Phenomenology of the Zee model for Dirac neutrinos and general neutrino interactions

Julián Calle, Diego Restrepo, Óscar Zapata

Instituto de Física, Universidad de Antioquia, Calle 70 # 52-21, Apartado Aéreo 1226, Medellín, Colombia

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Abstract

The Zee model for Dirac neutrinos is one of the simplest models featuring one-loop Dirac neutrino masses. The interactions between the new scalars (two singly-charged fields) and neutrinos induce general neutrino interactions (GNI) which, as a generalisation of the non standard neutrino interactions, constitute an additional tool to probe models beyond the SM like this. In this work, we consider a $U(1)_{B-L}$ gauge symmetry as the responsible for the Diracness of the neutrinos and the radiative character of the neutrino masses. We determine the viable parameter space consistent with neutrino oscillation data, leptonic rare decays and collider constraints, and establish the most relevant experimental prospects regarding lepton flavor violation searches and GNI in future solar neutrino experiments.

1 Introduction

The underlying mechanism behind the massiveness of neutrinos [1, 2] holds, after several decades of research dedicated to neutrino physics, as one of the unresolved conundrums of the Standard Model (SM). In addition to this, the yet-to-be determined question of whether neutrinos are different from their antiparticles [3] and the lack of signatures of lepton flavor violation (LFV) in the charged sector [4, 5] make the landscape even more unclear. Fortunately the very rich experimental program in the lepton sector is undoubtedly rather clear and promising, as for instance, due to the entrance in operation during the current decade of neutrino oscillation experiments such as JUNO [6], DUNE [7], Hyper-Kamiokande [8], among others, and the great experimental effort on searching for LFV with, in some cases, reaching a sensitivity improvement of several orders of magnitude [9–14].
The simplest way to generate neutrino masses is to add to the SM spectrum three right-handed neutrinos, $\nu_{Ri}$, coupled to SM Higgs through a term in the Lagrangian as \[15–19\]
\[
\mathcal{L}_{4\nu} = y_D \bar{L} \tilde{H} \nu_R + \text{h.c.},
\]
with $L$ the lepton doublet, $H$ the SM Higgs doublet and $y_D$ the matrix of neutrino Yukawa couplings. Since after the electroweak symmetry breaking (and assuming lepton number conservation) such a term leads to sub-eV neutrino Dirac masses for $|y_D| \lesssim 10^{-13}$, the explanation of the smallness of the neutrino mass scale is yet to be settled. A way out of this dilemma may be to forbid $\mathcal{L}_{4\nu}$ by imposing a $U(1)_{B-L}$ local symmetry but generate it at loop level via the 5-dimensional operator \[20–22\]
\[
\mathcal{L}_{5\nu} = y'_D \bar{L} \tilde{H} \nu_R S + \text{h.c.},
\]
with $S$ being the singlet scalar field responsible for the breaking of the $B-L$ symmetry. From the huge variety of one-loop Dirac neutrino models\footnote{For a recent review see Ref. [23].} the one including only two singlet charged scalars and three singlet right-handed neutrino \[24–26\] may be considered as the simplest one -we dub it the Zee model for Dirac neutrinos.

In this paper we consider the Zee model for Dirac neutrinos with a $B-L$ symmetry, a realization obtained following the approach discussed in Refs. [20, 21, 27, 28]. Thanks to this new symmetry (with the charge assignment given in Table 1) it is possible to forbid the tree level contribution of $\mathcal{L}_{5\nu}$ without any extra \textit{ad-hoc} symmetry\footnote{Notice that in Refs. [24, 25] such contribution is forbidden by a softly-broken $Z_2$ symmetry.}. As occurs in the original Zee model \[29, 30\], neutrino masses are generated at one loop with the SM charged leptons running inside the loop, which is closed thanks to a trilinear interaction term between the charged scalars and the Higgs field. Both charged singlets not only give rise to LFV processes and may leave signatures at the LHC but also induce general neutrino interactions (GNI) \[31–33\]. We will determine the viable parameter space consistent with neutrino oscillation observables, LFV and LHC constraints, and we will establish the most relevant experimental perspectives regarding LFV searches and GNI in future solar neutrino experiments.

The rest of the paper is organized as follows. In section 2 we present the model and discuss the neutrino mass generation, the LFV processes and LHC signatures. In section 3 we deduce the expressions for the general neutrino interactions, and the phenomenological analysis is presented in section 4. Finally, we conclude in section 5.

## 2 The Model

The realization of the Zee Dirac model through a $U(1)_{B-L}$ gauge symmetry and without new extra fermions besides the three right-handed neutrinos $\nu_{Rij}$ requires the introduction of two $SU(2)_L$-singlet charged scalars, $\sigma_1^\pm$ and $\sigma_2^\pm$, and the two bosonic fields associated to the new symmetry: $Z'_\mu$ and $S$ as the gauge boson and the scalar responsible for the symmetry breaking, respectively\footnote{If such a breaking is associated to a strong first order phase transition with the subsequent broadcasting of gravitational waves it may be possible to further test this model in the planned gravitational wave experiments (see Refs. [34–36] for related studies).}.
The most general Lagrangian contains the following Yukawa terms

\[-\mathcal{L} \supset f_{\alpha \beta} \overline{L}_{\alpha} L_{\beta} \sigma_1^+ + h_{\beta j} \overline{\nu}_{Rj} \nu_{Rj} \sigma_2^+ + \text{h.c.},\]  

