Dark Matter, Neutrino mass, Cutoff for Cosmic-Ray Neutrino, and Higgs Boson Invisible Decay from a Neutrino Portal Interaction

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Abstract

We study an effective theory beyond the standard model (SM) where either of 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 $O(10)\text{ MeV} - O(10)\text{ GeV}$, with a neutrino portal interaction which plays important roles not only in particle physics but also in cosmology and astronomy. We point out that the dark matter abundance can be thermally explained with (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 $O(100)\text{ MeV}$, the interaction can affect the neutrino propagation in the universe due to its annihilation with the cosmic background neutrino into the gauge singlets. Although can not be the dominant dark matter in this case, the singlets are produced by the invisible decay of the Higgs boson at a rate fully within the reach of the 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) neutrinos or non-observation of Glashow resonance, can be set. Interestingly, given the cutoff and the mass (range) of the WIMP, a neutrino mass can be “measured” kinematically.

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1 Introduction

Weakly Interacting Massive Particles (WIMPs) are promising candidates of dark matter [1–4]. However, the WIMPs with mass $6 \text{ GeV} - \mathcal{O}(10^2) \text{ TeV}$ has been severely constrained by the XENON, LUX and PandaX experiments [5–12]. This situation gives the motivation to investigate WIMPs lighter than GeVs. 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 a SM gauge singlet operator composed only of the SM fields (only of the hidden fields including the WIMP).

It is interesting to study a 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; $\phi_H$ 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 of 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 categorized into the last one. In particular, we will focus on the dark matter mass range between $\sim 10 \text{ MeV}$ and $\sim 10 \text{ GeV}$, which differs from the previous studies where the mass is greater than GeVs. 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,

$$\frac{\phi_H \cdot L \psi \phi}{M}, \quad (1)$$

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

We point out that $\psi$ and $\phi$ are restricted to nearly degenerate to satisfy the neutrino mass constraint 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 decay into $\psi, \phi$ and neutrino kinematically. Such an invisible decay rate will 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 dark matter abundance can be thermally explained with (co)annihilation of the WIMPs through the neutrino portal interaction. Furthermore, this dark matter, if heavier than around 2 GeV, can be tested in the future lepton colliders.

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

This paper is organized as follows. In Sec.2 we will explain the model with several constraints and show that $\phi$ or $\psi$ can explain the dark matter abundance thermally. In Sec.3 the propagation of the cosmic-ray neutrino with the neutrino portal interaction will be discussed. In Sec.4 the UV model will be discussed. The last section is devoted to conclusions and discussion.
2 A simple Effective Theory for WIMP

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

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

where the total Lagrangian is given by $\mathcal{L} = \mathcal{L}_{SM} + \delta \mathcal{L}$; $M_\psi (m_\phi)$ is the mass of $\psi$ ($\phi$); $V(\phi, \phi_H)$ 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, $\langle \partial^2 V / \partial \phi^2 \rangle = 0$, to $\phi$. 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 $\lambda_{H}$ is sufficiently small if $\lambda_{H} \lesssim \frac{m_\phi^2}{v^2}$, where $v = 174$ GeV is the VEV of the Higgs field. A small portal coupling larger than the order $\frac{1}{16\pi^2} (\frac{\Lambda_{c.o}}{M})^2$ is stable under quantum correction, where $\Lambda_{c.o}$ is the cut off scale of the model which could be smaller than $M$. The other dimension-five operators, $(\phi_H \cdot L)^2, F^\mu_{\nu} \bar{\psi} \sigma_{\mu\nu} \psi, |\phi_H|^2 \psi^2$, and $\phi^2 \psi^2$ are suppressed due to approximate lepton number symmetry under which $L$ and $\bar{\psi}$ have 1 while others 0. In particular, we suppose that a tree-level $(\phi_H \cdot L)^2$-term induces a neutrino mass smaller or of the same order of the physical one. Notice that the tree-level $(\phi_H \cdot L)^2$-term is not generated in a UV model if all the heavy particle masses and interaction preserve lepton number (see Sec.4.).

