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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

$$\mathcal{L}_{CC}^{\mathcal{N}} = -\frac{g}{2\sqrt{2}}\,\bar{\ell}\,\gamma_\alpha\,(RV)_{\ell k}(1-\gamma_5)\,\mathcal{N}_k W^\alpha + \text{h.c.}\,, \tag{1.1}$$

$$\mathcal{L}_{NC}^{\mathcal{N}} = -\frac{g}{4\,c_w}\,\overline{\nu_{\ell L}}\,\gamma_\alpha\,(RV)_{\ell k}\,(1-\gamma_5)\,\mathcal{N}_k Z^\alpha + \text{h.c.}\,, \tag{1.2}$$

$$\mathcal{L}_{H}^{\mathcal{N}} = -\frac{g\,M_k}{4\,M_W}\,\overline{\nu_{\ell L}}\,(RV)_{\ell k}\,(1+\gamma_5)\,\mathcal{N}_k h + \text{h.c.} \tag{1.3}$$

The couplings $(RV)_{\ell k}$ ($\ell = e,\mu,\tau$ 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]:

$$|(RV)_{\ell 1}|^2 = \frac{1}{2}\frac{y^2 v^2}{M_1^2}\frac{m_3}{m_2+m_3}\left|U_{\ell 3} + i\sqrt{m_2/m_3}\,U_{\ell 2}\right|^2, \quad \text{for normal hierarchy}\,, \tag{1.4}$$

$$|(RV)_{\ell 1}|^2 = \frac{1}{2}\frac{y^2 v^2}{M_1^2}\frac{m_2}{m_1+m_2}\left|U_{\ell 2} + i\sqrt{m_1/m_2}\,U_{\ell 1}\right|^2, \quad \text{for inverted hierarchy}\,, \tag{1.5}$$

$$(RV)_{\ell 2} = \pm i\,(RV)_{\ell 1}\sqrt{\frac{M_1}{M_2}}\,,\ \ell = e,\mu,\tau\,, \tag{1.6}$$

where U denotes the PMNS neutrino mixing matrix and $v \simeq 174$ 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 y in the expressions above represents the largest eigenvalue of the matrix of the neutrino Yukawa couplings: $y^2 v^2 = 2M_1^2\,(|(RV)_{e1}|^2 + \left|(RV)_{\mu 1}\right|^2 + |(RV)_{\tau 1}|^2)$.

Electroweak precision data (EWPD) provide an upper bound of the size of the Yukawa coupling y for a given seesaw scale, that is $y \lesssim 0.06\,(M_1/100\text{ GeV})$ [2]. However, the most stringent constraint on y comes from lepton flavour violating observables, in particular from the present upper limit on $\mu^+ \to e^+\gamma$ branching ratio reported by the MEG experiment [4]: $\text{BR}(\mu \to e\gamma) < 5.7 \times 10^{-13}$ at 90% confidence level. Indeed, in this case, taking the best fit values of the neutrino oscillation parameters [5], we get the bound: $y \lesssim 0.026$ for $M_1 = 100$ GeV.

We show in Fig. 1 all the relevant constraints on the effective couplings $(RV)_{\mu 1}$ and $(RV)_{e1}$ which come from the requirement of reproducing neutrino oscillation data, from the EWPD bounds

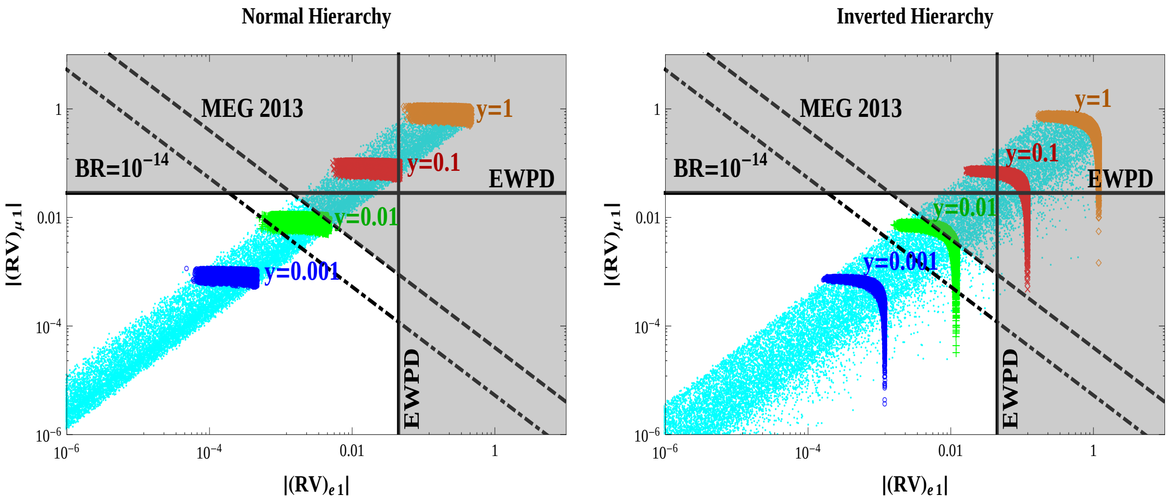

Figure 1: Correlation between $|(RV)_{e1}|$ and $|(RV)_{\mu 1}|$ for $M_1 = 100$ GeV in the case of normal (left panel) and inverted (right panel) light neutrino mass spectrum (see the text for details).

