Radiative Neutrino Mass via Fermion Kinetic Mixing

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We propose that the radiative generation of the neutrino mass can be achieved by incorporating the kinetic mixing of fermion fields which arises radiatively at 1 loop level. As a demonstrative example of the application of the mechanism, we will present the particular case of Standard Model extension by $U(1)_D$ symmetry. As a result, we show how neutrino masses can be generated via kinetic mixing portal instead of mass matrix with residual symmetries responsible for stability of multicomponent dark matter.

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I. INTRODUCTION

Since the late 1990’s there have been many observations of neutrino oscillation data [1–10] requiring neutrinos to have small non-zero masses and mixing among themselves. The first explanations of the natural smallness of neutrino masses came even before that in 1970-1980. In Standard Model (SM) of strong and electroweak interactions neutrinos are predicted to be massless but one unique dimension-5 effective operator [11] can be written which would generate Majorana neutrino masses

$$\mathcal{L}_5 = \frac{f_{ij}}{2\Lambda} (\nu_i \phi^0 \rho_j \rho^+ \nu_j \phi^0 \rho^+ \nu_j + \text{h.c.})$$ (1)

Then neutrino masses would be proportional to $v^2/\Lambda$, where $v = \langle \phi^0 \rangle$ is the SM Higgs doublet vacuum expectation value (VEV). There are only 3 tree level realizations of this operator: seesaw-I [12–15] requiring addition of SM singlet Majorana fermion $N_R$, seesaw-II [16] which extends SM scalar sector with electroweak triplet $\xi \sim (1,3,1)$ under $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ SM gauge group, and seesaw-III [20] which adds a Majorana fermion electroweak triplet $\Sigma \sim (1,3,0)$. There are also 3 one loop realizations of this dimension-5 operator: first one was done in 1980 and is known as Zee model [21], another one became well known in 2006 as Scotogenic model [22] (scotos from Greek meaning darkness), and the third one was realized in 2014 in the context of scotogenic inverse seesaw model [23]. All these Majorana neutrino mass realizations were systematically discussed back in 1998 [24].

On the other hand, neutrino might be Dirac in nature, in which case it would require additional symmetries and different high energy realizations for the smallness of neutrino masses. Three years ago a systematic study of small Dirac neutrino mass realizations [25] at tree and one loop level have been performed. It has been shown that there are only and only four possible cases at tree level and two possible realizations at one loop order. Neutrino mass generation has also been realized in the context of left-right symmetric models [26, 27].

The purpose of this work is to deviate from this canonical approaches to the neutrino mass problem which is achieved by relying on the new idea of fermion kinetic mixing we introduce and to show how it can be realized in the context of the neutrino mass generation with dark matter(DM), i.e. scotogenic scenario.

The paper is organized as follows: in section I

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we introduce and explain the idea of kinetic mixing, following with section II which describes the model and presents radiative neutrino mass generation via kinetic mixing of fermions, section III gives the generation of the fermion and boson masses, DM candidates are discussed in section IV, section V concludes.

II. FERMION KINETIC MIXING

In order to present the idea of kinetic mixing for further generation of neutrino masses, we introduce the minimal set of fields in the context of $U(1)_D$ gauge symmetry extension of the SM. The fermion, call it $A_L$ with $m_A = 0$, and the massive fermion $C_L$, with $m_C \neq 0$ are introduced. The Feynman diagram representing the mixing mechanism is shown in Fig. 1. As can be seen from the diagram, in order to complete the loop, we add the $\Psi_L$ fermion field and a pair of scalars, namely $s_7$ and $s_{11}$. $\phi$ serves the purpose of $U(1)_D$ gauge symmetry breaking in a specific manner, such that there is a residual symmetry left which stabilizes the DM. This residual symmetry is also needed in order to prevent the collapsing of the loop down to tree level, aka to prevent the $s_7$ and $s_{11}$ obtaining the vacuum expectation values (VEV’s). In our case it is achieved by choosing the specific $Q_D$ charge assignments for the fields in the loop. Quantum number assignments for these fields are presented in the table. The result of the diagram in Fig. 1 produces the effective kinetic mixing between two fermion fields which leads to the lagrangian kinetic term

$$\mathcal{L}_{\text{KinMix}} = \imath a A_L \partial C_L + \imath a^* C_L \partial A_L,$$

(2)

In order to bring kinetic part to the canonical form we consider all relevant kinetic terms

$$\mathcal{L} = \imath A_L \partial A_L + \imath C_L \partial C_L \quad \text{with}$$

$$+ \imath a A_L \partial C_L + \imath a^* C_L \partial A_L = \imath \left( A_L \partial C_L \right) \beta \left( 1 - \frac{a}{a^*} \right) \left( \frac{A_L}{C_L} \right).$$

(3)

In order to bring kinetic terms into canonical form, the first step is to rotate by $\pi/4$ so that the kinetic matrix can be diagonalized. When the kinetic terms are diagonal but still not properly normalized, we rescale or renormalize the corresponding fermion fields, so that the kinetic part becomes canonically structured.

