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Research Article

High Scale Type-II Seesaw, Dominant Double Beta Decay within Cosmological Bound and LFV Decays in SU(5)

M. K. Parida and Rajesh Satpathy

Centre for Excellence in Theoretical and Mathematical Sciences, Siksha 'O' Anusandhan (SOA) Deemed to be University, Khandagiri Square, Bhubaneswar 751030, India

Correspondence should be addressed to M. K. Parida; minaparida@soa.ac.in

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Very recently novel implementation of type-II seesaw mechanism for neutrino mass has been proposed in SU(5) grand unified theory with a number of desirable new physical phenomena beyond the standard model. Introducing heavy right-handed neutrinos and extra fermion singlets, in this work we show how the type-I seesaw cancellation mechanism works in this SU(5) framework. Besides predicting verifiable LFV decays, we further show that the model predicts dominant double beta decay with normal hierarchy or inverted hierarchy of active light neutrino masses in concordance with cosmological bound. In addition a novel right-handed neutrino mass generation mechanism, independent of type-II seesaw predicted mass hierarchy, is suggested in this work.

1. Introduction

Renormalizable standard model (SM) predicts neutrinos to be massless whereas oscillation experiments prove them to be massive [1–5]. All the generational mixings have been found to be much larger than the corresponding quark mixings. Theoretically [6–13] neutrino masses are predicted through various seesaw mechanisms [14–43]. In a minimal left-right symmetric [44–47] grand unified theory (GUT) like $SO(10)$ [48, 49] where parity (P) violation in weak interaction is explained along with fermion masses [50–55], a number of these seesaw mechanisms can be naturally embedded while unifying the three forces of the SM [56–72]. More recently precision gauge coupling unification has been successfully implemented in direct symmetry breaking of $SO(18) \longrightarrow SM$ which may have high potential for new physics [73].

The SO(10) model that predicts the most popular canonical seesaw as well as the type-II seesaw has also the ability to explain baryon asymmetry of the universe via leptogenesis through heavy RH neutrino [74] or Higgs triplet decays [75–77]. But because of underlying quark lepton symmetry [44], the type-I seesaw scale as well as RHν masses are so large that the model predicts negligible lepton flavor violating (LFV) decays like $\mu \longrightarrow e\gamma$, $\tau \longrightarrow \mu\gamma$, $\tau \longrightarrow e\gamma$, and $\mu \longrightarrow e\bar{e}e$. Similarly direct mediation of large mass of scalar triplet required for type-II seesaw gives negligible contribution to lepton number violating (LNV) and lepton flavor violating (LFV) decays. Ever since the proposal of left-right symmetry, extensive investigations continue in search of experimentally observable double beta decay [78–82] in the $W_R - W_R$ channel [83, 84]. Adding new dimension to such lepton number violating (LNV) process, the like-sign dilepton production has been suggested as a possible means of detection of W_R-boson at accelerator energies [85], particularly the LHC [43]. However, no such signals of TeV scale W_R− have been detected so far. Even if W_R mass and seesaw scales are large and inaccessible for direct verification, neutrinoless double beta decay ($0\nu2\beta$) in the $W_L - W_L$ channel [20, 86–91] is predicted close to observable limit with $\tau_{\beta\beta} \geq 10^{25}$ yrs provided light neutrino masses predicted by such high scale seesaw mechanisms are quasidegenerate (QD) with each mass $m_\nu \geq \mathcal{O}(0.2)$ eV [78] and their sum > 0.6 eV. But as noted by the recent Planck data such QD type masses violate the cosmological bound [92].

$$\Sigma_\nu \equiv \Sigma_i \widehat{m}_i \leq 0.23 \ \ (\text{eV}). \tag{1}$$

The fact that such QD type ν masses violate the cosmological bound may be unravelling another basic fundamental reason why detection of double beta decay continues to elude experimental observation for several decades. On the other hand, if neutrinos have smaller NH or IH type masses, there is no hope for detection of these LNV events in near future with RHν extended SM. In other words, predicting observable double beta decay in the $W_L - W_L$ channel with left-handed helicities of both the beta particles has been a formidable problem confronting theoretical and experimental physicists. However, it has been shown that in case of dynamical seesaw mechanism generating Dirac neutrinos the seesaw scale is accessible for direct experimental verification [93].

The path breaking discovery of inverse seesaw [25–31] with one extra singlet fermion per generation not only opened up the neutrino mass generation mechanism for direct experimental tests, but also lifted up lepton flavor violating (LFV) decays [94] from the abysmal depth of experimental inaccessibility of negligible branching ratios ($Br.(l_\alpha \longrightarrow l_\beta\gamma) \sim 10^{-50}$) to the illuminating salvation of profound observability (Br.$\simeq 10^{-8} - 10^{-16}$) [95–100] which has been discussed extensively [99, 101–118]. Despite inverse seesaw, observable double beta decay in the $W_L - W_L$ channel and the non-QD type neutrino masses remained mutually exclusive until both the RH neutrinos and singlet fermions (S_i) were brought into the arena of LFV and LNV conundrum through the much needed extension of the Higgs sector. The King-Kang [119] mechanism cancelled out the ruling supremacy of canonical seesaw which was profoundly exploited in SO(10) models with the introduction of both the SO(10) Higgs representations 16_H and $126^\dagger_H$ [13, 84, 120–125] with successful prediction of observable double beta decay in the $W_L - W_L$ channel [20, 86]. Very interestingly, even though high scale type-II seesaw can govern light neutrino masses of any hierarchy, possibility of observable LFV and double beta decay prediction in the $W_L - W_L$ channel irrespective of light neutrino mass hierarchies has been realized at least theoretically [13, 125].

The purpose of this work is to point out that there are new interesting physics realizations with suitable extension of a non-SUSY SU(5) GUT model proposed recently [126] where type-II seesaw, precision coupling unification, verifiable proton decay, scalar dark matter, and vacuum stability have been already predicted. However with naturally large type-II seesaw scale $> 10^{9.2}$ GeV, observable double beta decay accessible to ongoing experiments [78–82] is possible in this model too with QD type neutrinos only of common mass with $|m_0| \geq 0.2$ eV like many other high scale seesaw models as noted above. In this work we make additional prediction that dominant double beta decay in the $W_L - W_L$ channel can be realized with NH or IH type hierarchy consistent with much lighter neutrino masses $|m_i| << 0.2$ eV. Thus, this realization is consistent with cosmological bound of (1). Although such possibilities were realized earlier in SO(10) with TeV scale W_R or Z' bosons as noted above, in SU(5) without the presence of left-right symmetry and associated gauge bosons, we have shown here for the first time that the dominant double beta decay is mediated by a sterile neutrino (Majorana fermion singlet) of $\mathcal{O}(1)$ GeV mass of first generation. The model further predicts LFV decay branching ratios only 4–5 orders smaller than the current experimental limits. An additional interesting part of the present work is the first suggestion of a new mechanism for heavy RHν mass generation that permits these masses to have hierarchies independent of conventional type-II seesaw prediction. Thus highlights of the present model are as follows:

(i) first implementation of type-I seesaw cancellation mechanism leading to the dominance of type-II seesaw in SU(5);

(ii) prediction of verifiable LFV decays only 4 − 5 orders smaller than the current experimental limits;

(iii) prediction of dominant double beta decay in the $W_L - W_L$ channel close to the current experimental limits for light neutrino masses of NH or IH type in concordance with cosmological bound;

(iv) suggestion of a new right-handed neutrino mass generation mechanism independent of type-II predicted mass hierarchy;

(v) precision gauge coupling unification with verifiable proton decay which is the same as discussed in [126].

This paper is organised in the following manner. In Section 2 we briefly review the SU(5) model along with gauge coupling unification and predictions of the intermediate scales. In Section 3 we discuss how type-I seesaw formula for active neutrino masses cancels out giving rise to dominance of type-II seesaw and prediction of another type-I seesaw formula for sterile neutrino masses. Fit to neutrino oscillation data is discussed in Section 4. In Section 5 we suggest a new mechanism of RHν mass generation. Prediction on LFV decay branching ratios is discussed in Section 6. Lifetime prediction for double beta decay is presented in Section 7. In Section 8 we discuss the results of this work and state our conclusion. Block diagonalization procedure is explained in more detail in Appendix A.

2. A Non-Supersymmetric SU(5) Model

2.1. Extension of SU(5). As noted in [127], inclusion of the scalar $\kappa(3,0,8) \subset 75_H$ with mass $M_\kappa = 10^{9.23}$GeV in the extended non-SUSY SU(5) achieves precision gauge coupling unification. Then it has been shown in [126] that type-II seesaw ansatz for neutrino mass is realized by inserting the entire Higgs multiplet $15_H \subset SU(5)$ containing the LH Higgs triplet $\Delta_L(3,-1,1)$ at the same mass scale $M_{15_H} = 10^{12}$GeV.

$$\begin{gathered}\kappa(3,0,8) \subset 75_H, \quad M_\kappa = 10^{9.23}\text{GeV},\\ \Delta_L(3,-1,1) \subset 15_H, \quad M_{15_H} = 10^{12}\text{GeV},\\ \xi(1,0,1), \quad M_\xi \sim \mathcal{O}(1)\text{ TeV}.\end{gathered} \tag{2}$$

The scalar singlet $\xi(1,0,1)$ has played two crucial interesting roles of stabilising the SM scalar potential as well as serving as WIMP DM candidate. The introduction of 15_H at any scale

$> 10^{9.23}$ GeV in this model maintains precision coupling unification. In the present model we extend the model further by the inclusion of the following fermions and an additional scalar $\chi_S(1,0,1)$:

(i) three right-handed neutrino singlets N_i ($i = 1,2,3$), one for each generation, with masses to be fixed by this model phenomenology;

(ii) three left-handed Majorana fermion singlets S_i ($i = 1,2,3$), one for each generation, similar to those introduced in case of inverse seesaw mechanism [25–31];

(iii) a Higgs scalar singlet $\chi_S(1,0,1)$ to generate $S-N$ mixings through its VEV.

Being singlets under the SM gauge group, they do not affect precision gauge coupling unification of [126].

