that this is an on-shell renormalization. Since it is constrained by the CMB and BAO observations [21] as

$$m_\nu \lesssim 0.2\text{eV}(95\%\text{CL}), \qquad (4)$$

while it is also constrained by the double beta decay experiment for an electron neutrino [37], the neutrino mass crucially restricts the mass range of the two WIMPs. In Fig. 1, the contour plot of the generated neutrino mass and its constraint (gray shaded region) are represented in the $m_\phi - M$ plane with M_ψ=12 MeV. One sees that m_ϕ is restricted to be around M_ψ, and smaller the M, smaller is the difference $|m_\phi - M_\psi|$. Since one of ψ and ϕ is stable, M has an upper bound for sufficient (co-) annihilation of ϕ or ψ so as not to over-close the universe. Thus, ϕ and ψ are constrained to be nearly degenerate,

$$m_\phi \simeq M_\psi. \qquad (5)$$

Note that this constraint disappears when ϕ is replaced by a complex scalar field with only a bilinear mass term because the lepton number symmetry recovers. However, let us pursue with the simple real scalar case with $m_\phi \simeq M_\psi$, The following discussion will be qualitatively the same in a specific parameter region but with complex extension of the scalar field.

2.2   Heavy boson decays in colliders

Since $m_\phi + M_\psi \lesssim$20 GeV in our consideration, the anomalous decays of the Higgs, W and Z bosons into ψ, ϕ and a lepton are possible. Thus, in colliders this scenario is constrained and tested. In particular, the Higgs boson invisible decay is represented as

$$H \to \psi + \phi + \nu\,(\overline{\nu}). \qquad (6)$$

The decay width of the process is obtained as

$$\Gamma_{H\to inv} \simeq \frac{1}{1536\pi^3}\frac{m_H^3}{M^2}, \qquad (7)$$

where m_H is the Higgs boson mass and the decay products are approximated to be massless. Given the total decay width of the Higgs boson $\simeq$ 4 MeV, the branching ratio of this process is estimated as

$$Br_{H\to inv} \simeq 0.01\%\left(\frac{10\,\text{TeV}}{M}\right)^2, \qquad (8)$$

where the bound from the LHC is $Br_{H\to inv}$<25% (95%CL) [38, 39].

On the other hand, the decay rate of W boson to a charged lepton and missing energy (Z boson to missing energy) can be estimated as $\Gamma_{W^-\to l^-+\text{missing}} \simeq \Gamma^{\text{tree}}_{W^-\to l^-+\overline{\nu}_l}$ $\left(1 - \dfrac{1}{24\pi^2}\left(\dfrac{\text{v}}{M}\right)^2\right)\left(\Gamma_{Z\to\text{missing}} \simeq \Gamma^{\text{tree}}_{Z\to\overline{\nu}+\nu}\left(3 - \dfrac{1}{12\pi^2}\left(\dfrac{\text{v}}{M}\right)^2\right)\right)$ at the leading order of the anomalous decay[1], where $\Gamma^{\text{tree}}_{W^-\to l^-+\overline{\nu}_l}$ $\left(\Gamma^{\text{tree}}_{Z\to\nu_l+\overline{\nu}_l}\right)$ is the decay rate of the subscript at the tree-level in the SM. The branching ratio of W boson to lepton + missing energy (Z boson to missing energy) differs from the SM by $1 \times 10^{-6}\%\left(\dfrac{10\,\text{TeV}}{M}\right)^2\left(8 \times 10^{-7}\%\left(\dfrac{10\,\text{TeV}}{M}\right)^2\right)$. The corresponding LEP bound is given as 0.1% (0.06%) [13].

One finds that M can be as small as $O(100)$ GeV to be consistent with the current experiments. To be conservative, let us set a bound[2]

$$M \gtrsim 400\,\text{GeV}. \qquad (9)$$

This is represented as the horizontal black band in Fig. 2.

On the other hand, the Higgs invisible decay with

$$M \lesssim 5\,\text{TeV} \qquad (10)$$

can be tested in future lepton colliders, such as the CEPC, ILC, FCC-ee, and CLIC, where the branching ratio of the
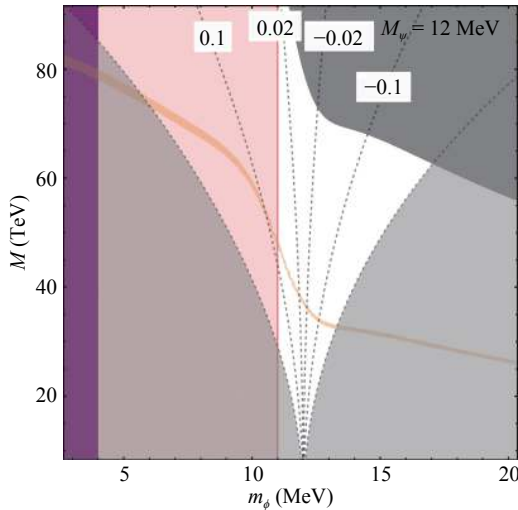

Fig. 1.   (color online) The contour plots of the radiatively generated neutrino mass [eV] with M_ψ = 12 MeV. The purple region may be excluded due to the neutrino effective number. In the gray region, the universe is over-closed. On the orange band, the thermal abundance of the lighter WIMP explains the dark matter. The pink region may be tested in future CMB/BAO observations.