After we have worked out the kinetic part, we need to diagonalize back the mass matrix of $A_L$ and $C_L$ fermions in the new basis. This rotation will differ from $\pi/4$ due to presence of rescaling. The final form of the relevant lagrangian is given by

$$\mathcal{L} = \imath \bar{F}_L \beta F_{aL} + \imath \bar{F}_L \beta F_{cL} - m_{F_L} F_{cL} F_{cL},$$

(4)

where $F_{aL}$ and $F_{cL}$ fermions correspond to now properly normalized mass eigenstates. The relation between $A_L$, $C_L$ and the canonically normalized mass eigenstates is given by

$$\left( \frac{A_L}{C_L} \right) = U(\pi/4, \Delta)^\dagger R_{cL} U(\alpha, 0)^\dagger \left( F_{aL}/F_{cL} \right)$$

(5)

and their mass eigenvalues are given by

$$m_{F_a} = 0,$$

(7)

$$m_{F_C} = \frac{m_C e^{12\Delta}}{1 - |\epsilon|^2}.$$  

(8)

Here $U(\Delta)$, unitary $2 \times 2$ transformation, and other relevant parameters given as shown below

$$U(\theta, \Delta) = \begin{pmatrix} c & \epsilon \Delta \\ -\epsilon \Delta & c \end{pmatrix},$$

(9)

$$\text{Im}[\epsilon] = \text{Im}[\epsilon e^{-2\Delta}] = 0 \rightarrow \text{Arg}[\epsilon] = 2\Delta,$$

(10)

$$\tan 2\alpha = -\frac{\sqrt{1 - |\epsilon|^2}}{|\epsilon|},$$

(11)

$$R_{cc} = \text{Diag}(\sqrt{1 - |\epsilon|^2}, \sqrt{1 + |\epsilon|^2}).$$

(12)

And finally $\alpha$ is given by

$$U(\theta, \Delta) = \begin{pmatrix} c & \epsilon \Delta \\ -\epsilon \Delta & c \end{pmatrix},$$

(9)

$$\text{Im}[\epsilon] = \text{Im}[\epsilon e^{-2\Delta}] = 0 \rightarrow \text{Arg}[\epsilon] = 2\Delta,$$

(10)

$$\tan 2\alpha = -\frac{\sqrt{1 - |\epsilon|^2}}{|\epsilon|},$$

(11)

$$R_{cc} = \text{Diag}(\sqrt{1 - |\epsilon|^2}, \sqrt{1 + |\epsilon|^2}).$$

(12)

And finally $a$ is given by
\[ a = \frac{1}{16\pi^2} Y^*_\chi Y_C, \]

where

\[ \gamma = s_R c_R G(y_{12R}, y_{31R}) + s_I c_I G(y_{11I}, y_{21I}), \]

\[ G(y_i, y_j) = \frac{1}{(1 - y_i)^2 (1 - y_j)^2} [y_i (y_i/2 - 1) \ln y_i - y_j (y_j/2 - 1) \ln y_j + y_i y_j/2 (y_i y_j + 4) \ln \left( \frac{y_i}{y_j} \right) + \frac{1}{2} (y_i (y_i - 1) - y_j (y_j - 1) + y_i y_j (y_i - y_j))] , \]

\[ y_i = \frac{m^2}{m*}, \]

FIG. 2: |\[\epsilon|\] vs \(y = \frac{m^2}{m*}\) plot with \(y_{12} = 5\), yukawa’s are taken to be 1 and mixing angles are set to \(\pi/4\).

where Yukawa’s are given with flavour indies suppressed. Here \(m*\) is the Dirac mass of the \(\Psi\) Dirac fermion, \(s\) and \(c\) stand for sinus and cosinus respectively, and correspond to the mixing of \(s_7\) with \(s_{11}\) scalars, which is proportional to \(\mu_3\) term in the lagrangian (see eq. (25)). The scalar mixing angles and mass eigenvalues, \(m^2_{s_i(h/t)}\), are given in section IV equation (13).