2.2. Coupling Unification, GUT Scale, and Proton Lifetime. As already discussed [126, 127] using renormalization group equations for gauge couplings and the set of Higgs scalars of (2), precision unification has been achieved with the PDG values of input parameters [128–130] on $\sin^2\theta_W(M_Z)$, $\alpha_S(M_Z)$ resulting in the following mass scales and the GUT fine-structure constant α_G:

$$\begin{aligned} M_U &= 10^{15.23}\,\text{GeV}, \\ M_\kappa &= 10^{9.23}\,\text{GeV}, \\ M_{\Delta_L} &= M_{15_H} = 10^{12}\,\text{GeV}, \\ \frac{1}{\alpha_G} &= 37.765. \end{aligned} \tag{3}$$

Using threshold effects due to superheavy Higgs scalars [131–138], proton lifetime prediction for $p \longrightarrow e^+\pi^0$ turns out to be in the experimentally accessible range [139]:

$$\begin{aligned} \tau_p\left(p \longrightarrow e^+\pi^0\right) = (1.01 \times 10^{34\pm0.44} \\ -\left(5.5 \times 10^{35\pm0.44}\right)\ \text{yrs}. \end{aligned} \tag{4}$$

Extensive investigations with number of SU(5) GUT extensions have been carried out with proton lifetime predictions consistent with experimental limits [140–149]. But implementation of type-II seesaw dominance due to type-I seesaw cancellation resulting in dominant LFV and LNV decays as discussed below is new especially in the context of non-SUSY SU(5).

3. Cancellation of Type-I and Dominance of Type-II Seesaw

Due to introduction of heavy RHνs in the present model which were absent in [126], it may be natural to presume a priori that besides type-II seesaw, type-I seesaw may also contribute substantially to light neutrino masses and mixings. But it has been noted that there is a natural mechanism to cancel out type-I seesaw contribution while maintaining dominance of inverse seesaw [84, 119, 120, 122–124] or type-II seesaw or even linear seesaw [13, 125] as the case may be. Briefly we discuss below how this cancellation mechanism operates in the present extended model resulting in type-II seesaw dominance even in the presence of heavy RHνs.

The SM invariant Yukawa Lagrangian of the model is

$$\begin{aligned} \mathscr{L}_{\text{Yuk}} = Y^{\ell}\overline{\psi}_L\psi_R\phi + f\psi_L^c\psi_L\Delta_L + y_\chi\overline{N}^C S\chi_s \\ + \left(\frac{1}{2}\right)M_N\overline{N}^C N + h.c. \end{aligned} \tag{5}$$

Using the VEVs of the Higgs fields and denoting $M = y_\chi\langle\chi_S\rangle = y_\chi V_\chi$, $M_D = Y\langle\phi\rangle$, a 9×9 neutral fermion mass matrix has been obtained which, upon block diagonalization, yields 3×3 mass matrices for each of the light neutrino (ν_α), the right-handed neutrino (N_α), and the sterile neutrino (S_α) [120, 124, 125]. The block diagonalization of 9×9 neutral fermion mass matrix was presented in useful format in [32] but without cancellation of type-I seesaw. Later on, this diagonalization procedure has been effectively utilized to study the type-I seesaw cancellation mechanism in SO(10) models [13, 120, 124, 125].

In this model the left-handed triplet Δ_L and RH neutrinos M_N being much heavier than the other mass scales with $M_{\Delta_L} \gg M_N \gg M \gg M_D$ are at first integrated out from the Lagrangian leading to

$$\begin{aligned} -\mathscr{L}_{\text{eff}} = \left(m_\nu^{II} + M_D\frac{1}{M_N}M_D^T\right)_{\alpha\beta}\nu_\alpha^T\nu_\beta \\ + \left(M_D\frac{1}{M_N}M^T\right)_{\alpha m}\left(\overline{\nu_\alpha}S_m + \overline{S_m}\nu_\alpha\right) \\ + \left(M\frac{1}{M_N}M^T\right)_{mn}S_m^T S_n, \end{aligned} \tag{6}$$

which, in the (ν, S) basis, gives the 6×6 mass matrix

$$\mathscr{M}_{\text{eff}} = \begin{pmatrix} M_D M_N^{-1} M_D^T + m_\nu^{II} & M_D M_N^{-1} M^T \\ M M_N^{-1} M_D^T & M M_N^{-1} M^T \end{pmatrix}, \tag{7}$$

while the 3×3 heavy RH neutrino mass matrix M_N is the other part of the full 9×9 neutrino mass matrix. This 9×9 mass matrix $\widetilde{\mathscr{M}}_{BD}$ which results from the first step of block diagonalization procedure as discussed above and in the Appendix is

$$\mathscr{W}_1^\dagger\mathscr{M}_\nu\mathscr{W}_1^* = \widetilde{\mathscr{M}}_{\text{BD}} = \begin{pmatrix} \mathscr{M}_{\text{eff}} & 0 \\ 0 & M_N \end{pmatrix}. \tag{8}$$

Defining

$$\begin{aligned} X &= M_D M^{-1}, \\ Y &= M M_N^{-1}, \\ Z &= M_D M_N^{-1}, \end{aligned} \tag{9}$$

the transformation matrix $\mathscr{W}_1$ has been derived as shown in (10) [120, 125]

$$\mathscr{W}_1 = \begin{pmatrix} 1-\frac{1}{2}ZZ^\dagger & -\frac{1}{2}ZY^\dagger & Z \\ -\frac{1}{2}YZ^\dagger & 1-\frac{1}{2}YY^\dagger & Y \\ -Z^\dagger & -Y^\dagger & 1-\frac{1}{2}\left(Z^\dagger Z+Y^\dagger Y\right) \end{pmatrix}. \tag{10}$$

After the second step of block diagonalization, the type-I seesaw contribution cancels out and gives in the (ν, S, N) basis

$$\mathscr{W}_2^\dagger \widetilde{\mathscr{M}}_{\mathrm{BD}} \mathscr{W}_2^* = \mathscr{M}_{\mathrm{BD}} = \begin{pmatrix} m_\nu & 0 & 0 \\ 0 & m_{\mathcal{S}} & 0 \\ 0 & 0 & m_{\mathcal{N}} \end{pmatrix}, \tag{11}$$

where $\mathscr{W}_2$ has been derived in (12) [120, 125]. We have used the bare mass of S_i and VEV of $\chi_L(2,-1/2,1)$ to be vanishing, i.e., $\mu_S = 0$, $<\chi_L>= 0$, to get the form suitable for this model building.

$$\mathscr{W}_2 = \begin{pmatrix} 1-\frac{1}{2}XX^\dagger & X & 0 \\ -X^\dagger & 1-\frac{1}{2}X^\dagger X & 0 \\ 0 & 0 & 1 \end{pmatrix} \tag{12}$$

In (11) the three 3×3 matrices are

$$m_\nu = m_\nu^{II} = f v_L \tag{13}$$

$$m_{\mathcal{S}} = -M M_N^{-1} M^T \tag{14}$$

$$m_{\mathcal{N}} = M_N, \tag{15}$$

the first of which is the well-known type-II seesaw formula and the second is the emergence of the corresponding type-I seesaw formula for the singlet fermion mass. The third of the above equations represents the heavy RHν mass matrix.

In the third step, m_ν, $m_{\mathcal{S}}$, and $m_{\mathcal{N}}$ are further diagonalized by the respective unitary matrices to give their corresponding eigenvalues:

$$\begin{aligned} U_\nu^\dagger m_\nu U_\nu^* &= \widehat{m}_\nu = \mathrm{diag}\left(m_1, m_2, m_3\right), \\ U_S^\dagger m_{\mathcal{S}} U_S^* &= \widehat{m}_S = \mathrm{diag}\left(m_{S_1}, m_{S_2}, m_{S_3}\right), \\ U_N^\dagger m_{\mathcal{N}} U_N^* &= \widehat{m}_N = \mathrm{diag}\left(M_{N_1}, M_{N_2}, M_{N_3}\right). \end{aligned} \tag{16}$$

The complete mixing matrix [32, 120] diagonalizing the above 9×9 neutrino mass matrix occurring in (8) and in (A.1) of the Appendix turns out to be

$$\mathscr{V} \equiv \begin{pmatrix} \mathscr{V}_{\alpha i}^{\nu\widehat{\nu}} & \mathscr{V}_{\alpha j}^{\nu\widehat{S}} & \mathscr{V}_{\alpha k}^{\nu\widehat{N}} \\ \mathscr{V}_{\beta i}^{S\widehat{\nu}} & \mathscr{V}_{\beta j}^{S\widehat{S}} & \mathscr{V}_{\beta k}^{S\widehat{N}} \\ \mathscr{V}_{\gamma i}^{N\widehat{\nu}} & \mathscr{V}_{\gamma j}^{N\widehat{S}} & \mathscr{V}_{\gamma k}^{N\widehat{N}} \end{pmatrix} \tag{17}$$

$$= \begin{pmatrix} \left(1-\frac{1}{2}XX^\dagger\right)U_\nu & \left(X-\frac{1}{2}ZY^\dagger\right)U_S & ZU_N \\ -X^\dagger U_\nu & \left(1-\frac{1}{2}\left\{X^\dagger X+YY^\dagger\right\}\right)U_S & \left(Y-\frac{1}{2}X^\dagger Z\right)U_N \\ 0 & -Y^\dagger U_S & \left(1-\frac{1}{2}Y^\dagger Y\right)U_N \end{pmatrix}, \tag{18}$$

as shown in the Appendix. In (18), $X = M_D M^{-1}$, $Y = MM_N^{-1}$, and $Z = M_D M_N^{-1}$.

The mass of the singlet fermion is acquired through a type-I seesaw mechanism:

$$m_S = -M\frac{1}{M_N}M^T \tag{19}$$

where M is the $N-S$ mixing mass term in the Yukawa Lagrangian (6).

4. Type-II Seesaw Fit to Oscillation Data

4.1. Neutrino Mass Matrix from Oscillation Data. Using diagonalization of neutrino mass matrix (m_ν) by the PMNS matrix U_{PMNS}

$$m_\nu = U_{\mathrm{PMNS}}\,\mathrm{diag}\left(m_1, m_2, m_3\right) U_{\mathrm{PMNS}}^T, \tag{20}$$

where m_i $(i = 1,2,3)$ denote the mass eigen values. For neutrino mixings we use the abbreviated cyclic notations $t_i = \sin\theta_{jk}$, $c_i = \cos\theta_{jk}$, where i, j, k are cyclic permutations of generational numbers 1, 2, 3. Following the standard parametrisation we denote the PMNS matrix [128–130]

$$U_{\mathrm{PMNS}} = \begin{pmatrix} c_3c_2 & t_3c_2 & t_2e^{-i\delta_D} \\ -t_3c_1-c_3t_1t_2e^{i\delta_D} & c_3c_1-t_3t_1t_2e^{i\delta_D} & t_1c_2 \\ t_3t_1-c_3c_1s_2e^{i\delta_D} & -c_3t_1-t_3c_1t_2e^{i\delta_D} & c_1c_2 \end{pmatrix} \cdot \mathrm{diag}\left(e^{i\alpha_M/2}, e^{i\beta_M/2}, 1\right), \tag{21}$$

where δ_D is the Dirac CP phase and (α_M, β_M) are Majorana phases.