(3)

where \( L_{\alpha} \) and \( \ell_{R\alpha} \) are lepton doublets and singlets, respectively, \( c \) is charge conjugation operator, \( f_{\alpha \beta} \) and \( h_{\beta j} \) are \( 3 \times 3 \) matrices in the flavor space, and \( \alpha, \beta = 1, 2, 3 \) are family indices. The coupling \( f \) is an antisymmetric complex matrix\(^4\) whereas all the non zero \( h_{\beta j} \) couplings are complex in general. Note that there are three vanishing couplings because the lepton number (\( L \)) assignment of right-handed neutrino of \( L \)-charge \(-5\): \( \nu_{R3} (\nu_{R1}) \) for the case of a normal (inverted) hierarchy of neutrino masses\(^5\). This massless chiral fermion, can either contribute to the effective number of relativistic degrees of freedom, \( N_{\text{eff}} \), in the early universe [20, 22], or becomes massive and be a good Majorana dark matter candidate after the introduction of an extra singlet scalar, \( S' \), of \( L \)-charge \(-10 \) [22, 38]. In the last case, the extra contribution to \( N_{\text{eff}} \) comes from the Goldstone boson associated to the imaginary part of the complex singlet scalar field \( S' \) [22]. The introduction of the extra scalar fields leads to the scalar potential

\[\mathcal{V}(H, S, \sigma_1, \sigma_2) = \mu_H^2 H^\dagger H + \lambda_H (H^\dagger H)^2 + \mu_S^2 S^\dagger S + \lambda_S (S^\dagger S)^2 + \mu_1^2 |\sigma_1^+|^2 + \mu_2^2 |\sigma_2^+|^2
\]
\[+ [\mu_3 \sigma_1^+ \sigma_2^- S + \text{h.c.}] + \lambda_1 |\sigma_1^+|^4 + \lambda_2 |\sigma_2^-|^4 + \lambda_5 (H^\dagger H)(S^\dagger S) + \lambda_4 |\sigma_1^+|^2 |\sigma_2^-|^2
\]
\[+ \lambda_6 (H^\dagger H)|\sigma_1^+|^2 + \lambda_7 (H^\dagger H)|\sigma_2^-|^2 + \lambda_8 (S^\dagger S)|\sigma_1^+|^2 + \lambda_9 (S^\dagger S)|\sigma_2^-|^2.\]  

(4)

We assume \( \mu_H^2 < 0 \) and \( \mu_S^2 < 0 \) and \( \text{Im}(\mu_3) = 0 \). Moreover, we assume \( \lambda_3 \ll 1 \) such that the scalar \( S \) and \( H \) do not mix, allowing us to identify the CP even scalar particle in \( H \) as the SM Higgs boson. To establish the scalar spectrum, we expand the scalar fields as

\[H = \left( \begin{array}{cc} 0 \\ \frac{1}{\sqrt{2}} (h + v_H) \end{array} \right), \quad S = \frac{1}{\sqrt{2}} (S_R + v_S),\]

\( ^4\)Notice that by performing a field redefinition on \( \sigma_1 \) it is possible to make one of such couplings real.

\( ^5\)If only one charged scalar is added then the right-handed and left-handed neutrinos would have the same charge under the \( B - L \) symmetry, thus rendering the Yukawa interaction term \( \bar{L}_\alpha H \nu_{R3} \) allowed.
with $v_H = 246.22$ GeV. Of the original ten scalar degrees of freedom in the model, the gauge bosons $W^\pm$, $Z^0$ and $Z'$ absorb four of them. Thus, the scalar spectrum contains two neutral CP-even states ($h$ and $S_R$) and two charged scalars ($s^\pm_1$ and $s^\pm_2$). Mass eigenstates for new charged scalars are defined through the mixing angle $\varphi$ as

$$
\begin{pmatrix}
    s^\pm_1 \\
    s^\pm_2
\end{pmatrix} = \begin{pmatrix}
    \cos \varphi & -\sin \varphi \\
    \sin \varphi & \cos \varphi
\end{pmatrix}
\begin{pmatrix}
    \sigma^\pm_1 \\
    \sigma^\pm_2
\end{pmatrix},
$$

where the mixing angle is given by $\sin(2\varphi) = \sqrt{2}\mu_3 v_S/(m^2_{s_2} - m^2_{s_1})$. By defining $m^2_{\sigma_1} = \mu_1^2 + \lambda_5 v_H^2/2 + \lambda_7 v_S^2/2$, $m^2_{\sigma_2} = \mu_2^2 + \lambda_6 v_H^2/2 + \lambda_8 v_S^2/2$, their masses take the form

$$
m^2_{\sigma_1,\sigma_2} = \frac{1}{2} \left( m^2_{\sigma_2} + m^2_{\sigma_1} \pm \sqrt{(m^2_{\sigma_2} - m^2_{\sigma_1})^2 + 8\mu_3^2 v_S^2} \right),
$$

with $s^\pm_1$ been the lightest one. Since the effect of $\lambda_{5,6,7,8}$ on $m^2_{\sigma_i}$ can be absorbed by re-scaling the $\mu_i^2$ parameters, hereafter we further assume $\lambda_{5,6,7,8} \ll 1$.

