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A new mechanism of neutrino mass generation in the NMSSM with broken lepton number

Asmaa Abada [a], Gautam Bhattacharyya [b,*], Grégory Moreau [c]

[a] *Laboratoire de Physique Théorique, Université de Paris-sud XI, Bâtiment 210, 91405 Orsay cedex, France*
[b] *Saha Institute of Nuclear Physics, 1/AF Bidhan Nagar, Kolkata 700064, India*
[c] *Centro de Física Teórica de Partículas, Departamento de Física, Instituto Superior Técnico, Av. Rovisco Pais 1, 1049-001 Lisboa, Portugal*

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Abstract

In the Minimal Supersymmetric Standard Model with bilinear R-parity violation, only one neutrino eigenstate acquires a mass at tree level, consequently experimental data on neutrinos cannot be accommodated at tree level. We show that in the next-to-minimal extension, where a gauge singlet superfield is added to primarily address the so-called μ-problem, it is possible to generate two massive neutrino states at tree level. Hence, the global three-flavour neutrino data can be reproduced at tree level, without appealing to loop dynamics which is vulnerable to model-dependent uncertainties. We give analytical expressions for the neutrino mass eigenvalues and present examples of realistic parameter choices.

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Keywords: Neutrino mass; R-parity violation; NMSSM

1. Introduction

In the Minimal Supersymmetric Standard Model (MSSM) with bilinear interactions in the superpotential explicitly violating the R-parity symmetry [1], a neutrino Majorana mass can be generated. Nevertheless, the rank 1 nature of the neutrino mass matrix suggests that only one eigenstate becomes massive at tree level, whereas the neutrino oscillation data require at least two non-zero mass eigenvalues [2]. The bilinear R-parity violating ($\not{R}_p$) soft masses induce one more non-zero eigenvalue, but only at one-loop order. What happens if one considers the next-to-minimal version of the MSSM, called 'NMSSM' [3], in the presence of bilinear $\not{R}_p$ terms in the superpotential? Here, the particle content is extended by one gauge singlet superfield. Our main result in this work is that *two* non-degenerate massive neutrino states now emerge at *tree level*. The upshot is that one can now reproduce the neutrino oscillation data with the superpotential parameters and gaugino masses just from *tree level* physics. On the other hand, in the usual MSSM with bilinear $\not{R}_p$ terms, the generation of the second neutrino mass eigenvalue relies on the soft supersymmetry breaking scalar masses which in turn bring more uncertainties from the supersymmetry breaking mechanism; furthermore, uncertainties from loop dynamics creep in too.

An increasingly important virtue of the NMSSM [3] (see [4] for phenomenological studies) is that it ameliorates the 'little hierarchy' problem originating from the requirement of large soft supersymmetry breaking scalar masses compared to the electroweak scale (in order to sufficiently push the lightest Higgs mass beyond the LEP limit). The NMSSM also provides a solution to the

* Corresponding author.
E-mail addresses: abada@th.u-psud.fr (A. Abada), gb@theory.saha.ernet.in (G. Bhattacharyya), greg@cftp.ist.utl.pt (G. Moreau).

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so-called μ-problem by arranging the vacuum expectation value (vev) of the gauge singlet scalar of the order of the supersymmetry breaking scale, so that the μ parameter turns out to be at the electroweak scale.

A recent paper [5], in the context of NMSSM with $\not R_p$ couplings, deals with the generation of neutrino masses where two eigenvalues arise at loop level. In another recent analysis [6], it has been shown that the NMSSM with bilinear $\not R_p$ terms offers a possible mechanism of neutrino mass suppression, thus significantly reducing the hierarchy between μ_i and μ (defined below). Besides, an alternative extension of the MSSM explicitly breaking the R-parity has been proposed in order to simultaneously address the μ-problem and the issue of small neutrino masses [7].

In Section 2, we present the superpotential of the model we consider. In Section 3, we discuss the effective neutrino mass matrix. We present our numerical results in Section 4. Finally, we conclude in Section 5.

