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Scalar dark matter search from the extended ν THDM

Seungwon Baek,^{a,b} Arindam Das^a and Takaaki Nomura^a

^a*School of Physics, KIAS,
Seoul 02455, Korea*

^b*Department of Physics, Korea University,
Seoul 02841, Korea*

E-mail: sbaek@korea.ac.kr, arindam@kias.re.kr, nomura@kias.re.kr

ABSTRACT: We consider a neutrino Two Higgs Doublet Model (ν THDM) in which neutrinos obtain *naturally* small Dirac masses from the soft symmetry breaking of a global $U(1)_X$ symmetry. We extended the model so the soft term is generated by the spontaneous breaking of $U(1)_X$ by a new scalar field. The symmetry breaking pattern can also stabilize a scalar dark matter candidate. After constructing the model, we study the phenomenology of the dark matter: relic density, direct and indirect detection.

KEYWORDS: Beyond Standard Model, Higgs Physics, Neutrino Physics

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Contents

1	Introduction	1
2	The Model	2
3	DM phenomenology	6
4	Conclusion	12

1 Introduction

The existence of the tiny neutrino mass can be naturally explained by the seesaw mechanism [1–7] which extends the Standard Model (SM) through Majorana type Right Handed Neutrinos (RHNs). As a result the SM light neutrinos become Majorana particles. Alternatively there is a simple model, neutrino Two Higgs Doublet Model (ν THDM) [8, 9], which can generate the Dirac mass term for the light neutrinos as well as for the other fermions in the SM. In this model we have two Higgs doublets; one is the same as the SM-like Higgs doublet and the other one is having a small VEV ($\mathcal{O}(1)$) eV to explain the tiny neutrino mass correctly. Due to this fact, the neutrino Dirac Yukawa coupling could be order 1. It has been discussed in [8] that a global softly broken $U(1)_X$ symmetry can forbid the Majorana mass terms of the RHNs; a hidden $U(1)$ gauge symmetry can be also applied to realize ν THDM as in ref. [10]. In this model all the SM fermions obtain Dirac mass terms via Yukawa interactions with the SM-like Higgs doublet (Φ_2) whereas only the neutrinos get Dirac masses through the Yukawa coupling with the other Higgs doublet (Φ_1). Another scenario of the generation of Dirac neutrino mass through a dimension five operator has been studied in [11]. The corresponding Yukawa interactions of the Lagrangian can be written as

$$\mathcal{L}_Y = -\bar{Q}_L Y^u \tilde{\Phi}_2 u_R - \bar{Q}_L Y^d \Phi_2 d_R - \bar{L}_L Y^e \Phi_2 e_R - \bar{L}_L Y^\nu \tilde{\Phi}_1 \nu_R + \text{H.c.} \quad (1.1)$$

where $\tilde{\Phi}_i = i\sigma_2 \Phi_i^*$ ($i = 1, 2$), Q_L is the SM quark doublet, L_L is the SM lepton doublet, e_R is the right handed charged lepton, u_R is the right handed up-quark, d_R is the right handed down-quark and ν_R are the RHNs. The Φ_1 and ν_R are assigned with the global charge 3 under the $U(1)_X$ group. The global symmetry forbids the Majorana mass term between the RHNs. In the original model [8], the global symmetry is softly broken by the mixed mass term between Φ_1 and Φ_2 ($m_{12}^2 \Phi_1^\dagger \Phi_2$) such that a small VEV is obtained by seesaw-like formulas

$$v_1 = \frac{m_{12}^2 v_2}{M_A^2}, \quad (1.2)$$

	Scalar Fields				New Fermion
	Φ_1	Φ_2	S	X	ν_R
$SU(2)_L$	2	2	1	1	1
$U(1)_Y$	$\frac{1}{2}$	$\frac{1}{2}$	0	0	0
$U(1)_X$	3	0	3	1	3

Table 1. Scalar fields and new fermion in our model.

where M_A is the pseudo-scalar mass in [8]. If $M_A \sim 100$ GeV and $m_{12} \sim \mathcal{O}(100)$ keV then v_1 can be obtained as $\mathcal{O}(1)$ eV. In the paper [12], the model is extended to include singlet scalar S which breaks the $U(1)_X$ symmetry. The soft term m_{12}^2 is identified with $\mu\langle S \rangle$ where μ is the Higgs mixing term, $\mu\Phi_1^\dagger\Phi_2 S + h.c.$. It has been studied in [12] that an SM singlet fermion being charged under $U(1)_X$ could be a potential DM candidate.

In this paper we extend the model with a natural scalar Dark Matter (DM) candidate (X). In this model the global $U(1)_X$ symmetry is spontaneously broken down to Z_2 symmetry by VEV of a new singlet scalar S . The remnant of the Z_2 symmetry makes the DM candidate stable. The Z_2 symmetry would be broken by quantum gravity effect and DM would decay via effective interaction [13]. This can be avoided if the $U(1)_X$ is a remnant of local symmetry at a high energy scale and we assume the Z_2 symmetry is not broken. A CP odd component of S becomes the Goldstone boson and hence we study the DM annihilation from this model and compare with the current experimental sensitivity.

