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The Higgs seesaw induced neutrino masses and dark matter

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ABSTRACT

In this paper we propose a possible explanation of the active neutrino Majorana masses with the TeV scale new physics which also provide a dark matter candidate. We extend the Standard Model (SM) with a local $U(1)'$ symmetry and introduce a seesaw relation for the vacuum expectation values (VEVs) of the exotic scalar singlets, which break the $U(1)'$ spontaneously. The larger VEV is responsible for generating the Dirac mass term of the heavy neutrinos, while the smaller for the Majorana mass term. As a result active neutrino masses are generated via the modified inverse seesaw mechanism. The lightest of the new fermion singlets, which are introduced to cancel the $U(1)'$ anomalies, can be a stable particle with ultra flavor symmetry and thus a plausible dark matter candidate. We explore the parameter space with constraints from the dark matter relic abundance and dark matter direct detection.

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1. Introduction

With the discovery of the Higgs-like scalar at the Large Hadron Collider (LHC), the Higgs mechanism of the SM for spontaneous breaking of the $SU(2)_L \times U(1)_Y$ gauge symmetry appears to be a correct description of the nature. In addition to explaining the spontaneous breaking of the electroweak symmetry, the Higgs boson is also responsible for the origin of fermion masses, via the Yukawa interactions. The minimal Higgs mechanism, however, is not able to address the fermion mass hierarchy problem, where the quark–lepton masses range from the top quark with mass of order electroweak scale, $M_t = 172$ GeV, down to electron of mass, $M_e = 0.511$ MeV, and the first order phase transition which is relevant for baryon asymmetry of the Universe. More precise measurement of Higgs boson properties will help determine whether there are new degrees of freedom that participate in electroweak symmetry breaking or otherwise involve in new Higgs boson interactions.

Furthermore, the discovery of the neutrino oscillation has confirmed the theoretical expectation that neutrinos are massive and lepton flavors are mixed [1], which provided the first piece of evidence for physics beyond the Standard Model (SM). In order to accommodate the tiny neutrino masses, one can extend the SM by introducing several right-handed neutrinos, which are taken to be singlets under the $SU(2)_L \times U(1)_Y$ gauge group. In this case, the gauge invariance allows right-handed neutrinos to have Majorana mass M_R, which is not subject to the electroweak symmetry breaking scale. Thus the effective mass matrix of three light Majorana neutrinos can be highly suppressed if M_R is much larger than the electroweak scale, which is the so-called canonical seesaw mechanism [2]. Two other types of tree-level seesaw mechanisms have also been proposed [3,4]. Despite its simplicity and elegance, the canonical seesaw mechanism is impossible to be tested in current collider experiments, especially at the LHC, due to its inaccessibility to the high right-handed Majorana mass scale. Heavy Majorana neutrinos can also give large radiative corrections to the SM Higgs mass, which leads to the seesaw hierarchy problem [5]. An alternative way to generate tiny Majorana neutrino masses at the TeV scale is the inverse seesaw mechanism [6,7], in which the neutrino mass m_ν is proportional to a small effective Majorana mass term μ. But there is no dynamical explanation of the smallness of μ. The argument is that neutrinos become massless in the limit of vanishing μ and the global lepton number, $U(1)_L$, is then restored, leading to a larger symmetry [8]. This argument, however, only works when we give the left-handed singlets (S_L) the same quantum(lepton) number as that of the right-handed heavy neutrinos (N_R). If the lepton number of S_L is zero, the argument above does not hold up any more. Besides, the lepton number is only an accidental symmetry of the SM and is explicitly broken by anomalies.

Since neutrino is the only neutral matter field in the SM, it is reasonable to conjecture that neutrinos are correlated with

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the dark matter, which provides another evidence of the new physics beyond the SM from the precise cosmological observations, through certain dark symmetry. The nature of the dark matter and the way it interacts with ordinary matter are still mysteries. The discovery of the Higgs boson opens up new ways of probing the world of the dark matter. The neutrino flux from the annihilation of the dark matter at the center of the dark matter halo also provides a way of indirect detecting the dark matter.