The \(|\epsilon|\) vs \(y = \frac{m^2}{m*}\) plot is shown in Fig. 2. Here \(\Delta y_{12} = y_1 - y_2\) indicates the mass splitting between \(s_{7,11}\) scalar eigenvalues, whereas scalar and pseudo-scalar splittings are labeled by \(\Delta y_{11} = y_R - y_I\).

Some self consitensy remarks: in the limit of \(a, \epsilon \to 0\) kinetic mixing is absent and \(\alpha \to -\pi/4, R_{sc} \to 1, F_{sL} \to A_L, F_{cL} \to C_L\) which means no rescaling, no mixing. Furthermore, the component of \(A_L \propto F_{cL}\) vanishes, leading to \(m_{loop} \to 0\) from eq. (26), which in turn will produce \(m_\nu \to 0\).

III. MODEL

In order to present a complete model for the kinetic mixing mechanism presented in section II, we extend the SM \(SU(3)_C \otimes SU(2)_L \otimes U(1)_Y\) symmetry group by \(U(1)_D\) dark sector gauge symmetry. Field content for \(SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_D\) gauge symmetry case is shown in table III.

We give all fermions as left handed fields. The last column in the table shows the number of copies besides the flavor count. \(A_L, C_L\) fermions and \(\eta_L, \eta_D\) scalars are added for the generation of Dirac radiative neutrino mass mechanism (see Fig. 3). \(N_L\) is needed as a Dirac mass partner for \(\nu_L\), aka to produce effective Dirac mass term \(m_{loop}\eta_L N_L\). Next, \(\Psi_L\) fermion and \(s_{7,11}\) scalars are needed for the kinetic mixing mechanism as was explained in section II. \(\phi\) is introduced for the spontaneous symmetry breaking of the \(U(1)_D\) dark gauge symmetry. It carries even \(Q_D\) charge in order to generate a residual \(\mathbb{Z}_2\) symmetry (as will be shown at the end of this section) upon breaking of \(U(1)_D\). \(\Psi_R\) is added to produce a Dirac mass for \(\Psi_L\). Lastly, \(N_R\), multiple copies of \(F_{2L}\) and \(C_L\) fermions are added to render the model chiral anomaly free. Even though, in general all five copies of \(C_L\) can kinetically mix with one copy of \(A_L\), without loss of generality and for the sake of simplicity we take the bases of \(C_L\) fermions in which only one particuliar \(C_L\) mixes with \(A_L\).
TABLE I: Particle content for $\mathcal{G}_{SM} \otimes U(1)_D$ case.

| Field | SU(2)$_L$ | SU(2)$_R$ | SU(3)$_c$ | SU(2)$_L$ | U(1)$_Y$ | U(1)$_D$ | Flavors | copies |
|-------|------------|------------|-----------|------------|-----------|-----------|---------|-------|
| $Q \sim (u, d)^T$ | $(\frac{1}{2}, 0)$ | $3$ | $2$ | $\frac{1}{6}$ | $0$ | $3$ | $1$ |
| $u^c$ | $(\frac{1}{2}, 0)$ | $3$ | $1$ | $-\frac{1}{2}$ | $0$ | $3$ | $1$ |
| $d^c$ | $(\frac{1}{2}, 0)$ | $3$ | $1$ | $\frac{1}{2}$ | $0$ | $3$ | $1$ |
| $L \sim (\nu, e)^T$ | $(\frac{1}{2}, 0)$ | $1$ | $2$ | $-\frac{1}{2}$ | $0$ | $3$ | $1$ |
| $e^c$ | $(\frac{1}{2}, 0)$ | $1$ | $1$ | $1$ | $0$ | $3$ | $1$ |
| $H \sim (H^+, H_0)^T$ | $(0, 0)$ | $1$ | $2$ | $\frac{1}{2}$ | $0$ | $1$ | $1$ |
| $A_L$ | $(\frac{1}{2}, 0)$ | $1$ | $1$ | $0$ | $3$ | $3$ | $1$ |
| $C_L$ | $(\frac{1}{2}, 0)$ | $1$ | $1$ | $0$ | $1$ | $3$ | $5$ |
| $N_L$ | $(\frac{2}{3}, 0)$ | $1$ | $1$ | $0$ | $-4$ | $3$ | $1$ |
| $N_R$ | $(\frac{2}{3}, 0)$ | $1$ | $1$ | $0$ | $4$ | $3$ | $1$ |
| $F_{2L}$ | $(\frac{2}{3}, 0)$ | $1$ | $1$ | $0$ | $-2$ | $3$ | $4$ |
| $\Psi_L$ | $(\frac{2}{3}, 0)$ | $1$ | $1$ | $0$ | $\frac{3}{2}$ | $3$ | $1$ |
| $\Psi_R$ | $(\frac{2}{3}, 0)$ | $1$ | $1$ | $0$ | $-\frac{3}{2}$ | $3$ | $1$ |
| $\eta \sim (\eta^0, \eta^{-1})^T$ | $(0, 0)$ | $1$ | $2$ | $-\frac{1}{2}$ | $3$ | $1$ | $1$ |
| $\eta_D$ | $(0, 0)$ | $1$ | $1$ | $0$ | $-1$ | $1$ | $1$ |
| $\phi$ | $(0, 0)$ | $1$ | $1$ | $0$ | $2$ | $1$ | $1$ |
| $\sigma_7$ | $(0, 0)$ | $1$ | $1$ | $0$ | $\frac{7}{2}$ | $1$ | $1$ |
| $s_{11}$ | $(0, 0)$ | $1$ | $1$ | $0$ | $-\frac{7}{2}$ | $1$ | $1$ |