### 2.1 Radiative neutrino masses

The interplay of the new Yukawa interactions and the trilinear interaction $\mu_3$ generates Dirac neutrino masses at one loop level as displayed in Fig. 1. The expression for the corresponding mass matrix can be cast as

$$
[M_\nu]_{\alpha i} = \kappa [f^T]_{\alpha \beta} [M_\ell]_{\beta \beta} h_{\beta i},
$$

where $M_\ell$ is the diagonal mass matrix of the charged leptons and

$$
\kappa = \frac{\sin(2\varphi)}{16\pi^2} \ln \frac{m^2_{s_2}}{m^2_{s_1}}.
$$

$M_\nu$ is diagonalized by the biunitary transformation

$$
\left( U_L^{(\nu)} \right)^\dagger M_\nu U_R^{(\nu)} = M_\nu^{\text{diag}} = \text{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3}),
$$

![Figure 1: B-L flux in the one-loop diagram leading to Dirac neutrino masses.](image)
where \( U_R^{(\nu)} \) and \( U_L^{(\nu)} \) are rotation \( 3 \times 3 \) matrices associated to the right-handed and left-handed neutrinos, respectively. Assuming \( U_R^{(\nu)} = 1 \) the lepton mixing matrix takes the form

\[
U_{PMNS} = U_L^{(\nu)}.
\]  

(10)

Due to the antisymmetric character of the Yukawa coupling matrix \( f \) one neutrino state remains massless, and from Eq. (7) one can express six of the nine non-zero Yukawa-couplings in terms of the neutrino oscillation observables, the masses of the charged leptons and scalars. The expressions for the non-zero couplings in the case of a normal neutrino mass hierarchy are:

\[
h_{22} = -\frac{(U_{13}U_{32} - U_{12}U_{33}) (f_{13}h_{12}\kappa m_e - m_2U_{32})}{f_{13}\kappa (U_{23}U_{32} - U_{22}U_{33}) m_\mu},
\]  

(11)

\[
h_{23} = \frac{(U_{12}U_{33} - U_{13}U_{32}) (f_{13}h_{13}\kappa m_e - m_3U_{33})}{f_{13}\kappa (U_{23}U_{32} - U_{22}U_{33}) m_\mu},
\]  

(12)

\[
h_{32} = \frac{f_{13}h_{12}\kappa m_e(U_{13}U_{22} - U_{12}U_{23}) + m_2U_{22}(-U_{13}U_{32} + U_{12}U_{33})}{f_{13}\kappa (U_{23}U_{32} - U_{22}U_{33}) m_\tau},
\]  

(13)

\[
h_{33} = \frac{f_{13}h_{13}\kappa m_e(U_{13}U_{22} - U_{12}U_{23}) + m_3U_{23}(-U_{13}U_{32} + U_{12}U_{33})}{f_{13}\kappa (U_{23}U_{32} - U_{22}U_{33}) m_\tau},
\]  

(14)

\[
f_{12} = \frac{f_{13} (U_{13}U_{22} - U_{12}U_{23})}{U_{13}U_{32} - U_{12}U_{33}},
\]  

(15)

\[
f_{23} = \frac{f_{13} (U_{23}U_{32} - U_{22}U_{33})}{U_{13}U_{32} - U_{12}U_{33}},
\]  

(16)

with \( U_{ij} = (U_{PMNS})_{ij} \), \( m_2 = \sqrt{\Delta m_{sol}^2} \) and \( m_3 = \sqrt{\Delta m_{atm}^2} \). The expressions for the case of an inverted neutrino mass hierarchy (III) are:

\[
h_{22} = \frac{(U_{11}U_{32} - U_{12}U_{31}) (f_{13}h_{12}\kappa m_e - m_2U_{32})}{f_{13}\kappa (U_{22}U_{31} - U_{21}U_{32}) m_\mu},
\]  

(17)

\[
h_{32} = \frac{f_{13}h_{12}\kappa m_e(U_{12}U_{21} - U_{11}U_{22}) + m_2U_{22}(-U_{12}U_{31} + U_{11}U_{32})}{f_{13}\kappa (U_{22}U_{31} - U_{21}U_{32}) m_\tau},
\]  

(18)

\[
h_{21} = -\frac{(U_{12}U_{31} - U_{11}U_{32}) (f_{13}h_{11}\kappa m_e - m_1U_{31})}{f_{13}\kappa (U_{22}U_{31} - U_{21}U_{32}) m_\mu},
\]  

(19)

\[
h_{31} = \frac{f_{13}h_{11}\kappa m_e(U_{12}U_{21} - U_{11}U_{22}) + m_1U_{21}(-U_{12}U_{31} + U_{11}U_{32})}{f_{13}\kappa (U_{22}U_{31} - U_{21}U_{32}) m_\tau},
\]  

(20)

\[
f_{12} = \frac{f_{13} (U_{12}U_{21} - U_{11}U_{22})}{U_{12}U_{31} - U_{11}U_{32}},
\]  

(21)

\[
f_{23} = \frac{f_{13} (U_{22}U_{31} - U_{21}U_{32})}{U_{12}U_{31} - U_{11}U_{32}},
\]  

(22)

with \( m_1 = \sqrt{\Delta m_{atm}^2} \) and \( m_2 = \sqrt{\Delta m_{sol}^2 + \Delta m_{atm}^2} \).