In this paper, we propose a possible explanation of the smallness of neutrino masses and a possible candidate of the dark matter. The discovery of the Higgs-like boson makes the Higgs mechanism more promising as a possible way to understand the origin of the fermion masses. We study the possibility of generating a small Majorana mass term with the help of the seesaw mechanism in the Higgs sector. We extend the SM with a local $U(1)'$ gauge symmetry, which is spontaneously broken by the vacuum expectation value (VEV) $\langle\varphi\rangle$ of an extra scalar singlet. Furthermore there is a seesaw mechanism in the scalar singlet sector: a second scalar singlet gets a tiny VEV $\langle\Phi\rangle$ in a way similar to that of the Higgs triplet in the type-II seesaw model [3]. $\langle\varphi\rangle$ is responsible for the origin of the dark matter mass and the Dirac neutrino mass term, while $\langle\Phi\rangle$ is responsible for the origin of a small Majorana neutrino mass term. The active neutrino mass matrix arises from the modified inverse seesaw mechanism. Compared with various existed inverse seesaw models [33–42], our studies are new in the following three aspects

- All the mass terms (Dirac and Majorana mass terms) originate from the spontaneous breaking of local gauge symmetries in our model.
- The smallness of the Majorana mass term (μ term in the traditional inverse seesaw model) is naturally explained by the so-called Higgs seesaw mechanism.
- The dark matter phenomenology is closely correlated with the neutrino physics via the $U(1)'$ gauge symmetry in our model.

We study constraints on the parameter space of this model from astrophysical observation and dark matter direct detections.

The paper is organized as follows. In Section 2 we describe our model, including the full Lagrangian, Higgs VEVs and mass spectrum. In Section 3 we study the neutrino masses and the effective lepton mixing matrix of the model. Section 4 is devoted to the study of the dark matter phenomenology. We summarize in Section 5.

2. The model

We extend the SM with three generations of right-handed neutrinos N_R and singlets S_L as in the inverse seesaw mechanism, together with two extra scalar singlets, φ and Φ, as well as a spontaneously broken $U(1)'$ gauge symmetry and a global $U(1)_D$ flavor symmetry. The quantum numbers of the fields are given in Table 1, where ℓ_L is left-handed lepton doublet, E_R is the right-handed charged lepton, H is the SM Higgs doublet, and χ_L and χ_R are the fermion singlet pair carrying the same $U(1)_D$ quantum number. Three generations of gauge singlets $\chi_{L,R}$ are needed to cancel anomalies [9–15] of the $U(1)'$ gauge symmetry. The lightest generation of $\chi_{L,R}$ is stable due to the global $U(1)_D$ flavor symmetry and thus plays the role of dark matter [30–32].

The Higgs potential of the model can be written as

$$
\begin{aligned}
V = & -m^2 H^\dagger H - m_1^2 \varphi^\dagger\varphi + m_2^2 \Phi^\dagger\Phi + \lambda(H^\dagger H)^2 \\
& + \lambda_1(\varphi^\dagger\varphi)^2 + \lambda_2(\Phi^\dagger\Phi)^2 \\
& + \lambda_3(H^\dagger H)(\varphi^\dagger\varphi) + \lambda_4(H^\dagger H)(\Phi^\dagger\Phi) + \lambda_5(\varphi^\dagger\varphi)(\Phi^\dagger\Phi) \\
& + \sqrt{2}\lambda_6\left(\Lambda\varphi^2\Phi^* + \text{h.c.}\right), \qquad (1)
\end{aligned}
$$

Table 1
Quantum numbers of the relevant fields under the local $U(1)'$ and the global $U(1)_D$ flavor symmetry.

	ℓ_L	E_R	N_R	S_L	χ_R	χ_L	H	φ	Φ
$U(1)'$	0	0	0	1	1	0	0	1	2
$U(1)_D$	0	0	0	0	1	1	0	0	0

where we define $H = (h^+, (h_0 + iA + v)/\sqrt{2})^T$, $\varphi = (\varphi_0 + i\delta + v_1)/\sqrt{2}$ and $\Phi = (\Phi_0 + i\rho + v_2)/\sqrt{2}$. After imposing the conditions of the global minimum, one has

$$
-m^2 v + \lambda v^3 + \frac{1}{2}v(\lambda_3 v_1^2 + \lambda_4 v_2^2) = 0\,, \qquad (2)
$$

$$
-m_1^2 v_1 + \lambda_1 v_1^3 + \frac{1}{2}v_1(\lambda_3 v^2 + \lambda_5 v_2^2) + 2\lambda_6\Lambda v_2 = 0\,, \qquad (3)
$$

$$
+m_2^2 v_2 + \lambda_2 v_2^3 + \frac{1}{2}v_2(\lambda_4 v^2 + \lambda_5 v_1^2) + \lambda_6\Lambda v_1^2 = 0\,. \qquad (4)
$$

