



Non-renormalizable operators for solar neutrino mass generation in Split SuSy with bilinear R-parity violation

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Abstract

The Minimal Supersymmetric Extension of the Standard Model (MSSM) is able to explain the current data from neutrino physics. Unfortunately Split Supersymmetry as low energy approximation of this theory fails to generate a solar square mass difference, including after the addition of bilinear R-Parity Violation. In this work, it is shown how one can derive an effective low energy theory from the MSSM in the spirit of Split Supersymmetry, which has the potential of explaining the neutrino phenomenology. This is achieved by going beyond leading order in the process of integrating out heavy scalars from the original theory, which results in non-renormalizable operators in the effective low energy theory. It is found that in particular a $d = 8$ operator is crucial for the generation of the neutrino mass differences.

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1. Introduction

The Standard Model (SM) of particle physics provides a theoretical explanation for the experimental evidence on weak interactions and the masses of the corresponding mediators. This last point has its climax with the discovery of the Higgs boson in 2012 [1,2]. Even though these

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evidences have been important in the historical developing of the SM, the theory is far from complete since neutrino masses and Dark Matter appear to be issues that cannot be explained by the original theory.

The proposal of Weinberg's $d = 5$ operator [3] points towards a way to extend the SM in order to include an explanation for neutrino masses, such that realizations of this operator at tree level and one loop can be achieved. Other proposals can be found in [4,5]. In this context, in order to obtain Majorana neutrino masses it is necessary to have some mechanism to provide lepton number violation as well. When exploring supersymmetric extensions of the SM, lepton number and baryon number can be incorporated in an *ad hoc* symmetry called *R-Parity*. Amongst the proposals in this area, we are interested exclusively in the addition of bilinear R-parity violation to the minimal supersymmetric extension of the SM (MSSM) as a way to study neutrino physics.

On top of that, within the working hypothesis of a supersymmetric extension of the SM, at the present time data from ATLAS and CMS have no direct evidence for the observation of *sparticles* [6]. This seems to indicate that the supersymmetric scalar particles may be too massive to be produced at the LHC. In this context it is appealing to study models like Split SuSy (SS) [7,8], where the claim is that the sfermions and one Higgs doublet are very heavy and therefore they are decoupled from the low-energy theory. In the original proposal, sfermions are pushed to live at scales in the order of 10^{10} GeV where the testing of sfermions at present colliders turns out to be very difficult. Moreover, it leads to an *unnatural* model since corrections to Higgs mass must be very fine tuned in order to give rise to a Higgs of 125 GeV. However, such problems can be avoided for Split SuSy if the sfermion masses lie around 10^4 GeV. Such a scale allows for a Higgs mass determined by experiments, and thus having the possibility to test effects of heavy sfermions in the present time experiments [8,9].

The MSSM plus BRpV is able to generate two neutrino mass eigenvalues at loop level. If we combine Split SuSy and BRpV in one single model, though, the decoupling of scalars retaining only the operators with lowest mass dimension (renormalizable), leaves a theory that cannot explain the observed neutrino oscillations, since it cannot produce a solar neutrino mass scale. Thus, one faces a problem, on the one hand it seems that Split SuSy is favored from collider phenomenology, while on the other hand it seems to be disfavored by neutrino oscillation data [10]. One alternative to solve this tension is to include new physics in terms of gravity motivated operators in order to achieve neutrino masses [11,12]. Nevertheless, it would be much more attractive to solve this problem within Split SuSy as a low energy approximation to the MSSM.

The origin of this tension between collider and neutrino phenomenology within the simplest version of Split SuSy can be understood in the light of the underlying supersymmetric theory. A quite general theorem states that within this theory the existence of a tree level neutrino Majorana mass implies the appearance of a B-L violating superpartner *Majorana*-like mass term and vice versa [13]. This is exactly what happens in the MSSM with RpV terms. Thus, within conventional Split SuSy it is actually not possible to accommodate both neutrino and collider phenomenology. The problem seems to be that after integrating out the scalars and retaining the resulting renormalizable operators only, the interplay between the neutrino Majorana mass and the superpartner scalars is lost, and with it, the generation of a second non-zero neutrino mass.