Considering triangular anomalies, $SU(3)_C \times SU(2)_L \times U(1)_Y$ anomalies are canceled in the same way as in canonical SM case. Since there are no fermions that transform non-trivialy under SM and dark sector simultaniously, any cross anomalies between SM and $U(1)_D$ are trivialy absent. The only anomalies to consider for cancelation are $U(1)_{DGrav}$ and $[U(1)_D]^3$. For this purpose, multiple copies of $C_L$ and $F_{2L}, N_R$, and $\Psi_R$ fields are added. Considering $U(1)_D$ sector anomalies, they cancel in the following way

\[ \sum_i Q_{Di} = 1 \times (3) + 5 \times (1) + 4 \times (-2) + 1 \times (-4) + 1 \times (4) + 1 \times \left( \frac{5}{2} \right) + 1 \times \left( -\frac{5}{2} \right) = 0 \]  \hspace{1cm} (18)

\[ \sum_i Q_{Di}^2 = 1 \times (3)^2 + 5 \times (1)^2 + 4 \times (-2)^2 + 1 \times (-4)^2 + 1 \times (4)^2 + 1 \times \left( \frac{5}{2} \right)^2 + 1 \times \left( -\frac{5}{2} \right)^2 = 0 \]  \hspace{1cm} (19)

The interaction terms of the lagrangian for the $\mathcal{G}_{SM} \otimes U(1)_D$ gauge symmetry case are given by
\[
-\mathcal{L}_{Yuk}^{SM} = \bar{u}_a Y^{ab}_L H_j e^\dagger_j + \bar{d}_R Y^{ab}_L H^\dagger_i + \bar{L}_a Y^\dagger_e e_{Rb} H_i + \text{h.c.} \\
-\mathcal{L}_{m_F}^{New} = \bar{\Psi}_L m^{ab}_L \Psi_{Rb} + \bar{N}_L m^{ab}_N N_{Rb} + \text{h.c.} \\
-\mathcal{L}_{Yuk}^{New} = L_a Y^{ab}_L A_{Lb} \eta^*_L + \bar{\Psi}_L Y^{ab}_L A_{Lb} s_{11} + \bar{\Psi}_L Y^{ab}_C A_{Lb} s^*_7 + A_{Lb} Y^{ab}_N N_{Lb} \eta^*_D \\
+ C_{Lba} Y^{a\beta}_C A_{Lb} \phi^* + A_{Lb} Y^{a\beta}_F \bar{F}_{2Lb} & \phi^* + C_{Lba} Y^{a\beta}_F \bar{F}_{2Lb} \eta^*_D \\
+ \bar{F}_{2Lao} Y^{a\beta}_N N_{Rb} \phi + \text{h.c.} \\
V_0 = \sum_{H,\phi,\eta_L,\eta_D, s_{11}, s^*_7} (-1)^{q_x} m^2 |x|^2 + \frac{\lambda_x}{2} |x|^4 \\
+ \sum_{H,\phi,\eta_L,\eta_D, s_{11}, s^*_7} \lambda_{xy} |x|^2 |y|^2 + \lambda_{HL2} (H^\dagger \eta_L) (\eta^*_L H) \\
V_3 = \mu_D \eta_D \phi + \mu_3 \phi s_{11} s^*_7 + \text{h.c.} \\
V_4 = \lambda_{H\eta_D} H_i \eta_{Lj} \eta_D \phi^* e^\dagger_j + \lambda_{s\eta_D} s^*_7 s^*_1 + \text{h.c.}
\]