Then the VEVs can be solved in terms of the parameters

$$
\begin{aligned}
& v^2 \approx \frac{2m_1^2\lambda_3 - 4m^2\lambda_1}{\lambda_3^2 - 4\lambda_1\lambda}\,, \qquad v_1^2 \approx \frac{2m^2\lambda_3 - 4m_1^2\lambda}{\lambda_3^2 - 4\lambda\lambda_1}\,, \\
& v_2 \approx -\frac{2\lambda_6\Lambda v_1^2}{2m_2^2 + \lambda_4 v^2 + \lambda_5 v_1^2}\,, \qquad (5)
\end{aligned}
$$

where v_2 is proportional to Λ and suppressed by m_2^2. Thus v_2 can be a small value given a large m_2^2 or small Λ.

In the basis $(h_0,\ \phi_0,\ \Phi_0)$, the mass matrix of the CP-even Higgs can be written as

$$
M^2_{\text{CP-even}} = \begin{pmatrix} 2v^2\lambda & vv_1\lambda_3 & vv_2\lambda_4 \\ vv_1\lambda_3 & 2\lambda_1 v_1^2 & 2\Lambda v_1\lambda_6 \\ vv_2\lambda_4 & 2\Lambda v_1\lambda_6 & 2v_2^2\lambda_2 - \lambda_6\Lambda v_1^2 v_2^{-1} \end{pmatrix}. \qquad (6)
$$

The mass eigenstates of the CP-even Higgs are then denoted as h_i including the SM-like Higgs h and two exotic Higgs, h_1 and h_2. There is no mixing between the SM CP-odd Higgs A, which is the Goldstone boson eaten by the Z gauge boson, and those of the Higgs singlets, i.e. δ and ρ. The mass matrix of the CP-odd Higgs singlets in the basis of (δ, ρ) is

$$
M^2_{\text{CP-odd}} = \begin{pmatrix} -4\Lambda v_2\lambda_6 & 2\Lambda v_1\lambda_6 \\ 2\Lambda v_1\lambda_6 & -\Lambda\lambda_6 v_1^2 v_2^{-1} \end{pmatrix}. \qquad (7)
$$

The massless eigenstate of the eq. (7) is the Goldstone boson eaten by the Z' and the nonzero mass eigenstate of the CP-odd scalar is then denoted as A', the mass squared of which can be written as $m^2_{A'} = -(4v_2 + v_1^2 v_2^{-1})\Lambda\lambda_6$.

Since the SM particles are not charged under $U(1)'$, there is no experimental constraint on the new symmetry. We assume that there is no kinetic mixing between Z and Z' at the tree level. Thus the mass and coupling of Z' are not constrained by current experiments either. We refer the reader to Ref. [16] for the discussion of Z–Z' mixing at the one-loop level. Notice that the SM Higgs, h, mainly mixes with the CP-even scalar singlet h_1. For the case $m_{h_1} < 1/2m_h$, where m_h is the mass eigenvalue of the SM Higgs and m_{h_1} is the mass eigenvalue of the CP-even scalar singlet, the SM Higgs can decay into h_1, providing enhancement to the Higgs to invisible decay width, which is thus disfavored by the LHC data. For the region $1/2m_h < m_{h_1}$, a global χ^2 fit to the current Higgs data from both ATLAS and CMS shows that the present 95% C.L.

limit on the mixing angle is $\cos\theta \geq 0.84$ [17]. For $m_{h_1} > 2m_h$, one can probe the extended scalar through resonant di-Higgs production [18]. We show in Fig. 2, contours of λ_3 (left-panel) and m_{h_1} (right-panel) in the λ–λ_1 plane by fixing the SM like Higgs mass to be 126 GeV.