2. Superpotential

The NMSSM superpotential contains two dimensionless couplings λ and κ in addition to the usual Yukawa couplings:

$$W_{\mathrm{NMSSM}} = Y^u_{ij} Q_i H_u U^c_j + Y^d_{ij} Q_i H_d D^c_j + Y^\ell_{ij} L_i H_d E^c_j + \lambda S H_u H_d + \frac{1}{3}\kappa S^3, \tag{1}$$

where $Y^{u,d,\ell}_{ij}$ are the Yukawa coupling constants (i, j, k are family indexes), and Q_i, L_i, U^c_i, D^c_i, E^c_i, H_u, H_d, S respectively are the superfields for the quark doublets, lepton doublets, up-type anti-quarks, down-type anti-quarks, anti-leptons, up Higgs, down Higgs, extra singlet under the standard model gauge group. An effective μ term, given by $\lambda\langle s\rangle H_u H_d$, is generated via the vev of the scalar component s of the singlet superfield S.

We now take note that in supersymmetric theories there is no deep underlying theoretical principle for the imposition of R-parity as a symmetry [8]. However, there exist strong constraints on the $\not R_p$ couplings coming from various phenomenological considerations [9,10]. Limits on neutrino masses and mixings have also been translated into tight upper bounds for $\not R_p$ couplings [11].

In the present Letter, we consider a generic NMSSM superpotential containing both the bilinear and trilinear $\not R_p$ terms:

$$W = W_{\mathrm{NMSSM}} + \mu_i L_i H_u + \lambda_i S L_i H_u, \tag{2}$$

where μ_i (λ_i) are the dimension-one (dimensionless) $\not R_p$ parameters. Actually, the contribution of trilinear term $\lambda_i S L_i H_u$ was studied in Ref. [5]. Admittedly, the most generic NMSSM superpotential also contains the other renormalizable trilinear $\not R_p$ interactions, namely, $\lambda_{ijk} L_i L_j E^c_k$, $\lambda'_{ijk} L_i Q_j D^c_k$ and $\lambda''_{ijk} U^c_i D^c_j D^c_k$, which are not relevant so long as we stick to tree level neutrino mass matrix.

Normally, in the NMSSM, only trilinear couplings with dimensionless parameters (like λ and κ) are kept in the superpotential, while dimensional parameters (like μ) are generated from the vev $\langle s\rangle$. In this Letter, the $\not R_p$ NMSSM superpotential (2) is assumed to arise in either one of the following three possible scenarios:

(1) All possible renormalizable terms are included in the superpotential. Then both bilinear ($\mu_0 H_u H_d$, $\mu_i L_i H_u$) and trilinear ($\lambda S H_d H_u$, $\lambda_i S L_i H_u$) terms are admitted. However, even if one may start with a term $\mu_0 H_u H_d$, it can be rotated away by a redefinition of fields through a rotation on $L_\alpha = (H_d, L_i)$ [$\alpha = 0, \dots, 3$], since H_d and L_i have the same gauge quantum numbers. There is no reason why this redefinition would remove also the λ_α terms ($\lambda_\alpha = (\lambda, \lambda_i)$), since the corresponding 4×4 rotation matrix depends on the μ_α parameters (the generic case is considered here, where μ_α and λ_α are *not* proportional). The coefficient μ_s of the S^2 term is assumed to be zero which can be considered as a possible natural value for a superpotential parameter. It should be noted that in the standard NMSSM with conserved R-parity there is an accidental Z_3 discrete symmetry whose spontaneous breaking causes the domain wall problem. In our version of the NMSSM, the $\not R_p$ bilinear $L_i H_u$ term explicitly breaks that Z_3 symmetry. In this scenario, we simply assume the existence of the dimensionful μ_i terms, which as we will see later will be constrained from neutrino data. But we do not advance any theoretical reason as to why μ_i would be small.

