Gauge and Yukawa Mediated Supersymmetry Breaking in the Triplet Seesaw Scenario

Filipe R. Joaquim[1,2] and Anna Rossi[1]

[1]Dipartimento di Fisica "G. Galilei", Università di Padova, I-35131 Padua, Italy

[2]Istituto Nazionale di Fisica Nucleare, Sezione di Padova, I-35131 Padua, Italy

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We propose a novel supersymmetric unified scenario of the triplet seesaw mechanism where the exchange of the heavy triplets generates both neutrino masses and soft supersymmetry breaking terms. Our framework is very predictive since it relates neutrino mass parameters, lepton-flavor-violation in the slepton sector, sparticle and Higgs spectra, and electroweak symmetry breakdown. The phenomenological viability and experimental signatures in lepton flavor-violating processes are discussed.

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Modern particle physics has been confronting the intriguing issue of neutrino mass generation and its phenomenological implications. The seesaw mechanism provides a natural explanation for the generation of neutrino masses and their suppression with respect to the other fermion masses of the standard model (SM). In its most popular versions, the seesaw mechanism is realized either by exchanging singlet fermions N [1], or a $SU(2)_W$ scalar triplet T with nonzero hypercharge [2], at a high scale M_L. An attractive feature of the supersymmetric version of the above scenarios is that lepton flavor-violating (LFV) processes (otherwise unobservable) can be enhanced through one-loop exchange of lepton superpartners if their masses do not conserve flavor. Regarding this aspect, most of the literature has been focusing on the most conservative scenario of universal sfermion masses at high energy, as in minimal supergravity or gauge-mediated supersymmetry (SUSY) breaking models. In such cases, flavor nonconservation in the sfermion masses arises from renormalization group (RG) effects due to flavor-violating Yukawa couplings [3–5]. We recall that in the triplet seesaw the flavor structure of the slepton mass matrix $\mathbf{m}_{\tilde{L}}^2$ after RG running can be univocally determined in terms of the low-energy neutrino parameters [5]. This is in contrast with the singlet seesaw where the structure of $\mathbf{m}_{\tilde{L}}^2$ cannot be unambiguously related to the neutrino parameters.

In this Letter we present a novel SUSY scenario of the triplet seesaw mechanism in which the soft SUSY-breaking (SSB) parameters in the minimal supersymmetric extension of the SM (MSSM) are generated at the decoupling of the heavy triplets and the mass scale of such SSB terms is fixed *only* by the triplet SSB bilinear term B_T. This scenario is highly predictive since it relates neutrino masses, LFV in the sfermion sector, sparticle and Higgs spectra, and electroweak symmetry breaking (EWSB).

The supersymmetric version of the triplet seesaw requires introducing the triplets as supermultiplets T, $\bar{T}$ in a vectorlike $SU(2)_W \times U(1)_Y$ representation, $T \sim (3, 1)$, $\bar{T} \sim (3, -1)$. In order to preserve successful gauge coupling unification, we embed our framework in a $SU(5)$ grand unified theory (GUT) [5] where the triplet states fit into the 15 representation $15 = S + T + Z$ transforming as $S \sim (6, 1, -\frac{2}{3})$, $T \sim (1, 3, 1)$, $Z \sim (3, 2, \frac{1}{6})$ under $SU(3) \times SU(2)_W \times U(1)_Y$ (the $\overline{15}$ decomposition is obvious). The SUSY-breaking mechanism is parametrized by a gauge singlet chiral supermultiplet X, whose scalar S_X and auxiliary F_X components are assumed to acquire a vacuum expectation value through some unspecified dynamics in the secluded sector. It is suggestive to assume that the $SU(5)$ model conserves the combination $B - L$ of baryon and lepton number. As a result, the relevant superpotential reads

$$W_{SU(5)} = \frac{1}{\sqrt{2}}(\mathbf{Y}_{15}\bar{5}\,15\,\bar{5} + \lambda 5_H \overline{15}\,5_H) + \mathbf{Y}_5 \bar{5}\,\bar{5}_H 10 + \mathbf{Y}_{10} 10\,10\,5_H + M_5 5_H \bar{5}_H + \xi X 15\,\overline{15}, \quad (1)$$