Inserting the best fit values \([39]\) for the mixing angles and squared mass differences along with
$\delta_{\text{CP}} = \pi(0)$ for a NH (IH), the non-free $f_{\alpha\beta}$ couplings show a mild hierarchy,

\begin{align}
  f_{12} &\approx 1.81 f_{13}, \quad f_{23} \approx 2.93 f_{13}, \quad \text{for NH}; \\
  f_{12} &\approx -0.88 f_{13}, \quad f_{23} \approx -0.20 f_{13}, \quad \text{for IH}. 
\end{align}

(23)\hspace{1cm}(24)

On the other hand, from the expressions for the non-free $h_{\beta i}$ Yukawa couplings it is highly expected that those associated to the tau ($h_{3i}$) turn to be suppressed in comparison to the ones associated to the muon ($h_{2i}$), and these in turn suppressed in comparison to the free $h_{\beta i}$ parameters ($h_{12}$ and $h_{13}$ for NH and $h_{11}$ and $h_{12}$ for IH). Introducing the benchmark point $m_{s1} = m_{s2}/2 = 300$ GeV, $\sin(2\varphi) = 0.5$, $f_{13} = 0.01$ and $h_{12} = h_{13} = 0.1 (h_{11} = h_{12} = 0.1)$ along the previous inputs, the non-free $h_{\beta i}$ couplings become

\begin{align}
  h_{22} &\approx -0.00027, \quad h_{23} \approx -0.00027, \quad h_{32} \approx 0.00001, \quad h_{33} \approx 0.00001, \\
  f_{12} &\approx 0.00738, \quad f_{23} \approx 0.01761, \quad \text{for NH}; \\
  h_{22} &\approx 0.00243, \quad h_{32} \approx 0.00013, \quad h_{21} \approx 0.00238, \quad h_{31} \approx 0.00012, \\
  f_{12} &\approx -0.00876, \quad f_{23} \approx -0.00202, \quad \text{for IH}. 
\end{align}

(25)\hspace{1cm}(26)

These results illustrate the strong hierarchy present in the $h_{\beta i}$ couplings.

Here a comment is in order concerning the leptonic CP phase, $\delta_{\text{CP}}$. Since only one $f_{ij}$ Yukawa coupling -say $f_{13}$- can be rendered real, it follows from the expression for $f_{12}$ and $f_{23}$ that the CP symmetry is conserved in the neutrino sector only in the scenario when those couplings are real. That is, such a scenario demands that the Dirac phase necessarily takes the values $\delta_{\text{CP}} = 0, \pi$. This scenario is in agreement to a large extent with the most recent analysis [39, 40] which give a best fit for the complex phase close to $\delta_{\text{CP}} = \pi$ for a normal hierarchy (NH), whereas $\delta_{\text{CP}} = 0$ is marginally allowed within the $3\sigma$ range for a inverted hierarchy (IH). This in turn makes such a scenario to be falsifiable at the upcoming neutrino experiments such as DUNE [7] and Hyper-Kamiokande [8].

### 2.2 LFV processes

The same Yukawa interactions involved in the neutrino mass generation also induce charged lepton flavor violating processes such as $\ell_i \to \ell_j \gamma$, $\ell_i \to 3\ell_j$ and $\mu - e$ conversion in nuclei. In this model
such processes are generated at one-loop level and are mediated by the charged scalars $s^+_k$ and neutrinos. The expressions for the branching ratios $B(\ell_i \to \ell_j \gamma)$ are given by

$$B(\mu \to e\gamma) = \frac{\alpha e Br(\mu \to ev_\mu \bar{v}_e)}{768\pi G_F^2} \left[16C_{\varphi 12}^2|f_{13}|^2|f_{23}|^2 + C_{\varphi 21}^2(|h_{12}h_{22} + h_{13}h_{23}|^2)\right],$$

(27)

$$B(\tau \to e\gamma) = \frac{\alpha_\tau Br(\tau \to ev_\tau \bar{v}_e)}{768\pi G_F^2} \left[16C_{\varphi 12}^2|f_{12}|^2|f_{23}|^2 + C_{\varphi 21}^2(|h_{12}h_{32} + h_{13}h_{33}|^2)\right],$$

(28)

$$B(\tau \to \mu\gamma) = \frac{\alpha_\tau Br(\tau \to \mu v_\mu \bar{v}_\mu)}{768\pi G_F^2} \left[16C_{\varphi 12}^2|f_{12}|^2|f_{13}|^2 + C_{\varphi 21}^2(|h_{22}h_{32} + h_{23}h_{33}|^2)\right],$$

(29)

for the case of a NH, and

$$B(\mu \to e\gamma) = \frac{\alpha e Br(\mu \to ev_\mu \bar{v}_e)}{768\pi G_F^2} \left[16C_{\varphi 12}^2|f_{13}|^2|f_{23}|^2 + C_{\varphi 21}^2(|h_{11}h_{21} + h_{12}h_{22}|^2)\right],$$

(30)

$$B(\tau \to e\gamma) = \frac{\alpha_\tau Br(\tau \to ev_\tau \bar{v}_e)}{768\pi G_F^2} \left[16C_{\varphi 12}^2|f_{12}|^2|f_{33}|^2 + C_{\varphi 21}^2(|h_{11}h_{31} + h_{12}h_{32}|^2)\right],$$

(31)

$$B(\tau \to \mu\gamma) = \frac{\alpha_\tau Br(\tau \to \mu v_\mu \bar{v}_\mu)}{768\pi G_F^2} \left[16C_{\varphi 12}^2|f_{12}|^2|f_{13}|^2 + C_{\varphi 21}^2(|h_{21}h_{31} + h_{22}h_{32}|^2)\right],$$

(32)

for a IH. Here we have defined $C_{\varphi ij} = \cos \varphi^2/m^2_{zi} + \sin \varphi^2/m^2_{zj}$. Notice that, in contrast to neutrino masses, these branching ratios are not suppressed by the scalar mixing angle $\varphi$, as is expected since the rare decays $\ell_i \to \ell_j \gamma$ do not depend on the massiveness of neutrinos.