(2) Our second scenario is based on the 't Hooft criteria of naturality: the parameters μ_i, μ_0 and μ_s are naturally small if the symmetry of the theory increases as these parameters are set to zero. For instance, one can assume that somehow a weak breaking (compared to the electroweak scale Q_{EW}) of some symmetry (like e.g. a U(1) symmetry forbidding the bilinear terms) generates the bilinear terms in the superpotential associated to μ_i, μ_0, $\mu_s \ll Q_{\mathrm{EW}}$. This small breaking would allow to address the μ-problem. Indeed, the main contribution to the dimension-one coefficient of $H_u H_d$ here comes from $\lambda\langle s\rangle$, as $\mu_0 \ll \mu = \lambda\langle s\rangle \sim Q_{\mathrm{EW}}$. The weak breaking of the symmetry is also responsible for the smallness of $\not R_p$ couplings and neutrino masses, since $\mu_i \ll \mu$. Thus, in such a scenario, the treatment of the mu-naturalness (à la NMSSM) and of the neutrino masses (à la $\not R_p$) are nicely connected via the weak breaking of a common symmetry. Admittedly, we do not provide any specific realization of this weak breaking. We only hint at such a possibility that the bilinear μ_i, μ_0, μ_s couplings may arise from powers of some small spurion vev ($\ll Q_{\mathrm{EW}}$).

(3) Finally, we propose a scenario where the trilinear λ_i terms in superpotential (2) are not present. This scenario relies on the Z_3 symmetry, where the chiral superfields transform by $\exp(i2\pi q/3)$, with the following charge assignments: $q = 0$ for U^c, D^c, E^c; $q = 1$ for S, H_u, H_d; and $q = 2$ for Q, L. Such a symmetry allows all couplings in Eqs. (1) and (2) except $S L_i H_u$. The other

terms S^2 and $H_u H_d$ are also eliminated by this symmetry. A spontaneous breaking of this Z_3 symmetry admittedly creates the domain wall problem, as happens in the standard NMSSM *with* R-parity. In this scenario, the μ term is created from the vev of S, but the μ_i terms are present in the superpotential from the beginning only to be constrained by neutrino data.

3. Neutrino mass matrix

Neutralino mass matrix. Within our framework, the neutralino mass terms read as,

$$\mathcal{L}^m_{\tilde{\chi}^0} = -\frac{1}{2}\Psi^{0^T}\mathcal{M}_{\tilde{\chi}^0}\Psi^0 + \text{h.c.} \tag{3}$$

in the basis $\Psi^{0^T} \equiv (\tilde{B}^0, \tilde{W}_3^0, \tilde{h}_d^0, \tilde{h}_u^0, \tilde{s}, \nu_i)^T$, where $\tilde{h}^0_{u,d}(\tilde{s})$ are the fermionic components of the superfields $H^0_{u,d}(S)$ and ν_i [$i = 1,2,3$] denote the neutrinos. In Eq. (3), the neutralino mass matrix is given, in a generic basis (where $\langle\tilde{\nu}_i\rangle \equiv v_i \neq 0$, $\mu_i \neq 0$ and $\lambda_i \neq 0$), by

$$\mathcal{M}_{\tilde{\chi}^0} = \begin{pmatrix} \mathcal{M}_{\text{NMSSM}} & \xi^T_{\not{R}_p} \\ \xi_{\not{R}_p} & 0_{3\times 3} \end{pmatrix}, \tag{4}$$

where $\mathcal{M}_{\text{NMSSM}}$ is the neutralino mass matrix corresponding to the NMSSM. While writing the latter mass matrix, we assume $v_i \ll v_{u,d}$ (the exact expression of M_Z being given by $v^2 = v_u^2 + v_d^2 + \sum_{i=1}^3 v_i^2 = 2c^2_{\theta_W} M_Z^2/g^2 \simeq (246/\sqrt{2}\ \text{GeV})^2$). Also, we use s and c to stand for sine and cosine, respectively.

$$\mathcal{M}_{\text{NMSSM}} = \begin{pmatrix} M_1 & 0 & -M_Z s_{\theta_W} c_\beta & M_Z s_{\theta_W} s_\beta & 0 \\ 0 & M_2 & M_Z c_{\theta_W} c_\beta & -M_Z c_{\theta_W} s_\beta & 0 \\ -M_Z s_{\theta_W} c_\beta & M_Z c_{\theta_W} c_\beta & 0 & -\mu & -\lambda v_u \\ M_Z s_{\theta_W} s_\beta & -M_Z c_{\theta_W} s_\beta & -\mu & 0 & -\lambda v_d + \sum_{i=1}^3 \lambda_i v_i \\ 0 & 0 & -\lambda v_u & -\lambda v_d + \sum_{i=1}^3 \lambda_i v_i & 2\kappa\langle s\rangle \end{pmatrix}. \tag{5}$$

Above, M_1 (M_2) is the soft supersymmetry breaking mass of the bino (wino), $\tan\beta = v_u/v_d = \langle h_u^0\rangle/\langle h_d^0\rangle$, and $\mu = \lambda\langle s\rangle$. We assume for simplicity that λ, κ and the soft supersymmetry breaking parameters are all real.