The paper is organized as follows. In section 2 we describe the model. In section 3 we discuss the DM phenomenology and finally in section 4 we conclude.

2 The Model

We discuss the extended version of the model in [8] with a scalar field (X). We write the scalar and the RHN sectors of the particle content in table 1. The gauge singlet Yukawa interaction between the lepton doublet (L_L), the doublet scalars (Φ_1, Φ_2) and the RHNs (ν_R) can be written as

$$\mathcal{L} \supset -Y_{ij}^e \bar{L}_{Li} \Phi_2 e_{Rj} - Y_{ij}^\nu \bar{L}_{Li} \tilde{\Phi}_1 \nu_{Rj} + \text{H.c.} \quad (2.1)$$

We assume that the Yukawa coupling constants Y_{ij}^e and Y_{ij}^ν are real. The scalar potential can be written by

$$\begin{aligned} V(\Phi_1, \Phi_2, S) = & -m_{11}^2 \Phi_1^\dagger \Phi_1 - m_{22}^2 \Phi_2^\dagger \Phi_2 - m_S^2 S^\dagger S + M_X^2 X^\dagger X - (\mu \Phi_1^\dagger \Phi_2 S + h.c.) \\ & + \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) + \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) \\ & + \lambda_S (S^\dagger S)^2 + \lambda_{1S} \Phi_1^\dagger \Phi_1 S^\dagger S + \lambda_{2S} \Phi_2^\dagger \Phi_2 S^\dagger S + \lambda_X (X^\dagger X)^2 + \lambda_{1X} \Phi_1^\dagger \Phi_1 X^\dagger X \\ & + \lambda_{2X} \Phi_2^\dagger \Phi_2 X^\dagger X + \lambda_{SX} S^\dagger S X^\dagger X - (\lambda_{3X} S^\dagger X X X + \text{H.c.}). \end{aligned} \quad (2.2)$$

The Dirac mass terms of the neutrinos are generated by the small VEV of Φ_1 . According to [8, 9] we assume that the VEV of Φ_1 is much smaller than the electroweak scale. The

vacuum stability analysis of a general scalar potential has been studied in [14]. Additionally, a remaining Z_3 symmetry is also involved when $U(1)_X$ is broken by non-zero VEV of S . Here X is the only Z_3 charged stable (scalar) particle and as a result X could be considered as a potential Dark Matter (DM) candidate. The mass term M_X of X in eq. (2.2) is positive definite which forbids X to get VEV and as a result the Z_3 symmetry promotes the stability of X as a DM candidate. It has already been discussed in [12] that a CP-odd component in S becomes massless Goldstone boson. Then we write scalar fields as follows

$$\Phi'_1 = \begin{pmatrix} \phi_1^+ \\ \frac{1}{\sqrt{2}}(v_1 + h_1 + ia_1) \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \frac{1}{\sqrt{2}}(v_2 + h_2 + ia_2) \end{pmatrix}, \quad (2.3)$$

$$X = X' e^{i \frac{a_S}{2v_S}}, \quad \Phi_1 = \Phi'_1 e^{i \frac{a_S}{v_S}}, \quad S = \frac{1}{\sqrt{2}} r_S e^{i \frac{a_S}{v_S}}, \quad (2.4)$$

where $r_S = \rho + v_S$. We assume X does not develop a VEV while the VEVs of Φ_1 , Φ_2 and S are obtained by requiring the stationary conditions $\partial V(v_1, v_2, v_S)/\partial v_i = 0$ following

$$\begin{aligned} -2m_{11}^2 v_1 + 2\lambda_1 v_1^3 + v_1(\lambda_{1S} v_S^2 + \lambda_3 v_2^2 + \lambda_4 v_2^2) - \sqrt{2} \mu v_2 v_S &= 0, \\ -2m_{22}^2 v_2 + 2\lambda_2 v_2^3 + v_2(\lambda_{2S} v_S^2 + \lambda_3 v_1^2 + \lambda_4 v_1^2) - \sqrt{2} \mu v_1 v_S &= 0, \\ -2m_S^2 v_S + 2\lambda_S v_S^3 + v_S(\lambda_{1S} v_1^2 + \lambda_{2S} v_2^2) - \sqrt{2} \mu v_1 v_2 &= 0. \end{aligned} \quad (2.5)$$

We then find that these conditions can be satisfied with $v_1 \simeq \mu \ll \{v_2, v_S\}$ and SM Higgs VEV is given as $v \simeq v_2 \simeq 246 \text{ GeV}$. From the first one of the eq. (2.5) we find that v_1 is proportional to and of the same order with μ such that

$$v_1 \simeq \frac{\sqrt{2} \mu v_2 v_S}{\lambda_{1S} v_S^2 + (\lambda_3 + \lambda_4) v_2^2 - 2m_{11}^2}. \quad (2.6)$$