As regards to the $\ell_i \to \ell_j \ell_j \ell_j$ processes there exist two main diagrams (see Fig. 2) contributing to the total amplitude: the photon- and $Z$-penguin diagrams (left panel) and the box diagram (right panel). The contribution of the Higgs-penguin diagram is subdominant due to the suppression coming from the Yukawa couplings associated to the first two families of charged leptons (the contribution of the processes involving tau leptons is not negligible but the corresponding limits are less restrictive) whereas the $Z'$-penguin contribution is also subdominant due to the large mass of the $Z'$ gauge boson required to surpass the LHC lower bound (see next section). The $\mu - e$ conversion diagrams are obtained when the pair of lepton lines attached to the photon and $Z$ boson in the penguin diagrams are replaced by a pair of light quark lines. The corresponding amplitudes for the $\ell_i \to \ell_j \ell_j \ell_j$, $\ell_i \to \ell_j \bar{\ell}_j \ell_k$ and $\mu - e$ processes are calculated by using the FlavorKit code [41].

Finally, when performing our numerical analysis in Section 4 we will consider the current experimental bounds and their future expectations shown in Table 2.

### 2.3 LHC observables

LHC searches for additional neutral gauge bosons have delivered stringent constraints on the fraction of the $Z'$ mass to the gauge coupling, $M_{Z'}/g'$. The recasting [51] of the latest ATLAS results for the search of dilepton resonances using 139 fb$^{-1}$ [52] for the $U(1)_{B-L}$ model gives

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6Higgs-penguin diagrams are again suppressed, in this case by the Yukawa couplings of light quarks.
Future sensitivity given by subsequent decays into a lepton and a neutrino involve precisely these Yukawa interactions. 

| Observable | Present limit | Future sensitivity |
|------------|---------------|--------------------|
| $\mathcal{B}(\mu \rightarrow e\gamma)$ | $4.2 \times 10^{-13}$ | $6.3 \times 10^{-14}$ |
| $\mathcal{B}(\tau \rightarrow e\gamma)$ | $3.3 \times 10^{-8}$ | $3 \times 10^{-9}$ |
| $\mathcal{B}(\tau \rightarrow \mu\gamma)$ | $4.4 \times 10^{-8}$ | $3 \times 10^{-9}$ |
| $\mathcal{B}(\mu \rightarrow eee)$ | $1.0 \times 10^{-12}$ | $10^{-16}$ |
| $\mathcal{B}(\tau \rightarrow eee)$ | $2.7 \times 10^{-8}$ | $3 \times 10^{-9}$ |
| $\mathcal{B}(\tau \rightarrow \mu\mu\mu)$ | $2.1 \times 10^{-8}$ | $10^{-9}$ |
| $\mathcal{B}(\tau \rightarrow e\mu\mu)$ | $2.7 \times 10^{-8}$ | $-$ |
| $\mathcal{B}(\tau \rightarrow \mu\mu\epsilon)$ | $1.8 \times 10^{-8}$ | $-$ |
| $R_{\mu e}(\text{Ti})$ | $4.3 \times 10^{-12}$ | $10^{-18}$ |
| $R_{\mu e}(\text{Au})$ | $7.3 \times 10^{-13}$ | $-$ |

Table 2: Current bounds [42–49] and projected sensitivities [9–11, 50] for the charged LFV observables.

$M_{Z'} \gtrsim 5$ TeV for $g' \sim 0.1$. Let us recall that the mass of the $Z'$ boson is given by $M_{Z'} = 3g'v_s$, which in turn implies $v_s \gtrsim 17$ TeV for $g' \sim 0.1$.

Regarding collider bounds on the charged scalar $s^\pm$, its main production mechanism at the LHC for small and intermediate Yukawa couplings ($f_{\alpha\beta}, h_{\beta i} \lesssim 0.1$) is via Drell-Yan processes. The subsequent decays into a lepton and a neutrino involve precisely these Yukawa interactions.

The expressions for decay width of the processes $\Gamma(s_1 \rightarrow \ell \sum_i \nu_i) \equiv \Gamma(s_1 \rightarrow \ell\nu)$ for NH are given by

\[
\Gamma(s_1 \rightarrow e\nu) = \frac{m_{s_1}}{16\pi} (4 \cos \varphi^2(|f_{12}|^2 + |f_{13}|^2) + \sin \varphi^2(|h_{12}|^2 + |h_{13}|^2)) ,
\]

\[
\Gamma(s_1 \rightarrow \mu\nu) = \frac{m_{s_1}}{16\pi} (4 \cos \varphi^2(|f_{12}|^2 + |f_{23}|^2) + \sin \varphi^2(|h_{22}|^2 + |h_{23}|^2))
\]

\[
\Gamma(s_1 \rightarrow \tau\nu) = \frac{m_{s_1}}{16\pi} (4 \cos \varphi^2(|f_{13}|^2 + |f_{23}|^2) + \sin \varphi^2(|h_{32}|^2 + |h_{33}|^2)) ,
\]