In Eq. (4), $\xi_{\not{R}_p}$ is the $\not{R}_p$ part of the matrix mixing neutrinos and neutralinos:

$$\xi_{\not{R}_p} = \begin{pmatrix} -\frac{g' v_1}{\sqrt{2}} & \frac{g v_1}{\sqrt{2}} & 0 & \mu_1 + \lambda_1\langle s\rangle & \lambda_1 v_u \\ -\frac{g' v_2}{\sqrt{2}} & \frac{g v_2}{\sqrt{2}} & 0 & \mu_2 + \lambda_2\langle s\rangle & \lambda_2 v_u \\ -\frac{g' v_3}{\sqrt{2}} & \frac{g v_3}{\sqrt{2}} & 0 & \mu_3 + \lambda_3\langle s\rangle & \lambda_3 v_u \end{pmatrix}, \tag{6}$$

g and g' being the SU(2) and U(1) gauge couplings.

Effective neutrino mass matrix. We restrict ourselves to the situation where $v_i/v_{u,d} \ll 1$ (as before), $|\mu_i/\mu| \ll 1$ and $|\lambda_i/\lambda| \ll 1$ so that (i) no considerable modifications of the NMSSM scalar potential are induced by the additional bilinear and trilinear term in superpotential (2), (ii) the neutrino–neutralino mixing is suppressed, leading to sufficiently small neutrino masses as shown later, and (iii) the effective neutrino mass matrix can be written to a good approximation by the following see-saw type structure,

$$m_\nu = -\xi_{\not{R}_p}\mathcal{M}^{-1}_{\text{NMSSM}}\xi^T_{\not{R}_p}. \tag{7}$$

From Eqs. (5)–(7), we deduce the analytical expression of the effective Majorana neutrino mass matrix:

$$m_{\nu_{ij}} = \frac{1}{|\mathcal{M}_{\text{NMSSM}}|}\big[\mu_i\mu_j\mathcal{F} + (\mu_i\Lambda_j + \mu_j\Lambda_i)\mathcal{G} + \Lambda_i\Lambda_j\mathcal{H}\big], \tag{8}$$

where $|\mathcal{M}_{\text{NMSSM}}|$ is the determinant of matrix (5), $\Lambda_i = \langle s\rangle(\lambda_i + \lambda\frac{v_i}{v_d})$ and

$$\mathcal{F} = \lambda^2 v_u^2 M_1 M_2 + \mathcal{X}, \qquad \mathcal{G} = \mathcal{X} + \left(\lambda v_d - \cos^2\beta\sum_{i=1}^3\lambda_i v_i\right)\mathcal{Y}, \qquad \mathcal{H} = \mathcal{X} + 2\cos^2\beta\left(\lambda v_d - \sum_{i=1}^3\lambda_i v_i\right)\mathcal{Y}, \tag{9}$$

with

$$\mathcal{X} = 2\cos^2\beta\kappa\langle s\rangle M_Z^2\big(\cos^2\theta_W M_1 + \sin^2\theta_W M_2\big), \qquad \mathcal{Y} = \frac{v_u}{\langle s\rangle}M_Z^2\big(\cos^2\theta_W M_1 + \sin^2\theta_W M_2\big).$$

One should notice that the neutrino mass matrix (8) arises entirely at the tree level.