The small order of $v_1 (\sim \mu)$ is required to keep v_2 and v_S in the electroweak scale. Considering the neutrino mass scale as $m_\nu \sim 0.1 \text{ eV}$, the value of μ/v_2 should be small such as $\mu/v_2 \sim \mathcal{O}(10^{-12})$ ensuring Y^ν as $\mathcal{O}(1)$ such that $m_e/v_2 \sim \mathcal{O}(10^{-6})$. Hence v_1 is considered to be smaller than the other VEVs. It also interesting to notice that $\mu = 0$ restores the symmetry of the Lagrangian hence a technically natural small value of μ is acceptable [15, 16]. It is also interesting to notice that $\mu = 0$ enhances the symmetry of the Lagrangian in the sense that we can assign arbitrary $U(1)_X$ charge to Φ_1 , which ensures the radiative generation of the μ -term is proportional to μ itself. Hence a small value of μ is technically natural [15, 16]. Now we identify mass spectra in the scalar sector.

Charged scalar: in this case we calculate the mass matrix in the basis $(\phi_1^\pm, \phi_2^\pm)$ where $\phi_1^\pm$ is approximately physical charged scalar while $\phi_2^\pm$ is approximately NG boson absorbed by $W^\pm$ boson. In the following we write physical charged scalar field as $H^\pm \simeq \phi_1^\pm$. The charged scalar mass matrix can be written as

$$M_{H^\pm}^2 = \begin{pmatrix} \frac{v_2(\sqrt{2}\mu v_S - \lambda_4 v_1 v_2)}{2v_1} & -\frac{1}{2}(\sqrt{2}\mu v_S - \lambda_4 v_1 v_2) \\ -\frac{1}{2}(\sqrt{2}\mu v_S - \lambda_4 v_1 v_2) & \frac{v_1(\sqrt{2}\mu v_S - \lambda_4 v_1 v_2)}{2v_2} \end{pmatrix} \simeq \begin{pmatrix} \frac{v_2(\sqrt{2}\mu v_S - \lambda_4 v_1 v_2)}{2v_1} & 0 \\ 0 & 0 \end{pmatrix}. \quad (2.7)$$

The charged Higgs mass can be written as

$$m_{H^\pm}^2 \simeq \frac{v_2(\sqrt{2}\mu v_S - \lambda_4 v_1 v_2)}{2v_1}. \quad (2.8)$$

CP-even neutral scalar: in the case of CP-even scalar all three components are physical. Hence the mass matrix can be written in the basis of (h_1, h_2, ρ) as

$$M_H^2 = \begin{pmatrix} 2\lambda_1 v_1^2 + \frac{\mu v_2 v_S}{\sqrt{2}v_1} & (\lambda_3 + \lambda_4)v_1 v_2 - \frac{\mu v_S}{\sqrt{2}} & \lambda_{1S} v_1 v_S - \frac{\mu v_2}{\sqrt{2}} \\ (\lambda_3 + \lambda_4)v_1 v_2 - \frac{\mu v_S}{\sqrt{2}} & 2\lambda_2 v_2^2 + \frac{\mu v_1 v_S}{\sqrt{2}v_2} & \lambda_{2S} v_2 v_S - \frac{\mu v_1}{\sqrt{2}} \\ \lambda_{1S} v_1 v_S - \frac{\mu v_2}{\sqrt{2}} & \lambda_{2S} v_2 v_S - \frac{\mu v_1}{\sqrt{2}} & 2\lambda_S v_S^2 + \frac{\mu v_1 v_2}{\sqrt{2}v_S} \end{pmatrix} \\ \simeq \begin{pmatrix} \frac{\mu v_2 v_S}{\sqrt{2}v_1} & 0 & 0 \\ 0 & 2\lambda_2 v_2^2 & \lambda_{2S} v_2 v_S \\ 0 & \lambda_{2S} v_2 v_S & 2\lambda_S v_S^2 \end{pmatrix}. \quad (2.9)$$

We find that all the masses of the mass eigenstates, $H_i (i = 1, 2, 3)$, are at the electroweak scale and the mixings between h_1 and other components are negligibly small while the h_2 and ρ can have sizable mixing. The mass eigenvalues and the mixing angle for h_2 and ρ system can be given by

$$m_{H_2, H_3}^2 = \frac{1}{2} \left[m_{22}^2 + m_{33}^2 \mp \sqrt{(m_{22}^2 - m_{33}^2)^2 + 4m_{23}^4} \right], \quad (2.10)$$

$$\tan 2\theta = \frac{-2m_{23}^2}{m_{22}^2 - m_{33}^2}, \quad (2.11)$$

$$m_{22}^2 = 2\lambda_2 v_2^2, \quad m_{33}^2 = 2\lambda_S v_S^2, \quad m_{23}^2 = \lambda_{2S} v_2 v_S. \quad (2.12)$$

Hence the mass eigenstates are obtained as

$$\begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix} \simeq \begin{pmatrix} 1 & 0 & 0 \\ 0 & \cos \theta & -\sin \theta \\ 0 & \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} h_1 \\ h_2 \\ \rho \end{pmatrix}. \quad (2.13)$$

Here H_2 is the SM-like Higgs, h , and $m_{H_2} \simeq m_h$ where the mixing angle θ between H_2 and H_3 is constrained as $\sin \theta \leq 0.2$ by the LHC Higgs data [17–19] using the numerical analyses on the Higgs decay followed by [20, 21].