Here we present numerical analyses within 3σ limit of the neutrino oscillation data in the type-II seesaw framework [126]. As we do not have any experimental information about Majorana phases, they are determined by means of random sampling; i.e., from the set of randomly generated values, each confined within the maximum allowed limit of 2π, only one set of values for (α_M, β_M) is chosen. Very recent analysis of the oscillation data has determined the 3σ and 1σ limits of Dirac CP phase δ_D [1]. The best fit values of δ_D in the normally ordered (NO) and invertedly ordered (IO) cases are near 1.2π and 1.5π, respectively, which we utilize for the sake of simplicity. A phenomenological model analysis has yielded $\delta_D = \pm 1.32\pi$ [150].

Global fit to the oscillation data [1] is summarized below including respective parameter uncertainties at 3σ level:

$$\begin{aligned}
\theta_{12}{}^{\circ} &= 34.5 \pm 3.25,\\
\theta_{23}{}^{\circ}\,(\mathrm{NO}) &= 41.0 \pm 7.25,\\
\theta_{23}{}^{\circ}\,(\mathrm{IO}) &= 50.5 \pm 7.25,\\
\theta_{13}{}^{\circ}\,(\mathrm{NO}) &= 8.44 \pm 0.5,\\
\theta_{13}{}^{\circ}\,(\mathrm{IO}) &= 8.44 \pm 0.5,\\
\frac{\delta_D}{\pi\,(\mathrm{NO})} &= 1.40 \pm 1.0,\\
\frac{\delta_D}{\pi\,(\mathrm{IO})} &= 1.44 \pm 1.0,\\
\Delta m_{21}^2 &= (7.56 \pm 0.545) \times 10^{-5}\mathrm{eV}^2,\\
\left|\Delta m_{31}\right|^2 (\mathrm{NO}) &= (2.55 \pm 0.12) \times 10^{-3}\mathrm{eV}^2,\\
\left|\Delta m_{31}\right|^2 (\mathrm{IO}) &= (2.49 \pm 0.12) \times 10^{-3}\mathrm{eV}^2.
\end{aligned} \tag{22}$$

We denote the cosmologically constrained parameter, the sum of the three active neutrino masses, as

$$\Sigma_\nu = \Sigma_i \widehat{m}_i. \tag{23}$$

For normally hierarchical (NH), inverted hierarchical (IH), and quasidegenerate (QD) patterns, the experimental values of mass squared differences have been fitted by the following values of light neutrino masses and the respective values of the cosmological parameter Σ_ν.

$$\begin{aligned}
&\widehat{m}_\nu = (0.00127, 0.008838, 0.04978)\ \mathrm{eV} \quad (\mathrm{NH})\\
&\Sigma_\nu = 0.059888\ \mathrm{eV},\\
&\widehat{m}_\nu = (0.04901, 0.04978, 0.00127)\ \mathrm{eV} \quad (\mathrm{IH})\\
&\Sigma_\nu = 0.059888\ \mathrm{eV},\\
&\widehat{m}_\nu = (0.2056, 0.2058, 0.2)\ \mathrm{eV} \quad (\mathrm{QD}),\\
&\Sigma_\nu = 0.6114\ \mathrm{eV}.
\end{aligned} \tag{24}$$

Using oscillation data and best fit values of the mixings, we have also determined the PMNS mixing matrix numerically:

$$U_{\mathrm{PMNS}} = \begin{pmatrix} 0.816 & 0.56 & -0.0199 - 0.0142i \\ -0.354 - 0.0495i & 0.675 - 0.0346i & 0.650 \\ 0.450 - 0.0568i & -0.485 - 0.0395i & 0.75 \end{pmatrix}. \tag{25}$$

4.2. Determination of Majorana Yukawa Coupling Matrix. Now inverting the relation $\widehat{m}_\nu = U^\dagger_{PMNS} \mathscr{M}_\nu U^*_{PMNS}$ where $\widehat{m}_\nu$ is the diagonalized neutrino mass matrix, we determine $\mathscr{M}_\nu$ for three different cases and further determine the corresponding values of the f matrix using $f = \mathscr{M}_\nu / v_L$ where we use the predicted value of $v_L = 0.1$ eV.

NH

$$f = \begin{pmatrix} 0.117 + 0.022i & -0.124 - 0.003i & 0.144 + 0.025i \\ -0.124 - 0.003i & 0.158 - 0.014i & -0.141 + 0.017i \\ 0.144 + 0.025i & -0.141 + 0.017i & 0.313 - 0.00029i \end{pmatrix} \tag{26}$$

IH

$$f = \begin{pmatrix} 0.390 - 0.017i & 0.099 + 0.01i & -0.16 + 0.05i \\ 0.099 + 0.01i & 0.379 + 0.02i & 0.176 + 0.036i \\ -0.16 + 0.05i & 0.176 + 0.036i & 0.21 - 0.011i \end{pmatrix} \tag{27}$$

QD

$$f = \begin{pmatrix} 2.02 + 0.02i & 0.0011 + 0.02i & -0.019 + 0.3i \\ 0.0011 + 0.02i & 2.034 + 0.017i & 0.021 + 0.21i \\ -0.019 + 0.3i & 0.021 + 0.21i & 1.99 - 0.04i \end{pmatrix} \tag{28}$$

Randomly chosen Majorana phases [126] $\alpha_M = 74.84^\circ$, $\beta_M = 112.85^\circ$, and the central value of the Dirac phase $\delta_D = 218^\circ$ have been used in this analysis. Using the well-known definition of the Jarlskog-Greenberg [151, 152] invariant,

$$J_{CP} = -t_3 c_3 t_2 c_2^2 t_1 c_1 \sin\delta_D, \tag{29}$$

and keeping δ_D at its best fit values, we have estimated the predicted allowed ranges of the CP-violating parameter in both cases:

$$\begin{aligned}
J_{CP} &= 0.0175 - 0.0212 \quad (\mathrm{NH})\\
J_{CP} &= 0.0302 - 0.0365 \quad (\mathrm{IH})
\end{aligned} \tag{30}$$

where the variables have been permitted to acquire values within their respective 3σ ranges of the oscillation data. Besides these there are nonunitarity contributions which have been discussed extensively in the literature.

4.3. Scaling Transformation of Solutions. In general there could be type-II seesaw models characterizing different seesaw scales and induced VEVs matching the given set of neutrino oscillation data represented by the same neutrino mass matrix. For two such models,

$$m_\nu = f^{(1)} v_L^{(1)} = f^{(2)} v_L^{(2)}, \tag{31}$$

Then the $f-$ matrix in one case is determined up to good approximation in terms of the other from the knowledge of the two seesaw scales.

$$f^{(1)} \simeq f^{(2)} \frac{M_{\Delta^{(1)}}}{M_{\Delta^{(2)}}} \tag{32}$$

At $M_{\Delta^{(1)}} = 10^{12}$ GeV our solutions are the same as in [126]. In view of this scaling relation, we can determine the values of the Majorana Yukawa matrix in the present case from the estimations of [126]. For example, if we choose $M_{\Delta^{(1)}} = 10^{10}$ GeV in the present case compared to $M_{\Delta^{(2)}} = 10^{12}$ GeV in [126], we rescale the solutions of [126] by a factor 10^{-2} to derive solutions in the present case. Thus graphical representations of solutions are similar to those of [126] for $M_{\Delta^{(2)}} = 10^{12}$ GeV which we do not repeat here. The values of magnitudes of f_{ij} at any new scale are obtained by rescaling them by the appropriate scaling factor while the phase angles remain the same as in [126].

4.4. Dirac Neutrino Mass Matrix. The Dirac neutrino mass matrix M_D plays crucial role in predicting LFV and LNV decays. In certain SO(10) models [50, 51, 125] this is usually determined by fitting the charged fermion masses at the GUT scale and equating it with the up-quark mass matrix. The fact that $M_D^0 \simeq M_u^0$ at the GUT scale follows from the underlying quark lepton symmetry [44] of SO(10). In SU(5) itself, however, there is no such symmetry to predict the structure of M_D in terms of quark matrices. Also in this SU(5) model we do not attempt any charged fermion mass fit at the GUT scale or above it. Since the Dirac neutrino mass matrix is not predicted by the SU(5) symmetry itself, for the sake of simplicity and to derive maximal effects on LFV and LNV decays, we assume M_D^0 to be equal to the up-quark mass matrix M_u^0 at the GUT scale. Noting that N is SU(5) singlet fermion, in the context of relevant Yukawa interaction Lagrangian

$$-\mathscr{L}_{\text{Yuk}} = \left[Y_N \overline{5}_F \cdot \mathbf{1}_F \cdot 5_H + Y_u 10_F \cdot 10_F \cdot 5_H + \cdots\right] + h.c., \tag{33}$$

this assumption is equivalent to alignment of the two Yukawa couplings:

$$Y_N \simeq Y_u. \tag{34}$$

This alignment is naturally predicted in SO(10) or SO(18) [73], but in the present SU(5) case it is assumed.

We realize this matrix M_D using renormalization group equations for fermion masses and gauge couplings and their numerical solutions [153–155] starting from the PDG values [128–130] of fermion masses at the electroweak scale. Following the bottom-up approach and using the down-quark diagonal basis, the quark masses and the CKM mixings are extrapolated from low energies using renormalization group (RG) equations [153–157]. After assuming the approximate equality $M_D^0 \simeq M_u^0$ at the GUT scale where M_u^0 is the up-quark mass matrix, the top-down approach is exploited to run down this mass matrix M_D^0 using RG equations [153]. Then the value of M_D near $1-10$ TeV scale turns out to be

$$M_D \simeq \begin{pmatrix} 0.014 & 0.04-0.01i & 0.109-0.3i \\ 0.04+0.01i & 0.35 & 2.6+0.0007i \\ 0.1+0.3i & 2.6-0.0007i & 79.20 \end{pmatrix} GeV. \tag{35}$$

As already noted above, although on the basis of SU(5) symmetry alone there may not be any reason for the rigorous validity of (35), in what follows we study the implications of this assumed value of M_D to examine maximum possible impact on LFV and LNV decays discussed in Sections 6 and 7. Another reason is that the present assumption on M_D may be justified in direct SO(10) breaking to the SM which we plan to pursue in a future work.