3. Neutrino masses

Now we investigate how to realize the neutrino masses in our model. The Yukawa interactions of the lepton sector can be given by

$$-\mathcal{L} = \overline{\ell_L} Y_E H E_R + \overline{\ell_L} Y_\nu \tilde{H} N_R + \overline{S_L} Y_N \varphi N_R + \overline{S_L} Y_S \Phi S_L^C + \overline{\chi_L} Y_\chi \varphi \chi_R + \text{h.c.} \tag{8}$$

where the first and second terms are the charged lepton and neutrino Yukawa interactions separately, the third and fourth terms are the Yukawa coupling of heavy neutrinos to the scalar singlets, and the last term is the Yukawa coupling of the additional fermions. We assume that there is no $\overline{N_R^C} M N_R$ type of mass term, which can be easily forbidden by an extra global $U(1)$ symmetry, in which all the right-handed fermions, H, χ_L and S_L are singly charged, Φ doubly charged and all other particles neutral. The symmetry is explicitly broken by the last term of the Higgs potential in Eq. (1). According to 't Hooft's naturalness criteria [8], which states that a parameter should be much smaller than unity if setting it to zero increases the symmetry of the theory, otherwise the theory is unnatural, the coupling λ_6 should be much smaller than unity, which naturally leads to a small VEV v_2 as required by our model. The same argument can be applied to the Majorana mass term of right-handed neutrinos, which thus should be zero or very much small, here we set it to be zero for simplicity. We can write down the mass matrix of neutrinos in the basis $(\nu_L, N_R^C, S_L)^T$:

$$\mathcal{M} = \begin{pmatrix} 0 & Y_\nu v & 0 \\ Y_\nu^T v & 0 & Y_N v_1 \\ 0 & Y_N^T v_1 & Y_S v_2 \end{pmatrix} \tag{9}$$

where v, v_1, v_2 are given in Eq. (5). Given $v_1 \sim 1$ TeV and $v_2 \sim 1$ MeV, the inverse seesaw mechanism is naturally realized. The matrix $\mathcal{M}$ can be diagonalized by the unitary transformation $\mathcal{U}^\dagger \mathcal{M} \mathcal{U}^* = \hat{\mathcal{M}}$; or explicitly,

$$\begin{pmatrix} A & B & C \\ D & E & F \\ G & H & I \end{pmatrix}^\dagger \begin{pmatrix} 0 & Y_\nu v & 0 \\ Y_\nu^T v & 0 & Y_N v_1 \\ 0 & Y_N^T v_1 & Y_S v_2 \end{pmatrix} \begin{pmatrix} A & B & C \\ D & E & F \\ G & H & I \end{pmatrix}^* = \begin{pmatrix} \hat{M}_\nu & 0 & 0 \\ 0 & \hat{M}_N & 0 \\ 0 & 0 & \hat{M}_S \end{pmatrix}, \tag{10}$$

where $\hat{M}_{\nu,N,S}$ are 3×3 diagonal matrices. The nine mass eigenstates correspond to three observed light neutrinos $\hat{\nu}$ and six heavy Majorana neutrinos $\hat{S}$ and $\hat{N}$, which pair up to form three pseudo-Dirac neutrinos.

Alternatively, the neutrino mass matrix can be block diagonalized and the effective Majorana mass matrix of the active neutrinos can be approximately written as

$$\mathcal{M}_\nu = M_D M_R^{-1} \mu M_R^{T-1} M_D^T = v^2 v_1^{-2} v_2 Y_\nu Y_N^{-1} Y_S Y_N^{T-1} Y_\nu^T . \tag{11}$$

The mass eigenvalues of the three pairs of heavy neutrinos are of the order M_R, and the mixing between $SU(2)_L$ singlets and doublets is suppressed by M_D/M_R. In the basis where the flavor eigenstates of the three charged leptons are identified with their mass eigenstates, the charged-current interactions between neutrinos and charged leptons turn out to be

$$-\mathcal{L}_{\rm CC} = \frac{g}{\sqrt{2}} \overline{\ell_L^\alpha} \gamma_\mu P_L \left(A_{\alpha i} \hat{\nu}_i + B_{\alpha i} \hat{N}_i + C_{\alpha i} \hat{S}_i \right) + \text{h.c.} \tag{12}$$

Obviously A describes the charged-current interactions of light Majorana neutrinos, while B and C are relevant to the charged currents of heavy neutrinos. The neutral current interactions between Majorana neutrinos and neutral gauge boson or Higgs can be also written down in a similar way.