2.1 Constraint from Neutrino Mass

At the broken phase of the electroweak symmetry, one obtains an interaction $\frac{\nu}{M} \bar{\nu} \psi \phi$.

It was pointed out that the neutrino mass is generated at the 1-loop level in this broken phase interaction [33–36]:

$$m_\nu = \frac{1}{16\pi^2} \frac{v^2}{M^2} K M_\psi + O((\frac{1}{16\pi^2})^2) \frac{v^2}{M^2} M_\psi, \quad (3)$$

where $K \equiv K (\frac{m_\phi^2}{M_\psi^2})$ with $K(x) = 1 - \frac{x}{x-1} \log (x)$ 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 is constrained by the CMB and BAO observations [21] as

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

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

2The 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.
while 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 the Fig.1, the contour plot of the generated neutrino mass and the constraint on it (gray shaded region) are represented in $m_\phi - M$ plane with $M_\psi = 12$ MeV. One sees that $m_\phi$ is restricted to be around $M_\psi$, and the smaller the $M$, the smaller the difference $|m_\phi - M_\psi|$. Since one of $\psi$ and $\phi$ is stable, $M$ has an upper bound for sufficient (co)annihilation of $\phi$ or $\psi$ not to over-close the universe. Thus, $\phi$ and $\psi$ are constrained to be nearly degenerated,

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

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

### 2.2 Heavy Boson Decays in Colliders

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

$$H \rightarrow \psi + \phi + \nu \ (\bar{\nu}). \quad (6)$$

The decay width of the process is obtained as

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

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 are estimated as

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

where the bound from the LHC is $Br_{H \rightarrow \text{inv}} < 25\%(95\%\text{CL})$ [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 $\Gamma_{W \rightarrow l^- + \text{missing}} \simeq \Gamma_{W \rightarrow l^- + \text{missing}}^{\text{tree}} (1-$
Fig. 1: The contour plots of the radiatively generated neutrino mass [eV] with $M_\psi = 12$ 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 the future CMB/BAO observations.

$$\frac{1}{24\pi^2} \left(\frac{\nu}{M}\right)^2 \left(\Gamma_{Z\to\text{missing}} \approx \Gamma_{Z\to\nu_\tau+\nu}^{\text{tree}} (3 - \frac{1}{12\pi^2} \left(\frac{\nu}{M}\right)^2)\right)$$

at the leading order of the anomalous decay$^3$, where $\Gamma_{Z\to\nu_\tau+\nu}^{\text{tree}} (\Gamma_{Z\to\nu_\tau+\nu}^{\text{tree}})$ is the decay rate of the subscript at the tree-level in the SM. The branching ratio of $W$-boson to lepton + missing ($Z$-boson to missing) differs from the SM one by $1 \times 10^{-6}\% \left(\frac{10}{M}\right)^2$ ($8 \times 10^{-7}\% \left(\frac{10}{M}\right)^2$).

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$^4$

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

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$^3$The processes with virtual $\phi, \psi$ emission and absorption are also included in the decay width.

$^4$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).
This is represented as the horizontal black band in Fig.2.

On the other hand, the Higgs invisible decay with

\[ M \lesssim 5 \text{ TeV} \quad (10) \]

can be tested in the future lepton colliders, such as the CEPC, ILC, FCC-ee, and CLIC, where the branching ratio of the invisible decay is planned to be measured at a precision around 0.1% (the purple shaded region in Fig.2.) [22–26].

### 2.3 Thermal Relic Abundance of WIMP

Now let us discuss the thermal relic abundance for the lighter of \( \phi \) or \( \psi \). The lighter one annihilates into (anti-)neutrinos through t(u)-channel,

\[ \phi + \phi \rightarrow \nu + \bar{\nu} \quad (m_\phi < M_\psi), \quad (11) \]
\[ \psi + \psi \rightarrow \nu + \bar{\nu}, \nu + \nu, \nu + \bar{\nu} \quad (M_\psi < m_\phi). \quad (12) \]