and the current upper limit on $\mu \to e\gamma$. The seesaw mass scale is fixed at the 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 $|(RV)_{e1}|$, 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 3σ 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 $\mathrm{BR}(\mu \to e\gamma) < 10^{-14}$. The scatter points correspond to different values of the maximum neutrino Yukawa coupling y: $y = 0.001$ (blue $\circ$), *ii*) $y = 0.01$ (green $+$), *iii*) $y = 0.1$ (red $\times$), *iv*) $y = 1$ (orange $\Diamond$) 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 $|U_{\mu 2} + i\,U_{\mu 1}|$ 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 (1.1-1.3) 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, ϕ. The fermion sector, instead, consists of 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

Field	L_α	$e_{R\alpha}$	N_1	N_2	H_1	H_2	ϕ
$\mathrm{SU(2)_L}$	**2**	**1**	**1**	**1**	**2**	**2**	**1**
$\mathrm{U(1)_Y}$	-1/2	-1	0	0	1/2	1/2	0
$\mathrm{U(1)_X}$	-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 ϕ 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 $\mathscr{V}_{\rm SB}$, invariant under the $\mathrm{SU(2)_L}\times\mathrm{U(1)_Y}\times[\mathrm{U(1)_X}]$ symmetry, is derived in [6]. Given the fields in Table 1, we have

$$\begin{aligned}\mathscr{V}_{\rm SB} \;=&\; -\mu_1^2\,(H_1^\dagger H_1)+\lambda_1\,(H_1^\dagger H_1)^2-\mu_2^2\,(H_2^\dagger H_2)+\lambda_2\,(H_2^\dagger H_2)^2-\mu_3^2\,\phi^*\phi+\lambda_3\,(\phi^*\phi)^2\\ &+\;\kappa_{12}\,(H_1^\dagger H_1)\,(H_2^\dagger H_2)+\kappa_{12}'\,(H_1^\dagger H_2)\,(H_2^\dagger H_1)+\kappa_{13}\,(H_1^\dagger H_1)\phi^*\phi+\kappa_{23}\,(H_2^\dagger H_2)\phi^*\phi\\ &-\;\frac{\mu'}{\sqrt{2}}\left((H_1^\dagger H_2)\phi+(H_2^\dagger H_1)\phi^*\right).\end{aligned} \tag{2.1}$$

The two $\mathrm{SU(2)_L}$ doublets $H_{1,2}$ and the singlet ϕ [6] take a non-zero vacuum expectation value (vev) $v_{1,2}$ and v_ϕ, respectively. In this case, the global $\mathrm{U(1)_X}$ 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, and 2 pseudo-scalars A^0 and J. The latter is the Goldstone boson associated with the breaking of the global $\mathrm{U(1)_X}$ 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 $Z_L \sim \sqrt{2}\,\mathrm{Im}\,(H_1^0)$, while the scalar mass eigenstates are to a good approximation: $H^\pm \sim H_2^\pm$, $h_A \sim \sqrt{2}\,\mathrm{Re}\,(H_2^0)$, $A_0 \sim \sqrt{2}\,\mathrm{Im}\,(H_2^0)$ and $J \sim \sqrt{2}\,\mathrm{Im}\,(\phi)$. Moreover, the two neutral scalars h^0 and H^0 arise from the mixing of $\sqrt{2}\,\mathrm{Re}\,(H_1^0)$ and $\sqrt{2}\,\mathrm{Re}\,(\phi)$. Typically, we have $v_2 \lesssim 10$ MeV [6], $v_1 \simeq 246.2$ GeV, while v_ϕ 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 [6] 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 $N_2^C \equiv C\overline{N_2}^T$ 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

$$\mathscr{L} \supset -m_N \overline{N_D} N_D - \left(Y^{\beta}_{\nu 1} \overline{N_D} \widetilde{H}^*_1 L_\beta + Y^{\gamma}_{\nu 2} \overline{N_D}^C \widetilde{H}^*_2 L_\gamma + \frac{\delta_N}{\sqrt{2}} \phi \overline{N_D} N^C_D + \text{h.c.} \right) \tag{2.2}$$

