Natural Seesaw Realization at the TeV Scale

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We summarize the phenomenological constraints on seesaw scenarios defined at the TeV scale and provide a simple extension of the Standard Model which naturally leads to a testable mechanism of neutrino mass generation.

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We review the most important phenomenological constraints on type I/inverse seesaw scenarios defined at the electroweak scale. We also consider a simple extension of the Standard Model (SM) which realizes a “testable” seesaw scenario, without imposing any fine-tuning of the neutrino Yukawa couplings in order to generate light active neutrino masses.

1. Phenomenological constraints on the TeV scale seesaw parameter space

We consider a phenomenological seesaw [1] extension of the SM with two heavy Majorana fermion singlets $\mathcal{N}_{1,2}$, which in principle can be tested in low energy and collider experiments for masses $M_{1,2} \sim (100 - 1000)$ GeV.

Following [2], $\mathcal{N}_{1,2}$ have charged-current (CC) and neutral-current (NC) interactions with the SM leptons and interact with the Higgs boson. They are given by

\begin{align}
\mathcal{L}_{CC}^N &= -\frac{g}{2\sqrt{2}} \ell^a \gamma^\alpha (RV)_{ik} (1 - \gamma_5) \mathcal{N}_k W^{\alpha} + \text{h.c.}, \\
\mathcal{L}_{NC}^N &= -\frac{g}{4c_w} \overline{\nu}_L \gamma^a (RV)_{ik} (1 - \gamma_5) \mathcal{N}_k Z^a + \text{h.c.}, \\
\mathcal{L}_H^N &= -\frac{g M_k}{4M_W} \overline{\nu}_L (RV)_{ik} (1 + \gamma_5) \mathcal{N}_k h + \text{h.c.}
\end{align}

The couplings $(RV)_{ik}$ ($\ell = e, \mu, \tau$ and $k = 1, 2$) arise from the mixing between heavy and light Majorana neutrinos and, therefore, are suppressed by the seesaw scale. They can be conveniently parametrized as follows [2]:

\begin{align}
\left|(RV)_{e1}\right|^2 &= \frac{1}{2} \frac{\mu^2}{m_1^2 - m_2^2} \left| \frac{m_3}{m_2 + m_3} U_{e3} + i \sqrt{m_2/m_3} U_{e2} \right|^2, \quad \text{for normal hierarchy,} \\
\left|(RV)_{\mu1}\right|^2 &= \frac{1}{2} \frac{\mu^2}{m_1^2 - m_2^2} \left| \frac{m_2}{m_1 + m_2} U_{\mu2} + i \sqrt{m_1/m_2} U_{\mu1} \right|^2, \quad \text{for inverted hierarchy,} \\
(RV)_{\ell 2} &= \pm i (RV)_{e1} \sqrt{\frac{M_1}{M_2}}, \quad \ell = e, \mu, \tau,
\end{align}

where $U$ denotes the PMNS neutrino mixing matrix and $\nu \simeq 174$ GeV. The relative mass splitting of the two heavy Majorana neutrinos must be very small, $|M_1 - M_2|/M_1 \ll 1$, due to the current upper limit on the effective Majorana mass probed in neutrinoless double beta decay experiments [2]. In this case, the flavour structure of the neutrino Yukawa couplings is fixed by the neutrino oscillation parameters [2,3] and the two heavy Majorana neutrinos form a pseudo-Dirac fermion. The parameter $\mu$ in the expressions above represents the largest eigenvalue of the matrix of the neutrino Yukawa couplings: $\mu^2 = 2M_1^2 (\left|(RV)_{e1}\right|^2 + \left|(RV)_{\mu1}\right|^2 + \left|(RV)_{\tau1}\right|^2)$.

Electroweak precision data (EWPD) provide an upper bound of the size of the Yukawa coupling $\mu$ for a given seesaw scale, that is $\mu \leq 0.06 (M_1/100$ GeV) [2]. However, the most stringent constraint on $\mu$ comes from lepton flavour violating observables, in particular from the present upper limit on $\mu^+ \rightarrow e^+ \gamma$ branching ratio reported by the MEG experiment [3]: $\text{BR}(\mu \rightarrow e \gamma) < 5.7 \times 10^{-13}$ at 90% confidence level. Indeed, in this case, taking the best fit values of the neutrino oscillation parameters [3], we get the bound: $\mu \lesssim 0.026$ for $M_1 = 100$ GeV.