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1) The processes with virtual ϕ, ψ emission and absorption are also included in the decay width.

2) For $M \lesssim 1$ TeV, one may care for the constraint for a heavy field in some UV models. The constraint in a UV model, which will be discussed Sec.4, is represented by a lower bound (34) similar to (9).

invisible decay is planned to be measured with a precision of around 0.1% (the purple shaded region in Fig. 2.) [22-26].
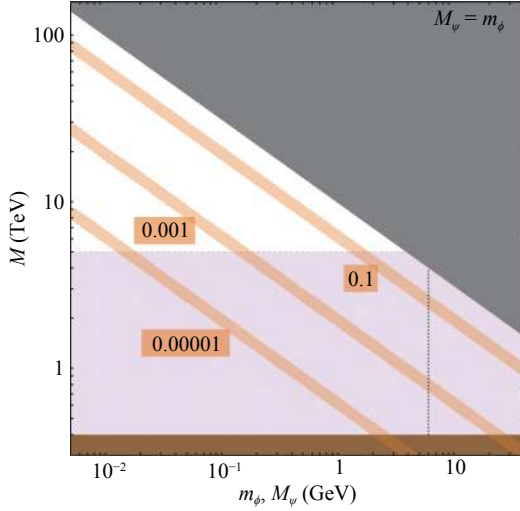

Fig. 2.   (color online) The contour plots of the WIMP relic abundance with $M_\psi \simeq m_\phi$. The brown region is excluded by heavy boson decays in colliders. The black dotted line denotes $M_\psi \simeq m_\phi = 6$ GeV. The light pink region could be tested in future lepton colliders by measuring the Higgs boson invisible decay rate.

2.3   Thermal relic abundance of WIMP

Let us now discuss the thermal relic abundance of the lighter of ϕ or ψ, which annihilates into (anti-)neutrinos through the $t(u)$-channel,

$$\phi + \phi \rightarrow \nu + \overline{\nu} \ (m_\phi < M_\psi), \tag{11}$$

$$\psi + \psi \rightarrow \nu + \nu, \overline{\nu} + \overline{\nu}, \nu + \overline{\nu} \ (M_\psi < m_\phi). \tag{12}$$

In the first row, one does not have $\phi + \phi \rightarrow \nu + \nu$ or $\overline{\nu} + \overline{\nu}$, because, by integrating out ψ, the corresponding effective vertex vanishes in the equation-of-motion for external neutrinos. The total annihilation cross-sections times the relative velocity at the tree-level are given as,

$$v_{\rm rel}\sigma_{\psi\psi}(s) \simeq \frac{1}{16\pi}\left(\frac{{\rm v}^2}{M^2}\right)^2 \frac{M_\psi^2}{\left(m_\phi^2 + M_\psi^2\right)^2}\left(1 + O\left(\frac{s}{4m_\phi^2}\right)\right),$$

and

$$v_{\rm rel}\sigma_{\phi\phi}(s) \simeq \frac{1}{2\pi}\left(\frac{{\rm v}^2}{M^2}\right)^2 \frac{M_\psi^2}{\left(m_\phi^2 + M_\psi^2\right)^2}\left(1 + O\left(\frac{s}{4M_\psi^2}\right)\right), \tag{13}$$

for annihilations of $\psi\psi$ and $\phi\phi$, respectively. The $O(s)$ terms are calculated by FeynRules and FormCalc [40, 41], which are also used to confirm all amplitude calculations in this paper. The dark matter abundance is es-

timated as

$$\Omega_{\phi,\psi}h^2 = 0.1\left(\frac{4 \times 10^{-26}{\rm cm}^3/{\rm s}}{\langle\sigma_{\rm eff}v\rangle}\right)\left(\frac{x_f\sqrt{g_*}}{5g_{*s}}\right), \tag{14}$$

where $\langle\sigma_{\rm eff}v\rangle$ is the thermally averaged annihilation cross-section given by

$$\langle\sigma_{\rm eff}v\rangle = \frac{\sum_i g_i^2 \int_{2m_i^2} {\rm d}s \sqrt{s}K_1(\sqrt{s}/T)\left(s/4 - m_i^2\right)\sigma_i(s)}{2T(\sum_i g_i m_i^2 K_2(m_i/T))^2}$$