where the matter multiplets are understood as $\bar{5} = (d^c, L)$, $10 = (u^c, e^c, Q)$ and the Higgs doublets fit with their colored partners, t, $\bar{t}$ like $5_H = (t, H_2)$, $\bar{5}_H = (\bar{t}, H_1)$. The $B - L$ quantum numbers are a combination of the hypercharges and the following charges: $Q_{10} = \frac{1}{5}$, $Q_{\bar{5}} = -\frac{3}{5}$, $Q_{5_H} = -\frac{2}{5}$, $Q_{\bar{5}_H} = \frac{2}{5}$, $Q_{15} = \frac{6}{5}$, $Q_{\overline{15}} = \frac{4}{5}$ and $Q_X = -2$. The form of $W_{SU(5)}$ implies that the 15, $\overline{15}$ states play the role of *messengers* of both $B - L$ and SUSY breaking to the visible (MSSM) sector thanks to the coupling with X. Namely, while $\langle S_X \rangle$ only breaks $B - L$, $\langle F_X \rangle$ breaks both SUSY and $B - L$. These effects are parametrized by the superpotential mass term $M_{15} 15\overline{15}$, where $M_{15} = \xi \langle S_X \rangle$, and the bilinear SSB term $-BM_{15} 15\overline{15}$, with $BM_{15} = -\xi \langle F_X \rangle$. Once $SU(5)$ is broken to the SM group we find [5], below the GUT scale M_G,

$$W = W_0 + W_T + W_{S,Z}$$

$$W_0 = \mathbf{Y}_e e^c H_1 L + \mathbf{Y}_d d^c H_1 Q + \mathbf{Y}_u u^c Q H_2 + \mu H_2 H_1$$

$$W_T = \frac{1}{\sqrt{2}}(\mathbf{Y}_T L T L + \lambda H_2 \bar{T} H_2) + M_T T \bar{T}$$

$$W_{S,Z} = \frac{1}{\sqrt{2}}\mathbf{Y}_S d^c S d^c + \mathbf{Y}_Z d^c Z L + M_Z Z \bar{Z} + M_S S \bar{S}. \quad (2)$$

Here, W_0 denotes the MSSM superpotential, W_T contains the triplet Yukawa and mass terms, and $W_{S,Z}$ includes the couplings and masses of the colored fragments S, Z. As in [5], we have relaxed the strict $SU(5)$ symmetry relations for the Yukawa interactions and mass terms by allowing $SU(5)$ breaking effects, induced, for example, by adjoint 24-insertions, such as $\mathbf{Y}_5 = \mathbf{Y}_5^{(0)} + \mathbf{Y}_5^{(1)} 24/\Lambda + \ldots$ with a cutoff scale $\Lambda > M_G$. These insertions are necessary to correct the relation $\mathbf{Y}_e = \mathbf{Y}_d^T$ and to solve the doublet-triplet splitting problem. For the sake of simplicity, we take $M_T = M_S = M_Z$ and $\mathbf{Y}_S, \mathbf{Y}_Z \ll \mathbf{Y}_T$ at M_G (possibly due to 24-insertions), which does not alter the major point of our discussion. The $SU(5)$ scenario with $\mathbf{Y}_S = \mathbf{Y}_Z = \mathbf{Y}_T$ implies correlations between LFV and quark flavor violation; this case has been considered in detail in [6]. In Eq. (2), W_T is responsible for the realization of the seesaw mechanism. Actually, at the scale M_T the triplets act as tree-level *messengers* of lepton number and flavor violation [7] via the symmetric Yukawa matrix $\mathbf{Y}_T$, generating the $d = 5$ effective operator $\lambda \mathbf{Y}_T (LH_2)^2/M_T$. Subsequently, at the electroweak scale the Majorana neutrino mass matrix is obtained

$$\mathbf{m}_\nu^{ij} = \lambda \langle H_2 \rangle^2 \mathbf{Y}_T^{ij}/M_T, \qquad i, j = e, \mu, \tau. \tag{3}$$

In the basis where $\mathbf{Y}_e$ is diagonal, it is apparent that all LFV is encoded in $\mathbf{Y}_T$. Namely, the nine independent parameters contained in $\mathbf{m}_\nu$ are directly linked to the neutrino parameters according to $\mathbf{m}_\nu = \mathbf{U}^* \mathbf{m}_\nu^D \mathbf{U}^\dagger$, where $\mathbf{m}_\nu^D = \mathrm{diag}(m_1, m_2, m_3)$ are the mass eigenvalues, and $\mathbf{U}$ is the leptonic mixing matrix.