whereas for IH they read

\[
\Gamma(s_1 \rightarrow e\nu) = \frac{m_{s_1}}{16\pi} (4 \cos \varphi^2(|f_{12}|^2 + |f_{13}|^2) + \sin \varphi^2(|h_{11}|^2 + |h_{12}|^2)) ,
\]

\[
\Gamma(s_1 \rightarrow \mu\nu) = \frac{m_{s_1}}{16\pi} (4 \cos \varphi^2(|f_{12}|^2 + |f_{23}|^2) + \sin \varphi^2(|h_{21}|^2 + |h_{22}|^2))
\]

\[
\Gamma(s_1 \rightarrow \tau\nu) = \frac{m_{s_1}}{16\pi} (4 \cos \varphi^2(|f_{13}|^2 + |f_{23}|^2) + \sin \varphi^2(|h_{31}|^2 + |h_{32}|^2)) .
\]

These direct leptonic decays lead to the collider signature of dileptons plus missing transverse momentum since

\[
pp \rightarrow \gamma^*/Z^* \rightarrow s^+_is^-i \rightarrow \ell^+_i \ell^-_j \nu
\]

\footnote{Notice that the decays involving gauge bosons are not present.}
Table 3: Current bounds and projected sensitivities for GNI [32].

Such a signature is analogous to the one of electroweak production of sleptons in the context of simplified supersymmetric scenarios [53, 54]. Assuming a 100% branching ratio into $\ell = e, \mu$ and an integrated luminosity of 139 fb$^{-1}$ ATLAS has set upper limits on the slepton pair production cross sections in such a way right-handed slepton masses are excluded up to 420 GeV for vanishing neutralino masses [53]. A recast of the excluded cross section for slepton pair production was done in Ref. [55], which allows to exclude $SU(2)_L$-singlet charged scalars decaying into a lepton plus a neutrino with masses up to 200 GeV (or even more in some cases). Moreover, the analysis in Ref. [56]$^8$ shows that singly charged scalars with masses below 500 GeV decaying mostly into electrons and muons may be excluded at the high luminosity phase of the LHC.

On the other hand, single production of charged scalars can also take place (for large Yukawa couplings) through the radiation from a lepton external leg in $s$-channel diagrams featuring a gauge boson ($\gamma, Z$ or $W$). The collider signatures for these topologies are one, two, or three leptons plus missing transverse momentum. Concerning to the final state with three leptons ($e$ or $\mu$), a recasting [59] of the ATLAS trilepton searches [60] excludes charged scalar not decaying into taus with masses $\sim 190$ GeV and Yukawa couplings of $\mathcal{O}(1)$.

3 General neutrino interactions in Zee Dirac model

The $h_{\beta i}$ and $f_{\alpha \beta}$ interactions also induce new effective four-fermion interactions between neutrinos and charged leptons with all the Lorentz-invariant structures: scalar ($S$), pseudoscalar ($P$), vector ($V$), axialvector ($A$) and tensor ($T$)$^9$. These new interactions can modify the SM prospect for the neutrino-electron elastic scattering, and therefore be constrained [32, 33, 61] by using the solar neutrinos experiment results such as Borexino [62], TEXONO [63] and CHARM-II [64].

$^8$See Refs. [57, 58] for related works.

$^9$There are also effective four-fermion interactions mediated by the $Z'$ boson, however, they turn to be heavily suppressed by $M_{Z'}$. 

| $\Gamma_{\mu\mu}/\Gamma_{\tau\tau}$ | $\Gamma_{\tau\tau}/\Gamma_{\tau\tau}$ | $\Gamma_{\mu\mu}/\Gamma_{\tau\tau}$ | $\Gamma_{\mu\mu}/\Gamma_{\tau\tau}$ |
|----------------|----------------|----------------|----------------|
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
| $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ | $\epsilon_{V_{ee}}$ | $\epsilon_{A_{ee}}$ |
The most general effective Lagrangian that describes the GNI can be cast as:

\[
\mathcal{L} = \frac{G_F}{\sqrt{2}} \sum_{a=S,P,V,A,T} (\bar{\nu}_a \Gamma^a \nu_\beta) \left[ \bar{\Gamma}^a \left( \epsilon_{\alpha\beta}^a + \epsilon_{\alpha\beta}^a i \gamma^5 \right) \Gamma^a \right],
\]

where \(\{\Gamma^S, \Gamma^P, \Gamma^V, \Gamma^A, \Gamma^T\} \equiv \{I, i\gamma^5, \gamma^\mu, \gamma^\mu \gamma^5, \sigma^{\mu\nu}\}\), \(\{i^S, i^P, i^V, i^A,i^T\} \equiv \{i, i, 1, 1, i\}\) and \(G_F\) is the Fermi constant. The 3 \(\times\) 3 matrices \(\epsilon\) and \(\tilde{\epsilon}\) -which are assumed to be real- parametrize the departures of the SM result.