5. Right-Handed Neutrino Mass in SU(5) vs SO(10)

5.1. RHν Mass in SO(10). The fermions responsible for type-I and type-II seesaw are the LH leptonic doublets and the RH fermionic singlets of three generations. In SO(10) the left-handed lepton doublet $(\nu, l)^T$ and the right-handed neutrino N are in the same spinorial representation 16_F.

$$(\nu, l)^T \subset 16_F, \quad N \subset 16_F. \tag{36}$$

The Higgs representation $126_H^\dagger \subset SO(10)$ contains both the left-handed and the right-handed triplets carrying $B-L = -2$,

$$126_H^\dagger = \Delta_L(3,1,-2,1) + \Delta_R(1,3,-2,1) + \cdots \tag{37}$$

where the quantum numbers are under the left-right symmetry group $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$. The common Yukawa coupling f_{10} in the Yukawa term

$$-\mathscr{L}_{\text{Yuk10}} = f_{10} 16_F 16_F 126_H^\dagger \tag{38}$$

generates the dilepton-Higgs triplet interactions both in the left-handed and right-handed sectors giving rise to type-I and type-II seesaw mechanisms. The RH neutrino mass is generated through the VEV of the neutral component of the of Δ_R

$$\mathscr{M}_N = f_{10} \left\langle \Delta_R^0 \right\rangle. \tag{39}$$

The type-II seesaw contribution to light neutrino mass is

$$\mathscr{M}_\nu = f_{10} v_L \tag{40}$$

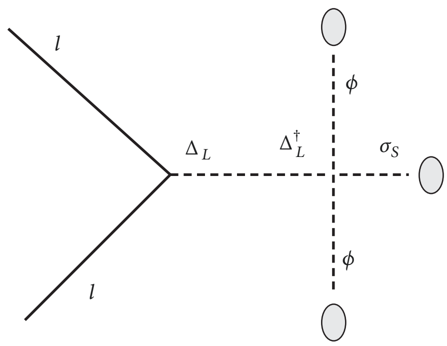

Figure 1: Feynman diagram representing type-II seesaw mechanism for neutrino mass generation in SU(5). Scalar fields ϕ, σ_S, and Δ_L represent SM Higgs doublet, singlet, and LH triplet as defined in the text. This diagram defines the trilinear coupling mass $\mu_\Delta = \lambda\langle\sigma_S\rangle$.

where v_L is the corresponding induced VEV of Δ_L

$$v_L = \lambda_{10}\frac{\left\langle \Delta_R^0 \right\rangle v_{ew}^2}{M_{\Delta_L}^2}. \tag{41}$$

Here λ_{10} is the quartic coupling in the part of the scalar potential

$$V_{10} = \lambda_{10}\Delta_L^\dagger \Delta_R \phi^\dagger \phi \subset \lambda_{10} 126_H^\dagger 126_H 10_H 10_H. \tag{42}$$

Thus with type-II seesaw dominance, the predicted heavy RH neutrino masses in SO(10) follow the same hierarchical pattern as the active light neutrino masses

$$M_{N_1} : M_{N_2} : M_{N_3} :: m_1 : m_2 : m_3. \tag{43}$$

5.2. RHν Mass in SU(5). Feynman diagram for type-II seesaw mechanism in the present SU(5) model is shown in Figure 1.

In contrast to SO(10) where the LH leptonic doublet and the RHν are in one and the same representation 16_F, in SU(5) they are in different representations

$$\begin{gathered}(\nu, l)^T \subset \overline{5}_F, \\ N \subset \mathbf{1}_F.\end{gathered} \tag{44}$$

In SU(5), while the dilepton-Higgs interaction is given by

$$-\mathscr{L}_{Yukll} = f\overline{5}_F\overline{5}_F 15_H, \tag{45}$$

the RH neutrino mass is generated through

$$-\mathscr{L}_{YukNN} = \left(\frac{1}{2}\right) f_N NN\sigma_S + h.c. \tag{46}$$

The fact that N is a singlet under SU(5) forces σ_S to be a singlet too. Further this singlet σ_S must carry $B - L = -2$ as its VEV generates the heavy Majorana mass

$$M_N = f_N \left\langle \sigma_S \right\rangle. \tag{47}$$

In sharp contrast to SO(10) where the LH triplet Δ_L and the RH triplet Δ_R scalars contained in the same representation $126_H^\dagger$ generate the type-II seesaw and M_N, the situation in SU(5) is different. Since LH triplet $\Delta_L(3, -1, 1)$ mediating type-II seesaw belongs to Higgs representation $15_H \subset$ SU(5) and σ_S belongs to a completely different representation (which is a singlet $\subset$ SU(5)), the two relevant Majorana type couplings in general may not be equal

$$f_N \neq f. \tag{48}$$

Also this assertion is further strengthened if we do not assume SU(5) to be a remnant of SO(10). Then the RH neutrino mass hierarchy can be decoupled from the type-II seesaw prediction. It is interesting to note that in SU(5)

$$v_L = \frac{\mu_\Delta v_{\mathrm{ew}}^2}{M_\Delta^2} \tag{49}$$

where μ_Δ is the trilinear coupling in the potential term

$$V_{tri} = \mu_\Delta \Delta_L \phi\phi + h.c. \tag{50}$$

The VEV of this singlet σ_S can explain the dynamical origin of such trilinear coupling through its VEV $v_\sigma = \langle\sigma_S\rangle$

$$\mu_\Delta = \lambda v_\sigma, \tag{51}$$

where λ is the quartic coupling in the potential term

$$V_{ql} = \lambda\sigma_S \Delta_L \phi\phi + h.c. \tag{52}$$

$$\subset \lambda\sigma_S 15_H 5_H 5_H + h.c \tag{53}$$

where the second line represents the SU(5) origin. For GUT scale $U(1)_{B-L}$ symmetry breaking driving VEV $v_\sigma \simeq M_{GUT}$ in the natural absence of any intermediate symmetry, it is possible to ensure $\mu_\Delta \simeq M_{\Delta_L}$ for

$$\lambda \simeq \frac{M_{\Delta_L}}{M_{GUT}}. \tag{54}$$

Thus the SU(5) model gives similar explanation for quartic coupling as in direct breaking case of SO(10). But the predicted hierarchy of RHν masses in SU(5) may not, in general, follow the same hierarchical pattern as given by SO(10) shown in (43). This is precisely so because (43) follows from the fact that the same matrix U_{PMNS} diagonalizes both the LH and the RH neutrino mass matrices which is further rooted in the fact that the same Majorana coupling f_{10} that generates the type-II seesaw mass term also generates M_N. But because of the general possibility in SU(5) that $f_N \neq f$, the RHνs may acquire a completely different pattern depending upon the value of f_N. Unlike SO(10), these masses emerging from SU(5) are also allowed to be quite different from the type-II seesaw scale.

Even if the value of v_σ may be needed to be near M_{Δ_L}, the value of M_N is allowed to be considerably lower by fine-tuning the value of f_N. Our LFV and LNV decay phenomenology as discussed below may need $M_N = 1 - 10$ TeV which is realizable using this new technique in SU(5). In contrast SO(10) needs $U(1)_R \times U(1)_{B-L}$ [84, 125, 157] or

$SU(2)_R \times U(1)_{B-L}$ gauge symmetry and hence new gauge bosons near the TeV scale to generate such RHν masses which should be detected at LHC [84, 125]. Thus a new mechanism for RHν mass emerges here by noting the coupling $f_N \neq f$ which has the potential to generate RHν masses over a wide range of values $M_N \sim 100 - 10^{15}$ GeV. Then the RHν mass predictions in the two GUTs in the presence of type-II seesaw dominance are as follows:

Type-II Seesaw Dominated SO(10)

$$M_{N_i} \simeq \frac{m_i M_{\Delta_L}^2}{v_{ew}^2}. \tag{55}$$

Type-II Seesaw Dominated SU(5)

$$M_{N_i} = \left[\mathcal{O}(10)\,\text{GeV} - \mathcal{O}\left(M_{\Delta_L}\right)\right]. \tag{56}$$

Here $m_i, i = 1, 2, 3$, are the three mass eigen values of light neutrinos. It is to be noted that m_i is absent in the RHS of (56) in the SU(5) case.

5.3. Realization of Mass Hierarchies. Here we discuss how the stated hierarchy in Section 3

$$M_{\Delta_L} \gg M_N \gg M \gg M_D, \tag{57}$$

for type-II seesaw dominance, is realized through fine tuning. At first noting that the mass squared term for $15_H \subset SU(5)$ in the scalar potential, $M_{15}^2 15_H^\dagger 15_H$, is SU(5) invariant, M_{15} can have any mass below the GUT scale subject to proton stability and gauge coupling unification. Since 15_H, unlike 5_H, does not have Yukawa interactions with SM fermions, the Higgs mediated proton decay is suppressed. We now explain why we have used $M_{15} = M_{\Delta_L} = 10^{12}$ GeV. In our model it is possible to assign any mass $M_{15} = M_{\Delta_L} > M_\eta$ where the mass of $\eta(3, 0, 8) \subset 75_H$ is $M_\eta = 10^{9.23}$ GeV. Because of the presence of $\eta(3, 0, 8)$ at such intermediate mass value, precision gauge coupling unification is achieved which has been discussed separately [126, 127]. Following the standard symmetry breaking SU(5) $\longrightarrow$ SM through the GUT scale VEV of the SM singlet scalar in the adjoint representation 24_H, $V_{GUT} = \langle 24_H^0 \rangle$, a SU(5) invariant scalar potential V_η gives the mass of $\eta(3, 0, 8)$

$$\begin{aligned} V_\eta &= M_{75}^2 75_H^2 + m_{(24,75)} 24_H 75_H^2 + \lambda_{(24,75)} 24_H^2 75_H^2 \\ &\supset \left[M_{75}^2 + m_{(24,75)} V_{GUT} + \lambda_{(24,75)} V_{GUT}^2\right] \eta^2. \end{aligned} \tag{58}$$