where \((-1)^{q_x} = -1\) for \(H\) and \(\phi\), and 1 for the other scalars. Here \(a, b, \text{etc.}\) indicate the flavour index and run over 1,2,3, whereas \(\alpha, \beta\) indicate the copy index, which exists only for \(C_L\) and \(F_2L\), and run over 1-5 and 1-4 for \(C_L\) and \(F_2L\), respectively. Here important points to notice are that \(L\) lepton doublet couples only to \(A_L\) new fermion field, \(\mu_3\) trilinear scalar term mixes \(s_7\) with \(s_{11}\) when \(\langle \phi \rangle \neq 0\) which is needed for the kinetic mixing, \(\lambda_{H\phi}\) quartic scalar term which mixes \(\eta^*_L\) with \(\eta_D\) needed to generate Dirac radiative neutrino mass, and that all terms which could give \(A_L\) fermion a mass, even at 1-loop order, are forbidden by symmetry and field content of the model.

Diagram representing radiative neutrino mass generation via kinetic mixing for \(G_{SM} \times U(1)_D\) gauge symmetry is shown in Fig. 3.

\[
\begin{align*}
\langle H \rangle & \quad \langle \phi \rangle \\
\eta^*_L & \quad \eta_D \\
\bar{\Psi}_L & \quad A_L \quad C_L \quad C_L \quad A_L \quad N_L \\
\langle \phi \rangle
\end{align*}
\]

FIG. 3: Radiative neutrino mass generation via fermion kinetic mixing in the \(G_{SM} \otimes U(1)_D\) gauge symmetry case. Crosses between \(A_L\) and \(C_L\) fields correspond to kinetic mixing given by diagram in Fig. 1.
The Dirac mass mixing of $\nu$ and $N_L$, leading to

$$m_{\text{loop}} = \frac{1}{16\pi^2} Y_N m_{F_C} Y_L \frac{\epsilon^2}{1 - \epsilon^2} \left[ s_{\eta R} c_{\eta R} F(x_{1R, x_{2R}}) + s_{\eta I} c_{\eta I} F(x_{2L, x_{1L}}) \right],$$

(26)

The plot showing how $m_{\text{loop}}$ depends on its parameters is depicted in Fig. 4. As can be seen the extra suppression from $\epsilon$ allows for wider range of masses, mass splittings and yukawa couplings. Here $\Delta m_{12} = m_1 - m_2$ indicates the mass splitting between $\eta_L, D$ scalar eigenvalues, whereas scalar and pseudo-scalar splittings are labeled by $\Delta m_{RI} = m_R - m_I$.

After the electroweak and dark symmetry breaking, due to symmetry and field content of the model, we obtain 2 residual dark $Z_2$ symmetries which are not ad hoc. The first $Z_2$ symmetry is analogous to the one from canonical Scotogenic model, but here it is obtained from $U(1)_D$ spontaneous symmetry breaking. The other $Z_2$ symmetry is now and present here due to fractional charge assignments of the particles involved in the kinetic mixing which will be explained in the section IV. This gives us opportunity for multicomponent DM case as will be discussed in section V. Fields odd, $(1)^{Q_D} = \text{odd}$, under the first $Z_2$ symmetry include $\eta_L, \eta_D, A_L, C_{L \alpha}$. The ones that are odd under the second $Z_2$ symmetry are the ones with fractional $Q_D$ charges, $s_7, s_{11}, \Psi_{L,R}$. This is summarized in table II.