Regarding the SSB term one has, in the broken phase: $-B_T M_T (T\bar{T} + S\bar{S} + Z\bar{Z}) + \mathrm{H.c.}$, where $B_T \equiv B_{15}$. These terms lift the tree-level mass degeneracy in the MSSM supermultiplets. Indeed, at the scale M_T, all the states T, T, $\bar{T}$, S, $\bar{S}$, and Z, $\bar{Z}$ are *messengers* of SUSY breaking to the MSSM sector via gauge interactions, as it happens in conventional gauge-mediation scenarios [8]. However, in our framework, the states T, $\bar{T}$ also transmit SUSY breaking via Yukawa interactions. Finite contributions for the trilinear couplings of the superpartners with the Higgs doublets, $\mathbf{A}_e$, $\mathbf{A}_u$, $\mathbf{A}_d$ the gaugino masses $M_a (a = 1, 2, 3)$, and the Higgs bilinear term $-B_H \mu H_2 H_1$ emerge at the one-loop level:

$$\begin{aligned} \mathbf{A}_e &= \frac{3B_T}{16\pi^2} \mathbf{Y}_e \mathbf{Y}_T^\dagger \mathbf{Y}_T, \qquad \mathbf{A}_u = \frac{3B_T}{16\pi^2} \mathbf{Y}_u |\lambda|^2, \\ \mathbf{A}_d &= 0, \qquad M_a = \frac{7B_T}{16\pi^2} g_a^2, \qquad B_H = \frac{3B_T}{16\pi^2} |\lambda|^2, \end{aligned} \tag{4}$$

(g_a are the gauge couplings). As for the SSB squared scalar masses, the leading $\mathcal{O}(F_X^2/M_T^2) = \mathcal{O}(B_T^2)$ contributions do not emerge at one-loop level [9], but instead at two-loop [10]:

$$\begin{aligned}
\mathbf{m}_{\tilde{L}}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \bigg[\frac{21}{10} g_1^4 + \frac{21}{2} g_2^4 - \left(\frac{27}{5} g_1^2 + 21 g_2^2 \right) \mathbf{Y}_T^\dagger \mathbf{Y}_T \\
&\quad + 3 \mathbf{Y}_T^\dagger \mathbf{Y}_e^T \mathbf{Y}_e^* \mathbf{Y}_T + 18 (\mathbf{Y}_T^\dagger \mathbf{Y}_T)^2 \\
&\quad + 3\,\mathrm{Tr}(\mathbf{Y}_T^\dagger \mathbf{Y}_T) \mathbf{Y}_T^\dagger \mathbf{Y}_T \bigg] \\
\mathbf{m}_{\tilde{e}^c}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \left[\frac{42}{5} g_1^4 - 6 \mathbf{Y}_e \mathbf{Y}_T^\dagger \mathbf{Y}_T \mathbf{Y}_e^\dagger \right] \\
\mathbf{m}_{\tilde{Q}}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \left[\frac{7}{30} g_1^4 + \frac{21}{2} g_2^4 + \frac{56}{3} g_3^4 - 3|\lambda|^2 \mathbf{Y}_u^\dagger \mathbf{Y}_u \right] \\
\mathbf{m}_{\tilde{u}^c}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \left[\frac{56}{15} g_1^4 + \frac{56}{3} g_3^4 - 6|\lambda|^2 \mathbf{Y}_u \mathbf{Y}_u^\dagger \right] \\
\mathbf{m}_{\tilde{d}^c}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \left[\frac{14}{15} g_1^4 + \frac{56}{3} g_3^4 \right] \\
m_{H_1}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \left[\frac{21}{10} g_1^4 + \frac{21}{2} g_2^4 \right] \\
m_{H_2}^2 &= \frac{|B_T|^2}{(16\pi^2)^2} \bigg[\frac{21}{10} g_1^4 + \frac{21}{2} g_2^4 - \left(\frac{27}{5} g_1^2 + 21 g_2^2 \right) |\lambda|^2 \\
&\quad + 9|\lambda|^2\,\mathrm{Tr}(\mathbf{Y}_u \mathbf{Y}_u^\dagger) + 21|\lambda|^4 \bigg].
\end{aligned} \tag{5}$$

The results (4) and (5) can be obtained either by diagrammatic computations or from generalization of the wave function renormalization method proposed in [11].