On the other hand, the effective four-fermion Lagrangian for GNI contributions in the present model is given by

\[
\mathcal{L}_{\text{eff}} = -2 f_{\alpha\rho} f_{\beta\sigma}^* \left[(\bar{\ell}_\beta \gamma^\mu P_L \ell_\alpha)(\bar{\nu}_\sigma \gamma_\mu P_L \nu_\rho)\right] C_{\phi12} - \frac{1}{2} h_{\alpha\rho} h_{\beta\sigma}^* \left[(\bar{\ell}_\beta \gamma^\mu P_R \ell_\alpha)(\bar{\nu}_\sigma \gamma_\mu P_R \nu_\rho)\right] C_{\phi21}
\]

\[
+ f_{\alpha\rho} h_{\beta\sigma}^* \left[(\bar{\ell}_\beta P_L \ell_\alpha)(\bar{\nu}_\sigma P_L \nu_\rho) - \frac{1}{4} \bar{\ell}_\beta \sigma_{\mu\nu} P_L \ell_\alpha)(\bar{\nu}_\sigma \sigma_{\mu\nu} P_L \nu_\rho)\right] D_{\phi12}
\]

\[
+ f_{\beta\sigma}^* h_{\alpha\rho} \left[(\bar{\ell}_\beta P_R \ell_\alpha)(\bar{\nu}_\sigma P_R \nu_\rho) - \frac{1}{4} \bar{\ell}_\beta \sigma_{\mu\nu} P_R \ell_\alpha)(\bar{\nu}_\sigma \sigma_{\mu\nu} P_R \nu_\rho)\right] D_{\phi21},
\]

where \(D_{\phi ij} = \cos \varphi \sin (1/m^2_{s_j} - 1/m^2_{s_i})\). It follows that the factors in the second and third lines in the previous expression become suppressed for a small scalar mixing angle, while the two factors in the first line are suppressed for small Yukawa couplings correspondingly.

For electron neutrinos \(\sigma = \rho = e\) the corresponding GNI read

\[
\epsilon_{ee}^V = \tilde{\epsilon}_{ee}^V = \epsilon_{ee}^A = \tilde{\epsilon}_{ee}^A = -\frac{\sqrt{2}}{8G_F} |h_{11}|^2 C_{\phi21},
\]

\[
\epsilon_{ee}^S = \tilde{\epsilon}_{ee}^S = \epsilon_{ee}^P = \tilde{\epsilon}_{ee}^P = \epsilon_{ee}^T = \tilde{\epsilon}_{ee}^T = 0.
\]

Notice that all GNI associated to electron neutrinos for the NH are zero since \(h_{11} = 0\). For muon neutrinos \(\sigma = \rho = \mu\) it follows that

\[
\epsilon_{\mu\mu}^V = \epsilon_{\mu\mu}^A = -\frac{\sqrt{2}}{8G_F} (|h_{12}|^2 C_{\phi21} + 4 |f_{12}|^2 C_{\phi12})
\]

\[
\epsilon_{\mu\mu}^S = -\epsilon_{\mu\mu}^P = -4 \epsilon_{\mu\mu}^T = \frac{\sqrt{2} D_{\phi12}}{4G_F} (f_{12} h_{12}^* + f_{12}^* h_{12}),
\]

\[
\epsilon_{\mu\mu}^S = -\epsilon_{\mu\mu}^P = -4 \epsilon_{\mu\mu}^T = \frac{i \sqrt{2} D_{\phi12}}{4G_F} (f_{12} h_{12}^* - f_{12}^* h_{12}).
\]

\[^{10}\text{We closely follow the notation introduced in Ref. [32]}\]

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Figure 3: General neutrino interactions for the IH case.

For the case of tau neutrinos $\sigma = \rho = \tau$ the GNI are

$$\epsilon_{\tau\tau}^V = \epsilon_{\tau\tau}^A = -\frac{\sqrt{2}}{8G_F} (|h_{13}|^2C_{\varphi21} + 4|f_{13}|^2C_{\varphi12}),$$  \hspace{1cm} (48)

$$\tilde{\epsilon}_{\tau\tau}^V = \tilde{\epsilon}_{\tau\tau}^A = -\frac{\sqrt{2}}{8G_F} (|h_{13}|^2C_{\varphi21} - 4|f_{13}|^2C_{\varphi12}),$$  \hspace{1cm} (49)

$$\epsilon_{\tau\tau}^S = -\epsilon_{\tau\tau}^P = -4\epsilon_{\tau\tau}^T = \frac{\sqrt{2}D_{\varphi12}}{4G_F} (f_{13}h_{13}^* + f_{13}^*h_{13}),$$  \hspace{1cm} (50)

$$\tilde{\epsilon}_{\tau\tau}^S = -\tilde{\epsilon}_{\tau\tau}^P = -4\tilde{\epsilon}_{\tau\tau}^T = \frac{i\sqrt{2}D_{\varphi12}}{4G_F} (f_{13}h_{13}^* - f_{13}^*h_{13}).$$  \hspace{1cm} (51)

Because $h_{i3} = 0$ in a inverted neutrino mass ordering these GNI reduce to

$$\epsilon_{\tau\tau}^V = \epsilon_{\tau\tau}^A = -\tilde{\epsilon}_{\tau\tau}^V = \tilde{\epsilon}_{\tau\tau}^A = -\frac{\sqrt{2}}{2G_F} |f_{13}|^2C_{\varphi12},$$  \hspace{1cm} (52)

$$\epsilon_{\tau\tau}^S = \epsilon_{\tau\tau}^P = \epsilon_{\tau\tau}^T = \epsilon_{\tau\tau}^S = \epsilon_{\tau\tau}^P = \epsilon_{\tau\tau}^T = 0.$$  \hspace{1cm} (53)
Figure 4: General neutrino interactions for the NH case.

By using the Borexino results, the authors of Ref. [32] have derived 90% C.L. constraints on general neutrino interactions for all neutrino flavors. The most constraining bounds that apply to the current model are displayed in Table 3.