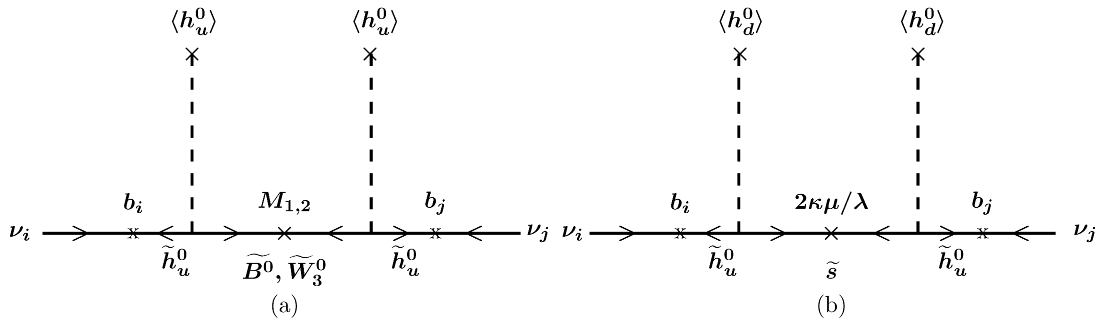

Fig. 1. Tree level Feynman diagrams in the $R\!\!\!/_p$ NMSSM generating Majorana neutrino masses proportional to $b_i b_j$. The effective $R\!\!\!/_p$ bilinear parameter b_i stands for either μ_i or $\lambda_i\langle s\rangle$. The mass-dimensional couplings appearing at the two vertices are of the type $m = M_Z t(\theta_W)\sin\beta$ [$t(\theta_W) = \sin\theta_W$ for $\tilde{B}^0$ and $t(\theta_W) = -\cos\theta_W$ for $\tilde{W}_3^0$] in (a) and $m = -\lambda v_d$ in (b). A cross indicates either a mass insertion or a vev. The arrows show the flow of momentum for the associated propagators.

The emergence of the specific mass matrix structure of Eq. (8) at tree level is the primary result of our analysis. Accordingly to this particular structure, if the two (effective) quantities μ_i and Λ_i (characteristic of the NMSSM) take simultaneously non-vanishing values, then the mass matrix ceases to be of rank 1, even though the determinant is still zero. In this situation, we get two non-zero neutrino mass eigenvalues. It is worth comparing the situation with what happens in the MSSM with bilinear $R\!\!\!/_p$ violation. In the latter case, we get a similar kind of analytic structure of the mass matrix from the simultaneous consideration of the μ_i as well as the corresponding soft B_i terms. While the $\mu_i\mu_j$ contributions originate at tree level, the $\mu_i B_j$ and $B_i B_j$ contributions arise at one-loop order through Grossman–Haber diagrams [12] which proceed through slepton-Higgs and neutrino–neutralino mixings (for a series of analysis in a three-flavour framework, see [13]). The Grossman–Haber loops would still contribute in our scenario, but now that we have two tree level masses, those loop-suppressed contributions are not so crucial for generating a viable neutrino mass spectrum.

Neutrino mass eigenvalues at tree level. The eigenvalues of the effective neutrino mass matrix (8) turn out to be $\{0, m_\nu^-, m_\nu^+\}$ with

$$m_\nu^\pm = \frac{1}{2|\mathcal{M}_{\mathrm{NMSSM}}|}\Bigg[\left(\sum_{i=1}^3 \mu_i^2\right)\mathcal{F} + \left(\sum_{i=1}^3 \Lambda_i^2\right)\mathcal{H} + 2\left(\sum_{i=1}^3 \mu_i\Lambda_i\right)\mathcal{G}$$
$$\pm\left\{\left[\left(\sum_{i=1}^3 \mu_i^2\right)\left(\sum_{i=1}^3 \Lambda_i^2\right) - \left(\sum_{i=1}^3 \mu_i\Lambda_i\right)^2\right]\mathcal{I} + \left[\left(\sum_{i=1}^3 \mu_i^2\right)\mathcal{F} + \left(\sum_{i=1}^3 \Lambda_i^2\right)\mathcal{H} + 2\left(\sum_{i=1}^3 \mu_i\Lambda_i\right)\mathcal{G}\right]^2\right\}^{1/2}\Bigg], \tag{10}$$

where $\mathcal{F}, \mathcal{G}, \mathcal{H}$ are defined in Eq. (9), and $\mathcal{I} = 4(\mathcal{G}^2 - \mathcal{F}\mathcal{H})$.