Notice that the generation of all gaugino masses requires a complete 15 representation. More specifically, M_1, M_2, and M_3 arise from the exchange of (T, S, Z), (T, Z), and (S, Z), respectively. Equations (4) and (5) hold at the decoupling scale M_T and therefore are meant as boundary conditions for the SSB parameters which then undergo (MSSM) RG running to the low-energy scale μ_{SUSY}. Observe that the Yukawas $\mathbf{Y}_T$ induce LFV to $\mathbf{A}_e$, to $\mathbf{m}_{\tilde{L}}^2$ and to a much less extent in $\mathbf{m}_{\tilde{e}^c}^2$. This makes the present scenario different from pure gauge-mediated models [8] where flavor violation is naturally suppressed (for other examples of Yukawa mediated SUSY breaking, see, e.g, [9,12]). We suppose that possible gravity mediated contributions $\sim F/M_{\mathrm{pl}}$ (where $F^2 = \langle |F_X|^2 \rangle + \ldots$ is the sum of F terms in the secluded sector) are negligible. This is the case if $M_T \ll 10^{16}$ GeV$\xi \langle F_X \rangle / F$. Furthermore, it is necessary that $\xi \langle F_X \rangle < M_T^2$ (or $B_T < M_T$) to avoid tachyonic scalar messengers.

It is worth stressing that here the LFV entries $(\mathbf{m}_{\tilde{L}}^2)_{ij}$ ($i \neq j$)show up as finite radiative contributions induced by B_T at M_T, and they are not essentially modified by the (MSSM) RG evolution to low-energy. This is different from a previous work [5] where a common SSB scalar mass $m_0 \sim \mathcal{O}(100\ \mathrm{GeV})$ was assumed at M_G and the dominant LFV contributions to $\mathbf{m}_{\tilde{L}}^2$ were generated by RG evolution from M_G down to the decoupling scale M_T. In such a case, finite contributions like those in Eqs. (4) and (5), also emerge at M_T, but they are subleading with respect to the RG corrections, since B_T is of the same order as m_0. Instead, in the present picture, there is a hierarchy between the SSB parameter B_T and the remaining ones [see Eqs. (4) and (5)], $B_T^2 \gg (B_T g^2/16\pi^2)^2 \sim m_0^2$. However, in both scenarios the flavor structure of $\mathbf{m}_{\tilde{L}}^2$ is proportional to $\mathbf{Y}_T^\dagger \mathbf{Y}_T$ and can be written by using Eq. (3) in terms of the neutrino parameters (the terms $\propto g^2 \mathbf{Y}_T^\dagger \mathbf{Y}_T$ are generically the leading ones):

$$(\mathbf{m}_{\tilde{L}}^2)_{ij} \propto (\mathbf{Y}_T^\dagger \mathbf{Y}_T)_{ij} \sim \left(\frac{M_T}{\lambda \langle H_2 \rangle^2}\right)^2 [\mathbf{U}(\mathbf{m}_\nu^D)^2 \mathbf{U}^\dagger]_{ij}. \quad (6)$$

Consequently, the relative size of LFV in the different leptonic families can be univocally predicted as

$$\frac{(\mathbf{m}_{\tilde{L}}^2)_{\tau\mu}}{(\mathbf{m}_{\tilde{L}}^2)_{\mu e}} \approx \frac{\rho s_{23} c_{23}}{s_{12} c_{12} c_{23}} \sim 40, \quad \frac{(\mathbf{m}_{\tilde{L}}^2)_{\tau e}}{(\mathbf{m}_{\tilde{L}}^2)_{\mu e}} \approx -\frac{s_{23}}{c_{23}} \sim -1, \quad (7)$$

where $\rho = (m_3/m_2)^2$, θ_{12} and θ_{23} are lepton mixing angles and $\theta_{13} = 0$ is taken (here $c_{ij} = \cos\theta_{ij}, \ldots$). A hierarchical neutrino mass spectrum is considered and the best-fit values for the parameters are used [13]. Taking the present upper limit on $\sin\theta_{13} = 0.2$, the above ratios become 3 and 0.8, respectively, while varying the other neutrino parameters within their experimental range affects these ratios by less than 10% (see also [6]). The above relations imply that also the branching ratios (BR) of LFV processes such as $l_i \to l_j \gamma$ can be predicted

$$\frac{\mathrm{BR}(\tau \to \mu\gamma)}{\mathrm{BR}(\mu \to e\gamma)} \sim 300, \qquad \frac{\mathrm{BR}(\tau \to e\gamma)}{\mathrm{BR}(\mu \to e\gamma)} \sim 10^{-1}. \quad (8)$$