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

\[ v_{\text{rel}} \sigma_{\psi\psi}(s) \simeq \frac{1}{16\pi} \left( \frac{v^2}{M^2} \right)^2 \frac{M_\psi^2}{(m_\phi^2 + M_\psi^2)^2} \left( 1 + \mathcal{O} \left( \frac{s}{4m_\phi^2} \right) \right), \quad (13) \]

and

\[ v_{\text{rel}} \sigma_{\phi\phi}(s) \simeq \frac{1}{2\pi} \left( \frac{v^2}{M^2} \right)^2 \frac{M_\psi^2}{(m_\phi^2 + M_\psi^2)^2} \left( 1 + \mathcal{O} \left( \frac{s}{4M_\psi^2} \right) \right), \]

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

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

where \( \langle \sigma_{\text{eff}} v \rangle \) is the thermal averaged annihilation cross section given by

\[ \langle \sigma_{\text{eff}} v \rangle = \frac{\sum_i g_i^2 \int_{2m_i} ds \sqrt{s} K_1(\sqrt{s}/T) (s/4 - m_i^2) \sigma_i(s)}{2T(\sum_i g_i m_i^2 K_2(m_i/T))^2} \]

(\( K_j(x) \) is the modified Bessel function of the \( j \)-th kind) [42], where the coannihilation effect is included; \( g_s(g_{ss}) \) is the degree of freedom for the energy (entropy) density.
of the radiation which is typically around $10-100$ for $\mathcal{O}(10)$ MeV < $m_{DM}$ < $10$ GeV;
$x_f = \frac{m_{DM}}{T_f}$ is the freeze-out temperature in the unit of $m_{DM} = \min(m_\phi, M_\psi)$ 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.\footnote{\textsuperscript{5}The width of a band in the figure is obtained by using the largest and smallest $x_f\sqrt{\frac{\pi}{g^{*}}} \sqrt{\frac{\pi}{g^{*}}}$ on the band.} The gray regions in both figures denote the over-closure of the universe, $\Omega_{\phi,\psi} \gtrsim 1$. In this figure, one finds that the thermal dark matter can be tested in the future lepton colliders with mass

$$m_\phi \simeq M_\psi \gtrsim 2 \text{ GeV}. \quad (15)$$

If the lighter of $\phi$ or $\psi$ is part of the dark matter, the testable mass range increases.

In particular, the mass greater than 6 GeV is now testing in Xenon1T, LUX, and PandaX \cite{43–45} (This boundary is represented as the black dotted line in Fig.2.), and it is interesting that we can have a cross-check if the dark matter is detected in the direct-detection experiments.

### 2.4 $N_{\text{eff}}$ and BBN

The mass of $\phi$ or $\psi$ should be larger than MeVs, otherwise the created neutrinos from the annihilation of them or themselves will change $N_{\text{eff}}$ by $\mathcal{O}(1)$ and could spoil the BBN \cite{46–49}. According to \cite{47},

$$m_\phi > 5 \text{ MeV}, M_\psi > 7 \text{ MeV} (|m_\phi - M_\psi| \gtrsim \text{MeVs})$$

$$m_\phi, M_\psi > 9 \text{ MeV} (m_\phi \simeq M_\psi) \quad (16)$$

is obtained from the bound for $N_{\text{eff}}$ \cite{21}.

On the other hand, the viable region with

$$m_\phi \text{ or } M_\psi \lesssim 11 \text{ MeV} \quad (17)$$

has a slightly larger $N_{\text{eff}}$ and may be tested by several future CMB observations such as the PIXIE and CMB-S4 experiments, as well as the BAO observation \cite{50–52}. 
Fig. 2: The contour plots of the WIMP relic abundance with $M_\psi \simeq m_\phi$. The brown region is excluded by heavy boson decays in the colliders. The black dotted line denotes $M_\psi \simeq m_\phi = 6$ GeV. The light pink region is testable in the future lepton colliders by measuring the Higgs boson invisible decay rate.

3 Propagation of Cosmic-Ray Neutrino with Neutrino Portal Interaction

Now we 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 with other dark matter components: a WIMP with neutrino portal interaction of different generation (see Sec.5 and footnote 9), a superpartner$^6$, an inflaton [76–85], etc. This region is

$^6$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.
interesting because it would affect the propagation of the neutrino in the universe.