The explicit expression of A can be obtained by integrating out heavy neutrinos and performing the normalization to the light neutrino wave functions. So the effective lepton-mixing matrix can be written as

$$A_{\alpha i} = \left(\delta_{\alpha\beta} - \frac{1}{2} \left| M_D M_R^{-1} \mu (M_R^T)^{-1} \right|^2_{\alpha\beta} - \frac{1}{2} \left| M_D M_R^{-1} \right|^2_{\alpha\beta} \right) U_{\beta i} , \tag{13}$$

where U is the standard PMNS matrix. Obviously the effective neutrino mixing matrix is not unitary. The deviation of A from a unitary matrix is proportional to $|M_D M_R^{-1}|^2$. Constraints on the elements of the leptonic mixing matrix, combining data from neutrino oscillation experiments and weak decays was studied in Ref. [19]. So far neutrino mixing angles have all been measured to a good degree of accuracy, and a preliminary hint for a nontrivial value of δ has also been obtained from a global analysis of current neutrino oscillation data. But the constraint on the non-unitarity of the lepton mixing matrix still need to be improved and the future neutrino factory can measure this effect through the "zero-distance" effect and extra CP violations. The Daya Bay [20] reactor neutrino experiment has measured a nonzero value for the neutrino mixing angle θ_{13} with a significance of 5.2 standard deviations. For this case, even though the neutrino mixing matrix U, which diagonalizes the active neutrino mass matrix, takes the well-known lepton mixing pattens, such as Tri-Bimaximal [21], Bimaximal [22] and Democratic [23] pattens, where θ_{13} is exactly zero, it is still possible to get relatively large θ_{13} from the non-unitarity factors in eq. (13). One can also check the non-unitary effect from the lepton-flavor-violating (LFV) SM Higgs decay, which, interesting and important but beyond the scope of this paper, will be shown in somewhere else. We refer to Ref. [24] for the study of LFV in the inverse seesaw model and Ref. [16] for the LFV effects induced by the TeV-scale neutrinos. It should be emphasized that the low energy phenomenology of neutrinos in our model is the same as that in the conventional inverse seesaw model. To avoid redundancy, in this work we only provide the new dynamics of generating the light Majorana mass term ("μ" mass term in the inverse seesaw model) without repeating the canonical discussion about the neutrino phenomenology.

4. Dark matter

Precise cosmological observations have confirmed the existence of the non-baryonic cold dark matter. The lightest generation of $\chi_{L,R}$, the only odd particles under the global $U(1)$ symmetry, can be a stable dark matter candidate. In order to produce the dark matter relic abundance observed today $\Omega_{DM}h^2 = 0.1187 \pm 0.017$ [25], the thermally averaged annihilation rate $\sigma_A v$ should approximately be 3×10^{-27} cm^3 s$^{-1}/\Omega_{DM}h^2$. Interactions relevant to dark matter phenomenology can be written as

$$\text{I.} \quad \bar{\chi} \gamma^\mu P_R \chi Z'_\mu , \tag{14}$$

$$\text{II.} \quad \overline{\nu_L^C} F^2 \gamma^\mu Z'_\mu \nu_L^C , \tag{15}$$

$$\text{III.} \quad Y_\chi/\sqrt{2} \bar{\chi}_L (\cos\theta h_1 - \sin\theta h) \chi_R , \tag{16}$$

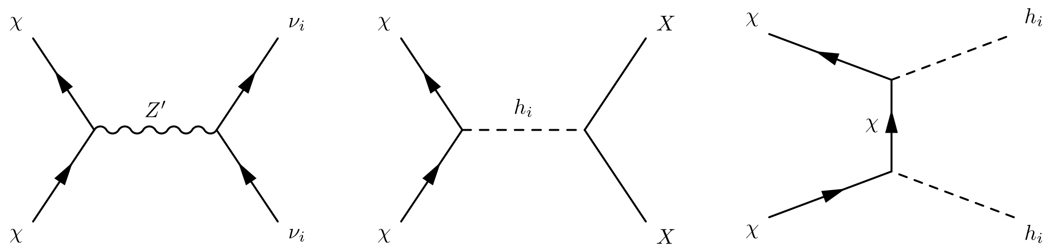

Fig. 1. Feynman diagrams relevant for the annihilations of the dark matter, X denotes the SM particles.