Other LFV processes and related correlations [14] have been considered in [6]. (Connections between neutrino parameters and other observables can arise also in different scenarios, see, e.g., [15]). Without loss of generality, we take B_T to be real since its phase has not physical effect. However, a different approach was considered in [16] where a complex B_T could generate sizable electric dipole moments for quarks and leptons since there was a relative phase between $\mathbf{A}_{e,d,u}$ and M_a shown in Eq. (4). Moreover, B_T could play a role in generating the baryon asymmetry of the Universe within resonant leptogenesis [17].

We shall now discuss the phenomenological viability taking $M_T > 10^7$ GeV so that the gauge couplings remain perturbative up to M_G. Our approach follows a bottom-up perspective where, for a given ratio M_T/λ and $\tan\beta$, $\mathbf{Y}_T$ is determined at M_T according to Eq. (3) using the low-energy neutrino parameters. The Yukawa matrices $\mathbf{Y}_{e,u,d}$ are determined by the related charged fermion masses, modulo $\tan\beta$. Although the μ-parameter is not predicted by the underlying theory, it is nevertheless determined with $\tan\beta$ by correct EWSB conditions. Therefore, we end up with only three free parameters, B_T, M_T, and λ. In Fig. 1 we show the (λ, M_T) parameter space allowed by the perturbativity and EWSB requirements, the experimental lower bound on the lightest Higgs mass [18] m_h and the upper bound on $\mathrm{BR}(\mu \to e\gamma)$, for $B_T = 20(50)$ TeV in the left (right) panel. First notice the light-gray regions excluded by the perturbativity requirement which are independent of B_T. For each value of M_T there is a minimum value of λ, which scales as $\sim 2 \times 10^{-4}(M_T/10^{11}$ GeV), below which the couplings $\mathbf{Y}_T$ reach the Landau pole below M_G. Similarly, there is a maximum value of λ beyond which λ itself blows up below M_G. The EWSB constraint excludes a region for $\lambda \sim 1 - 1.2$ and $M_T \gtrsim 10^{12}$ GeV (independently of B_T), which is limited by the least achievable value of $\tan\beta$, $\tan\beta \sim 2.5$. As for the μ-parameter (dashed lines), it slightly increases with increasing M_T due to the large RG factor which affects $m_{H_2}^2(\mu_{\mathrm{SUSY}})$ in the minimization condition, $\mu^2(\mu_{\mathrm{SUSY}}) \approx -m_{H_2}^2(\mu_{\mathrm{SUSY}})$, covering the range $\mu \sim$ 450–550 (1000–1200) GeV for $B_T = 20(50)$ TeV. We observe that $\lambda < 0.6(0.7)$ for $B_T = 20(50)$ TeV is required by the constraint $m_h > 110$ GeV. The related contour lies on the correspondent minimum value of $\tan\beta \sim 5(3.5)$ for $B_T = 20(50)$ TeV. When $B_T = 50$ TeV, the sparticle spectrum is heavier. Hence the radiative corrections $\sim \log(\frac{\mu_{\mathrm{SUSY}}}{m_t})$ to m_h are larger and in the tree-level contribution $\sim M_Z|\cos 2\beta|$ smaller $\tan\beta$ can be tolerated.