The aim of this paper is to investigate the possibility to accommodate neutrino oscillation physics within Split SuSy in a self consistent way (i.e. with no addition of operators motivated from physics outside the MSSM as a UV complete theory) when going beyond the leading term after integrating out the scalar fields. Other alternatives within this context have been investigated (see for instance [14] and [15]), though, we will stick ourselves to original proposal of

Split SuSy and BRpV. Clearly it would be favorable to be able to explain the solar mass difference without the incorporation of new physics and new couplings. In this spirit it is shown that already retaining next to leading operators from the original MSSM Lagrangian gives rise to non-renormalizable higher dimensional operators in the effective low energy theory, which can reconcile the Split SuSy with neutrino oscillation data.

2. Split Supersymmetry from minimal integration of scalars

In Split Supersymmetry [7] all sfermions together with one heavy Higgs doublet have a large mass, which is taken for simplicity to be degenerate at the scale \tilde{m} .

In this framework, since the effects of sfermions appear at very large scale, the hierarchy problem is not solved. It is argued that gauge unification is still achieved, and a Dark Matter candidate is still present (the neutralino) when R-Parity is conserved. In addition, large flavor and CP violation are avoided. This type of model is supported by the fact the LHC collaborations do not see sfermions up to a mass near the TeV scale [6]. Split SuSy can be obtained from the MSSM which includes two Higgs doublets H_u and H_d [16,17]. The Higgs potential in the MSSM is

$$V = \left(m_{H_u}^2 + |\mu|^2\right) H_u^\dagger H_u + \left(m_{H_d}^2 + |\mu|^2\right) H_d^\dagger H_d - B_\mu \left(H_u^T i\sigma_2 H_d + h.c.\right) + \frac{1}{8}(g^2 + g'^2) \left(H_u^\dagger H_u - H_d^\dagger H_d\right)^2 + \frac{1}{2}g^2 |H_u^\dagger H_d|^2, \quad (1)$$

where $m_{H_u}^2$ and $m_{H_d}^2$ are soft mass parameters corresponding to the H_u and H_d bosons, μ is the supersymmetric Higgs mass parameter, and B_μ is the soft Higgs mixing parameter (with units of mass square).

It is assumed that one Higgs doublet, H_{SM} , remains light and resembles the SM Higgs doublet. The second one, H_{SS} , is heavy and for simplicity its components have a degenerate mass equal to \tilde{m} . This doublet comes from a mixing of H_u and H_d , given by

$$\begin{bmatrix} H_{SM} \\ i\sigma_2 H_{SS}^* \end{bmatrix} = \begin{bmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{bmatrix} \begin{bmatrix} H_u \\ i\sigma_2 H_d^* \end{bmatrix}, \quad (2)$$

where the mixing angle α is reminiscent of the MSSM. In the decoupling limit of the MSSM (equivalent to the Split Supersymmetric case) we have $\sin \alpha = -\cos \beta$ and $\cos \alpha = \sin \beta$, under the usual conventions for the angles in the MSSM [17]. In SS, the angle $\alpha = \beta - \pi/2$ is such that there is no mixing between H_{SM} and H_{SS} . The mass terms for the two doublets are $m_h \approx 125$ GeV and \tilde{m} respectively.