This leads to $M_\eta^2 = M_{75}^2 + m_{(24,75)} V_{GUT} + \lambda_{(24,75)} V_{GUT}^2$. Here $M_{75} \sim m_{(24,75)} \sim V_{GUT} \sim M_{GUT}$. Thus fine-tuning any one of these four parameters can give $M_\eta = 10^{9.23}$ GeV. The presence of 15_H below M_η destabilizes unification but protects it for $M_{15} > M_\eta$. This has led to the chosen value of $M_{\Delta_L} = M_{15} = 10^{12}$ GeV. We have noted in the following section that the value of $M_3 < 1$ TeV violates the observed bound on the nonunitarity parameter $\eta_{\tau\tau} < 2.7 \times 10^{-3}$ leading to the lower bound on M in the degenerate case

$$M_1 = M_2 = M_3 \geq 1250\ \text{GeV}. \tag{59}$$

Here $M = \text{diag}(M_1, M_2, M_3)$. Noting the definition

$$M = y_\chi \left\langle \chi_S (1, 0, 1) \right\rangle = y_\chi V_\chi, \tag{60}$$

we now argue that even for GUT scale mass of χ_S and its VEV $V_\chi = V_{GUT}$, it is possible to satisfy (59). For $V_\chi \sim V_{GUT} \sim 10^{15}$ GeV we need a small fine-tuned value of Yukawa coupling

$$y_\chi > 10^{-12}, \tag{61}$$

which satisfies $M \geq 1$ TeV but does not affect any known fermion mass. This shows the interesting possibility that even without having a low mass nonstandard Higgs χ_S, it is possible to realize the extended seesaw with type-II seesaw dominance. However, if we insist on $y_\chi \leq 1$, we need $V_\chi \geq 1$ TeV and $M_{\chi_S} \geq 1$ TeV which is realizable as χ_S is a SU(5) scalar singlet. As we have assumed $M_D = M_u$, this gives at GUT scale on extrapolation [153]

$$M_{D_{33}} \sim m_{\text{top}} \simeq 85\ \text{GeV}. \tag{62}$$

Here m_{top} is the top-quark mass. Thus achieving precision unification and type-II seesaw dominance has already given us $M_{Delta} > M_\eta$ whereas fine-tuning the Majorana coupling, $f_N > 10^{-11}$, has yielded $M_N > 10^4$ GeV. Combining these with (59) and (62) gives the hierarchical relation of (57).

6. Lepton Flavor Violations

Using SM extensions there has been extensive investigation of lepton flavor violating phenomena $l_\alpha \longrightarrow l_\beta + \gamma, (\alpha \neq \beta)$ and other processes like $\mu \longrightarrow e\bar{e}e$ including unitarity violations [101–108, 110–116]. In the flavor basis we use the standard charged current Lagrangian

$$\mathscr{L}_{\text{CC}} = -\frac{1}{\sqrt{2}} \sum_{\alpha=e,\mu,\tau} \left[g_{2L} \bar{\ell}_{\alpha L} \gamma_\mu \nu_{\alpha L} W_L^\mu\right] + \text{h.c.} \tag{63}$$

In predicting the LFV branching ratios we have used the relevant formulas of [101] and assumed a simplifying diagonal structure for M,

$$M = \text{diag.}\left(M_1, M_2, M_3\right). \tag{64}$$

Then (64), in combination with (35), gives the elements of the $\nu - S$ mixing matrix

$$\mathscr{V}^{(lS)} = \begin{pmatrix} \dfrac{M_{De1}}{M_1} & \dfrac{M_{De2}}{M_2} & \dfrac{M_{De3}}{M_3} \\ \dfrac{M_{D\mu1}}{M_1} & \dfrac{M_{D\mu2}}{M_2} & \dfrac{M_{D\mu3}}{M_3} \\ \dfrac{M_{D\tau1}}{M_1} & \dfrac{M_{D\tau2}}{M_2} & \dfrac{M_{D\tau3}}{M_3} \end{pmatrix}. \tag{65}$$

The $S - N$ mixing matrix,

$$V^{(SN)} = \frac{M}{M_N}, \tag{66}$$

is relatively damped out since $M_N \gg M$. In fact the type-I cancellation condition $M_N >> M >> M_D$ ensures this damping. Noting that the physical neutrino flavor state ν_α is a mixture of $\widehat{\nu}$, $\widehat{S}$, and $\widehat{N}$,

$$\nu_\alpha = U_{\alpha i}\widehat{\nu}_i + V_{\alpha i}^{lS}\widehat{S} + V_{\alpha i}^{(SN)}\widehat{N}_i. \quad (67)$$

Here $U \sim U_{PMNS}$ and the other two mixings violate unitarity. For large $M_N \gg M$ the third term in the RHS of (67) can be dropped leading to the unitarity violation parameter η

$$U' \simeq (1-\eta)\, U_{PMNS}. \quad (68)$$

Here

$$\eta_{\alpha\beta} = \left(\frac{1}{2}\right)\left(X \cdot X^\dagger\right)_{\alpha\beta}, \qquad X = \frac{M_D}{M}. \quad (69)$$

There has been extensive discussion on the constraint imposed on this parameter [108, 111]. The largest out of these is $\eta_{\tau\tau} \leq 0.0027$. Theoretically

$$\frac{1}{2}\left[\sum_i \frac{M_{D_{\tau i}} \cdot M_{D_{\tau i}}^*}{M_i^2}\right] \leq 0.0027. \quad (70)$$

In the completely degenerate case of $S-N$ mixing, $M_1 = M_2 = M_3 = M$, we get the following.

$$M \geq 1250\mathrm{GeV} \quad (71)$$

The RH neutrinos in the present model being degenerate with masses $M_{N_i} \gg m_{S_i}$ have much less significant contributions than the singlet fermions. The predicted branching ratios being only few to four orders less than the current experimental limits [95] are verifiable by ongoing searches,

$$BR(\mu \longrightarrow e\gamma) = 1.19 \times 10^{-16}, \quad BR(\tau \longrightarrow e\gamma) = 1.69 \times 10^{-14}, \quad BR(\tau \longrightarrow \mu\gamma) = 1.8 \times 10^{-12}. \quad (72)$$

For the sake of completeness we present the variation of LFV decay branching ratios as a function of the lightest neutrino mass in Figure 2.

In this approach the LFV decay rate mediated by the W_L boson in the loop depends predominantly upon $N-S$ mixing matrix M and the Dirac neutrino mass matrix M_D, although subdominantly upon the RHν mass matrix M_N. However in the high scale type-II seesaw ansatz followed here LFV decay rate is independent of light neutrino masses. This behavior of LFV decay rates is clearly exhibited in Figure 2 where the three branching ratios have maintained constancy with the variation of m_ν.

7. Dominant W_L–W_L Channel Double Beta Decay within Cosmological Bound

7.1. Double Beta Decay Mediation by Sterile Neutrinos. In the absence of W_R bosons and right-handed $\Delta_R^{\pm\pm}$ in SU(5), there is

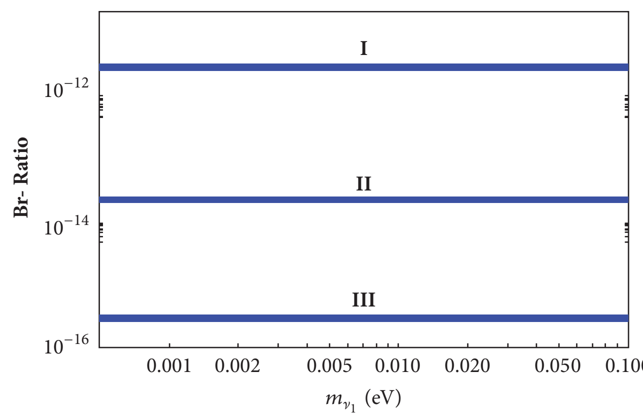

Figure 2: Variation of LFV decay branching ratios as a function of the lightest neutrino mass. Colored horizontal lines represent I : $BR(\tau \longrightarrow \mu\gamma)$, II : $BR(\tau \longrightarrow e\gamma)$, and III : $BR(\mu \longrightarrow e\gamma)$.

no contribution to right-handed double beta production. The gauge coupling unification constraint sets the lower bound on the masses of left-handed doubly charged Higgs bosons $\Delta_L^{\pm\pm}$ to be $M_{\Delta_L} \simeq M_{15_H} > 10^{9.23}$ GeV. As such they have negligible contributions for direct mediations of $0\nu\beta\beta$ process with left-handed electrons. Thus the only significant contributions in the $W_L - W_L$ channel could be through the mediation of ν, S, and N. Feynman diagrams for $0\nu\beta\beta$ decay amplitude in the $W_L - W_L$ channel due to the exchanges of Majorana fermions ν, S, and N are shown in Figure 3. In Figure 4 we also present Feynman diagram for $0\nu\beta\beta$ decay amplitude due to the sterile neutrino exchange where its mass insertion has been explicitly indicated. Mass eigen values of different sterile neutrinos for different sets of (M_1, M_2, M_3) consistent with constraints on unitarity violating parameters $\eta_{\alpha\beta}$ are presented in Table 1. We have used the singlet fermion mass seesaw formula of (15) and $M_{N_1} = M_{N_2} = M_{N_3} = 4000$ GeV. These solutions are displayed in Figure 5.