### IV. MASS SPECTRUM

**Fermion masses**

Masses of SM fermions are generated identical to SM case, so we focus on mass generations of new particles. We need to consider 3 different sectors that do not mix with each another. First, the $(\nu_L, N_L, N_R, F_{2L})$ sector, the $Z_2^{1,2} \sim (++, +)$ even sector. Next, $Z_2^2 \sim (-)$ odd sector, similar to the one present in the canonical Scotogenic paper $22$, $(A_L, C_{L \alpha})$ fields. Lastly $Z_2^2 \sim (-)$ odd fermions, special for this model, due to the presence of kinetic mixing mechanism, aka $\Psi$ fermion. Starting with $Z_2$ even fermions we have, considering the La-
In the case if \( m_{\text{loop}} \) does not provide enough suppression for the neutrino masses, the ratio of \( m_N/\nu \) can provide extra suppression for the neutrino masses. For example if \( m_{\text{loop}} \sim 10^{-4} \text{ GeV} \) then \( m_N/\nu \sim 10^6 \) would give \( m_\nu \sim 0.1 \text{eV} \). The other 3 \( F_{2L} \) states, orthonormal to the \( F_{2L} \) state coupled to \( N_R \) fermion, obtain their masses radiatively through the diagram shown in Fig. 5.

Next, considering the \( Z_2 \) odd sector, \( A_L \) remains massless till after neutrinos generate their masses, such that \( A_L \) mass is generated through neutrino mass \( (m_{\text{loop}}) \) as shown in Fig. 6. \( m_A \) is then given by

\[
m_A = \frac{1}{16\pi^2} Y_N m_{\text{loop}} Y_L \left[ s_\nu c_\nu F(x_{1R}, x_{2R}) + s_\nu c_\nu F(x_{1L}, x_{1L}) \right],
\]

where \( F(x_i, x_j) \) and \( x \) are given in eqs. (27) and (28). Fig. 7 shows the dependence of \( m_A \) on its parameters.

This is important point because if \( A_L \) obtained its mass in some other way, neutrinos would generate their masses through \( A_L \)'s mass and the kinetic mixing would contribute in the sub-leading order and be unnessesary. One more important point to mention is that, this predicts one dark fermion to be naturally lighter than the neutrino, since its mass is one loop suppressed with respect to the Dirac neutrino mass as can be seen from Fig. 6.

The five copies of \( C_L \) dark fermions obtain their masses through \( \langle \phi \rangle \) at tree level by incorporating...
the diagonalization of the $5 \times 5$ mass matrix in $C_L$ basis.

Lastly, considering the $Z_2 \sim (-)$ odd sector, we have only $\Psi$ fermion which is vector-like and has an invariant mass of $m_\Psi$.

**Boson masses**

The only new vector gauge boson, corresponding to $U(1)_D$ gauge symmetry of dark sector, gets its mass through a canonical higgs mechanism during spontaneous symmetry breaking of $U(1)_D$ gauge symmetry in the dark sector. Mass of dark $U(1)_D$ gauge boson is given by

$$m^2_{A_D} = \sum_i Q_D^i g^2 D_i v^2 = 2 g^2 D v^2,$$
(34)

the corresponding would-be Nambu-Goldstone boson is $\text{Im} [\phi]$. Due to absence of scalars, with non-zero VEV, that simultaneously transform under $G_{SM}$ and dark $U(1)_D$ gauge symmetry, there is no tree level mixing between $A^0_D$ and SM neutral gauge bosons. Mixing will appear at one loop order through $\eta_L^{±,0}$ running in the loop but it is loop suppressed and we will ignore the mixing here. The rest of gauge bosons obtain their masses just like in SM.

Coming to scalar sector. The charged higgs scalar from SM, $H^\pm$, corresponds to would-be Nambu-Goldstone boson and gets eaten up by $W^\pm$. The other electrically charged scalar, $\eta^+_L$, part of $\eta_L$ doublet needed for the neutrino mass generation, does not mix with $H^\pm$ due to presence of $Z_2^D$ under which $\eta_L \sim -$ and $H \sim +$. The charged scalar mass is given by

$$m^2_{\eta_L^+} = m^2_{\eta_L} + \frac{1}{2} \left( [\lambda_{HL} + \lambda_{HL2}] v^2 + \lambda_L \phi v^2_\phi \right).$$
(35)

Besides that, the real components of $H^0$ and $\phi$ mix with each other and after using scalar potential minimization conditions, $\partial V / \partial [H/\phi] = 0$ to eliminate $m^2_{\eta_L}$ and $m^2_\phi$, their corresponding 2x2 mass square matrix is given by

$$m^2_{H_{\phi},\phi} = \begin{pmatrix} \lambda_H v^2 & \lambda_H v v_\phi \\ \lambda_H v v_\phi & \lambda_\phi v^2_\phi \end{pmatrix},$$
(36)

with their corresponding eigenvalues and mixing angle given as

$$m^2_{1,2} = \frac{1}{2} \left[ \lambda_H v^2 + \lambda_\phi v^2_\phi \pm \sqrt{\left( \lambda_H v^2 - \lambda_\phi v^2_\phi \right)^2 + 4 \lambda^2_H \phi v^2_\phi} \right],$$
(37)

$$\tan 2\theta_{H\phi} = \frac{2 \lambda_H v v_\phi}{\lambda_H v^2 - \lambda_\phi v^2_\phi}.$$  
(38)

Mass matrix corresponding to $H^0_i$ and $\phi_i$ is given by zero, meaning that they become the longitudinal degrees of freedom of $Z$ and $A_D$ gauge bosons, as expected.