The transition from the MSSM to Split SuSy is achieved by integrating out the heavy scalars by solving their equation of motion. For instance, the heavy Higgs field can be replaced by

$$H_{SS} \approx -\frac{1}{\tilde{m}^2} \frac{c_\alpha}{\sqrt{2}} \left[g(\sigma \cdot \tilde{W}) - g' \tilde{B} \right] \tilde{H}_d + \frac{1}{\tilde{m}^2} \frac{s_\alpha}{\sqrt{2}} i\sigma_2 \left[g(\sigma^* \cdot \tilde{W}) + g' \tilde{B} \right] \tilde{H}_u + \dots \quad (3)$$

In the original proposal of Split SuSy this scalar was also taken to be heavy, therefore the limit $\tilde{m} \rightarrow \infty$ was invoked, which decouples completely the field H_{SS} . After replacing (3) in the Lagrangian of the MSSM and after decoupling the heavy Higgs, the R-Parity conserving part of the Split Supersymmetric Lagrangian reads [7–9]

$$\begin{aligned} \mathcal{L}_{RpC}^{split} = & \mathcal{L}_{kinetic}^{split} + m^2 H^\dagger H - \frac{\lambda}{2} (H^\dagger H)^2 - \left[Y_u \bar{q}_L u_R i\sigma_2 H^* + Y_d \bar{q}_L d_R H + \right. \\ & + Y_e \bar{l}_L e_R H + \frac{M_3}{2} \tilde{G} \tilde{G} + \frac{M_2}{2} \tilde{W} \tilde{W} + \frac{M_1}{2} \tilde{B} \tilde{B} + \mu \tilde{H}_u^T i\sigma_2 \tilde{H}_d + \\ & \left. + \frac{1}{\sqrt{2}} H^\dagger (\tilde{g}_u \sigma \tilde{W} + \tilde{g}'_u \tilde{B}) \tilde{H}_u + \frac{1}{\sqrt{2}} H^T i\sigma_2 (-\tilde{g}_d \sigma \tilde{W} + \tilde{g}'_d \tilde{B}) \tilde{H}_d + \text{h.c.} \right], \end{aligned} \tag{4}$$

where for simplicity we have called H the SM Higgs doublet. This Lagrangian contains the mass and quartic parameters of the light Higgs potential, the Yukawa couplings, the gaugino mass parameters, the supersymmetric higgsino mass parameter, and four Higgs–gaugino–higgsino couplings \tilde{g} , by-products of the decoupling of the scalars.

The original Split Supersymmetric model can be extended by adding bilinear R-Parity violating interactions,

$$\mathcal{L}_{RpV}^{split} = \epsilon_i \tilde{H}_u^T i\sigma_2 L_i - \frac{1}{\sqrt{2}} a_i H^T i\sigma_2 (-\tilde{g}_d \sigma_j \tilde{W}^j + \tilde{g}'_d \tilde{B}) L_i + \text{h.c.} \tag{5}$$

The terms proportional to ϵ_i come from the usual bilinear interaction, while the terms proportional to a_i appear after integrating out the sleptons and retaining renormalizable operators [10]. At this point, since higher order terms have been neglected, all the terms of the effective low energy Lagrangian are renormalizable (indicated by the subscript r). The neutralinos and neutrinos mix themselves by means of the following mass matrix,

$$\mathcal{M}_r^{SS} = \begin{bmatrix} \mathbf{M}_{\chi^0,r}^{SS} & (m_r^{SS})^T \\ m_r^{SS} & 0 \end{bmatrix}, \tag{6}$$

where the neutralino submatrix is given by

$$\mathbf{M}_{\chi^0,r}^{SS} = \begin{bmatrix} M_1 & 0 & -\frac{1}{2} \tilde{g}'_d v & \frac{1}{2} \tilde{g}'_u v \\ 0 & M_2 & \frac{1}{2} \tilde{g}_d v & -\frac{1}{2} \tilde{g}_u v \\ -\frac{1}{2} \tilde{g}'_d v & \frac{1}{2} \tilde{g}_d v & 0 & -\mu \\ \frac{1}{2} \tilde{g}'_u v & -\frac{1}{2} \tilde{g}_u v & -\mu & 0 \end{bmatrix}, \tag{7}$$

and the mixing between neutralinos and neutrinos is defined by the submatrix,

$$m_r^{SS} = \begin{bmatrix} -\frac{1}{2} \tilde{g}'_d a_1 v & \frac{1}{2} \tilde{g}_d a_1 v & 0 & \epsilon_1 \\ -\frac{1}{2} \tilde{g}'_d a_2 v & \frac{1}{2} \tilde{g}_d a_2 v & 0 & \epsilon_2 \\ -\frac{1}{2} \tilde{g}'_d a_3 v & \frac{1}{2} \tilde{g}_d a_3 v & 0 & \epsilon_3 \end{bmatrix}. \tag{8}$$