We use normalizations necessary for different contributions [158–164] due to exchanges of light neutrinos, sterile neutrinos, and the heavy RH neutrinos in the $W_L - W_L$ channel. They lead to the inverse half-life [120, 124, 125],

$$\begin{aligned}\left[T_{1/2}^{0\nu}\right]^{-1} &\simeq G_{01}\left|\frac{\mathcal{M}_\nu^{0\nu}}{m_e}\right|^2 \left|\mathbf{M}_\nu^{ee} + \mathbf{M}_S^{ee} + \mathbf{M}_N^{ee}\right|^2, \\ &= K_{0\nu}\left|\mathbf{M}_\nu^{ee} + \mathbf{M}_S^{ee} + \mathbf{M}_N^{ee}\right|^2, \\ &= K_{0\nu}\left|\mathbf{M}_{\mathrm{eff}}\right|^2.\end{aligned} \quad (73)$$

Here $G_{01} = 0.686 \times 10^{-14}\mathrm{yrs}^{-1}$, $\mathcal{M}_\nu^{0\nu} = 2.58 - 6.64$, and $K_{0\nu} = 1.57 \times 10^{-25}\mathrm{yrs}^{-1}\mathrm{eV}^{-2}$. In (73) the three effective mass parameters have been defined as

$$\mathbf{M}_\nu^{ee} = \sum_i \left(\mathcal{V}_{ei}^{\nu\nu}\right)^2 m_{\nu_i} \quad (74)$$

$$\mathbf{M}_S^{ee} = \sum_i \left(\mathcal{V}_{ei}^{\nu S}\right)^2 \frac{|p|^2}{\widehat{m}_{S_i}} \quad (75)$$

$$\mathbf{M}_N^{ee} = \sum_i \left(\mathcal{V}_{ei}^{\nu N}\right)^2 \frac{|p|^2}{M_{N_i}}, \quad (76)$$

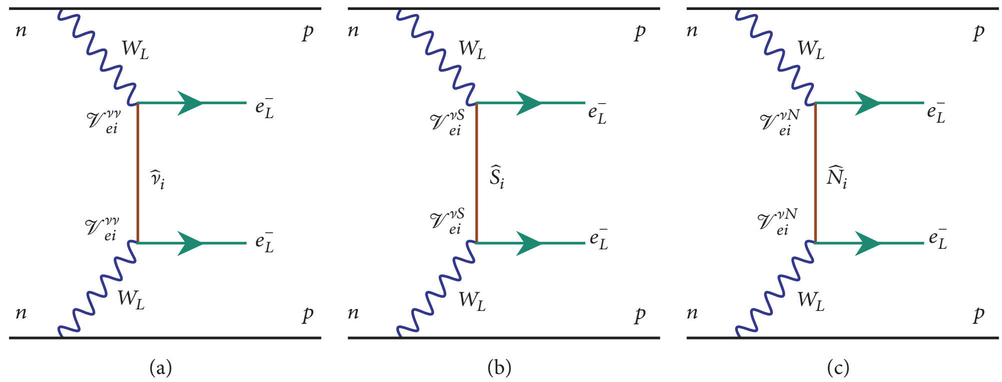

FIGURE 3: Feynman diagrams representing neutrinoless double beta decay due to exchanges of all three types of Majorana fermions ν, S, and N.

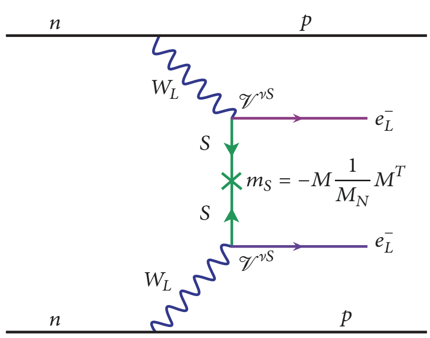

FIGURE 4: Feynman diagram representing neutrinoless double beta decay amplitude due to exchanges of singlet fermions S_i $(i = 1 - 3)$ with explicit mass insertion $m_s = -M(1/M_N)M^T$.

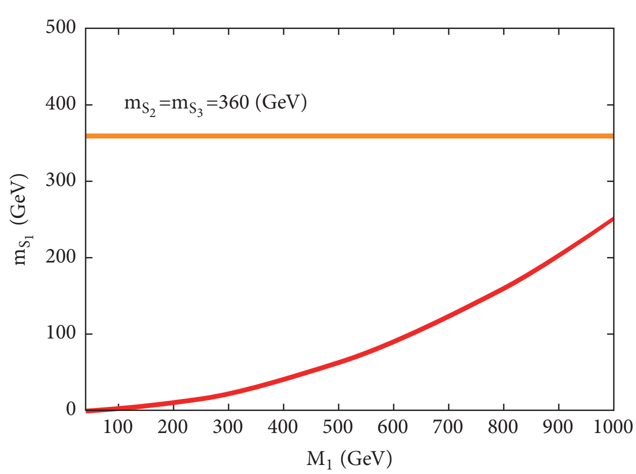

FIGURE 5: Prediction of singlet fermion mass eigen values as a function of $N - S$ mixing mass parameters M_i $(i = 1, 2, 3)$ for $M_{N_i} = 4$ TeV $(i = 1, 2, 3)$. The horizontal red colored line represents solutions for two other eigen values for $M_2 = M_3 = 1200$ GeV.

with

$$\mathbf{M}_{\mathrm{eff}} = \mathbf{M}_{\nu}^{ee} + \mathbf{M}_{S}^{ee} + \mathbf{M}_{N}^{ee}. \tag{77}$$

TABLE 1: Prediction of singlet fermion masses for different values of (M_1, M_2, M_3) where we have used $M_{N_1} = M_{N_2} = M_{N_3} = 4000$ GeV.

M (GeV)	$\widehat{m}_s$ (GeV)
(60, 1200, 1200)	(0.9, 360, 360)
(70, 1200, 1200)	(1.22, 360, 360)
(80, 1200, 1200)	(1.60, 360, 360)
(90, 1200, 1200)	(2.00, 360, 360)
(100, 1200, 1200)	(2.50, 360, 360)
(110, 1200, 1200)	(3.00, 360, 360)
(120, 1200, 1200)	(3.60, 360, 360)
(130, 1200, 1200)	(4.22, 360, 360)
(140, 1200, 1200)	(4.90, 360, 360)
(150, 1200, 1200)	(5.62, 360, 360)

The quantity $\widehat{m}_{S_i}$ is the i-th eigen value of the $S-$ fermion mass matrix m_S. The magnitude of neutrino virtuality momentum $|p|$ has been estimated to be in the allowed range $|p| = 120$ MeV-200 MeV [158–164]. The contributions of the RHνs, being much heavier than the singlet fermions, have been neglected.

7.2. Singlet Fermion Assisted Enhanced Double Beta Decay Rate. We use neutrino oscillation data to estimate M_{ν}^{ee} for NH and IH cases with the values of Dirac phase and Majorana phases as discussed above. We further use the values of M_i from Table 1 and Figure 5 and the Dirac neutrino mass matrix from (35) to estimate M_S^{ee} while treating the RHν mass at its assumed degenerate value of $M_{N_i} = 4$TeV $(i = 1, 2, 3)$. The variation of effective parameter m_{ee} as a function of lightest neutrino mass is shown in Figure 6 when $m_{s_1} = 2$ GeV.

As noted from the analytic formulas, the effective mass parameter in the singlet fermion dominated case, being inversely proportional to m_{s_1}, will proportionately decrease with increasing value of the mediating particle mass. This feature has been shown in Figure 7. We present predictions of double beta decay half-life as a function of the singlet fermion mass in Figure 8. It is clear that while for $m_{s_1} = 2$ GeV the half-life saturates the current experimental limit, for larger values of m_{s_1} the half-life is found to increase.

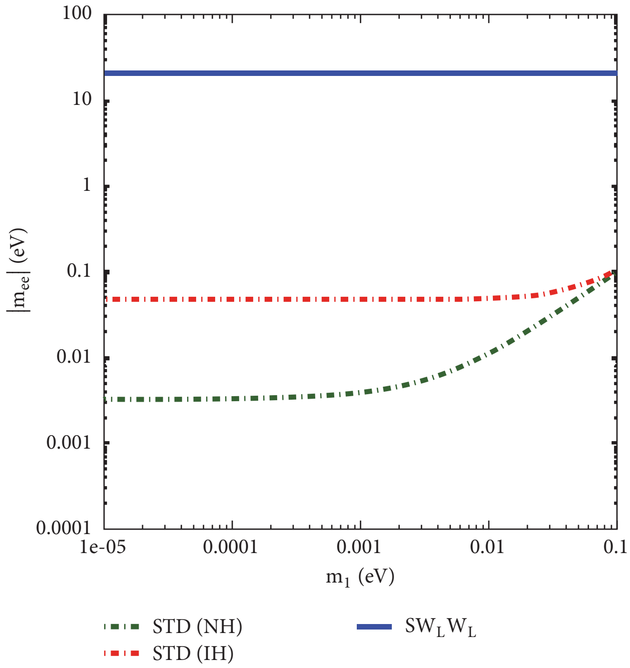

FIGURE 6: Variation of effective mass parameter as a function of lightest active neutrino mass m_1 for $m_{s_1} = 2$ GeV. For comparison, predictions in the standard model supplemented by light neutrino masses of NH type are shown by green dot-dashed curve. For IH pattern of mass hierarchy the standard prediction is shown by red dot-dashed curve.

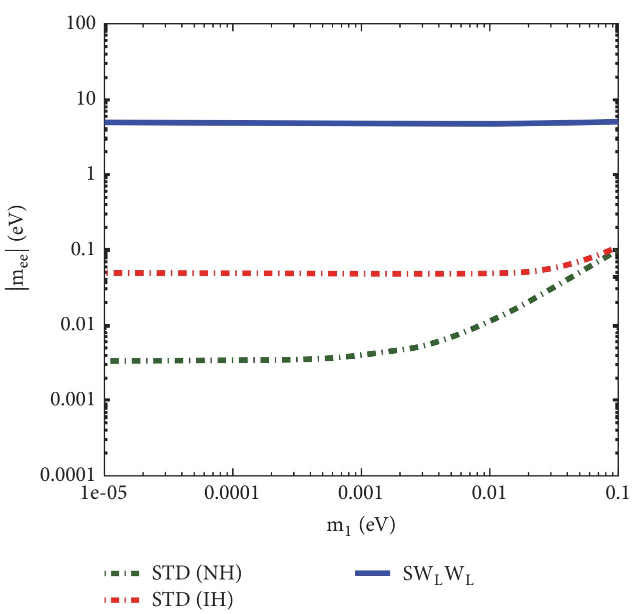

FIGURE 7: Same as Figure 6 but for $m_{s_1} = 4.0$ GeV.

Neglecting heavy RHν contributions but including those due to the lightest sterile neutrino and the IH type light neutrinos, our predictions of half-life as a function of the lightest sterile neutrino mass are shown in Figure 9.