($K_j(x)$ is the modified Bessel function of the j-th kind) [42], where the co-annihilation effect is included; $g_*(g_{*s})$ is the degree of freedom of the energy (entropy) density of the radiation, which is typically around 10-100 for $O(10)\,{\rm MeV} < m_{\rm DM} \lesssim 10$ GeV; $x_f = \frac{m_{\rm DM}}{T_f}$ is the freeze-out temperature in units of $m_{\rm DM} = \min(m_\phi, M_\psi)$, which is around 15-20; h=0.678. The region satisfying $\Omega_{\phi,\psi}h^2 \simeq 0.1$ is represented by the orange band in Fig. 1. In Fig. 2, the contours of $\Omega_{\phi,\psi}h^2$ are shown (orange bands) at the limit $M_\psi = m_\phi$. The width of the orange bands denotes the ambiguity of our calculation[1]. The gray regions in both figures denote the over-closure of the universe, $\Omega_{\phi,\psi} \gtrsim 1$. From the figures, one finds that the thermal dark matter can be tested in future lepton colliders with mass

$$m_\phi \simeq M_\psi \gtrsim 2\,{\rm GeV}. \tag{15}$$

If the lighter of ϕ or ψ is part of the dark matter, the mass range where it can be tested increases.

In particular, masses greater than 6 GeV are now being tested in the Xenon1T, LUX, and PandaX experiments [43-45] . The above boundary is represented as the black dotted line in Fig. 2. It is interesting to note that we can have a cross-check if the dark matter is detected in the direct-detection experiments.

2.4 $N_{\rm eff}$ and BBN

The mass of ϕ or ψ should be larger than MeV, otherwise the created neutrinos from their annihilations will change $N_{\rm eff}$ by $O(1)$ and could spoil the BBN [46-49]. According to [47],

$$\begin{aligned}&m_\phi > 5\,{\rm MeV},\ M_\psi > 7\,{\rm MeV}\ \left(\left|m_\phi - M_\psi\right| \gtrsim {\rm MeVs}\right)\\&m_\phi,\ M_\psi > 9\,{\rm MeV}\ (m_\phi \simeq M_\psi)\end{aligned} \tag{16}$$

is obtained from the bound for $N_{\rm eff}$ [21].

On the other hand, the viable region with

$$m_\phi\ {\rm or}\ M_\psi \lesssim 11\,{\rm MeV} \tag{17}$$

has a slightly larger $N_{\rm eff}$ and may be tested by several future CMB observations such as the PIXIE and CMB-S4 experiments, as well as the BAO observation [50-52].

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1) The width of a band in the figure is obtained by using the largest and smallest $\frac{x_f\sqrt{g_*}}{\sqrt{g_{*s}}}$ on the band.

3 Propagation of cosmic-ray neutrino with neutrino portal interaction

We now focus on the region where the lighter WIMP composes a fraction of the dark matter, $\Omega_{\phi,\psi}h^2 < 0.1$, i.e. the region with sufficiently strong neutrino portal interaction. The observed dark matter abundance can be explained by the other dark matter components: WIMP with neutrino portal interaction of different generation (see Section 5 and footnote 4), superpartner[1], inflaton [76-85], etc. This region is interesting because it would affect the neutrino propagation in the universe.

Although more statistics is needed, up to now no cosmic-ray neutrino event above several PeV has been detected in the IceCube experiment [27, 28], and the Glashow resonance around 6 PeV has also not been observed [32]. Despite several detections of cosmic-ray events of other kinds of particles up to $\sim 10^2$ EeV, this fact implies that there may be a special cutoff for cosmic-ray neutrino. In particular, if some of the observed cosmic-rays around 10^2 EeV are protons, cosmic-ray neutrinos of $O(\text{EeV})$ should also be observed. Protons of energy larger than $O(10^2)$ EeV interact with CMB photons and produce pions via the δ resonance, and hence lose energy before cosmic-ray neutrinos reach Earth. This scattering sets a GZK cutoff at an energy $O(10^2)$ EeV for protons [86, 87], which explains the observed cutoff for high energy cosmic-ray events. In the GZK cutoff scenario, GZK neutrinos of energy $O(\text{EeV})$ are produced from the decay of these pions [88] and should be detected at $O(0.1)-O(10)$ events/year in the IceCube neutrino observatory [29-31].

The absence of such energetic neutrino events can be explained from a viewpoint of particle physics[2]. Thanks to the neutrino portal interaction, annihilation between cosmic-rays and cosmic background neutrinos,

$$\nu/\bar{\nu} + \nu(\text{C}\nu\text{B}) \rightarrow \psi + \psi, \phi + \phi, \qquad (18)$$

is enhanced for sufficiently small M, so that before the neutrinos reach Earth they turn into WIMPs[3]. Namely, we propose that the neutrino portal interaction can set a cutoff for cosmic-ray neutrinos.