Here, v is the vacuum expectation value of the light Higgs field. If one defines $\lambda_i = a_i \mu + \epsilon_i$ and block-diagonalizes the matrix in eq. (6) one obtains the well known effective neutrino mass matrix,

$$\mathbf{M}_v^{eff} = \frac{v^2}{4 \det \mathbf{M}_{\chi^0,r}^{SS}} \left(M_1 \tilde{g}_d^2 + M_2 \tilde{g}_d'^2 \right) \begin{bmatrix} \lambda_1^2 & \lambda_1 \lambda_2 & \lambda_1 \lambda_3 \\ \lambda_2 \lambda_1 & \lambda_2^2 & \lambda_2 \lambda_3 \\ \lambda_3 \lambda_1 & \lambda_3 \lambda_2 & \lambda_3^2 \end{bmatrix}, \tag{9}$$

which gives the atmospheric neutrino mass scale. For notational purposes we call the matrix elements of this 3×3 neutrino mass matrix $\mathbf{M}_{v,ij}^{eff} = A \lambda_i \lambda_j$. The situation with Split Supersymmetry is that one loop corrections give $\Delta \mathbf{M}_{v,ij}^{eff} = (\Delta A) \lambda_i \lambda_j$, this is, a term proportional to the

tree level value. Thus, loop contributions in SS do not induce a second neutrino mass scale. In the MSSM with BRpV something different happens, since loop contributions additionally induce a term proportional to $\epsilon_i \epsilon_j$ [18] and in accordance with the result in ref. [13]. In order to obtain a second mass scale in Split Supersymmetry one alternative is to include extra contributions which go beyond the original SS or MSSM field content. For example an extra contribution from gravity has been studied in [12]. We consider here a different alternative, namely the effect of neglected contributions from non-renormalizable operators that arise from the original MSSM Lagrangian.

Before working out the relevant non-renormalizable terms, we mention the fact that in Split Supersymmetry the value of the light Higgs mass near 125 GeV implies an upper bound on the scale \tilde{m} of the order of 10^5 – 10^6 GeV [8,9], in other words, the value of the light Higgs mass implies a rather low value of the Split SuSy scale.

3. Split Supersymmetry beyond minimal integration of scalars

Like in the previous section, the starting point is the MSSM Lagrangian with the Higgs potential (1). However, now the integration of heavy scalars will not be truncated at the first contribution. For example, for the heavy Higgs the equation of motion is given by eq. (3). Instead of sending $\tilde{m} \rightarrow \infty$ which corresponds to simply erasing the H_{SS} from the Lagrangian we now retain a finite (but large) value of \tilde{m} and replace (3) back into the Lagrangian (1). In this way, new operators arise as a consequence of the finite mass of the heavy scalars. One of those terms is

$$\mathcal{L}^{nr} \ni \frac{1}{8}(g^2 + g'^2) \frac{c_\alpha^2}{\tilde{m}^4} \left\{ H_{SM}^T i \sigma_2 [g (\sigma \cdot \tilde{W}) - g' \tilde{B}] \tilde{H}_d \right\}^2, \quad (10)$$

which is a non-renormalizable dimension 8 operator (from now on, the script *nr* refers to “non-renormalizable”). The value of the coupling is given at \tilde{m} , and it leads to a neutrino mass contribution that will be developed in the next section. Notice that the bilinear R-parity violating mixing in matrix (6) allows us to convert from a \tilde{H}_d to a neutrino after block diagonalization [18,19]. Analogously one finds the following $d = 7$ operator when integrating out down squarks

$$\mathcal{L}^{nr} \ni -\frac{A_d Y_d^3}{m_Q^2 m_d^2} \left(c_\alpha H_{SM}^T i \sigma_2 \tilde{H}_d d_R \right) (\bar{Q}_i \sigma_2 \tilde{H}_d^c) + h.c. \quad (11)$$