Note that for either all $\mu_i = 0$ or all $\Lambda_i = 0$, the eigenvalue m_ν^- vanishes as expected since in this limit we recover the rank 1 form. An inspection of the form of Eq. (10) reveals that the coefficient of $\mathcal{I}$ can be written as $\sum_{i\neq j}(\mu_i\Lambda_j - \mu_j\Lambda_i)^2$, which indicates that the misalignment between μ_i and Λ_i is crucial in creating a non-vanishing m_ν^-.

Therefore, the condition for generating two non-vanishing and non-degenerate eigenvalues is to ensure $\mu_i \neq 0$ and $\Lambda_i \neq 0$ simultaneously. In other words, to achieve two non-zero eigenvalues, μ_i has to be non-zero always, but we can go to a basis of L_α fields where $v_i = 0$ (then generally $\lambda_i \neq 0$), or to the other extreme where $\lambda_i = 0$ (but $v_i \neq 0$), or any other basis in between basically maintaining $\Lambda_i \neq 0$. On the contrary, if $\mu_i = 0$, only one neutrino eigenstate gets a mass different from zero, as was also found by the authors of Ref. [5]. But all the scenarios we considered in this Letter, according to the above arguments, will yield two non-zero neutrino masses at tree level. We mention that within the first scenario, a rotation on the L_α fields has already been performed, whereas in the third one, no rotation is possible.

In Figs. 1 and 2, we present the Feynman diagrams contributing to the Majorana neutrino mass (8). All these diagrams proceed through the tree level exchange of the neutralinos (gauginos, higgsinos and singlino). In these figures, we have considered the basis corresponding to $v_i = 0$ for simplicity.

4. Numerical results

Thus the present model predicts a hierarchical neutrino mass spectrum at tree level. This hierarchical pattern could be modified by the loop level contributions to neutrino masses. At least, the massless state acquires a mass once the loop contributions are considered.

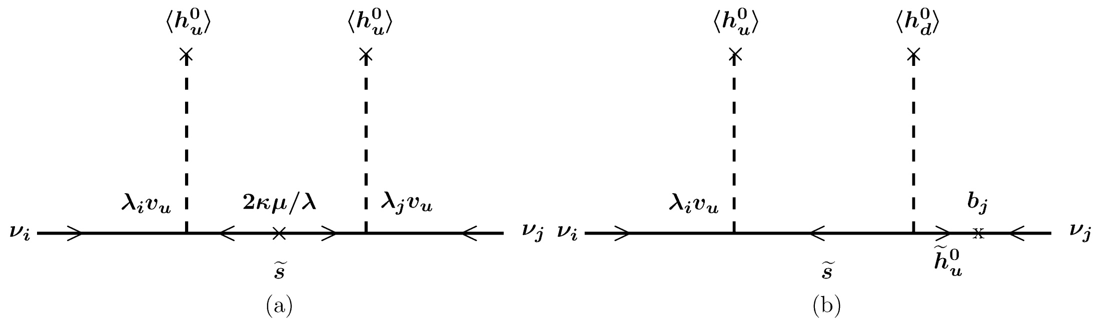

Fig. 2. Tree level Feynman diagrams in the $\not{R}_p$ NMSSM generating Majorana neutrino masses proportional to $\lambda_i\lambda_j v_u^2$ and $\lambda_i v_u b_j$. The effective $\not{R}_p$ bilinear parameter b_i, as in Fig. 1, stands for either μ_i or $\lambda_i\langle s\rangle$. The mass-dimensional coupling appearing at the right vertex in (b) is of the type $m = -\lambda v_d$.

It is not our aim to give detailed numerical fits in this short Letter. We would just like to numerically demonstrate that our scenario can reproduce the neutrino data from the tree level neutrino mass matrix, with a choice of NMSSM parameters that corroborate the μ-naturalness.

As an example, we take the following NMSSM parameters:

$$\lambda = 0.4, \qquad \kappa = 0.3, \qquad \mu = -200\ \text{GeV}, \qquad \tan\beta = 30, \qquad M_1 = 350\ \text{GeV}, \qquad M_2 = 500\ \text{GeV}. \tag{11}$$

These parameters satisfy the general NMSSM constraints described below. These constraints are not expected to be significantly modified by the presence of $\not{R}_p$ interactions in superpotential (2) as we work under the assumption $|\mu_i/\mu| \ll 1$ and $|\lambda_i/\lambda| \ll 1$.