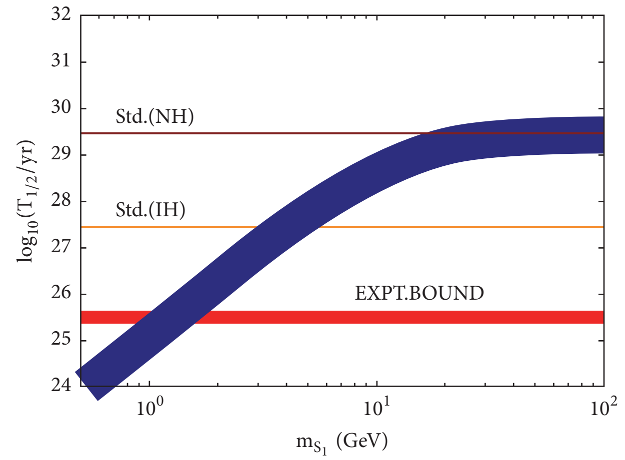

FIGURE 8: Prediction of double beta decay half-life as a function of sterile neutrino mass m_{s_1} GeV (blue shaded curve) where the NH type light neutrino and the sterile neutrino exchange contributions have been included. Effects of much larger masses $(m_{S_2}, m_{S_3}) \gg m_{S_1}$ have been neglected. The spread in the curve reflects uncertainty in the virtuality momentum $p = 120 - 190$ MeV. For comparison, the standard prediction with NH and IH pattern of light neutrino mass hierarchies is shown by the two respective horizontal lines. The bottom most thick red horizontal line closest to the X-axis represents overlapping experimental bounds from different groups [78–82].

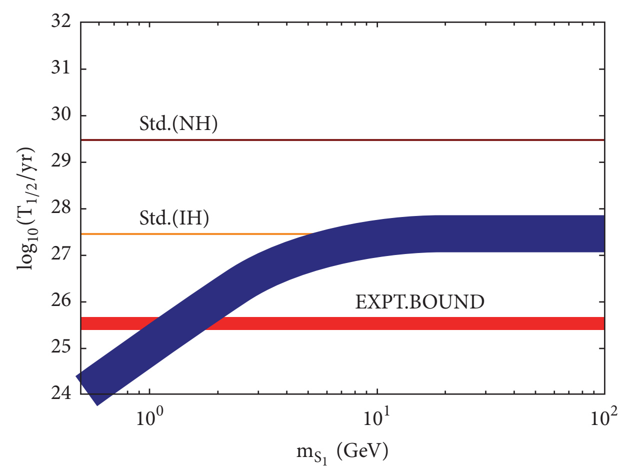

FIGURE 9: Same as Figure 8 but with contributions of IH type light neutrinos combined with lightest sterile neutrinos.

Predicted lifetimes are seen to decrease with increasing sterile neutrino mass. The sterile neutrino exchange contribution completely dominates over light neutrino exchange contributions for $m_{S_1} = 1.3 - 7$ GeV in case of IH but for $m_{S_1} = 1.5$–20 GeV in case of NH. At $m_{S_1} \simeq 1.5 GeV$ both types of solutions saturate the current laboratory limits reached by different experimental groups.

8. Summary, Discussion, and Conclusion

A recently proposed scalar extension of minimal non-SUSY SU(5) GUT has been found to realize precision gauge coupling unification, high scale type-II seesaw ansatz for neutrino masses, and prediction of a WIMP scalar DM candidate that also completes vacuum stability of the scalar potential. But the LFV decays are predicted to have negligible rates inaccessible to ongoing searches in foreseeable future. Likewise experimentally verifiable double beta decay rates measurable by different search experiments are possible only for quasidegenerate neutrino mass spectrum with large common mass scale $|m_0| > 0.2$ eV or $\sum_i m_i > 0.6$ eV. This violates the recently measured cosmological bound $\sum_i m_i \leq 0.23$ eV. In order to remove these theoretical shortcomings in the context of SU(5), we have extended this model by the addition of three RHνs, three extra Majorana fermion singlets S_i $(i = 1, 2, 3)$, and a scalar singlet $\chi_S(1, 0, 1)$ that generates $N - S$ mixing mass term through its vacuum expectation value. In the original theory of type-I seesaw cancellation mechanism, although the choice of particles is the same as N_i, S_i, and $\chi_S(1, 0, 1)$, the neutrino mass is given by double seesaw [119]. Further, there is no grand unification of gauge couplings or prediction of proton decay in this model [119], and the scalar potential of the model has vacuum instability. In addition the N_i are not gauged. The model does not predict dominant contributions to double beta decay for NH or IH type neutrino masses. In non-SUSY SO(10) models of unification of three forces implementing the cancellation of type-I seesaw [84, 120, 124], the TeV scale RH neutrinos are gauged but the neutrino masses are controlled by inverse seesaw. But in [13, 125] the RHνs are gauged and the neutrino mass formula is linear seesaw or type-II seesaw [13]. In all type-II seesaw dominated SO(10) models, the RHν masses have the same hierarchy as the left-handed neutrino masses: $M_{N_1} : M_{N_2} : M_{N_3} :: m_1 : m_2 : m_3$. This happens precisely because the left-handed and the right-handed dilepton Yukawa interactions originate from the same SO(10) invariant term: $f 16_F 16_F 126^{\dagger}$. In SU(5), however, as the LH triplet $\Delta_L(3, -1, 1)$ generating type-II seesaw and the singlet $\sigma_S(1, 0, 1)$ generating RHνs belong to different scalar representations, $15_H \subset SU(5)$ and $\mathbf{1}_H \subset SU(5)$, respectively, they can possess different Majorana couplings in their respective Yukawa interactions: $f l l \Delta_L{}^C$ and $f_N \sigma_S N N$. Because of this reason the generated RHν masses through $M_N = f_N \langle \sigma_s \rangle$ no longer follows the predicted type-II seesaw predicted hierarchical pattern. Then the allowed fine-tuning $|f_N| << |f|$ permits RH neutrino mass scale $M_N \sim \mathcal{O}(1 - 10)$ TeV even though, unlike SO(10) models, there are no possibilities of low mass W_R or Z' bosons at this scale. The apprehension of unacceptably large active neutrino mass generation through type-I seesaw mechanism is rendered inoperative through the well-established procedure of cancellation mechanism that is also shown to operate profoundly in this SU(5) model. Such RHνs generating $N - S$ mixing mass $M \simeq \mathcal{O}(100 - 1000)$ GeV now reproduce the well-known results on LFV decay branching ratios only 4 – 5 orders lower than the current experimental limit as well as the extensively investigated nonunitarity effects. Through the sterile neutrino canonical seesaw formula emerging from this cancellation mechanism (in the presence of N_i), $m_S = -M(1/M_N)M^T$, this mechanism predicts their masses over a wide range of values, $m_{s_1} = \mathcal{O}(1 - 100)$ GeV and $m_{s_2}, m_{s_3} \sim \mathcal{O}(10 - 1000)$ GeV. The lightest sterile neutrino mass m_{s_1} now predicts dominant double beta decay in the $W_L - W_L$ channel through the $\nu - S$ mixing close to the current experimental limits even though the light neutrino masses are of NH or IH type ($m_i << |0.2|$ eV) which satisfy the cosmological bound. For larger values of m_{S_1} the predicted decay rate decreases and the sterile neutrino contribution becomes negligible for $m_{s_i} >> 50$ GeV. In the limiting case when all the singlet fermion masses have such large values, the double beta decay rates asymptotically approach the respective standard NH or IH type contributions. The new mechanism of RHν mass generation also allows the second and the third generation sterile neutrino masses to be quasidegenerate (QD) near 1–10 TeV scale while keeping $m_{S_1} \sim 1 - 10$ GeV suitable for dominant double beta decay mediation. There is a possibility that such TeV scale QD masses while maintaining observable predictions on LFV decays can effectively generate baryon asymmetry of the universe via resonant leptogenesis [125]. A scalar singlet DM can be easily accommodated as discussed in [126, 165, 166] while resolving the issue of vacuum stability. Irrespective of scalar DM, the model can also accommodate a Majorana fermion singlet dark matter [167] which can emerge from the additional fermionic representation $24_F \subset SU(5)$.

The prediction of new fermions has an additional advantage over scalars as these masses are protected by leptonic global symmetries [168]. Also the prediction of such Majorana type sterile neutrinos can be tested by high energy and high luminosity accelerators through their like-sign dilepton production processes [169]. For example, at LHC they can mediate the process $pp \longrightarrow W_L X \longrightarrow l^{\pm} l^{\pm} jjX$ where the jets could be manifested as mesons. It would be quite interesting to examine emergence of such SU(5) theory as a remnant of SO(10) or E_6 GUTs.

We conclude that even in the presence of SM as effective gauge theory descending from a suitable SU(5) extension, it is possible to predict experimentally accessible double beta decay rates in the W_L–W_L channel satisfying the cosmological bound on active neutrino masses as well as verifiable LFV decays. The RHν masses can be considerably different from those constrained by conventional type-I or type-II seesaw frameworks which are instrumental in predicting interesting physical phenomena even if there are no nonstandard heavy gauge bosons anywhere below the GUT scale.

Appendix

A. Diagonalization, Masses, and Mixings

The purpose of this Appendix is to provide certain details of mixings among the fermions ν, S, and N and also derive their masses by block diagonalization of the resulting 9×9 neutral fermion mass matrix discussed in Section 3. We write the complete 9×9 mass matrix in the flavor basis $\{\nu_L, S_L, N_R^C\}$ after the effect of Δ_L is integrated out

$$\mathcal{M}_\nu = \begin{pmatrix} m^{(II)} & 0 & M_D \\ 0 & 0 & M^T \\ M_D^T & M & M_N \end{pmatrix} \tag{A.1}$$

where the type-II seesaw contribution has been noted as $m^{(II)} = f v_L$. The flavor basis to mass basis transformation and diagonalization of $\mathcal{M}_\nu$ is achieved by a unitary transformation matrix $\mathcal{V}$ defined below

$$|\psi\rangle_f = \mathcal{V}\,|\psi\rangle_m \tag{A.2}$$

$$\text{or,}\quad \begin{pmatrix} \nu_\alpha \\ S_\beta \\ N^C_\gamma \end{pmatrix} = \begin{pmatrix} \mathcal{V}^{\nu\nu}_{\alpha i} & \mathcal{V}^{\nu S}_{\alpha j} & \mathcal{V}^{\nu N}_{\alpha k} \\ \mathcal{V}^{S\nu}_{\beta i} & \mathcal{V}^{SS}_{\beta j} & \mathcal{V}^{SN}_{\beta k} \\ \mathcal{V}^{N\nu}_{\gamma i} & \mathcal{V}^{NS}_{\gamma j} & \mathcal{V}^{NN}_{\gamma k} \end{pmatrix} \begin{pmatrix} \widehat{\nu}_i \\ \widehat{S}_j \\ \widehat{N}_k \end{pmatrix} \tag{A.3}$$

$$\text{and}\quad \mathcal{V}^\dagger \mathcal{M}_\nu \mathcal{V}^* = \widehat{\mathcal{M}}_\nu = \operatorname{diag}\left(\widehat{m}_{\nu_i}; \widehat{m}_{\mathcal{S}_j}; \widehat{m}_{\mathcal{N}_k}\right) \tag{A.4}$$

where subscripts f, m denote the flavor and mass basis, respectively. Also $\mathcal{M}_\nu$ is the mass matrix in flavor basis with α, β, γ running over three generations of light neutrinos, sterile neutrinos, and right-handed heavy-neutrinos. Here $\widehat{\mathcal{M}}_\nu$ is the diagonal mass matrix with $(i, j, k = 1, 2, 3)$ running over corresponding mass states at the sub-eV, GeV, and TeV scales, respectively.