CP-odd neutral scalar: calculating the mass matrix of the pseudo-scalars in a basis (a_1, a_2, a_S) we get the mass matrix as

$$M_A^2 = \frac{\mu}{\sqrt{2}} \begin{pmatrix} \frac{v_2 v_S}{v_1} & -v_S & -v_2 \\ -v_S & \frac{v_1 v_S}{v_2} & v_1 \\ -v_2 & v_1 & \frac{v_1 v_2}{v_S} \end{pmatrix} \simeq \begin{pmatrix} \frac{\mu v_2 v_S}{\sqrt{2}v_1} & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad (2.14)$$

using $S \simeq \frac{v_S + \rho + i a_S}{\sqrt{2}}$. In the last step we used the approximation, $v_1 (\sim \mu) \ll v_2, v_S$. We find three mass eigenstates,

$$A = a_1 - \frac{v_1}{v_2} a_2 - \frac{v_1}{v_S} a_S, \\ G^0 = \frac{v_1}{v_2} a_1 + a_2, \\ a = \frac{v_1}{v_S} a_1 - \frac{v_1^2}{v_2 v_S} a_2 + \left(1 + \frac{v_1^2}{v_2^2} \right) a_S, \quad (2.15)$$

up to normalization. They correspond to massive pseudo-scalar, the massless Nambu-Goldstone (NG) mode which is absorbed by the Z boson, and a massless physical Goldstone boson associated with the $U(1)_X$ breaking, respectively. Hence the mass of A is given by

$$m_A^2 = \frac{\mu(v_1^2 v_2^2 + v_1^2 v_S^2 + v_2^2 v_S^2)}{\sqrt{2} v_1 v_2 v_S} \simeq \frac{\mu v_2 v_S}{\sqrt{2} v_1}, \quad (2.16)$$

which is at the electroweak scale. It can be shown [12] that the Goldstone boson, a , is safe from the phenomenological constraints such as $Z \rightarrow H_i a (i = 1, 2, 3)$ decay, stellar cooling from the interaction $a \bar{e} \gamma_5 e$, etc., because it interacts with the SM particles only via highly-suppressed ($\sim v_1/v_{2,S}$) mixing with the SM Higgs. Note that, in our analysis below, we approximate pseudo-scalars as $A \simeq a_1$, $G^0 \simeq a_2$ and $a \simeq a_S$ since we assume $v_1 \ll v_2, v_S$ in realizing small neutrino mass. Here we also discuss decoupling of the physical Goldstone boson from thermal bath where we assume it is thermalized via Higgs portal interaction. The interactions $\rho \partial_\mu a_S \partial^\mu a_S / v_S$, $\lambda_{2S} v_S v_2 \rho h_2$ and the SM Yukawa interactions generate the effective interaction among the Goldstone boson a and the SM fermions

$$- \frac{\lambda_{2S} m_f}{2 m_{H_3}^2 m_{H_2}^2} \partial_\mu a \partial^\mu a \bar{f} f, \quad (2.17)$$

where m_f is the mass of the SM fermion f , and we used $a_S \simeq a$. The temperature, T_a , at which a decouples from thermal bath is roughly estimated by [22]

$$\frac{\text{collision rate}}{\text{expansion rate}} \simeq \frac{\lambda_{2S}^2 m_f^2 T_a^5 m_{PL}}{m_{H_2}^4 m_{H_3}^4} \sim 1, \quad (2.18)$$

where m_{PL} denotes the Planck mass and m_f should be smaller than T_a so that f is in thermal bath. The decoupling temperature is then calculated by

$$T_a \sim 2 \text{ GeV} \left(\frac{m_{H_3}}{100 \text{ GeV}} \right)^{\frac{4}{5}} \left(\frac{\text{GeV}}{m_f} \right)^{\frac{2}{5}} \left(\frac{0.01}{\lambda_{2S}} \right)^{\frac{2}{5}}. \quad (2.19)$$

Thus Goldstone boson a can decouple from thermal bath sufficiently earlier than muon decoupling and does not contribute to the effective number of active neutrinos¹ [23]. Note that the Goldstone boson should be in thermal bath at temperature below that of freeze-out of DM when we consider the relic density of DM, X , is explained by the process, $X \bar{X} \rightarrow aa$, in our analysis below. Taking minimum DM mass as $\sim 100 \text{ GeV}$ freeze-out temperature T_f is larger than $\sim 100/x_f \text{ GeV} \sim 4 \text{ GeV}$ where $x_f = m_{\text{DM}}/T_f \sim 25$. Therefore we can get $T_f > T_a$ even with small $\lambda_{2S} (= 0.01)$ as long as m_{H_3} is not much heavier than the electroweak scale.