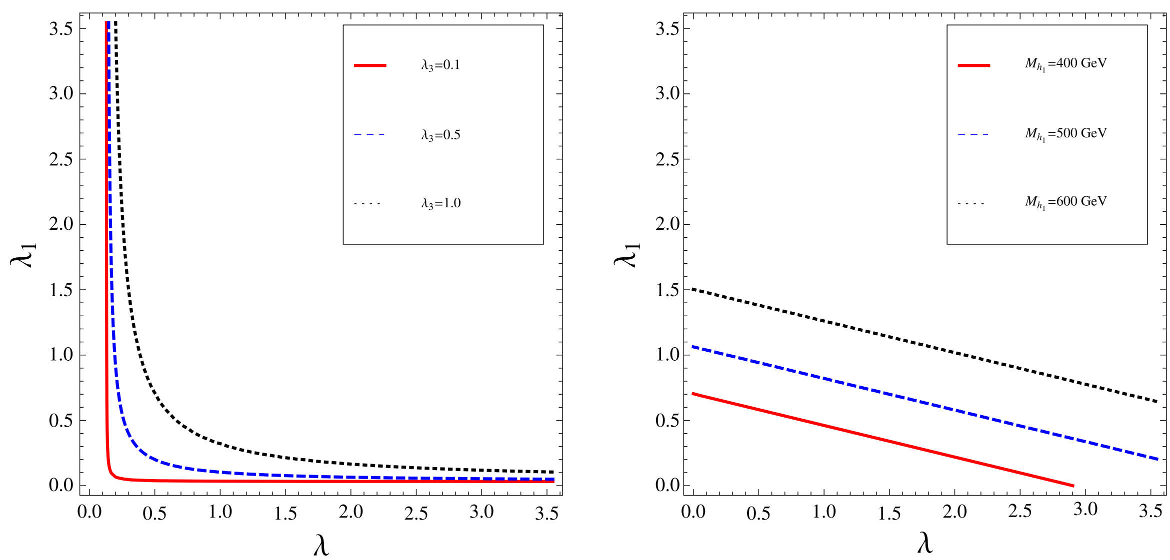

Fig. 2. Contours of the λ_3 (left panel) and the CP-even scalar singlet mass m_{h_1} (right-panel) in the λ–λ_1 plane, with the SM-like Higgs mass fixed at 126 GeV, which helps to reduce one free parameter of the model.

where θ is the mixing angle between the SM Higgs boson and the Higgs singlet. It's the 1–2 mixing angle of matrix given in Eq. (6). F is either D or G, the 21 and 31 entry in $\mathcal{U}$. The expressions of D and G can be written as

$$G \approx (M_R^{-1})^* M_D^\dagger U \,, \tag{17}$$

$$D \approx (M_R^*)^{-1} \mu^\dagger (M_R^\dagger)^{-1} M_D^\dagger \,, \tag{18}$$

from which it's easily seen that the active neutrinos mainly mix with S_L, while the mixing with N_R^C is highly suppressed by the factor μM_R^{-1}. The major contributions to the annihilation cross section come from two types of channels,

$$\chi\bar{\chi} \to Z' \to 2\nu \qquad \chi\bar{\chi} \to h_i \to 2X, \tag{19}$$

where X represents the SM fields including h_i but other than neutrinos. The relevant Feynman diagrams for dark matter annihilation are given in Fig. 1. Obviously the dark matter in our model is the hybrid of neutrino portal and Higgs portal.

To investigate the viability of this model of providing a good dark matter candidate, we fix those parameters irrelevant to the dark matter properties and vary the others. Without loss of generality we also simplify the calculation by taking diagonal Yukawa coupling matrices, which are relevant for the generation of neutrino mixing but irrelevant for the dark matter phenomenology. The typical input parameters are given in the Table 2. The relics density and direct detection cross section are calculated with *micrOMEGAs* [26], which solves the Boltzmann equations numerically and utilizes *CalcHEP* [27] to calculate the relevant cross section. We show in the left panel of Fig. 2 contours of λ_3 in the λ–λ_1 plane by choosing $m_h = 126$ GeV. We also show in the right panel of Fig. 2 contours of the CP-even exotic Higgs mass M_{h_1} in the λ–λ_1 plane, with the SM-like Higgs mass fixed at 126 GeV, which shows that the mass of the exotic CP-even Higgs is in the range of 300–700 GeV.

Table 2
Input parameters at the benchmark point. The parameters in the right part of the table do not change the DM relic density. λ_3 is calculated by imposing the condition $M_{h_0} = 126$ GeV. The choice of parameter space ensures v_1 is of TeV and v_2 is of GeV to generate the right neutrino mass scale.