Similar operators appear after integrating out sleptons and the procedure invoked is the same (although with small differences). These two kinds of operators (10), (11) will induce corrections that have been claimed to be, jointly with the neutralino–neutrino corrections, the largest contributions to the solar mass difference [19,20]. Notice that in order to write down a similar expression for the sleptons we just need to do the replacement: $d_R \rightarrow e_R$, $Q_i \rightarrow L_i$, and the subindex $d \rightarrow e$ (but not for \tilde{H}_d).

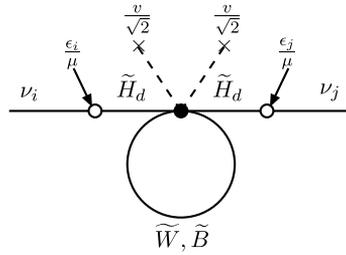
In order to write down this last coupling, we perform the decoupling of squarks in the soft trilinear sector of the MSSM instead the Higgs potential, as we did with the $d = 8$ operator (10). In the MSSM this term allows us to write the down–down loop. Nevertheless, equation (11) (and the sleptonic analogue) introduce a dependence on the trilinear scalar couplings A_d (A_e , for sleptons).

The non-renormalizable operators given in eqs. (10) and (11) lay the seed for the generation of the needed neutrino masses due to quantum corrections. This mechanism will be explained in the following section.

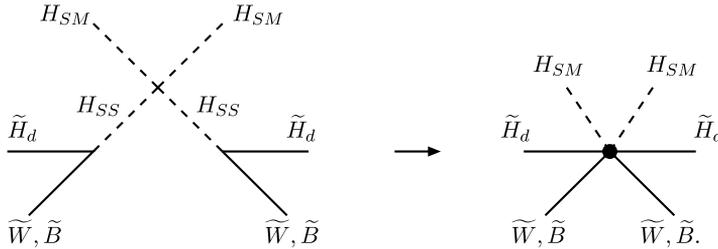
4. Loop corrections

4.1. Calculating the loops

The non-renormalizable terms in the Lagrangian given in eq. (10) lead to the following graph contributing to the neutrino mass matrix,



The non-renormalizable vertex is represented in the graph by a full circle at the center. The open circles represent the mixing between the neutrinos and the down type higgsino due to R-Parity violation. The central vertex of this diagram arises from the corresponding diagram in the underlying supersymmetric theory shown in the figure below

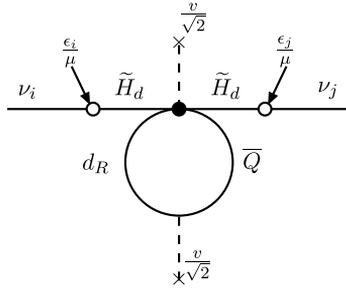


The SM Higgs field acquires a vacuum expectation value v . Inside the loop we can have winos or bins. The mass contributions of those two loops can be summarized as,

$$\Delta M_v^{nr} = \frac{1}{256\pi^2} (g^2 + g'^2) \frac{v^2 s_\beta^2}{\tilde{m}^4} \left[g^2 M_2 \tilde{A}_0(M_2^2) + g'^2 M_1 \tilde{A}_0(M_1^2) \right] \frac{\epsilon_i \epsilon_j}{\mu^2}, \tag{12}$$

where we see the dimension 8 vertex, the two vevs of the light Higgs field, the mixing between neutrinos and the down type higgsino, the loop factor and the finite Veltman’s function \tilde{A}_0 . Since the solar neutrino mass scale is very small, the contribution from (12) can be important, even though it is suppressed by a factor of $1/\tilde{m}^4$. An estimation of the orders of magnitudes in eq. (12) confirms that this contribution can do the job. In the next subsection this contribution will be evaluated numerically.