As the phenomenology of the Higgs sector has been discussed in [8, 12, 24, 25], we concentrate on the DM phenomenology in the following analysis.

¹If $m_{H_3} \approx 500 \text{ MeV}$ and $\lambda_{2S} \approx 0.005$, then a can make sizable contribution: $\Delta N_{\text{eff}} = 4/7$ [22].

Figure 1. Diagrams in (I), (II), (III) and (IV) correspond to DM annihilation process in scenario-I, II, III and IV.

3 DM phenomenology

In this section, we discuss DM physics of our model such as relic density, direct and indirect detections which are compared with experimental constraints. Since the Higgs portal interaction is strongly constrained by DM direct detection [26–29], we consider the case of small mixing so that $h_1 \simeq H_1$, $h_2 \simeq H_2$ and $\rho \simeq H_3$; here H_2 is the SM-like Higgs in our DM analysis.

Dark matter interaction. Firstly masses of dark matter candidates X is given by [27]

$$m_X^2 = M_X^2 + \frac{\lambda_{1X}}{2}v_1^2 + \frac{\lambda_{2X}}{2}v_2^2 + \frac{\lambda_{SX}}{2}v_S^2 \quad (3.1)$$

where the real and imaginary part of X has the same mass and X is taken as a complex scalar field; this is due to remnant Z_3 symmetry. The interactions relevant to DM physics are given by

$$\begin{aligned} \mathcal{L} \supset & \frac{1}{v_S} \partial_\mu a (X \partial^\mu X^* - X^* \partial^\mu X) + \frac{1}{4v_S^2} \partial_\mu a \partial^\mu a X^* X \\ & + \frac{\lambda_{1X}}{2} \left(H^+ H^- + \frac{1}{2} H_1^2 + A^2 \right) X^* X + \frac{\lambda_{2X}}{4} (2v_2 H_2 + H_2^2) X^* X \\ & + \frac{\lambda_{SX}}{4} (2v_S H_3 + H_3^2) X^* X + \frac{\lambda_{3X}}{2} (v_S + H_3) (X X X + c.c.) \\ & - \mu_{SS} H_3^3 + \frac{1}{v_S} H_3 \partial_\mu a \partial^\mu a - \mu_{1S} H_3 \left(H^+ H^- + \frac{1}{2} (H_1^2 + A^2) \right) - \frac{\mu_{2S}}{2} H_3 H_2^2, \end{aligned} \quad (3.2)$$

where we ignored terms proportional to v_1 since the value of VEV is tiny, $\mu_{SS} \equiv m_{H_3}^2/(2v_S)$, $\mu_{1S} \equiv \lambda_{1S}v_S$, $\mu_{2S} \equiv \lambda_{2S}v_S$, and omitted scalar mixing $\sin \theta (\cos \theta)$ assuming $\cos \theta \simeq 1$ and $\sin \theta \ll 1$. Thus relevant free parameters to describe DM physics are summarized as;

$$\{m_X, m_{H_1}, m_{H_3}, m_A, m_{H^\pm}, v_S, \lambda_{1X}, \lambda_{2X}, \lambda_{SX}, \lambda_{3X}, \mu_{1S}, \mu_{2S}\}, \quad (3.3)$$

Figure 2. Scatter plot for parameters on m_X - v_S plane under the DM relic abundance bound in Scenario-I.

where we choose $\mu_{1S,2S}$ as free parameter instead of $\lambda_{1S,2S}$ and we use $\mu_{SS} = m_{H_3}^2/(2v_S)$. In our analysis, we focus on several specific scenarios for DM physics by making assumptions for model parameters to illustrate some particular processes of DM annihilations. These scenarios are given as follows:

- Scenario-I: $100 \text{ GeV} < v_S < 2000 \text{ GeV}$, $\{\lambda_{1X}, \lambda_{2X}, \lambda_{SX}, \lambda_{3X}, \mu_{1S}/v\} \ll 1$.
- Scenario-II: $v_S \gg v$, $\{\lambda_{SX}, \mu_{1S}/v\} \gg \{\lambda_{1X}, \lambda_{2X}, \lambda_{3X}, \mu_{1S}/v\}$.
- Scenario-III: $v_S \gg v$, $\lambda_{1X} \gg \{\lambda_{2X}, \lambda_{SX}, \lambda_{3X}, \mu_{1S}/v\}$.
- Scenario-IV: $v_S \gg v$, $\lambda_{X3} \gg \{\lambda_{1X}, \lambda_{2X}, \lambda_{SX}, \mu_{1S}/v\}$.