Parameters	Values or range	Parameters	Values or range
v_1/GeV	500	m_h/GeV	126
Y_ν, Y_N	0.5	Y_S	0.5
λ, λ_1	$(0,\ \sqrt{4\pi}]$	λ_2, λ_4, λ_5	0.5
$M_{Z'}$/GeV	200, 1000	M_χ/GeV	[10, 2000]

As can be seen from Fig. 1, the dark matter in this model is hybrid-portals: vector portal and Higgs portal. For $M_\chi \lesssim M_W$, the dark matter pair annihilate mostly into neutrino pair and bottom quark pair, the amplitudes of which are suppressed by the mixing between the heavy neutrinos and light active neutrinos and the bottom quark–Yukawa coupling, respectively. As a result, the relic density will be too large (except the region near the SM like-Higgs resonance) and overclose the Universe. For $M_W \lesssim M_\chi \lesssim M_{Z'}$, the annihilation channels to dibosons including W, Z and different scalars will be open at various values of M_χ. For $M_\chi \gtrsim M_{Z'}$ all the annihilation channels are open and thus the dark matter relic density will be smaller. Fig. 3 shows dark matter relic density $\Omega_\chi h^2$ as a function of dark matter mass for $M_Z' = 200$ GeV in black and for $M'_Z = 1$ TeV in red. We have set $\lambda = 0.25$, $\lambda_1 \approx \lambda_3 = 0.5$, which results in $m_{h_1} \approx 514$ GeV. The black curve has dips at 53 GeV, 100 GeV and 257 GeV, which come from the SM like-Higgs, Z' and the second CP-even Higgs resonances, respectively. The red curve has dips at 53 GeV, 257 GeV and 500 GeV, which come from the SM like-Higgs, the second CP-even Higgs and the Z' resonances, respectively.

Dark matter is further constrained by direct detection experiments such as LUX [28] and XENON 100 [29]. The dark matter–quark interactions in the effective models naturally induce the dark matter–nucleus interactions. The effective Hamiltonian in our model can be written as

$$H_{\rm eff} = \sum_q c_\theta s_\theta \frac{m_\chi}{v_1} (\bar\chi \chi) \left(\frac{1}{M_h^2} - \frac{1}{M_{h1}^2} \right) \frac{m_q}{v} \bar q q \, , \tag{20}$$

where $c_\theta = \cos\theta$ and $s_\theta = \sin\theta$. Parameterizing the nucleonic matrix element as $\langle N \sum_q m_q \bar q q \rangle = f_N m_N$, where m_N is the proton or neutron mass and f_N are the nucleon form factors. We refer to [30–32] for explicit values of $f^{p,n}$. The cross section for the DM scattering elastically from a nucleus is given by

$$\sigma^{\rm SI} = \frac{\mu^2}{\pi} \left[\frac{c_\theta s_\theta m_\chi}{v v_1} \left(\frac{1}{M_h^2} - \frac{1}{M_{h1}^2} \right) \right]^2 [Z m_p f^p + (A - Z) m_n f^n]^2 \tag{21}$$

where $\mu = m_\chi m_N / (m_\chi + m_N)$ is the reduced mass of the WIMP-nucleon system, with m_N the target nucleus mass. Z and $(A - Z)$ are the numbers of protons and neutrons in the nucleus. One can see from (21) that the scattering cross section is sensitive to λ_3, which determines the mixing angle, θ, between the SM-like Higgs and the Heavier scalar singlet. The direct detection cross section gets bigger when λ_3 increases. For a given VEV of the extra scalar, λ_3 is solely determined in the λ_1–λ plane, as can be seen from the left panel of Fig. 2. In Fig. 4 we show the allowed parameter space of this model for two different Z' masses, $M_{Z'} = 200$ GeV and $M_{Z'} = 1000$ GeV respectively. For every λ and λ_1 in Fig. 4 we find the right dark matter mass which gives the right amount of dark matter relic abundance and we show the contours of the dark matter mass in dotted lines. The shaded region is excluded by the direct detection limit from LUX. One can conclude that for a light Z', a light dark matter is available in large λ and small λ_1 region and a heavy dark matter is available in a small λ region. For a heavy Z', the parameter space in the large λ_1 is available.