The method of complete diagonalization will be carried out in two steps: (**1**) the full neutrino mass matrix $\mathcal{M}_\nu$ has to be reduced to a block diagonalized form as $\mathcal{M}_{\mathrm{BD}}$; (**2**) this block diagonal form is further diagonalized to give physical masses of the neutral leptons $\widehat{\mathcal{M}}_\nu$.

(1) Determination of $\mathcal{M}_{\mathrm{BD}}$. We shall follow the parametrisation of the type given in [32] to determine the form of the diagonalizing matrices $\mathcal{W}_1$ and $\mathcal{W}_2$. We define their product as

$$\mathcal{W} = \mathcal{W}_1 \mathcal{W}_2 \tag{A.5}$$

where $\mathcal{W}_1$ and $\mathcal{W}_2$ satisfy

$$\begin{aligned} \mathcal{W}_1^\dagger \mathcal{M}_\nu \mathcal{W}_1^* &= \widehat{\mathcal{M}}_{\mathrm{BD}}, \\ \text{and}\ \mathcal{W}_2^\dagger \widehat{\mathcal{M}}_{\mathrm{BD}} \mathcal{W}_2^* &= \mathcal{M}_{\mathrm{BD}} \end{aligned} \tag{A.6}$$

where $\widehat{\mathcal{M}}_{\mathrm{BD}}$ and $\mathcal{M}_{\mathrm{BD}}$ are the intermediate block diagonal and full block diagonal mass matrices, respectively.

$$\widehat{\mathcal{M}}_{\mathrm{BD}} = \begin{pmatrix} \mathcal{M}_{eff} & 0 \\ 0 & m_{\mathcal{N}} \end{pmatrix} \tag{A.7}$$

$$\text{and}\ \mathcal{M}_{\mathrm{BD}} = \begin{pmatrix} m_\nu & 0 & 0 \\ 0 & m_{\mathcal{S}} & 0 \\ 0 & 0 & m_{\mathcal{N}} \end{pmatrix} \tag{A.8}$$

(2) Determination of $\mathcal{W}_1$. We need to first integrate out the heavy state (N_R), being heavier than other mass scales in our theory, such that up to the leading order approximation the analytic expression for $\mathcal{W}_1$ is

$$\mathcal{W}_1 = \begin{pmatrix} 1 - \frac{1}{2} B^* B^T & B^* \\ -B^T & 1 - \frac{1}{2} B^T B^* \end{pmatrix}, \tag{A.9}$$

where the matrix B is 6×3 dimensional and is described as

$$B^\dagger = M_N^{-1}\left(M_D^T, M^T\right) = \left(Z^T, Y^T\right) \tag{A.10}$$

where $X = M_D M^{-1}$, $Y = M M_N{}^{-1}$, and $Z = M_D M_N^{-1}$ so that $Z = X \cdot Y \neq Y \cdot X$.

Therefore, the transformation matrix $\mathcal{W}_1$ can be written purely in terms of dimensionless parameters Y and Z

$$\mathcal{W}_1 = \begin{pmatrix} 1 - \frac{1}{2} Z Z^\dagger & -\frac{1}{2} Z Y^\dagger & Z \\ -\frac{1}{2} Y Z^\dagger & 1 - \frac{1}{2} Y Y^\dagger & Y \\ -Z^\dagger & -Y^\dagger & 1 - \frac{1}{2}\left(Z^\dagger Z + Y^\dagger Y\right) \end{pmatrix} \tag{A.11}$$

while the light and heavy states can be now written as follows.

$$\mathcal{M}_{eff} = -\begin{pmatrix} M_D M_N^{-1} M_D^T & M_D M_N^{-1} M^T \\ M M_N^{-1} M_D^T & M M_N^{-1} M^T \end{pmatrix} \tag{A.12}$$

$$m_{\mathcal{N}} = M_N + \cdots \tag{A.13}$$

Determination of $\mathcal{W}_2$. From the above discussion, it is quite clear now that the eigenstates $\mathcal{N}_i$ are eventually decoupled from others and the remaining mass matrix $\mathcal{M}_{eff}$ can be block diagonalized using another transformation matrix

$$\mathcal{S}^\dagger \mathcal{M}_{\mathrm{eff}} \mathcal{S}^* = \begin{pmatrix} m_\nu & 0 \\ 0 & m_{\mathcal{S}} \end{pmatrix} \tag{A.14}$$

such that

$$\mathcal{W}_2 = \begin{pmatrix} \mathcal{S} & 0 \\ 0 & 1 \end{pmatrix}. \tag{A.15}$$

In a simplified structure,

$$\mathcal{M}_{eff} = \begin{pmatrix} m_\nu^{II} + M_D Z^T & M_D Y^T \\ Y M_D^T & M Y^T \end{pmatrix}. \tag{A.16}$$

Under the assumption at the beginning $Z << Y$ and of course $M_D << M$, this structure is similar to type-(I+II) seesaw. Therefore we immediately get the light neutrino masses as follows.

$$\begin{aligned} m_\nu &= -M_D Z^T + m_\nu^{II} + M_D Y^T \left(M Y^T\right)^{-1} Y M^T \\ &= -M_D Z^T + M_D Z^T + m_\nu^{II} = m_\nu^{II} \end{aligned} \tag{A.17}$$

$$m_{\mathcal{S}} = -M M_N^{-1} M^T \tag{A.18}$$

We see that in addition to $m_{\mathcal{N}}$, the $m_{\mathcal{S}}$ is also almost diagonal if M and M_N are assumed to be diagonal. The transformation matrix S is

$$S = \begin{pmatrix} 1 - \frac{1}{2} A^* A^T & A^* \\ -A^T & 1 - \frac{1}{2} A^T A^* \end{pmatrix} \tag{A.19}$$

such that

$$A^{\dagger} = \left(MY^{T}\right)^{-1} YM_D^{T} \simeq \left(MY^{T}\right)^{-1} YM_D^{T} = X^{T}. \quad \text{(A.20)}$$

The 3×3 block diagonal mixing matrix $\mathscr{W}_2$ has the following form.

$$\mathscr{W}_2 = \begin{pmatrix} S & \mathbf{0} \\ \mathbf{0} & \mathbf{1} \end{pmatrix} = \begin{pmatrix} 1 - \frac{1}{2}XX^{\dagger} & X & 0 \\ -X^{\dagger} & 1 - \frac{1}{2}X^{\dagger}X & 0 \\ 0 & 0 & 1 \end{pmatrix} \quad \text{(A.21)}$$

A.1. Physical Neutrino Masses from Complete Diagonalization. The block diagonal matrices m_ν, $m_{\mathcal{S}}$, and $m_{\mathcal{N}}$ can further be diagonalized to give physical masses for all neutral leptons by a unitary matrix $\mathscr{U}$ as

$$\mathscr{U} = \begin{pmatrix} U_\nu & 0 & 0 \\ 0 & U_S & 0 \\ 0 & 0 & U_N \end{pmatrix}. \quad \text{(A.22)}$$

Here the unitary matrices U_ν, U_S, and U_N satisfy the following.

$$\begin{aligned} U_\nu^{\dagger} m_\nu U_\nu^{*} &= \widehat{m}_\nu = \operatorname{diag}\left(m_{\nu_1}, m_{\nu_2}, m_{\nu_3}\right), \\ U_S^{\dagger} m_{\mathcal{S}} U_S^{*} &= \widehat{m}_S = \operatorname{diag}\left(m_{S_1}, m_{S_2}, m_{S_3}\right), \\ U_N^{\dagger} m_{\mathcal{N}} U_N^{*} &= \widehat{m}_N = \operatorname{diag}\left(m_{N_1}, m_{N_2}, m_{N_3}\right) \end{aligned} \quad \text{(A.23)}$$

With this discussion, the complete mixing matrix is as follows.

$$\begin{aligned} \mathscr{V} &= \mathscr{W} \cdot \mathscr{U} = \mathscr{W}_1 \cdot \mathscr{W}_2 \cdot \mathscr{U} \\ &= \begin{pmatrix} 1 - \frac{1}{2}ZZ^{\dagger} & -\frac{1}{2}ZY^{\dagger} & Z \\ -\frac{1}{2}YZ^{\dagger} & 1 - \frac{1}{2}YY^{\dagger} & Y \\ -Z^{\dagger} & -Y^{\dagger} & 1 - \frac{1}{2}\left(Z^{\dagger}Z + Y^{\dagger}Y\right) \end{pmatrix} \begin{pmatrix} 1 - \frac{1}{2}XX^{\dagger} & X & 0 \\ -X^{\dagger} & 1 - \frac{1}{2}X^{\dagger}X & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} U_\nu & 0 & 0 \\ 0 & U_S & 0 \\ 0 & 0 & U_N \end{pmatrix} \\ &= \begin{pmatrix} 1 - \frac{1}{2}XX^{\dagger} & X - \frac{1}{2}ZY^{\dagger} & Z \\ -X^{\dagger} & 1 - \frac{1}{2}\left(X^{\dagger}X + YY^{\dagger}\right) & Y - \frac{1}{2}X^{\dagger}Z \\ 0 & -Y^{\dagger} & 1 - \frac{1}{2}Y^{\dagger}Y \end{pmatrix} \cdot \begin{pmatrix} U_\nu & 0 & 0 \\ 0 & U_S & 0 \\ 0 & 0 & U_N \end{pmatrix} \end{aligned} \quad \text{(A.24)}$$

It is straightforward to verify that the matrix product of the right-hand side of (A.24) agrees with (18) of Section 3.

Data Availability

The data used to support the findings of this study are included within the article.

Conflicts of Interest

The authors declare that there are no conflicts of interest regarding the publication of this manuscript.

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