As it was indicated in the previous section, there are two other contributions that we studied in detail: A loop with down quarks, and a loop of leptons which are produced after the decoupling of down squarks and sleptons. In these two cases, the contribution can be related in a similar way to the diagram:



The corresponding $d = 7$ operator [see eq. (11)] gives rise to a numerically irrelevant contribution to the neutrino mass matrix. This happens despite of the fact that the bottom-sbottom loops are in many cases important in the usual MSSM+BRpV models. The reason is that when sbottom quarks are decoupled with mass \tilde{m} , the contribution from the corresponding non-renormalizable operator becomes irrelevant due to the smallness of the bottom quark mass compared to the mass of the gauginos.

In summary, in order to write down the corrections for neutrino masses, we use the notation for the 3×3 neutrino corrected mass matrix in the form of

$$\mathbf{M}_\nu^{eff} = A\lambda_i\lambda_j + C\epsilon_i\epsilon_j \tag{13}$$

with,

$$A = \frac{v^2}{4 \det M_{\chi^0, r}^{SS}} \left(M_1 \tilde{g}_d^2 + M_2 \tilde{g}'^2 \right)$$

$$C = \frac{1}{256\pi^2} (g^2 + g'^2) \frac{v^2 s_\beta^2}{\tilde{m}^4 \mu^2} \left[g^2 M_2 \tilde{A}_0(M_2^2) + g'^2 M_1 \tilde{A}_0(M_1^2) \right] \tag{14}$$

as can be read from eqs. (9) and (12). It is by virtue of a the second term that one can expect an additional non-vanishing neutrino mass scale.

Notice that the contribution in eq. (12) is in principle not finite (if we imagine replacing \tilde{A}_0 by A_0), as opposite to the renormalizable case. However, there is no reason for concern here, since the inclusion of a counterterm that absorbs this infinity must be considered after the decoupling of the heavy degrees of freedom. A similar case, where loops finite in a renormalizable theory, are turned into infinite in the effective theory can be found at the reference [21]. In addition, it is worthwhile to mention that value of the scale \tilde{m} between 10^3 – 10^5 GeV helps to realize the radiative mechanism for neutrino masses when one wants to compute the Wilson coefficients for Weinberg operator via new degrees of freedom added to the SM [22].

4.2. Numerical results

Our intention is to explore whether the non-renormalizable operators account for neutrino physics or their scale suppression makes the neutrino phenomenology unfeasible. Therefore, if a positive result is found, it is sufficient to perform a scan in a window of the parameter space. In order to see how a solar squared mass difference arises from the effect of the loop, we computed the neutrino squared mass differences and the corresponding mixing angles by using the equation (13). The atmospheric square mass difference can be generated via the λ parameters as it happens in the models that consider bilinearly violated R-parity.

Table 1
Ranges for parameters with mass units.

Value	Min. [GeV]	Max. [GeV]
M_χ	500	1500
M_1	M_χ	$1.5 M_\chi$
M_2	M_χ	$1.5 M_\chi$
M_3	$10^{3.5}$	10^5
μ	M_χ	$1.5 M_\chi$
ϵ_1	-10.0	10.0
ϵ_2	-10.0	10.0
ϵ_3	-10.0	10.0
\tilde{m}	$10^{3.5}$	10^5

Table 2
Ranges for parameters without units.