Real and imaginary components of $\eta^0_L$ and $\eta_D$ mix with each other, respectively. Their mass squared matrix is given by
The eigenvalues and mixing angles corresponding to $s$ is analogous to the kinetic mixing and generated in this case by operator from $[22]$. Here, the important point is the splitting of masses due to the kinetic mixing mechanism and the second one is analogous to the canonical Scotogenic model due to $2 \times Z_2$ sectors that do not mix with one another. This is due to $2 \times Z_2$ symmetries present here. The first one is analogous to the canonical Scotogenic model and plays the same role here, whereas the second $Z_2$ symmetry in this case is unique and is present here due to the kinetic mixing mechanism and the fractional charges of the particles involved. It can be thought of as a dark stabilizing symmetry for the dark sector within dark sector of Scotogenic model. In this way, corresponding neutral scalars can be categorized as $(H^0, \phi) \in \{+, +\}$, $(\eta^0_L, \eta_D) \in \{-, +\}$, and $(s_7, s_11) \in \{-, -\}$ under the $2 \times Z_2$ symmetries.

\[ m^2_{nLD(R,I)} = \left( m^2_{nL} + \frac{1}{2} \left( \lambda_{HHL} v^2 + \lambda_{LH} v^2 \right) \right) \pm \frac{1}{2} \lambda_{HE} v^2 \pm \frac{1}{2} \lambda_{HD} v^2 \pm \sqrt{2} \mu_\phi v^2. \]  

(39)

The corresponding eigenvalues and mixing angles are given by

\[ m^2_{1,2} = \frac{1}{2} \left[ m^2_{nL} + \frac{1}{2} \left( \lambda_{HHL} v^2 + \lambda_{LH} v^2 \right) + m^2_D + \frac{1}{2} \left( \lambda_{HE} v^2 + \lambda_{HD} v^2 \right) \right] \pm \frac{1}{2} \lambda_{HE} v^2 \pm \frac{1}{2} \lambda_{HD} v^2 \pm \sqrt{2} \mu_\phi v^2 \]

\[ \tan^2 \theta_{n[R,I]} = \frac{\pm 2 \frac{1}{2} \lambda_{HE} v^2 \pm 2 \frac{1}{2} \lambda_{HD} v^2 \pm \sqrt{2} \mu_\phi v^2}{m^2_{nL} + \frac{1}{2} \left( \lambda_{HHL} v^2 + \lambda_{LH} v^2 \right) - m^2_D - \frac{1}{2} \left( \lambda_{HE} v^2 + \lambda_{HD} v^2 \right) \pm \sqrt{2} \mu_\phi v^2}. \]  

(40)

The mass splitting, which is needed for the non-zero neutrino mass matrix, of real and imaginary parts is accomplished by $\lambda_{HE} v^2 H^0_L \eta_D \phi^*$ operator, which is analogous to the $(\eta^0 H)^* \phi$ operator from $[22]$. As the last piece, the $s_7$ and $s_{11}$ mix with each other, real and imaginary parts, respectively. Their mass matrix is given by

\[ m^2_{s_{7,11}[R,I]} = \left( m^2_{s_7} + \frac{1}{2} \left( \lambda_{HHL} v^2 + \lambda_{LH} v^2 \right) \right) \pm \frac{1}{2} \lambda_{HE} v^2 \pm \frac{1}{2} \lambda_{HD} v^2 \pm \sqrt{2} \mu_\phi v^2 \]

(42)