Here we set $v \equiv v_2 \simeq 246 \text{ GeV}$ since $v_1 \ll v_2$. In scenario-I DM mainly annihilates into $a_S a_S$ and $a_S H_3$ final state as shown in figure 1-(I). In scenario-II DM annihilates via H_3 portal interaction as figure 1-(II). In scenario-III DM annihilates into components of Φ_1 through contact interaction with coupling λ_{1X} as shown figure 1-(III). Finally scenario-IV represents semi-annihilation processes $XX \rightarrow XH_3$ as shown in figure 1-(IV). In our analysis, we assumed $\lambda_{2S} \ll \mathcal{O}(1)$ so that we can neglect the case of DM annihilation via the SM Higgs portal interaction since it is well known and constraints from direct detection experiments are strong.

Relic density. Here we estimate the thermal relic density of DM for each scenario given above. The relic density is calculated numerically with `micrOMEGAs 4.3.5` [30] to solve the Boltzmann equation by implementing relevant interactions. In numerical calculations we apply randomly produced parameter sets in the following parameter ranges. For all

Figure 3. Scatter plot for parameters on m_X - λ_{SX} and μ_{1S} - λ_{SX} planes in left and right panels under the DM relic abundance bound in Scenario-II.

Figure 4. Left: scatter plot for parameters on m_X - λ_{1X} plane under the DM relic abundance bound in Scenario-III. Right: that for parameters on m_X - λ_{12X} in Scenario-IV.

scenarios we apply parameter settings as

$$\begin{aligned}
 m_X &\in [50, 500] \text{ GeV}, \quad \mu_{2S} = 1 \text{ GeV}, \quad M_{H_1} = M_A = M_{H^\pm} \in [100, 1000] \text{ GeV}, \\
 \lambda_{2X} &\ll 1,
 \end{aligned} \tag{3.4}$$

where the setting for λ_{2X} is to suppress the SM Higgs portal interactions and small value of μ_{2S} is to suppress scalar mixing. Then we set parameter region for each scenarios as follows:

$$\begin{aligned}
 \text{Scenario-I: } v_S &\in [100, 2000] \text{ GeV}, \quad \lambda_{SX, 1X, 3X} \in [10^{-8}, 10^{-4}], \\
 \mu_{1S} &\in [0.001, 0.1] \text{ GeV}, \quad M_{H_3} \in [10, 30] \text{ GeV},
 \end{aligned} \tag{3.5}$$

$$\begin{aligned}
 \text{Scenario-II: } & v_S \in [3000, 10000] \text{ GeV}, & \lambda_{SX} \in [10^{-3}, 1], & \lambda_{1X,3X} \in [10^{-8}, 10^{-4}], \\
 & \mu_{1S} \in [100, 1000] \text{ GeV}, & M_{H_3} \in [150, 2000] \text{ GeV}, & \\
 \text{Scenario-III: } & v_S \in [3000, 10000] \text{ GeV}, & \lambda_{1X} \in [10^{-3}, 1], & \lambda_{SX,3X} \in [10^{-8}, 10^{-4}], \\
 & \mu_{1S} \in [0.001, 0.1] \text{ GeV}, & M_{H_3} \in [150, 2000] \text{ GeV}, & \\
 \text{Scenario-IV: } & v_S \in [3000, 10000] \text{ GeV}, & \lambda_{3X} \in [10^{-3}, 1], & \lambda_{SX,1X} \in [10^{-8}, 10^{-4}], \\
 & \mu_{1S} \in [0.001, 0.1] \text{ GeV}, & M_{H_3} \in [50, m_X] \text{ GeV}. &
 \end{aligned}
 \tag{3.6}$$

$$\tag{3.7}$$

$$\tag{3.8}$$

Then we search for the parameter sets which can accommodate with observed relic density. Here we apply an approximated region [31]

$$0.11 \lesssim \Omega h^2 \lesssim 0.13. \tag{3.9}$$

In figure 2, we show parameter points on m_X - v_S plane which can explain the observed relic density of DM in Scenario-I. In this scenario, relic density is mostly determined by the cross section of $XX \rightarrow a_S a_S$ process which depends on m_X/v_S via second term of the Lagrangian in eq. (3.2). Thus preferred value of v_S becomes larger when DM mass increases as seen in figure 2. In left and right panel of figure 3, we respectively show parameter points on m_X - λ_{SX} and μ_{1S} - λ_{SX} planes satisfying correct relic density in Scenario-II. In this scenario, the region $m_X \lesssim 100$ GeV requires relatively larger λ_{SX} coupling since scalar boson modes $\{H_3 H_3, H_1 H_1, AA, H^\pm H^\mp\}$ are forbidden by our assumption for scalar boson masses. On the other hand the region $m_X > 100$ GeV allow wider range of λ_{SX} around $0.01 \lesssim \lambda_{SX} \lesssim 1.0$ since DM can annihilate into other scalar bosons if kinematically allowed. In left (right) panel of figure 4, we show parameter region on m_X - $\lambda_{1X}(\lambda_{3X})$ satisfying the relic density in Scenario-III(IV). In scenario-III, DM mass should be larger than ~ 100 GeV to annihilate into scalar bosons from Φ_1 and required value of the coupling is $0.2 \lesssim \lambda_{1X} \lesssim 1.0$ for $m_X \leq 500$ GeV. In scenario-IV, the required value of the coupling λ_{3X} has similar behavior as λ_{1X} in the scenario-III for $m_X > 100$ GeV but slightly larger value. This is due to the fact that semi-annihilation process require larger cross section than that of annihilation process.