Value	Min.	Max.
$\tan \beta$	1	45
λ_1	-1.0	1.0
λ_2	-1.0	1.0
λ_3	-1.0	1.0

We also implemented the condition that a Higgs mass is within the zone allowed by the experiments [23]. This point is crucial since, as it was indicated above, the Higgs mass is one of the most stringent parameters in order to define the scale where new physics appears. It turns out that this condition restricts the SuSy mass scale to be of the order of $\tilde{m} \sim 10^3\text{--}10^5$ GeV in order to avoid an extreme fine tuning. We also imposed the condition that the spectrum of supersymmetric particles fulfills the experimental constraints from the supersymmetric searches [24]. The ranges where we varied the parameters of the model are depicted on Tables 1 and 2.

From the Table 1, we can highlight that the value of M_3 , this is, the Gluino mass, can be safely put at high scales, since the Gluino does not interact with SM particles in this context (the squarks have been decoupled), there is no reason to keep it at low scales. Two comments about the λ_i parameters are at place here. First, on the contrary to the notation of other BRpV models, in this case the λ_i have no units. Second, the values of λ_i have been chosen in order to fulfill the atmospheric square mass difference. This was achieved by choosing the parameters once the A factor on eq. (13) has been computed. When performing the scan, we obtained the full spectrum of particles which is in agreement with the constraints from the searches for supersymmetric particles [1]. The limiting values of the obtained spectra are shown in the Table 3. As mentioned above, it is imposed to have a spectrum where the masses of neutralinos and charginos are below the scale \tilde{m} , otherwise it would be inconsistent to integrate out other particles with masses of the order of \tilde{m} .

The main result of this section is that the parameter space allows to meet the neutrino and Higgs requirement and the collider bounds on SuSy masses. One notices a strong correlation between the parameters, but the good scenarios are not accumulated around an isolated point in the parameter space. Fig. 1 shows $|\tilde{\epsilon}|$ as a function of $\tan \beta$. One notes that the points indicating allowed parameter space show a balance between $|\tilde{\epsilon}|$ and $\tan \beta$. This can be understood from the fact that the presence of s_β^2 in eq. (12) implies that larger values of $|\tilde{\epsilon}|$ are needed at low $\tan \beta$ to obtain the correct solar mass scale. Notice that in this analysis, the ϵ values lie upon a range

Table 3
Some values for the parameters obtained in the search.

Parameter	Min. value	Max. value	Units	
Lightest neutralino	$5.10 \cdot 10^2$	$1.98 \cdot 10^3$	GeV	
Higgs mass	$1.24 \cdot 10^2$	$1.26 \cdot 10^2$	GeV	
Gluino mass	$3.46 \cdot 10^3$	$6.54 \cdot 10^4$	GeV	
$\tan\beta$	$2.10 \cdot 10^0$	$4.88 \cdot 10^1$	–	
Neutrino physics	Δm_{\odot}^2	$6.04 \cdot 10^{-5}$	$8.93 \cdot 10^{-5}$	eV ²
	Δm_{atm}^2	$2.44 \cdot 10^{-3}$	$2.48 \cdot 10^{-3}$	eV ²
	$\sin^2 \theta_{\odot}$	$2.81 \cdot 10^{-1}$	$3.72 \cdot 10^{-1}$	–
	$\sin^2 \theta_{atm}$	$5.39 \cdot 10^{-1}$	$5.86 \cdot 10^{-1}$	–
	$\sin^2 \theta_{rea}$	$1.80 \cdot 10^{-2}$	$2.77 \cdot 10^{-2}$	–

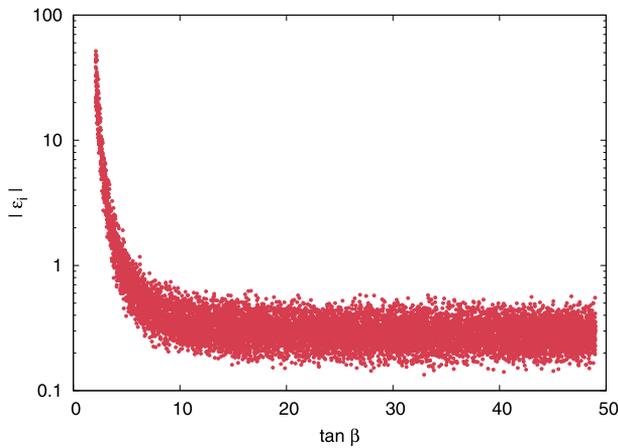


Fig. 1. Modulus of $\vec{\epsilon}$ as a function of $\tan\beta$. The strong correlation between the parameters, as the accumulation of points for $0.1 \leq |\vec{\epsilon}| \leq 0.5$ at $\tan\beta \geq \sim 10$ is given by neutrino physics (see body of the text).