Although more statistics are needed, up to now no cosmic-ray neutrino event above several PeVs is detected in the IceCube experiment [27,28], and the Glashow resonance around 6 PeV is also not observed [32]. Despite several detections of cosmic-ray events of other kind particles up to \( \sim 10^2 \) EeV, this fact implies that there may be a special cutoff for the cosmic-ray neutrino. In particular, if some of the observed cosmic-rays around \( 10^2 \) EeV are protons, cosmic-ray neutrinos of \( \mathcal{O}(\text{EeV}) \) should also be observed. Protons of energy larger than \( \mathcal{O}(10^2) \) EeV interacts with a CMB-photon and produces pions via \( \Delta \)-resonance, and hence loses energy before the cosmic-ray neutrino reaches the earth. This scattering sets a GZK cutoff at energy \( \mathcal{O}(10^2) \) 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 \( \mathcal{O}(\text{EeV}) \) are produced from the decay of these pions [88] and should be detected at \( \mathcal{O}(0.1) - \mathcal{O}(10) \) events/year in the IceCube neutrino observatory [29–31].

The non-detection of such energetic neutrino events can be explained from a viewpoint of particle physics.\(^7\) Thanks to the neutrino portal interaction, annihilation between the cosmic-ray and the cosmic background neutrinos,

\[
\nu / \bar{\nu} + \nu (\text{C\nu B}) \rightarrow \psi + \psi, \phi + \phi, \tag{18}
\]

is enhanced with sufficiently small \( M \) so that before the neutrino reaches the earth it turns into the WIMPs.\(^8\) 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 become greater than the threshold, \( 2M_\psi \) or \( 2m_\phi \), at \( E_\nu^{\text{cutoff}} \), as

\[
\frac{E_{cm}}{2} = \frac{1}{\sqrt{2}} \sqrt{m_\nu^2 + E_\nu E_{\text{C\nu B}} (1 - \cos \theta)} \gtrsim M_\psi \text{ or } m_\phi \tag{19}
\]

\(^7\)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 \( \mathcal{O}(10^2) \) EeV cosmic-rays observed are composed of heavy nuclei such as iron, the absence of the GZK neutrino event can be explained [90].

\(^8\)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].
where \( E_{\nu}^{cutoff} \) is defined at the equality

\[
M_\psi \text{ or } m_\phi \sim \left( \frac{E_{\nu}^{cutoff}}{6 \text{ PeV}} \frac{E_{\mathrm{CvB}}}{0.2 \text{ eV}} \right)^{\frac{1}{2}} 35 \text{ MeV}.
\]

(20)

Here \( E_{\mathrm{CvB}} \simeq \max [T_\nu, m_\nu] \) is the typical energy of cosmic background neutrinos with temperature \( T_\nu \simeq 2 \times 10^{-4} \text{ eV} \), and \( \theta \) is the angle between the momenta of two neutrinos.

Secondly, the mean free path, \( d(E_\nu) \), imposed by the annihilation should be shorter than the distance to the neutrino source. To discuss this, let us neglect for simplicity the neutrino oscillation.\(^9\) Following [94], one obtains the mean free path of a neutrino given by

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

(21)

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

\[
\sigma_{\nu\nu}(E_{cm}) \simeq \frac{v^4}{M_\phi^4} \sqrt{(\frac{E_{cm}}{2})^2 - m_\phi^2} + E_{cm} \log \left( \frac{\sqrt{(E_{cm})^2 - m_\phi^2} + E_{cm}}{m_\phi} \right).
\]

(22)

Then the neutrino flux from the source at \( L \) distant place is weakened by a factor of

\[
\kappa(E_\nu) = e^{-\frac{L}{\kappa_{(E_\nu)}}}.
\]

(23)

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

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

\[
\Phi(E_\nu)E_\nu^2 = 1.5 \times 10^{-8} \left( \frac{E_\nu}{10^5 \text{ GeV}} \right)^{-0.3} + V_{GZK}(E_\nu).
\]

(24)