To set a cutoff, there are two conditions. First, the annihilation channel should be turned on at $E_\nu > E_\nu^{\text{cutoff}}$, and hence the center-of-mass energy of the neutrino-(anti)

neutrino system, E_{cm}, should be greater than the threshold, $2M_\psi$ or $2m_\phi$, at E_ν^{cutoff}, as

$$\frac{E_{\text{cm}}}{2} \equiv \frac{1}{\sqrt{2}}\sqrt{m_\nu^2 + E_\nu E_{\text{C}\nu\text{B}}(1-\cos\theta)} \gtrsim M_\psi \text{ or } m_\phi, \quad (19)$$

where E_ν^{cutoff} is defined by

$$M_\psi \text{ or } m_\phi \sim \left(\frac{E_\nu^{\text{cutoff}}}{6\,\text{PeV}} \frac{E_{\text{C}\nu\text{B}}}{0.2\,\text{eV}}\right)^{\frac{1}{2}} 35\,\text{MeV}. \qquad (20)$$

Here, $E_{\text{C}\nu\text{B}} \simeq \max[T_\nu, m_\nu]$ is the typical energy of cosmic background neutrinos with temperature $T_\nu \simeq 2\times10^{-4}$ eV, and θ is the angle between the momenta of two neutrinos.

Secondly, the mean free path, $\text{d}(E_\nu)$, imposed by annihilation, should be shorter than the distance to the neutrino source. To discuss this, let us neglect for simplicity the neutrino oscillation[4]. Following [94], one obtains the mean free path of a neutrino given by

$$\text{d}(E_\nu) \simeq \int \frac{\text{d}^3\vec{p}}{(2\pi)^3} \sigma_{\nu\nu}(E_{\text{cm}}(\vec{p},E_\nu)) f_{\text{C}\nu B}(\vec{p}). \qquad (21)$$

Here, $f_{\text{C}\nu B}(\vec{p}) = 2\left(e^{|\vec{p}|/T_\nu}+1\right)^{-1}$ is the neutrino distribution function for cosmic background neutrinos, and $\sigma_{\nu\nu}(E_{\text{cm}})$ is the helicity averaged neutrino-(anti)neutrino annihilation cross-section, which is, with $m_\phi \simeq M_\psi$, approximated as

$$\sigma_{\nu\nu}(E_{\text{cm}}) \simeq$$

$$\frac{v^4}{M^4} \frac{\sqrt{\left(\frac{E_{\text{cm}}}{2}\right)^2 - m_\phi^2} + E_{\text{cm}}\log\left(\frac{\sqrt{\left(\frac{E_{\text{cm}}}{2}\right)^2 - m_\phi^2} + \frac{E_{\text{cm}}}{2}}{m_\phi}\right)}{16\pi E_{\text{cm}}^3}. \qquad (22)$$

The neutrino flux from the source at a distance L is then reduced by a factor

$$\kappa(E_\nu) = e^{-\frac{L}{d(E_\nu)}}. \qquad (23)$$

Here, we have neglected the effect of the redshift of E_ν due to the expansion of the universe, which would reduce the observed E_ν in IceCube by $O(10)\%$ with $L \sim O(1)$ Gpc.

For instance, the predicted neutrino flux is represented in Fig. 3 by assuming a neutrino flux before annihilation of

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1) There are several typical lightest superpartners (LSPs) which might be the dominant dark matter, depending on SUSY breaking scenarios: gravitino LSP in gauge mediation [53], bino-like LSP with SUSY breaking in a gauge unified manner, wino-like LSP in anomaly mediation and simple SUSY breaking scenarios based on the anomaly mediation [54-62], N=2 superpartners in N=2 partial breaking [63-75], etc.

2) There are also astronomical explanations for the absence of neutrino events above several PeVs. For example, if the neutrino is originated from the galaxy clusters or starburst galaxies the non-observation of the Glashow resonance can be accounted for [89]. If the $O(10^2)$ EeV cosmic-rays observed are composed of heavy nuclei such as iron, the absence of the GZK neutrino event can be explained [90].

3) The explanation of the neutrino events especially for absorption lines in the observed neutrino flux at IceCube, in the light of the interaction between a cosmic-ray neutrino and a cosmic background neutrino, was discussed in several recent studies. [91-99].

4) Given the neutrino oscillation, all kinds of the neutrinos share the strongest neutrino interaction and the mean free path for each neutrino should be Eq.(21) times a factor $\sim O(1)$. Thus, in the multi-generation extension of the neutrino portal interaction, one does not need all the interactions to be strong to set the cutoff for different kinds of neutrinos, and this allows one of the WIMPs becomes the dominant dark matter.
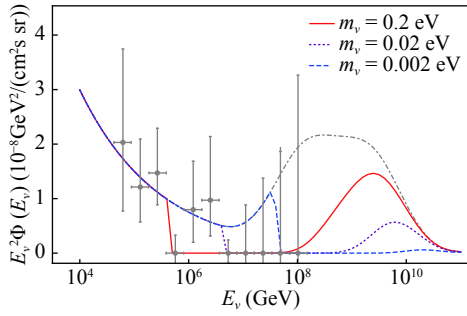
Fig. 3.   (color online) The predicted neutrino flux with several m_ν. $L = 300$ Mpc, $M_\psi = m_\varphi = 9$ MeV and $M = 450$ GeV are fixed. The red solid, purple dotted, and blue dashed lines represent the flux with $m_\nu = 0.2$ eV, 0.02 eV, and 0.002 eV, respectively. The gray points represent the IceCube observations obtained from [28] , while the gray dot-dashed line represents the flux distribution before annihilation Eq.(24).