4 Results and discussion

We have performed a random scan for each neutrino mass ordering, varying just a subset of the free parameters of the model with the aim of making the analysis simpler. The set of free parameters of the model relevant for our analysis has been varied as

\begin{align}
10^{-5} & \leq |f_{13}| \leq 3 ; \\
0.1 & \leq |h_{12}, h_{13} h_{11}, h_{12}| \leq 3 ; \\
10^{-6} & \leq \phi \leq \pi / 2 ; \\
80 \text{GeV} & \leq m_{s1} \leq 500 \text{GeV} ; m_{s2} = [m_{s1}, 1000 \text{GeV}].
\end{align}
Figure 5: LFV processes for IH (left) and NH (right). The dotted and dashed lines represent the sensitivity limit expected for the future searches for $R_{\mu e}(\text{Ti})$ and $\mathcal{B}(\mu \to 3e)$, respectively.

The magnitude of the non-free Yukawa couplings are restricted to be within the range $[10^{-5}, 3]$, and the other parameters of the scalar potential were fixed as $\lambda_i = 10^{-4}, i = 1, ..., 8$. The parameters of the new gauge boson are set to $M_{Z'} = 6$ TeV and $g' = 0.5$ in such a way $\nu_S = 4$ TeV.\footnote{With these values it follows that $\mu_3 < 200$ GeV, in order to obtain charged scalars below the TeV mass scale.} In our numerical analysis, all the viable benchmark points satisfy the current neutrino oscillation data within the 3$\sigma$ level \cite{40}, both for inverted and normal hierarchies, with $\delta_{\text{CP}} = \pi(0)$ for a NH (IH). Likewise, they satisfy the LFV upper bounds reported in Table 2 and the GNI constraints in Table 3. Finally, we impose the collider bounds on the charged scalar coming from searches for final states with an oppositely charged lepton pair ($e^-e^+$ and $\mu^-\mu^+$) and missing transverse energy \cite{55}. In order to obtain the particle spectrum and low energy observables we have used SPheno \cite{65, 66} and the FlavorKit \cite{41} of SARAH \cite{67, 68}.

In the Fig. 3 (4) are displayed the results for the general neutrino interactions assuming a inverted (normal) neutrino mass hierarchy, and the future sensitivity (dashed line) that may be reached in future solar neutrino experiments \cite{32}, with the color bar denoting the mass of the heaviest charged scalar. Only the results for the parameters $\tilde{\epsilon}_{ee}, \tilde{\epsilon}_{\mu\mu}, \epsilon_{ee}$ and $\epsilon_{\mu\mu}$ ($\tilde{\epsilon}_{\mu\tau}, \epsilon_{\mu\tau}$ and $\epsilon_{\tau\tau}$) are displayed since they are the ones with the most appreciable restrictions on future experiments. As is expected, the heavier the $s_2^\pm$, the lower the intensity of the general neutrino interactions.

The correlation between the most constrained LFV processes, namely $\mathcal{B}(\mu \to e\gamma), \mathcal{B}(\mu \to e\gamma)$ and $\mu - e$ conversion, is shown in the top panels of the Fig. 5. The corresponding prospects for the future sensitivities (dashed and dotted lines) point out that a large portion of the parameter space will be explored. On the other hand, the current limits impose mild restrictions on the Yukawa couplings, in particular $\mathcal{B}(\mu \to e\gamma)$ establishes that $f_{13}$ is restricted to values less than 0.04, furthermore, future searches for these processes restrict the value of $f_{13}$ around to $10^{-3}$. Finally, the Fig. 6 shows the results of the search for charged scalars in the LHC. The current limits impose that the mass of the lightest scalar must be greater than $\sim 260$ GeV and $\sim 300$ GeV.
for normal and inverse hierarchy, respectively [55]. The dashed line shows the future exclusion limits for an integrated luminosity of 300 fb\(^{-1}\) and the solid line for 3 ab\(^{-1}\) [56]. Notice that in principle it is possible to explore the whole parameter space in the high luminosity phase of the LHC.

Since the relevant GNI parameters and the decay rates for \(s_1 \rightarrow \ell \nu\) depend only on the magnitude of the Yukawa couplings it is understandable to expect that the above results hold in the scenario where the Yukawa couplings are complex. In order to check this we have performed an additional scan for the scenario where the Dirac CP phase vary within the 3σ level [40]. The results from this scan concerning the magnitude of the GNI interactions and the reach of LHC searches do not differ from those obtained in the scenario with CP conservation. Moreover, the LFV observables continue exhibiting a strong correlation as the one displayed in Fig. 5.

## 5 Summary

The simplest model leading to Dirac neutrino masses at one loop level is obtained by adding a pair of charged scalars to the SM along with three right-handed neutrinos. In this work we have assumed that the Diracness of the neutrinos is protected by only one extra \(U(1)_{B-L}\) gauge symmetry, with the anomalies canceled by the SM leptons and the three right-handed neutrinos. We studied the interesting phenomenological features of the model such as the neutrino mass generation, charged LFV processes and signatures at the LHC associated to the charged scalars and the new \(B-L\) gauge boson. Furthermore, we reported the expressions for the general neutrino-electron interactions and identified the regions of parameter space that may be explored in future solar neutrino experiments. Remarkably, future searches for charged scalars at the LHC will probe the entire parameter space considered in this analysis.
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