\(^9\) 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.
The first term is the best-fit power law in [28] while the second term represents a toy GZK neutrino flux,
\[
10^8 V_{\text{GZK}}(E_\nu) = \left( e^{\frac{3}{8} \cos \left( \frac{\pi}{2.5} \log_{10} \left( \frac{E_\nu}{\text{GeV}} \right) - 8.75 \right)} - \frac{1}{24} \cos \left( 3 \pi \log_{10} \left( \frac{E_\nu}{\text{GeV}} \right) - 8.75 \right) \right) - e^{-\frac{1}{3}}
\]
for \( E_\nu > 10^{6.25} \text{ GeV} \), otherwise 0. (The realistic ones for GZK neutrino are given in [29–31].) The neutrino flux in our scenario is approximated as
\[
\kappa(E_\nu) \Phi(E_\nu) E_\nu^2.
\]

In Fig.3, one finds the neutrino flux does get a cutoff or an absorption band through the \( t \)-channel annihilation. Notice that the cut-off is less efficient in a model where there is a significant \( s \)-channel annihilation/scattering process. In fact, the \( s \)-channel process itself does not contribute like a “cutoff” but 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 \( s \)-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.

**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 place with
\[
L > \mathcal{O}(10) \text{ Mpc}
\]
can be affected with the neutrino portal interaction.

As in the left panel of Fig. 4, we have checked that to obtain \( d(E_\nu) \) smaller than the scale of particle horizon size \( \sim 10 \text{ Gpc} \), i.e. when the neutrino propagation in the universe could be affected, \( M \) should be smaller than \( \sim 2 \text{ TeV} \). From the right panel, where \( M \) is at around the lower bound (9), one reads the upper bound of \( M_\psi \simeq m_\phi \) to be around \( \mathcal{O}(100) \text{ MeV} \). This upper bound becomes smaller 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,
\[
M \simeq 0.4 - 2 \text{ TeV} \text{ and } m_\phi \simeq M_\psi \sim 9 - \mathcal{O}(100) \text{ MeV}.
\]
Since the upper bound of \( M \) satisfies (10), the following is predicted: if the high energy neutrino flux in the IceCube is affected by the neutrino portal interaction,
Fig. 3: The predicted neutrino flux with several $m_\nu$. $L = 300$ Mpc, $M_\psi = m_\phi = 9$ 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 observation arranged from [28] while the gray dot-dashed line represents the flux distribution before the annihilation Eq.(24).

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

**Measuring Physical Neutrino Mass Range**

The neutrino flux carries the information of the 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 $\phi$ and $\psi$ are given in someway. This implies a mass scale can be obtained for one of the neutrinos if its mass is greater than $T_\nu$.

Even just given the mass range of $\phi, \psi$, one can predict the neutrino mass range.
Fig. 4: 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.

Since $E_{C\nu B} E_\nu^{cutoff} \sim m_\phi^2 \simeq M_\psi^2$, with the cutoff scale fixed, one obtains

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

which implies $E_{C\nu B}$ ($m_\phi \simeq M_\psi$) 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 neutrino mass lower bound becomes greater. Thus the neutrino mass range is predicted as

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

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 0.05 eV.

If the GZK neutrino source is originated from $\simeq O(1)$Gpc away from the earth, $m_\phi$ and $M_\psi$ 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 the GZK neutrino (See the blue dashed line in Fig.3), the neutrino mass range can be estimated as

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

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

4 A UV model

In the previous sections, we have studied a dimension 5 operator with two additional gauge singlets. It is questioned that whether there is a UV model, and if is, whether constraints for heavy particles in the UV model restrict our scenario especially for $M \lesssim \text{TeV}$.