We show in Fig. 1 all the relevant constraints on the effective couplings $(RV)_{\mu1}$ and $(RV)_{e1}$ which come from the requirement of reproducing neutrino oscillation data, from the EWPD bounds.
and the current upper limit on $\mu \to e\gamma$. The seesaw mass scale is fixed at benchmark value $M_1 = 100$ GeV. From Fig. 1, it is manifest that the allowed ranges of the right-handed (RH) neutrino couplings $|(RV)_{\mu 1}|$ and $|(RV)_{e 1}|$, in the case of normal (left panel) and inverted (right panel) light neutrino mass spectrum, are confined in a small strip of the overall representative plane. This corresponds to the scatter plot of the points which are consistent with the $3\sigma$ allowed ranges of the neutrino oscillation parameters [5]. The region of the parameter space which is allowed by the EWPD is marked with solid lines. The region allowed by the current bound on the $\mu \to e\gamma$ decay rate is indicated with a dashed line, while the dot-dashed line shows the exclusion limit for $\text{BR}(\mu \to e\gamma) < 10^{-14}$. The scatter points correspond to different values of the maximum neutrino Yukawa coupling $y$: $y = 0.001$ (blue ◊), $ii$) $y = 0.01$ (green +), $iii$) $y = 0.1$ (red ×), $iv$) $y = 1$ (orange ◊) and $v$) an arbitrary value of the Yukawa coupling $y \leq 1$ (cyan points).

As depicted in Fig. 1, in the case of a light neutrino mass spectrum with inverted hierarchy a strong suppression of the $\mu \to e\gamma$ decay rate is possible for specific values of the measured neutrino oscillation parameters. This is due to cancellations in the $\mu - e$ transition amplitude proportional to $|U_{\mu 2} + iU_{\mu 1}|$ which are possible in the case of neutrino mass spectrum with inverted hierarchy for $\sin \theta_{13} \gtrsim 0.13$ and CP conserving phases of the neutrino mixing matrix (see [2] for a detailed discussion).

2. A natural model realization of the TeV scale seesaw scenario

We discuss a simple model that naturally provides a testable seesaw scenario where the RH neutrino interactions in (1.1-1.2) are “sizable” and give rise to observable effects. We extend the scalar sector of the SM with an additional Higgs doublet, $H_2$, and a complex singlet, $\phi$. The fermion sector, instead, consists of the two RH neutrinos $N_{1,2}$, which allow to implement the seesaw mechanism in order to explain active neutrino masses and mixing. The resulting model Lagrangian is invariant under a global $U(1)_X$ symmetry, where $X$ corresponds to a generalization of the lepton
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| Field   | L_\alpha | e_{R\alpha} | N_1 | N_2 | H_1 | H_2 | \phi |
|---------|----------|-------------|-----|-----|-----|-----|------|
| SU(2)_L | 2        | 1           | 1   | 1   | 2   | 2   | 1    |
| U(1)_Y  | -1/2     | -1          | 0   | 0   | 1/2 | 1/2 | 0    |
| U(1)_X  | -1       | -1          | -1  | +1  | 0   | 2   | -2   |

Table 1: Charge assignment of the fields.

number. The particle content of the model and the charge assignment of the fields are reported in Table 1.

In this scenario, the presence of H_2 and \phi is motivated by the requirement of generating light neutrino masses through the type I/inverse/linear seesaw mechanism at the TeV scale. The most general scalar potential \mathcal{V}_{SB}, invariant under the SU(2)_L \times U(1)_Y \times [U(1)_X] symmetry, is derived in [6]. Given the fields in Table 1, we get

\begin{align}
\mathcal{V}_{SB} = -\mu_1^2 (H_1^\dagger H_1) + \lambda_1 (H_1^\dagger H_1)^2 - \mu_2^2 (H_2^\dagger H_2) + \lambda_2 (H_2^\dagger H_2)^2 - \mu_3^2 \phi^* \phi + \lambda_3 (\phi^* \phi)^2 \\
+ \kappa_{12} (H_1^\dagger H_1) (H_2^\dagger H_2) + \kappa_{13} (H_1^\dagger H_1) \phi^* \phi + \kappa_{23} (H_2^\dagger H_2) \phi^* \phi \\
- \frac{\mu'}{\sqrt{2}} \left( (H_1^\dagger H_2) \phi + (H_2^\dagger H_1) \phi^* \right).
\end{align}