Direct detection. Here we briefly discuss constraints from direct detection experiments estimating DM-nucleon scattering cross section in our model. Then we focus on our scenario-III since DM can have sizable interaction with nucleon via H_2 and H_3 exchange and investigate upper limit of mixing $\sin \theta$. The relevant interaction Lagrangian with mixing effect is given by

$$\mathcal{L} \supset \frac{\lambda_{SX} v_S}{2} X^* X (c_\theta H_3 - s_\theta H_2) + \sum_q \frac{m_q}{v} \bar{q} q (s_\theta H_3 + c_\theta H_2), \tag{3.10}$$

where q denote the SM quarks with mass m_q , and we assumed $\mu_X \ll \lambda_{SX} v_S$ as in the relic density calculation. We thus obtain the following effective Lagrangian for DM-quark interaction by integrating out H_2 and H_3 ;

$$\mathcal{L}_{\text{eff}} = \sum_q \frac{\lambda_{SX} v_S m_q s_\theta c_\theta}{2v} \left(\frac{1}{m_h^2} - \frac{1}{m_{H_3}^2} \right) X^* X \bar{q} q, \tag{3.11}$$

where $m_{H_2} \simeq m_h = 125 \text{ GeV}$ is used. The effective interaction can be rewritten in terms of nucleon N instead of quarks such that

$$\mathcal{L}_{\text{eff}} = \frac{f_N \lambda_{SX} v_S m_N s_\theta c_\theta}{v} \left(\frac{1}{m_h^2} - \frac{1}{m_{H_3}^2} \right) X^* X \bar{N} N, \quad (3.12)$$

where m_N is nucleon mass and f_N is the effective coupling constant given by

$$f_N = \sum_q f_q^N = \sum_q \frac{m_q}{m_N} \langle N | \bar{q} q | N \rangle. \quad (3.13)$$

The heavy quark contribution is replaced by the gluon contributions such that

$$\sum_{q=c,b,t} f_q^N = \frac{1}{m_N} \sum_{q=c,b,t} \langle N | \left(-\frac{\alpha_s}{12\pi} G_{\mu\nu}^a G^{a\mu\nu} \right) | N \rangle, \quad (3.14)$$

which is obtained by calculating the triangle diagram for heavy quarks inside a loop. Then we write the trace of the stress energy tensor as follows by considering the scale anomaly;

$$\theta_\mu^\mu = m_N \bar{N} N = \sum_q m_q \bar{q} q - \frac{7\alpha_s}{8\pi} G_{\mu\nu}^a G^{a\mu\nu}. \quad (3.15)$$

Combining eqs. (3.14) and (3.15), we get

$$\sum_{q=c,b,t} f_q^N = \frac{2}{9} \left(1 - \sum_{q=u,d,s} f_q^N \right), \quad (3.16)$$

which leads

$$f_N = \frac{2}{9} + \frac{7}{9} \sum_{q=u,d,s} f_q^N. \quad (3.17)$$

Finally we obtain the spin independent X - N scattering cross section as follows;

$$\sigma_{\text{SI}}(XN \rightarrow XN) = \frac{1}{8\pi} \frac{\mu_{NX}^2 f_N^2 m_N^2 \lambda_{SX}^2 v_S^2 s_\theta^2 c_\theta^2}{v^2 m_X^2} \left(\frac{1}{m_h^2} - \frac{1}{m_{H_3}^2} \right)^2, \quad (3.18)$$

where $\mu_{NX} = m_N m_X / (m_N + m_X)$ is the reduced mass of nucleon and DM. Here we consider DM-neutron scattering cross section for simplicity where that of DM-proton case gives almost similar result. In this case, we adopt the effective coupling $f_n \simeq 0.287$ (with $f_u^n = 0.0110$, $f_d^n = 0.0273$, $f_s^n = 0.0447$) in estimating the cross section. In figure 5, we show DM-nucleon scattering cross section as a function of $\sin \theta$ we take $m_X = 300 \text{ GeV}$, $m_{H_3} = 300 \text{ GeV}$, $v_S = 5000 \text{ GeV}$, and $\lambda_{SX} = 0.5(0.01)$ for red(blue) line as reference values. We find that some parameter region is constrained by direct detection when λ_{SX} is relatively large and $\sin \theta > 0.01$. More parameter region will be tested in future direct detection experiments.