The μ-naturalness forces one to restrict to $\langle s\rangle \lesssim 10\text{TeV}$, which translates into the condition $|\mu|[\text{GeV}] \times 10^{-4} \lesssim \lambda$. Furthermore, the absence of Landau singularities, for λ, κ, the top and bottom Yukawa coupling constants below the GUT energy scale, imposes [4] the typical bounds: $\lambda \lesssim 0.75$, $|\kappa| \lesssim 0.65$ and $1.7 \lesssim \tan\beta \lesssim 54$. Finally, the LEP bound on the lightest chargino mass, namely $m_{\tilde{\chi}_1^+} > 103.5$ GeV [14], translates into $|\mu| \gtrsim 100$ GeV.

Together with the values in Eq. (11), we take the following $\not{R}_p$ effective couplings,

$$\begin{aligned} &\mu_1 = 1 \times 10^{-4}\ \text{GeV}, \qquad \mu_2 = 2 \times 10^{-4}\ \text{GeV}, \qquad \mu_3 = 2 \times 10^{-4}\ \text{GeV}, \\ &\Lambda_1/\langle s\rangle = 1.9 \times 10^{-5}, \qquad \Lambda_2/\langle s\rangle = 1.4 \times 10^{-5}, \qquad \Lambda_3/\langle s\rangle = 1.5 \times 10^{-5}. \end{aligned} \tag{12}$$

This set of parameters yield the following three neutrino mass eigenvalues at tree level:

$$m_{\nu_1} = 0, \qquad m_{\nu_2} = 0.0095\ \text{eV}, \qquad m_{\nu_3} = 0.058\ \text{eV}. \tag{13}$$

These values are in agreement with the three-flavour analyzes including results from solar, atmospheric, reactor and accelerator oscillation experiments which lead to (4σ level): $6.8 \leqslant \Delta m_{21}^2 \leqslant 9.3\ [10^{-5}\ \text{eV}^2]$ and $1.1 \leqslant \Delta m_{31}^2 \leqslant 3.7\ [10^{-3}\ \text{eV}^2]$ [2]. Besides, the neutrino mass eigenvalues in (13) satisfy the bound extracted from WMAP and 2dFGRS galaxy survey (depending on cosmological priors): $\sum_{i=1}^3 m_{\nu_i} \lesssim 0.7$ eV [15]. Finally, these eigenvalues are perfectly compatible with the limits extracted from the tritium beta decay experiments (95% C.L.): $m_\beta \leqslant 2.2$ eV [Mainz] and $m_\beta \leqslant 2.5$ eV [Troitsk] [16], this effective mass being defined as $m_\beta^2 = \sum_{i=1}^3 |U_{ei}|^2 m_{\nu_i}^2$ where U_{ei} is the leptonic mixing matrix.

Although we have chosen a particular set of input parameters for illustration, solutions exist over a wide range of parameter space. More realistic estimates can be obtained by switching on the soft scalar terms $B_i\tilde{\ell}_i h_u + \text{h.c.}$ plus the superpotential terms $\lambda_{ijk} L_i L_j E_k^c$ and $\lambda'_{ijk} L_i Q_j D_k^c$. All these terms contribute to the neutrino mass matrix at one-loop order. A combined fit of all these parameters is beyond the scope of this Letter.

5. Conclusion

In the NMSSM, which is a gauge singlet extension of the MSSM addressing the μ-naturalness, *two non-vanishing neutrino mass eigenvalues can arise at tree level* when the lepton number violating bilinear terms $\mu_i L_i H_u$ are present. One can then explain the neutrino oscillation data without essentially depending on the loop-generated masses which otherwise bring in more uncertainties from unknown dynamics. This result is in contrast with any other supersymmetric $\not{R}_p$ scenario, as those scenarios do not generate more than one massive neutrino eigenstate at tree level (except the scenario proposed in [7] where 3 right-handed neutrinos are added to the field content).

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