The two SU(2)_L doublets H_{1,2} and the singlet \phi [3] take a non-zero vacuum expectation value (vev) v_{1,2} and v_{\phi}, respectively. In this case, the global U(1)_X is spontaneously broken down to a Z_2 discrete symmetry. The scalar mass spectrum of the model consists of: 1 charged scalar H^\pm, 3 CP-even neutral scalars h^0, H^0 and h_A, 2 pseudo-scalars A^0 and J. The latter is the Goldstone boson associated with the breaking of the global U(1)_X symmetry and is usually dubbed Majoron in theories with spontaneously broken lepton charge. Since it is a massless particle, strong constraints apply on its couplings to the SM fermions: a hierarchical pattern for the vevs of the scalar fields, namely v_2 \ll v_1, v_{\phi}, is required in order to satisfy the astrophysical constraints on the Majoron phenomenology. As discussed in detail in [6], a suppressed value of v_2 \propto \mu' is naturally realized from the minimization of the potential (2.1), due to the residual symmetries of the model.

In the limit of negligible v_2, the longitudinal gauge boson components are W_{L}^\pm \sim H_1^\pm and Z_L \sim \sqrt{2} \text{Im}(H_1^0), while the scalar mass eigenstates are to a good approximation: H^\pm \sim H_{2}^\pm, h_A \sim \sqrt{2} \text{Re}(H_2^0), \phi_0 \sim \sqrt{2} \text{Im}(H_2^0) and J \sim \sqrt{2} \text{Im}(\phi). Moreover, the two neutral scalars h^0 and H^0 arise from the mixing of \sqrt{2} \text{Re}(H_2^0) and \sqrt{2} \text{Re}(\phi). Typically, we have v_2 \lesssim 10 \text{ MeV} [3], v_1 \approx 246.2 \text{ GeV}, while v_{\phi} is free. Recalling that only H_1 has Yukawa couplings to SM fermions (cf. Table 1), h_A, A^0 and H^\pm couple to the SM sector only through gauge interactions and via the scalar quartic couplings, while h^0 and H^0 can have a priori sizable Yukawa couplings to SM fermions (see [3] for a discussion of the collider constraints on the scalar sector of the model).

Neutrino mass generation

We introduce for convenience a Dirac fermion field, N_D \equiv P_R N_1 + P_L N_2^C, where P_{L,R} are the usual chiral projectors and N_2^C \equiv C N_2^T is the conjugate of the N_2 RH neutrino field. The most
general interaction Lagrangian of \( N_D \) invariant under the global \( U(1)_X \) symmetry is

\[
\mathcal{L} \supset -m_N \overline{N_D} N_D - \left( Y^\beta_{\nu_1} \overline{N_D} H^*_1 L_\beta + Y^\gamma_{\nu_2} \overline{N_D} C \tilde{H}^*_2 L_\gamma + \frac{\delta_N}{\sqrt{2}} \phi \overline{N_D} N_D^C + \text{h.c.} \right)
\]  

(2.2)

where \( N_D^C \equiv C N_D^T \) and \( \tilde{H}_k \equiv -i \sigma_2 H^*_k \) \((k = 1, 2)\). The parameter \( \delta_N \) is made real through a phase transformation. The Yukawa interactions \( \propto Y_{\nu_1} (Y_{\nu_2}) \) couple \( N_1 (N_2) \) to the SM leptons. Therefore, after the Higgs doublets acquire a nonzero vev, the SM lepton number (i.e. the generalized \( X \) charge) is explicitly violated by \( Y_{\nu_2} \) mediated interactions. Furthermore, while the Dirac type mass \( m_N \) conserves the lepton number, the term proportional to \( \delta_N \) provides, after \( \phi \) takes a nonzero vev, a Majorana mass term for both \( N_1 \) and \( N_2 \). In the case \( m_N \gg \delta_N v_\phi \) we have a low energy realization of the type I/inverse seesaw scenario \([6]\).

In the chiral basis \((\nu_L, (N_1^C)_L, (N_2^C)_L)\), the \( 5 \times 5 \) symmetric neutrino mass matrix reads:

\[
\mathcal{M}_\nu = \begin{pmatrix}
0_{3 \times 3} & y_1 v_1 & y_2 v_2 \\
y_1^T v_1 & \delta_N v_\phi & m_N \\
y_2^T v_2 & m_N & \delta_N v_\phi \\
\end{pmatrix}.
\]

(2.3)