$$\phi(E_\nu)E_\nu^2 = 1.5 \times 10^{-8} \left(\frac{E_\nu}{10^5 \text{ GeV}} \right)^{-0.3} + V_{\text{GZK}}(E_\nu). \qquad (24)$$

The first term is the best-fit power law in [28] , while the second term represents a toy GZK neutrino flux,

$$10^8 V_{\text{GZK}}(E_\nu) = \left(e^{\left(\frac{3}{8} \cos\left(\pi \frac{\log_{10}(E_\nu/\text{GeV})-8.75}{2.5} \right) - \frac{1}{24} \cos\left(3\pi \frac{\log_{10}(E_\nu/\text{GeV})-8.75}{2.5} \right) \right)} - e^{-\frac{1}{3}} \right)$$

for $E_\nu > 10^{6.25}$ GeV, otherwise it is 0. (Realistic ones for GZK neutrino are given in [29-31].) The neutrino flux in our scenario is approximated as

$$\kappa(E_\nu)\phi(E_\nu)E_\nu^2. \qquad (25)$$

In Fig. 3, one finds that the neutrino flux does get a cutoff or an absorption band through the t-channel annihilation. Note that the cutoff is less efficient in a model where there is a significant s-channel annihilation/scatter-

ing process. In fact, the s-channel process itself does not contribute like a "cutoff", but is an absorption line due to the quick decrease of the cross-section when the center-of-mass energy exceeds the threshold. Moreover, the scattering process between neutrino and (anti)neutrino is at tree-level if the s-channel process exists. This is constrained by the CMB observation [100] and the efficiency of the cutoff is bounded. In our case, the scattering process is 1-loop suppressed and this bound is much looser than the heavy boson decay.

3.1   Relation with heavy boson decay

The contour plot of $d(E_\nu)$ is represented in Fig. 4. From the left panel, one finds that the neutrino flux originating from a distance

$$L > O(10) \text{ Mpc} \qquad (26)$$

can be affected by the neutrino portal interaction.

As shown in the left panel of Fig. 4, we have checked that to obtain $d(E_\nu)$ smaller than the scale of particle horizon size ~ 10 Gpc, i.e. where neutrino propagation in the universe could be affected, M should be smaller than ~ 2 TeV. From the right panel, where M is around the lower bound (9), one reads the upper bound of $M_\psi \simeq m_\phi$ to be around $O(100)$ MeV. This upper bound decreases with larger M due to the scaling of the cross-section. Hence, one obtains the parameter range where the neutrino propagation in the universe is affected,

$$M \simeq 0.4 - 2 \text{ TeV and } m_\phi \simeq M_\psi \sim 9 - O(100) \text{ MeV}. \qquad (27)$$

Since the upper bound of M satisfies (10), the following is predicted: if the high energy neutrino flux in IceCube is affected by the neutrino portal interaction, the Higgs invisible decay is fully within reach of future lepton colliders. Let us emphasize again that M required here is
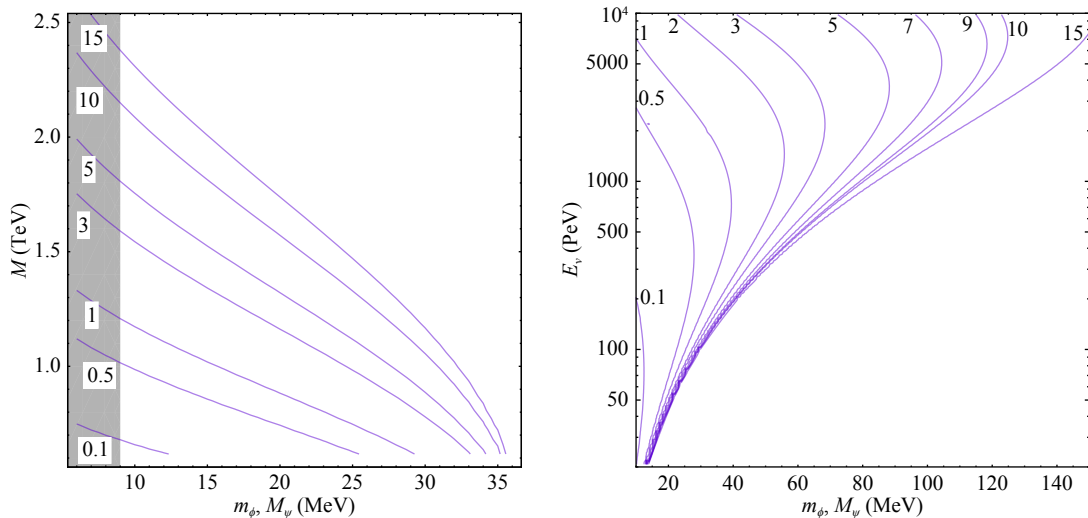