where $N^C_D \equiv C\overline{N}^T_D$ and $\widetilde{H}_k \equiv -i\sigma_2 H^*_k$ ($k = 1,2$). The parameter δ_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 δ_N provides, after ϕ 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 $\left(\nu_{\mathbf{L}}, (N^C_1)_{\mathbf{L}}, (N^C_2)_{\mathbf{L}}\right)$, the 5×5 symmetric neutrino mass matrix reads:

$$\mathscr{M}_\nu = \begin{pmatrix} \mathbf{0_{3\times 3}} & \mathbf{y_1}\, v_1 & \mathbf{y_2}\, v_2 \\ \mathbf{y_1}^{\mathrm{T}} v_1 & \delta_N v_\phi & m_N \\ \mathbf{y_2}^{\mathrm{T}} v_2 & m_N & \delta_N v_\phi \end{pmatrix}. \tag{2.3}$$

In the previous expression $\mathbf{0_{3\times 3}}$ denotes the null matrix of dimension 3 and we introduce the short-hand notation: $\mathbf{y_k} \equiv \left(Y^e_{\nu k}, Y^\mu_{\nu k}, Y^\tau_{\nu k}\right)^{\mathrm{T}} / \sqrt{2}$. The neutrino sector, therefore, consists of one massless neutrino, two massive light Majorana neutrinos and two heavy Majorana neutrinos $\mathscr{N}_{1,2}$. The latter are quasidegenerate, with masses $M_{1,2} = m_N \mp \delta_N v_\phi$, and form a pseudo-Dirac pair for $m_N \gg \delta_N v_\phi$ [6]. 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 $(RV)_{\ell 1}$ are proportional to the lepton number conserving Yukawa couplings $Y^\ell_{\nu k}$. The resulting effective light neutrino mass matrix is

$$(M_\nu)^{\alpha\beta} \simeq -\frac{v_1 v_2}{m_N} \left(\mathbf{y_1}^\alpha \mathbf{y_2}^\beta + \mathbf{y_2}^\alpha \mathbf{y_1}^\beta \right) + v_\phi\, \delta_N \frac{v_1^2}{m_N^2} \left(\mathbf{y_1}^\alpha \mathbf{y_1}^\beta + \mathbf{y_2}^\alpha \mathbf{y_2}^\beta \frac{v_2^2}{v_1^2} \right). \tag{2.4}$$

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_ϕ 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 ϕ, 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 [6]. Taking $\mu_N \equiv (\delta_N v_\phi)/m_N \ll 1$, the neutrino masses result

$$\begin{aligned} m^\pm_\nu &\simeq \frac{1}{m_N} \left(\sqrt{y_1^2 y_2^2 - \mu_N (y_1^2 + y_2^2)\,\mathrm{Re}(y_{12})} \pm \sqrt{|y_{12}|^2 - \mu_N (y_1^2 + y_2^2)\,\mathrm{Re}(y_{12})} \right) \\ &\simeq \frac{1}{m_N} (y_1 y_2 \pm |y_{12}|) \times \left(1 \mp \frac{\mu_N}{2} \frac{(y_1^2 + y_2^2)\,\mathrm{Re}(y_{12})}{y_1 y_2 |y_{12}|} \right), \end{aligned} \tag{2.5}$$

with $y_i \equiv \sqrt{\mathbf{y_i}^\dagger \cdot \mathbf{y_i}}\, v_i$, $y_{12} \equiv \mathbf{y_1}^\dagger \cdot \mathbf{y_2}\, v_1 v_2$ and $\eta_{12} \equiv \mathbf{y_1} \times \mathbf{y_2}\, v_1 v_2$. Notice that if the neutrino Yukawa vectors $\mathbf{y_1}$ and $\mathbf{y_2}$ are proportional, m^-_ν is exactly zero. For a normal hierarchical spectrum,

$m_\nu^+ = \sqrt{|\Delta m^2_{\rm A}|}$ and $m_\nu^- = \sqrt{\Delta m^2_\odot}$, while in the case of inverted hierarchy we have $m_\nu^+ = \sqrt{|\Delta m^2_{\rm A}|}$ and $m_\nu^- = \sqrt{|\Delta m^2_{\rm A}| - \Delta m^2_\odot}$, $|\Delta m^2_{\rm A}|$ 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 μ_N, the neutrino mass parameters satisfy the relation [6]: $|\mathbf{y_1} \times \mathbf{y_2}|\, v_1\, v_2/m_N \cong (\Delta m^2_\odot\, |\Delta m^2_{\rm A}|)^{1/4}$. Hence, barring accidental cancellations, the size of the neutrino Yukawa couplings is typically

$$|\mathbf{y_1}||\mathbf{y_2}| \approx 10^{-4}\,(m_N/1\ \text{TeV})\,(1\ \text{KeV}/v_2)\,. \tag{2.6}$$

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 is 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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