In the previous expression \( 0_{3 \times 3} \) denotes the null matrix of dimension 3 and we introduce the shorthand notation: \( y_k \equiv (Y^e_{\nu_k}, Y^\mu_{\nu_k}, Y^\tau_{\nu_k}) / \sqrt{2} \). The neutrino sector, therefore, consists of one massless neutrino, two massive light Majorana neutrinos and two heavy Majorana neutrinos \( \mathcal{M}_{1,2} \). The latter are quasi-degenerate, with masses \( M_{1,2} = m_N \mp \delta_N v_\phi \), and form a pseudo-Dirac pair for \( m_N \gg \delta_N v_\phi \) \([3]\). They have, therefore, naturally “sizable” CC and NC interactions with the SM leptons, eqs. (1.1) and (1.2), where in this case the mixing matrix elements \((RV)_{11}\) are proportional to the lepton number conserving Yukawa couplings \( Y^e_{\nu_k} \). The resulting effective light neutrino mass matrix is

\[
(M_\nu)^{\alpha\beta} \simeq -\frac{v_1 v_2}{m_N} \left( y_1^{\alpha} y_2^{\beta} + y_2^{\alpha} y_1^{\beta} \right) + v_\phi \delta_N \frac{v_1^2}{m_N} \left( y_1^{\alpha} y_1^{\beta} + y_2^{\alpha} y_2^{\beta} \frac{v_2^2}{v_1^2} \right).
\]

(2.4)

The first term in (2.4) acts as a linear seesaw contribution and its suppression originates from the small vev \( v_2 \). The second term is typical of inverse seesaw scenarios, where the small ratio \( v_\phi \delta_N/m_N \) suppresses the neutrino mass scale. Notice that, with only two RH neutrinos the linear seesaw contribution alone (i.e. neglecting \( v_\phi \) in (2.4)) allows to fit all current neutrino oscillation data, while if \( v_2 = 0 \) and \( v_\phi \neq 0 \) the inverse seesaw scenario can only account for one massive light neutrino. Therefore, the complex scalar field \( \phi \), with vev \( v_\phi \neq 0 \), is not mandatory in order to obtain two massive light neutrinos through the (linear) seesaw mechanism. On the other hand, \( v_\phi \neq 0 \) is a necessary condition to set a hierarchy between the Higgs doublet vevs, \( v_2 \ll v_1 \), without fine-tuning of the parameters \([8]\). Taking \( \mu_N \ll 1 \), the neutrino masses have a simple expression:

\[
m^\pm_\nu \simeq \frac{1}{m_N} \left( \sqrt{\frac{v_1^2}{v_2^2} - \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12})} \pm \sqrt{|y_{12}|^2 - \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12})} \right) \\
\simeq \frac{1}{m_N} \left( |y_1 y_2| \pm |y_{12}| \right) \times \left( 1 \mp \frac{\mu_N (y_1^2 + y_2^2) \text{Re}(y_{12})}{|y_1 y_2|^2 |y_{12}|} \right),
\]

(2.5)

with \( y_{12} \equiv \sqrt{y_1^2 y_2^2}, y_{12}^2 \equiv y_1^2 \cdot y_2^2, y_1 y_2 \equiv y_1 \cdot y_2, y_{12} \equiv y_1 \times y_2, v_1 v_2 \), \( \eta_{12} \equiv y_1 \times y_2 v_1 v_2 \) and \( \mu_N = (\delta_N v_\phi)/m_N \). Notice that if the neutrino Yukawa vectors \( y_1 \) and \( y_2 \) are proportional, \( m_\nu \) is exactly zero. For a normal hierarchical
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spectrum, \( m_\nu^+ = \sqrt{|\Delta m^2_A|} \) and \( m_\nu^- = \sqrt{|\Delta m^2_\odot|} \), while in the case of inverted hierarchy we have \( m_\nu^+ = \sqrt{|\Delta m^2_\odot|} \) and \( m_\nu^- = \sqrt{|\Delta m^2_A|} \), \( |\Delta m^2_A| \) and \( |\Delta m^2_\odot| \) being the atmospheric and solar neutrino mass square differences, respectively. It is easy to show that at leading order in \( \mu_N \), the neutrino mass parameters satisfy the relation \([6]\): \([y_1 \times y_2] v_1 v_2 / m_N \approx (|\Delta m^2_\odot| / |\Delta m^2_A|)^{1/4} \). Hence, barring accidental cancellations, the size of the neutrino Yukawa couplings is typically

\[
|y_1| |y_2| \approx 10^{-4} (m_N / 1 \text{ TeV}) (1 \text{ KeV} / v_2).
\]

Finally we point out that with the addition of a scalar field odd under \( U(1)_X \) it is possible to realize a viable dark matter candidate in the model, which results naturally stable due to the presence of the remnant \( Z_2 \) symmetry. A variation of leptogenesis in this case is possible at the TeV scale \([6]\).

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