The eigenvalues and mixing angles corresponding to $s$ scalars are given by

\[ m^2_{s_{1,2}} = \frac{1}{2} \left[ m^2_{s_7} + \frac{1}{2} \left( \lambda_{HHL} v^2 + \lambda_{LH} v^2 \right) + m^2_{s_{11}} + \frac{1}{2} \left( \lambda_{LH1} v^2 + \lambda_{S11} v^2 \right) \right] \pm \frac{1}{2} \lambda_{HE} v^2 \pm \frac{1}{2} \lambda_{HD} v^2 \pm \sqrt{2} \mu_\phi v^2 \]

\[ \tan^2 \theta_{s[R,I]} = \frac{\pm \sqrt{2} \mu_\phi v^2}{m^2_{s_7} + \frac{1}{2} \left( \lambda_{HHL} v^2 + \lambda_{LH} v^2 \right) - m^2_{s_7} - \frac{1}{2} \left( \lambda_{LH1} v^2 + \lambda_{S11} v^2 \right) \pm \sqrt{2} \mu_\phi v^2}. \]  

(44)

Here, the important point is the splitting of masses of real and imaginary parts which is needed for the kinetic mixing and generated in this case by operator $\mu_3 \phi s_{11} s_7$.

As can be seen above, there are 3 separate scalar sectors that do not mix with one another. This is due to $2 \times Z_2$ symmetries present here. The first one is analogous to the canonical Scotogenic model and plays the same role here, whereas the second $Z_2$ symmetry in this case is unique and is present here due to the kinetic mixing mechanism and the fractional charges of the particles involved. It can be thought of as a dark stabilizing symmetry for the dark sector within dark sector of Scotogenic model. In this way, corresponding neutral scalars can be categorized as $(H^0, \phi) \in \{+, +\}$, $(\eta^0_L, \eta_D) \in \{-, +\}$, and $(s_7, s_{11}) \in \{-, -\}$ under the $2 \times Z_2$ symmetries.

V. DARK MATTER

The $G_{SM} \otimes U(1)_D$ model discussed in section $[11]$ can accommodate multicomponent DM scenario. The lightest of the particles that transform as $Z_2 \sim$
(-, +) is one component, the stability is provided by the $\mathbb{Z}_2$ symmetry which is exact. Assuming $m_\Psi > m_{a_{11}}$, the second component is the lightest eigenstate of the $s_7, s_{11}$ sector, which transforms as $\mathbb{Z}_2^1 \sim (+, -)$. In this case $\Psi \sim (-, -)$ under $\mathbb{Z}_2^1$ would decay into lighter $s_7, s_{11}$ eigenstates $\sim (+, -)$ and $A_L \sim (-, +)$ through $Y_{A,C}$ yukawa couplings. Assuming that the lightest stable particle (LSP) of $(-, -)$ sector is $s$, the dominant contribution to its relic abundance would come from effective $ss \rightarrow HH$ diagram, where $H$ is the SM Higgs scalar field. The quartic $\lambda$ coupling between $s$ and $H$ would need to be suppressed to avoid elastic scattering between $s$ and nuclei, which is mediated by $sHH$ trilinear coupling. However, the trilinear coupling $ss\phi$ and $\phi HH$ which are proportional to $v_\phi$ are not suppressed and would contribute to $ss \rightarrow HH$, assuming the mixing between $\phi$ and $H$ is small enough to avoid the same elastic scattering off nuclei. The detailed analysis for the relic density is beyond the scope of this work.

VI. CONCLUSIONS

We have presented the mechanism of neutrino mass generation via kinetic mixing in the context of the anomaly free $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y \otimes U(1)_D$ gauge symmetry. In this case Dirac neutrino mass is generated after electroweak and $U(1)_D$ symmetry breaking via the kinetic mixing of 2 fermions in the dark sector. As a consequence the neutrino mass is naturally suppressed by the radiative nature of the generation mechanism(similar to Scotogenic scenario). This model includes 2 dark sectors, first one is the same as in Scotogenic scenario and the second is unique to this mechanism, present here due to the kinetic mixing mechanism, which allows for the multiparticle DM scenario.

Despite presenting the particular example with $G_{SM} \otimes U(1)_D$ gauge symmetry, the kinetic mixing idea is more general and can be realized in cases with other gauge symmetries as well. In principle, the kinetic mixing of fermions does not need to be carried out in the dark sector and we could kinetically mix neutrino with other fermion but this would require to include sterile neutrinos into the model, further more in this case in order to increase the neutral fermion mass matrix rank(give neutrino a mass) one would need to follow the scenario like: neutrino mixes with another neutral fermion leading to mass generation of dark sector particle and then using the same diagram neutrino would get a mass from this same dark sector particle. These and other prospects and phenomenology is among further possible research directions of this work.

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