Fig. 4.   (color online) The contour plot of the mean free path $d(E_\nu)$ [Gpc] for neutrino with $M_\psi \simeq m_\phi$. In the left panel, $E_\nu = 6$ PeV and $m_\nu = 0.2$ eV is fixed. The vertical shaded region represents the constrained region from the neutrino effective number. In the right panel, $M = 630$ GeV and $m_\nu = 0.01$ eV are fixed.

much smaller than the one for Eq. (14), and we can not provide dominant dark matter whose interaction affects cosmic-ray neutrinos. However, to explain the neutrino oscillation, an extension with several flavors of ψ or ϕ is needed. (See conclusions and discussion.) In this case, some of the flavors could be the dominant dark matter, while some could affect the spectra of cosmic-ray neutrinos. We note that in this case the dark matter should be lighter than the particles relevant for the cutoff, and thus the dark matter mass is lighter than $O(100)$ MeV.

3.2 Measuring the physical neutrino mass range

Neutrino flux carries information on annihilation during its propagation. In particular, as in a collider, one can measure $E_{C\nu B}$ kinematically once the cutoff and the masses of ϕ and ψ are given. This implies that a mass scale can be obtained for one of the neutrinos if its mass is greater than T_ν.

Even if just the mass range of ϕ, ψ is given, one can predict the neutrino mass range. Since $E_{C\nu B} E_\nu^{cutoff} \sim m_\phi^2 \simeq M_\psi^2$, with the cutoff scale fixed, one obtains

$$E_{C\nu B} \propto m_\phi^2 \simeq M_\psi^2, \tag{28}$$

which implies that $E_{C\nu B}$ ($m_\phi \simeq M_\psi$) has a lower bound corresponding to Eq.(16) ($E_{C\nu B} \gtrsim T_\nu$). If $E_{C\nu B} > T_\nu$ at the lower bound of Eq.(16), the lower bound of one of neutrino masses is predicted.

For instance, for a cutoff at 6 PeV, one obtains a lower bound of neutrino mass $m_\nu \gtrsim 1.4 \times 10^{-2}$ eV. With a cutoff $\lesssim 6$ PeV, which may explain the non-observation of the Glashow resonance, the lower bound for neutrino mass increases. Thus the neutrino mass range is predicted as

$$m_\nu \simeq 0.01 - 0.2 \text{ eV} \quad (\text{for } E_\nu^{cutoff} \lesssim 6 \text{ PeV}). \tag{29}$$

The neutrino flux around the lower limit is illustrated by the purple dotted line in Fig. 3. This mass range covers the atmospheric neutrino scale of 0.05 eV.

If the GZK neutrino source originates from $\simeq O(1)$ Gpc away from Earth, m_ϕ and M_ψ should be smaller than $O(100)$ MeV. This can be found in the right panel of Fig. 4, because for any $E_\nu m_\nu$ this is almost satisfied. Then, for a cutoff around 10 PeV, which may explain the non-detection of GZK neutrino (See the blue dashed line in Fig. 3), the neutrino mass range can be estimated as

$$0.008 \text{ eV} \lesssim m_\nu \lesssim 0.1 \text{ eV} \left(\text{for } E_\nu^{cutoff} \simeq 10 \text{ PeV}\right). \tag{30}$$

The lower bound, where annihilation is most efficient, is close to the solar neutrino scale of 0.009 eV.

4 UV model

In the previous sections, we studied a dimension 5 op-

erator with two additional gauge singlets. It is questionable whether there is a UV model, and if there is, whether the constraints for heavy particles in the UV model restrict our scenario, especially for $M \lesssim$ TeV.

To suppress the $(\phi_H \cdot L)^2$ term in order to satisfy the constraint (4) at the tree-level, the UV model should also have an approximate lepton number conservation. One of such UV models is given by

$$\mathcal{L} = -y\phi_H LN - \tilde{M}SN - \frac{\tilde{M}}{2f}\phi S\psi - \frac{M_N}{2}NN - \frac{M_\psi}{2}\psi\psi$$

$$+ h.c. - \frac{m_\phi^2}{2}\phi^2 - V(\phi, \phi_H), \tag{31}$$

where S and N are gauge singlet Weyl fermions with lepton number 1 and -1, respectively, and we have omitted the kinetic terms. For later convenience, we introduce the Yukawa coupling y, the decay constant f, the order parameter $\tilde{M}$, and the mass parameter M_N satisfying $M_N, M_\psi \ll \tilde{M}$ due to the approximate lepton number symmetry. We have forbidden the $M_S SS$ term at the tree-level by imposing Z_4 symmetry, under which S, N, ψ and the spurion $\tilde{M}$ have a charge 1/4, 1/2, 1/2 and 1/4, respectively. Leptons can carry a charge 1/2, so that Yukawa couplings are allowed. Thus, this symmetry is identified to be spontaneously broken down due to the VEV $\tilde{M}$ of some scalar field.