To suppress $(\phi_H \cdot L)^2$-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

$$\mathcal{L} = -y\phi_H L N - \frac{M}{2f} \phi S \psi - \frac{M_N}{2} NN - \frac{M_\psi}{2} \psi \psi + h.c. - \frac{m_\phi^2}{2} \phi^2 - V(\phi, \phi_H)$$

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 a 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 $M_S S S$ term at the tree-level by imposing $Z_4$ symmetry under which $S, N, \psi$ and the spurion $\tilde{M}$ are charged by $1/4, 1/2, 1/2$ and $1/4$, respectively. Leptons can carry $1/2$ 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 of $S \rightarrow S - \frac{y}{\tilde{M}} \phi_H L$, one finds that the neutrino portal interaction appears as

$$\mathcal{L} \rightarrow \frac{1}{M} \phi_H \cdot L \phi \psi - \tilde{M} S N - \frac{M}{2f} \phi S \psi - \frac{M_N}{2} NN - \frac{M_\psi}{2} \psi \psi + h.c. - \frac{m_\phi^2}{2} \phi^2 - V(\phi, \phi_H),$$

with

$$\frac{1}{M} = -\frac{y}{2f}. \quad (33)$$
The neutrino portal term (1st term) is decoupled from the heavy fields, $S$ and $N$. Moreover, $(\phi_H \cdot L)^2$ does not appear by integrating out the heavy fields up to 1loop level, because $\phi_H \cdot L$ does not directly couple to the heavy field. This fact can be also checked by integrating the heavy fields out in terms of Eq.(31) after diagonalizing the fermion mass matrix. Interestingly, with this UV model, the neutrino mass is purely generated radiatively through the neutrino portal interaction. The Higgs portal term, $|\phi_H|^2 \phi^2$, could be suppressed if $\phi$ is a pseudo-Nambu-Goldstone boson with breaking scale $f$ (See discussion for a concrete non-linear sigma model.)\textsuperscript{10}

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

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

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 $\mathcal{O}(10) \text{ MeV} - \mathcal{O}(10) \text{ GeV}$, where 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 the WIMPs plus a neutrino and this invisible decay can be searched for in the future lepton colliders, CEPC, ILC, FCC-ee and CLIC.

We showed that the neutrino portal interaction can successfully (co)annihilates the lightest WIMP and the WIMP relic abundance can explain the observed one for dark matter. Such a neutrino portal dark matter is tested in the future lepton colliders for the mass $\gtrsim 2 \text{ GeV}$.

\textsuperscript{10}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 $m_\phi$ could be the size of the explicit breaking.
When the WIMP explains a small fraction of the dark matter abundance, the neutrino propagation in the universe can be significantly affected. We pointed out this region can be fully tested in the 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 event or the non-observation of the Glashow resonance in the IceCube. Moreover, a neutrino mass can be “measured” kinematically from the scale of the cutoff and a WIMP mass.

Using a UV model, we have justified our set up and showed that a neutrino mass can be dominantly generated from 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 that with 3 generation cases, e.g. \( \phi_H \cdot L^i Y_{ij} \psi^i \phi \cdot \psi^j \cdot M^{-1} \), where \( i, j \) denotes the generation and \( Y_{ij} \) is the dimensionless coupling in the mass basis of \( \psi^i \) (\( \phi^j \)). The neutrino mass matrix is generated with \( m_{\nu ij} \sim \frac{\sum_k Y_{ik} M_{\psi k} Y_{kj}}{16 \pi^2} \left( \frac{\sum_k Y_{ik} K_{ik} M_{\psi k}}{16 \pi^2} \right) \) where \( K_j \) is \( K \) in Eq.(3) but with \( M_{\psi} \) (\( m_\phi \) ) to be replaced by the mass of \( \psi_j \) (\( \phi_j \)). In this extension, several parameter regions previously discussed can be simultaneously realized with the neutrino portal interactions of different generations.

Now let us provide a natural realization of the UV model Eq.(31) for our relevant parameter ranges where \( m_\phi \) and \( M \) are sufficiently small. A light scalar \( \phi \) suggests a naturalness problem. One of the solutions to this problem is to identify \( \phi \) as a pseudo-Nambu-Goldstone boson. Now consider the spontaneously 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 the 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 \( \pi^+ \), is \( Z_2 \) odd and contains \( \phi \) as \( \pi^+ = \frac{\phi + i \tilde{\phi}}{\sqrt{2}} \). This possibility not only explains the smallness of \( m_\phi \), but also allows a rather small decay constant, \( f \), for the composite scalar \( \phi \), like the pion decay constant in QCD.