The Higgs portal interaction can be also tested by collider experiments. The interaction can be tested via searches for invisible decay of the SM Higgs for $2m_X < m_h$

Figure 5. DM-Nucleon scattering cross section as a function of $\sin \theta$ in Scenario-II where we take $m_X = 300$ GeV, $m_{H_3} = 300$ GeV, $v_S = 5000$ GeV and $\lambda_{SX} = 0.5(0.01)$ for red(blue) line as reference values. The current bounds from XENON1T [32] and PandaX-II [33].

while collider constraint is less significant compared with direct detection constraints for $2m_X > m_h$ [34–36]. Furthermore DM can be produced via heavier Higgs boson H_3 if $2m_X < m_{H_3}$ and the possible signature will be mono-jet with missing transverse momentum as $pp \rightarrow H_3 j \rightarrow XXj$. However the production cross section will be small when the mixing effect $\sin \theta$ is small as we assumed in our analysis. Such a process would be tested in future LHC with sufficiently large integrated luminosity while detailed analysis is beyond the scope of this paper.

Indirect detection. Here we discuss possibility of indirect detection in our model by estimating thermally averaged cross section in current Universe with `micrOMEGAs 4.3.5` using allowed parameter sets from relic density calculations. Since $a_S a_S$ final state is dominant in scenario-I, we focus on the other scenarios in the following.

Figure 6 shows DM annihilation cross section in current Universe as a function of m_X where left and right panels correspond to Scenario-II and Scenario-III/IV. In Scenario-II, the cross section is mostly $\sim O(10^{-26})\text{cm}^3/\text{s}$ while some points give smaller(larger) values corresponding to the region with $2m_X \gtrless M_{H_3}$ as a consequence of resonant effect. The annihilation processes in the scenario provide the SM final state via decay of H_3 and $\{H_1, H^\pm, A\}$ where H_3 decay gives mainly $b\bar{b}$ via mixing with the SM Higgs and the scalar bosons from second doublet gives leptons. This cross section would be tested via γ -ray observation like Fermi-LAT [37] as well as high energy neutrino search such as IceCube [38, 39], especially when the cross section is enhanced. In Scenario-III, the cross section is mostly $\sim O(10^{-26})\text{cm}^3/\text{s}$ and the final states from DM annihilation include

Figure 6. Left: the current DM annihilation cross section in Scenario-II as a function of m_X . Right: that for Scenario-III and IV represented by red and blue points.

components of Φ_1 that are $\{H_1, H^\pm, A\}$. Thus DM mainly annihilate into neutrinos via the decay these scalar bosons while little amount of charged lepton appear from $H^\pm$. Therefore constraints from indirect detection is weaker in this scenario. In Scenario-IV, the values of cross section is relatively larger due to the nature of semi-annihilation scenario. In this case final states from DM annihilation give mostly $b\bar{b}$ via decays of H_3 in the final state. Then it would be tested by γ -ray search and neutrino observation as in the scenario-II.

4 Conclusion

We consider a neutrino Two Higgs Doublet Model (ν THDM) in which small Dirac neutrino masses are explained by small VEV, $v_1 \sim \mathcal{O}(1)$ eV, of Higgs H_1 associated with neutrino Yukawa interaction. A global $U(1)_X$ symmetry is introduced to forbid seesaw mechanism. The smallness of v_1 proportional to soft $U(1)_X$ -breaking parameter m_{12}^2 is technically natural.

We extend the model to introduce a scalar dark matter candidate X and scalar S breaking $U(1)_X$ symmetry down to discrete Z_2 symmetry. Both are charged under $U(1)_X$. The lighter state of X is stable since it is the lightest particle with Z_2 odd parity. The soft parameter m_{12}^2 is replaced by $\mu\langle S\rangle$. The physical Goldstone boson whose dominant component is pseudoscalar part of S is shown to be phenomenologically viable due to small ratio ($\sim \mathcal{O}(10^{-9})$) of v_1 compared to electroweak scale VEVs of the SM Higgs and S .

We study four scenarios depending on dark matter annihilation channels in the early Universe to simplify the analysis of dark matter phenomenology. In Scenario I, Goldstone modes are important. Scenario II is H_3 portal. In Scenario III, the dark matter makes use of the portal interaction with Φ_1 which generates Dirac neutrino masses. In Scenario IV the dominant interaction is $\lambda_{3X}S^\dagger X X X + \text{h.c.}$ which induces semi-annihilation process of our dark matter candidate. In Scenario II, the dark matter scattering cross section with

neutrons can be sizable and detected at next generation direct detection experiments. We calculated indirect detection cross section in Scenarios II, III, and IV, which can be tested by observing cosmic γ -ray and/or neutrinos.

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