By making a shift $S \rightarrow S - \frac{y}{\tilde{M}}\phi_H L$, one finds that the neutrino portal interaction appears as

$$\mathcal{L} \rightarrow \frac{1}{M}\phi_H \cdot L\phi\psi - \tilde{M}SN - \frac{\tilde{M}}{2f}\phi S\psi - \frac{M_N}{2}NN - \frac{M_\psi}{2}\psi\psi$$

$$+ h.c. - \frac{m_\phi^2}{2}\phi^2 - V(\phi, \phi_H), \tag{32}$$

with

$$\frac{1}{M} = -\frac{y}{2f}. \tag{33}$$

The neutrino portal term (1st term) is decoupled from the heavy fields, S and N. Moreover, $(\phi_H \cdot L)^2$ does not appear after integrating out the heavy fields up to 1-loop level, because $\phi_H \cdot L$ does not directly couple to the heavy field. This fact can be also checked by integrating out the heavy fields in the terms of Eq. (31) after diagonalizing the fermion mass matrix. Interestingly, with this UV model, the neutrino mass is generated purely radiatively through the neutrino portal interaction. The Higgs portal term, $|\phi_H|^2 \phi^2$, could be suppressed if ϕ is a pseudo-Nambu-Goldstone boson with breaking scale f (See discussion for a concrete non-linear sigma model)[1].

There are several constraints for N, because it be-

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1) At the tree-level, the portal coupling is of order $\sim m_\phi^2/f^2$ since a Nambu-Goldstone boson for an exact global symmetry does not have potential and m_ϕ could be the size of the explicit breaking.

haves as a right-handed neutrino with Yukawa coupling y and mass $\tilde{M}$ [101, 102]. If we adopt the constraint in [102] for a heavy right-handed neutrino, which dominantly mixes with τ neutrino, $\left|\frac{y\mathrm{v}}{\tilde{M}}\right| \lesssim 0.1$ is required. On the other hand, this effective theory should have $\frac{\tilde{M}}{2f} \lesssim \sqrt{4\pi}$ from the viewpoint of perturbative unitarity, and $\frac{\tilde{M}}{f} \lesssim 2\pi$ when ϕ is identified as a pion-like field, which, respectively, turn out to be

$$M \gtrsim \frac{\mathrm{v}}{0.1 \times \sqrt{4\pi}} \sim 500\,\mathrm{GeV} \text{ and } \frac{\mathrm{v}}{0.1 \times 2\pi} \sim 300\,\mathrm{GeV}. \quad (34)$$

5   Conclusions and discussion

In this paper, a simplest neutrino portal interaction for WIMPs was investigated, especially for the lightest WIMP mass in the range of $\mathcal{O}(10)\,\mathrm{MeV} - \mathcal{O}(10)\,\mathrm{GeV}$, where it is not yet severely constrained by direct detections. Neutrino portal interaction is interesting because it can affect not only collider physics for the Higgs boson but also neutrino physics.

We pointed out that the constraint for radiatively generated neutrino mass seriously restricts the parameter space so that the two WIMPs are nearly degenerate. Due to this restriction, the Higgs boson can decay into WIMPs plus a neutrino, and this invisible decay can be searched for in future lepton colliders, CEPC, ILC, FCC-ee and CLIC.

We showed that the neutrino portal interaction can successfully (co-)annihilate the lightest WIMP, and that the WIMP relic abundance can explain the observed abundance of dark matter. Such a neutrino portal dark matter could be tested in future lepton colliders for mass $\gtrsim 2\,\mathrm{GeV}$.

In the case that WIMPs explain a small fraction of the dark matter abundance, the neutrino propagation in the universe can be significantly affected. We pointed out that this region can be fully tested in future lepton colliders. In particular, this region can set a cutoff for cosmic-ray neutrino and can explain the non-detection of the GZK neutrino events, and the non-observation of the Glashow resonance in IceCube. Moreover, a neutrino mass can be "measured" kinematically from the scale of the cutoff and WIMP mass.

Using a UV model, we have justified our set up and showed that neutrino mass can be dominantly generated by the neutrino portal interaction.

Since there are generations in the SM, it is natural to make an extension of the neutrino portal interaction (1) to the case of 3 generations, e.g. $\phi_H \cdot L^i Y_{ij} M^{-1} \psi^j \phi \left(\phi_H \cdot L^i Y_{ij} M^{-1} \psi \phi^j\right)$, where i, j denotes the generation and Y_{ij} is the dimensionless coupling in the mass basis of ψ^j (ϕ^j). The neutrino mass matrix is generated with $m_{vij} \simeq \frac{\sum_k Y_{ik} M_{\psi k} Y_{kj} K_k}{16\pi^2} \frac{\mathrm{v}^2}{M^2} \left(\frac{\sum_k Y_{ik} Y_{kj} K_k M_\psi}{16\pi^2} \frac{\mathrm{v}^2}{M^2}\right)$ where K_j is K from Eq. (3) but with M_ψ (m_ϕ) replaced by the mass of ψ_j (ϕ_j). In this extension, several parameter regions previously discussed can be simultaneously realized by the neutrino portal interactions of different generations.