which is taken to be natural, since it should remain small in order to give rise to neutrino masses. Notice further that in the sector of large $\tan\beta$ the value of ϵ saturates around 10^{-1} . This can be understood from the observation that $\sin\beta < 1$.

The Fig. 2 shows $|\vec{\epsilon}|$ as a function of the SS scale \tilde{m} . There, one sees that the allowed parameter space follows a nice linear relation in the logarithmic plot, where a growth of one order of magnitude in \tilde{m} is compensated by a growth of two orders of magnitude in $|\vec{\epsilon}|$. This can be understood from relation (12), since $\Delta M_{\nu}^{nr} \sim \epsilon_i \epsilon_j / \tilde{m}^4$. The fact that this correlation between the $|\vec{\epsilon}|$ and \tilde{m} does not extend to arbitrary values of \tilde{m} is due the bounds imposed from the Higgs sector.

5. Summary

This paper is dedicated to the tension between Split SuSy on the one hand, a good candidate for a low energy MSSM in terms of collider physics, and the same model on the other hand, not doing well for neutrino physics since it cannot account for the solar mass difference of neutrinos. We explore the possibility that the problem of Split SuSy with the solar mass difference originates

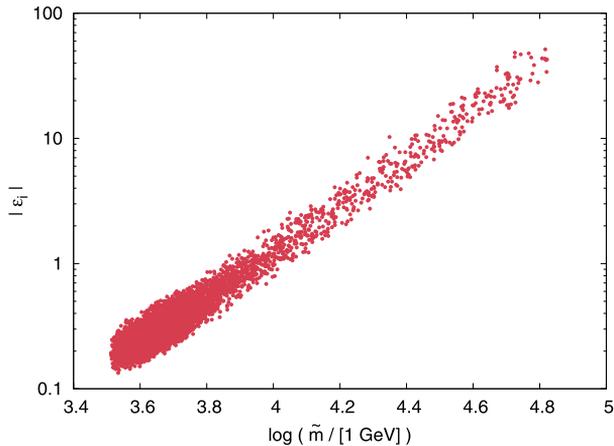


Fig. 2. Modulus of $\vec{\epsilon}$ as a function of \tilde{m} . The accumulation of points between $3.5 \leq \tilde{m} \leq 4.0$ obeys to neutrino physics, which pushes to have $|\vec{\epsilon}| \sim 0.25$.

from the fact that in the transition from the MSSM to Split SuSy the scalar fields are integrated out by a leading order approximation only.

In order to study this working hypothesis we include further terms in the integrating out procedure, which have been previously neglected. Those terms appear in the effective low energy Split SuSy Lagrangian in the form of non-renormalizable operators. With those inclusions we calculate quantum corrections and it is found that indeed a non-trivial contribution to the neutrino mass matrix is generated after spontaneous symmetry breaking. It is found that in particular the contribution coming from (12) has the potential of generating the observed solar mass difference. Finally, it is shown that this extended version of Split SuSy is indeed capable of reproducing the observed neutrino oscillations by simultaneously avoiding a strong fine tuning of the Higgs mass, if the mass scale \tilde{m} is rather moderate $\sim 10^4$ GeV.

The finding of good neutrino phenomenology at such a moderate scale \tilde{m} further supports our working hypothesis, namely, that terms that have been previously neglected in the integration out procedure can actually re-conciliate Split SuSy with the solar neutrino mass difference.

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