The present bound on $\mathrm{BR}(\mu \to e\gamma)$ provides a lower bound on λ for each value of M_T. This stems from the fact that the LFV entries $(\mathbf{m}_{\tilde{L}}^2)_{ij}$ scale as $(M_T/\lambda)^2$ [Eq. (5)]. Consequently, the allowed λ-range is wider for lower values of M_T and, comparing the two panels, the whole parameter space is larger for $B_T = 50$ TeV. In the allowed regions, the lightest MSSM sparticle is typically a charged slepton with mass around 100–200(300–450) GeV for $B_T = 20(50)$ TeV, although for small $\tan\beta$ there could be a mass degeneracy with the lightest neutralino. However, either the lightest slepton or neutralino would decay into the gravitino which is most likely the lightest

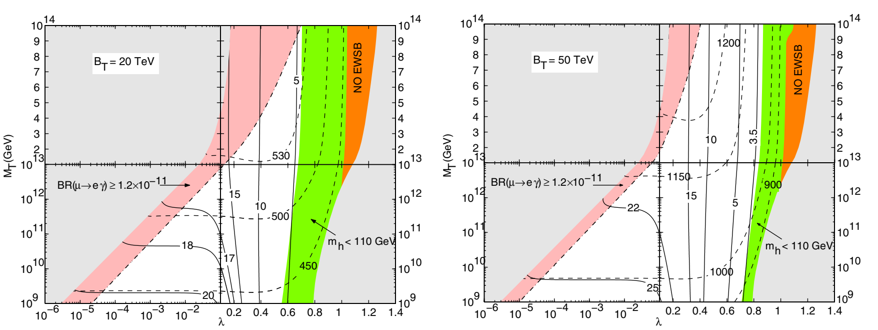

FIG. 1 (color online). The parameter space constrained by the perturbativity requirement (light-gray), correct EWSB from the one-loop corrected Higgs potential, lower bound on m_h and the upper bound on $\mathrm{BR}(\mu \to e\gamma)$, for $B_T = 20(50)$ TeV in the left (right) panel. We display the isocontours of $\tan\beta$ (solid line) and μ (dashed line). Here, the top pole mass is fixed at $m_t = 174$ GeV.

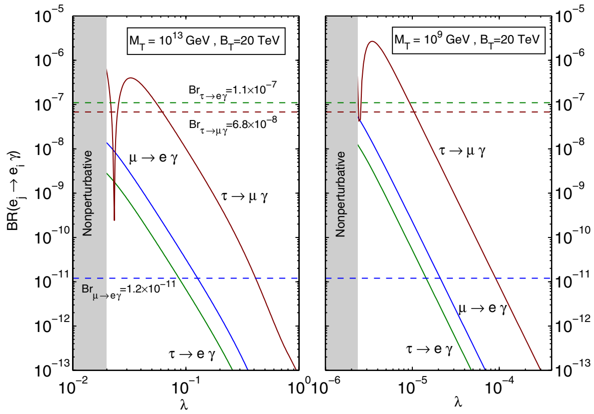

FIG. 2 (color online). Branching ratios of the lepton radiative decays. The horizontal lines indicate the present bound on each BR [21].

supersymmetric particle in our framework. Finally, we have checked that values of $B_T < 10$ TeV are phenomenologically unacceptable. In Fig. 2 we display the branching ratios $\mathrm{BR}(l_j \to l_i \gamma)$ as a function of λ for $B_T = 20$ TeV and $M_T = 10^{13}(10^9)$ GeV in the left (right) panel. The behavior of these branching ratios is in remarkable agreement with the estimates of Eq. (8). Hence, the relative size of LFV does not depend on the detail of the model, i.e., on λ, B_T, or M_T. This feature is not present for a very narrow range of λ where $\mathrm{BR}(\tau \to \mu\gamma)$ is strongly suppressed due to a conspiracy of the various contributions in $(\mathbf{m}^2_{\tilde{L}})_{\tau\mu}$ which mutually cancel out [see Eq. (5)].

Before concluding, we briefly mention that the tree-level exchange of the T, $\bar{T}$ states also generates the L-violating SSB operator $\lambda \mathbf{Y}_T B_T (\tilde{L} H_2)^2 / M_T$ which induces a sneutrino/antisneutrino mass splitting $\Delta\mathbf{m}^2_{\tilde{\nu}} = B_T \mathbf{m}_\nu$ at the weak scale. Since B_T is much larger than the weak scale, this could render interesting effects for sneutrino oscillations [20].

In conclusion, we have suggested a unified picture of the supersymmetric type-II seesaw where the triplets, besides being responsible for neutrino mass generation, communicate SUSY breaking to the observable sector through gauge and Yukawa interactions. We have performed a phenomenological analysis of the allowed parameter space emphasizing the role of LFV processes in testing our framework.

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