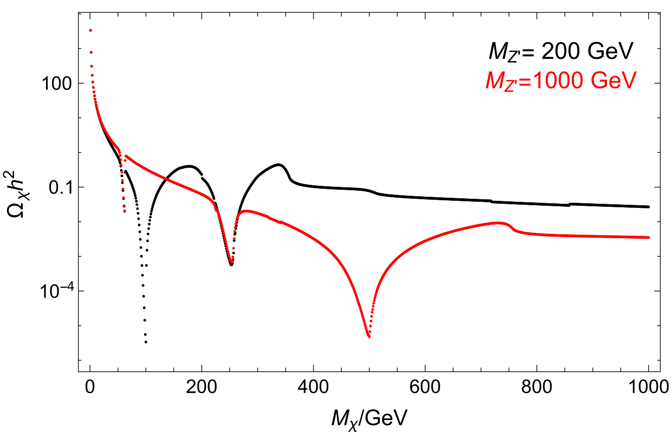

Fig. 3. Dark matter relic density $\Omega_\chi h^2$ as a function of χ mass for $M_{Z'} = 200$ GeV in black and $M_{Z'} = 1000$ GeV in red with $\lambda = 0.25$, $\lambda_1 = \lambda_3 \approx 0.5$ and other values taken according to Table 2. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

5. Concluding remarks

In this paper we extend the SM with a local and a global $U(1)$ symmetry. The smallness of active neutrino Majorana masses is explained by the modified inverse seesaw mechanism. Extra fermion singlets introduced to cancel anomalies of the model can play the role of dark matter. Constraints on the model parameter space from dark matter relic density as well as dark matter direct searches are studied. All the fermion masses arise from the spontaneous breaking of local gauge symmetries, which is a very appealing feature of the model in the era of Higgs physics.

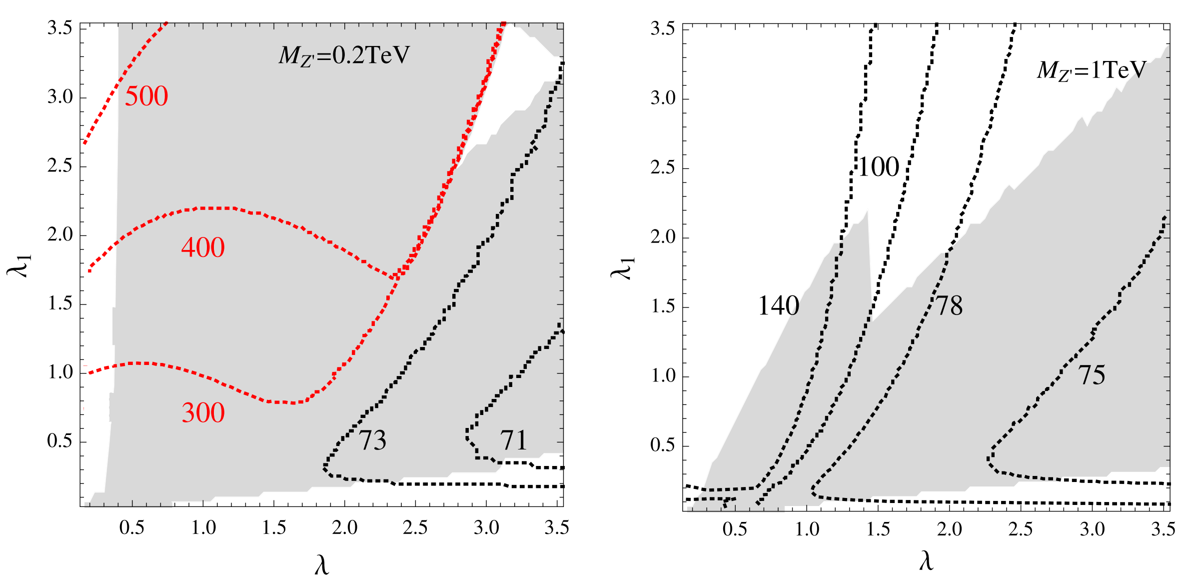

Fig. 4. Parameter space for $M_{Z'} = 200$ GeV and $M_{Z'} = 1000$ GeV respectively. The shaded region is excluded by the direct detection limit from LUX. The dotted lines in both plots are contours of dark matter mass in GeV. The red contours simply means that the major contribution to the dark matter annihilation are from the channels only open when the dark matter mass is above the Higgs mass, contrary to the ones in black. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

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