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

\[ L \sim \frac{\sum_l Y_{ik} M_{\psi k} Y_{lj}}{16 \pi^2} \left( \frac{\sum_k Y_{ik} K_{ik} M_{\psi k}}{16 \pi^2} \right) \]

\[ m_{\nu ij} \sim \frac{\sum_k Y_{ik} M_{\psi k} Y_{kj}}{16 \pi^2} \left( \frac{\sum_k Y_{ik} K_{ik} M_{\psi k}}{16 \pi^2} \right) \]

\[ K_j \]

\[ M_{\psi} \]

\[ m_\phi \]

\[ f \]

\[ \phi \]

\[ \tilde{\phi} \]

\[ \sqrt{2} \]

\[ Z_2 \]

\[ SU(2) \times U(1) \]

\[ \text{QCD} \]

\[ m_\phi \]

\[ M \]

\[ f \]

\[ \phi \]

\[ \tilde{\phi} \]

\[ \sqrt{2} \]

\[ Z_2 \]

\[ SU(2) \times U(1) \]

\[ \text{QCD} \]

\[ m_\phi \]

\[ M \]

\[ f \]

\[ \phi \]

\[ \tilde{\phi} \]

\[ \sqrt{2} \]

\[ Z_2 \]

\[ SU(2) \times U(1) \]

\[ \text{QCD} \]
pions,
\[ \mathcal{L}_{UV} = \mathcal{L}_{\text{sym}} + \mathcal{L}_{\text{exb}} \]  
(35)
\[ \mathcal{L}_{\text{sym}} = \overline{\mathbf{N}} \sigma^\mu \partial_\mu \mathbf{N} - \langle \Phi \rangle \cdot e^{i \frac{\Theta}{2f}} \cdot \overline{\mathbf{N}} \mathbf{S} + \text{h.c.}, \]  
(36)
\[ \mathcal{L}_{\text{exb}} = -\frac{M_\psi}{2} \psi \psi - \frac{m_\phi^2}{2} \phi^2 - \frac{M_N}{2} N N - \frac{\bar{m}_\phi^2}{2} \phi^2 - \frac{m_0^2}{2} \phi_0^2 - y \phi_H \cdot \mathbf{L} \mathbf{N} + \text{h.c.} \]  
(37)
Here, \( \mathcal{L}_{\text{sym}} \) is SU(2) \( \times \) U(1) symmetric Lagrangian, while terms, which explicitly break SU(2) \( \times \) U(1), are collected in \( \mathcal{L}_{\text{exb}} \); \( \overline{\mathbf{N}} = (N, \psi) \) is a matter doublet with U(1) charge \(-1/2\) and lepton number \(-1\), while the fermion \( S \) carries a lepton number 1; \( \langle \Phi \rangle = (\bar{M}, 0) \) 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 \( M_\psi, m_\phi, \bar{m}, m_0 \) are the explicit breaking terms of the SU(2) \( \times \) U(1) symmetry, and can be smaller than \( \bar{M} \) and \( f \) naturally. In particular, the unbroken U(1) is explicitly broken down to exact Z\(_2\) symmetry by \( M_\psi \) and \( \bar{m}_\phi^2 - m_\phi^2 \). Since \( \psi \) can be also charged under lepton number instead of \( \bar{\psi} \), the 1-loop neutrino mass is suppressed by an additional factor of \( \frac{\bar{m}_\phi^2 - m_\phi^2}{\bar{m}_\phi^2 + m_\phi^2} \) and could reduce the tuning between \( M_\psi \) and \( m_\phi \) to satisfy the neutrino mass constraint.

In this model, with these additional light particles which are assumed to be lighter than the Higgs boson, the testability in the future lepton colliders is even increased. Although the neutrino mass constraint is alleviated and \( M_\psi \) can deviate from \( m_\phi \simeq m_\phi \), for a given values of the mass and the cross section for dark matter-dark matter (neutrino-(anti)neutrino), the increase of \( \text{max} (M_\psi, m_\phi \simeq m_\phi) \) 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 affecting the propagation of the cosmic-ray neutrino.

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