Let us now provide a natural realization of the UV model Eq. (31) for our relevant parameter ranges where m_ϕ and M are sufficiently small. A light scalar ϕ suggests a naturalness problem. One of the solutions of this problem[1] is to identify ϕ as a pseudo-Nambu-Goldstone boson. Consider now the spontaneous breaking of an approximate $SU(2) \times U(1)$ global symmetry to $U(1)$ by some non-perturbative effect, in analogy with the chiral symmetry breaking in QCD. If all explicit breaking terms of $SU(2) \times U(1)$ can be identified as spurions with even charges under the residual $U(1)$, this residual symmetry contains an exact Z_2 symmetry. The $U(1)$ charged pion, say π_+, is Z_2 odd and contains ϕ as $\pi_+ = \frac{\phi + \mathrm{i}\tilde{\phi}}{\sqrt{2}}$. This possibility not only explains the smallness of m_ϕ, but also allows a rather small decay constant, f, for the composite scalar ϕ, like the pion decay constant in QCD.

To be concrete, let us consider the following non-linear realized Lagrangian for pions,

$$\mathcal{L}_{\mathrm{UV}} = \mathcal{L}_{\mathrm{sym}} + \mathcal{L}_{\mathrm{exb}}, \quad (35)$$

$$\mathcal{L}_{\mathrm{sym}} = \overline{\vec{N}}\vec{\sigma}^\mu \partial_\mu \vec{N} - \langle \vec{\phi} \rangle \cdot \mathrm{e}^{\mathrm{i}\frac{\pi^a \sigma_a}{2f}} \cdot \vec{N} S + \mathrm{h.c.}, \quad (36)$$

$$\mathcal{L}_{\mathrm{exb}} = -\frac{M_\psi}{2}\psi\psi - \frac{m_\phi^2}{2}\phi^2 - \frac{M_N}{2}NN$$
$$- \frac{\tilde{m}^2}{2}\tilde{\phi}^2 - \frac{m_0^2}{2}\pi_0^2 - y\phi_H \cdot LN + \mathrm{h.c.}. \quad (37)$$

Here, $\mathcal{L}_{\mathrm{sym}}$ is $SU(2) \times U(1)$ symmetric Lagrangian, while the terms that explicitly break $SU(2) \times U(1)$ are collected in $\mathcal{L}_{\mathrm{exb}}$; $\vec{N} = (N, \psi)$ is a matter doublet with $U(1)$ charge -1/2 and lepton number -1, while the fermion S carries lepton number 1; $\langle \vec{\phi} \rangle = (\tilde{M}, 0)$ is the VEV of an $SU(2)$ doublet operator with $U(1)$ charge -1/2, and the second term of Eq. (36) turns out to be the second and third terms in Eq. (31).

The mass parameters $M_\psi, m_\phi, \tilde{m}, m_0$ are the explicit breaking terms of the $SU(2) \times U(1)$ symmetry, and can be smaller than $\tilde{M}$ and f naturally. In particular, the un-

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1) Alternatively, this may indicate that a SUSY extension of the SM has a SUSY breaking soft scale around MeVs in the Z_2 odd sector, while that in the SM sector is above TeV to survive the experimental constraints. A candidate is a gauge mediation scenario [53], where sparticles charged under the SM gauge group gain weight via gauge interaction, while a singlet scalar acquires a highly suppressed mass either from higher order correction or the gravity effects.

broken $U(1)$ is explicitly broken down to the exact Z_2 symmetry by M_ψ and $\tilde{m}^2 - m_\phi^2$. Since ψ can also be charged under lepton number instead of ψ, the 1-loop neutrino mass is suppressed by an additional factor of $\frac{\tilde{m}^2 - m_\phi^2}{\tilde{m}^2 + m_\phi^2}$ and could reduce the tuning between M_ψ and m_ϕ to satisfy the neutrino mass constraint.

In this model, with additional light particles which are assumed to be lighter than the Higgs boson, the possibility of testing in future lepton colliders is even increased. Although the neutrino mass constraint is alleviated, and M_ψ can deviate from $m_{\tilde\phi} \simeq m_\phi$, for given values of the mass and the cross-section for dark matter-dark matter (neutrino-(anti)neutrino), the increase of $\max(M_\psi, m_{\tilde\phi} \simeq m_\phi)$ leads to the increase of the neutrino portal coupling $1/M$. Thus, the Higgs invisible decay rate is even enhanced for the regions of thermal dark matter and affects the propagation of cosmic-ray neutrino.

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