

Neutrino Masses in Split Supersymmetry

Marco Aurelio Díaz¹, Pavel Fileviez Pérez^{2,3}, and Clemencia Mora⁴

¹ *Departamento de Física, Universidad Católica de Chile,
Avenida Vicuña Mackenna 4860, Santiago, Chile*

² *Department of Physics, University of Wisconsin-Madison,
1150 University Avenue, Madison, WI 53706, USA*

³ *CFTP, Instituto Superior Tecnico, Ave. Rovisco Pais 1, 1049-001, Lisbon, Portugal*

⁴ *Département de Physique Nucléaire et Corpusculaire ,
Université de Genève, 24 Quai Ernest-Ansermet, 1211 Genève 4, Suisse*

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We investigate the possibility to generate neutrino masses in the context of Split supersymmetric scenarios where all sfermions are very heavy. All relevant contributions coming from the R-parity violating terms to the neutrino mass matrix up to one-loop level are computed, showing the importance of the Higgs one-loop corrections. We conclude that it is not possible to generate all neutrino masses and mixings in Split SUSY with bilinear R-Parity violating interactions. In the case of Partial Split SUSY the one-loop Higgs contributions are enough to generate the neutrino masses and mixings in agreement with the experiment. In the context of minimal SUSY SU(5) we find new contributions which help us to generate neutrino masses in the case of Split SUSY.

I. INTRODUCTION

Supersymmetric extensions of the Standard Model (SM) have been considered as one of the most appealing candidates for physics beyond the Standard Model (SM). Recently, different supersymmetric scenarios have been studied extensively. We mention low-energy SUSY [1], where the supersymmetric scale is around TeV, and Split SUSY where all the scalars, except for one Higgs doublet, are very heavy [2]. In both supersymmetric scenarios mentioned above it is possible to achieve unification of the gauge interactions at the high scale and the lightest supersymmetric particle (LSP) could be a natural candidate to describe the Cold Dark Matter of the Universe once the so-called R-parity is imposed as an exact symmetry of the theory. In SPLIT SUSY scenarios, ignoring the hierarchy problem, most of the unpleasant aspects of low-energy SUSY, such as excessive flavour and CP violation, and very fast dimension 5 proton decay, are eliminated.

It is very-well known that in general interactions which break the lepton or baryon number (or R-parity) are present in any SUSY extension of the SM. Therefore, we have the possibility to generate the neutrino masses and mixing [3], and we have to understand the predictions for proton stability [4]. For several phenomenological aspects of R-parity violating interactions see Ref. [5]. The possibility to describe the

neutrino properties with R-parity violating interactions in the context of the Minimal Supersymmetric Standard Model has been studied in detail by several groups in the context of low energy supersymmetry (See for example Refs. [6] and [7]). In the context of SPLIT SUSY the possibility to describe the masses and mixing of neutrinos has been studied in Ref. [8], where the authors concluded that it is not possible to use the R-parity bilinear terms alone to describe the neutrino properties.

In this work we re-examine the possibility to describe the properties of neutrinos using the R-parity violating interactions in the context of split supersymmetric scenarios. We agree with the results presented in Ref. [8] that in Split Supersymmetry, where only one Higgs doublet remains at the weak scale, it is not possible to generate the neutrino masses in agreement with the experiments and explain the reasons in detail. We study an alternative Split SUSY scenario where only the sfermions are very heavy while all Higgses can be light. We refer to this scenario as “Partial SPLIT SUSY”. Notice that in this scenario we can keep the nice features of SPLIT SUSY such as the suppression of proton decay due to R-parity violation and unification of gauge couplings at the high-scale. In this SUSY scenario we show that it is possible to generate the neutrino masses using all relevant interactions once the heavy sfermions are integrated out. Computing all contributions up to one-loop level, we find an example solution where it is shown that all constraints coming from neutrino experiments on the R-parity violating interactions are satisfied. In this scenario even if R-Parity is broken one could have the gravitino as a possible cold dark matter candidate.

We conclude that in Partial SPLIT SUSY it is possible to generate all neutrino masses and mixing in agreement with the experiments using the bilinear terms alone and the trilinear R-Parity violating (TRpV) couplings are essentially irrelevant. The key element is that the symmetry of the neutrino mass matrix at tree level is broken by the Higgs bosons loops together with neutralinos and charginos. The terms that break the symmetry of neutrino mass matrix vanish in the decoupling limit, making not possible the description of the neutrino masses in the “Standard” Split SUSY scenario. We study the same issue in the context of the minimal supersymmetric SU(5) where one finds new contributions which help us to generate neutrino masses in agreement with the experiments in the case of SPLIT SUSY.

II. R-PARITY VIOLATION AND NEUTRINO MASSES IN SPLIT SUSY

As we know in any supersymmetric extension of the Standard Model there are interactions terms which break the so-called R-parity. The R-parity is defined as $R = (-1)^{3(B-L)+2S}$, where L , B , and S are the lepton and baryon number, and the spin, respectively. Usually this symmetry is considered as an exact symmetry of the minimal supersymmetric extension of the Standard Model (MSSM) in order to avoid the dimension four contributions to proton decay and at the same time there is a possibility to have the lightest

supersymmetric particle as a good candidate for the Cold Dark Matter of the Universe.

In the context of the MSSM the so-called R-parity violating terms are given by

$$\mathcal{W}_{NR} = \alpha_{ijk} \hat{Q}_i \hat{L}_j \hat{D}_k^C + \beta_{ijk} \hat{U}_i^C \hat{D}_j^C \hat{D}_k^C + \gamma_{ijk} \hat{L}_i \hat{L}_j \hat{E}_k^C + \epsilon_i \hat{L}_i \hat{H}_u, \quad (1)$$

where $\beta_{ijk} = -\beta_{ikj}$ and $\gamma_{ijk} = -\gamma_{jik}$. As it is well-known due to the presence of the first and second terms in the above equation one has the so-called the dimension four contributions to the decay of the proton. In this case in order to satisfy the experimental bounds on the proton decay lifetime one has to assume that the multiplication of the couplings α_{ijk} and β_{ijk} is of the order 10^{-21} when the susy scale is at electroweak scale. In order to avoid these very small couplings in the theory one imposes by hand the R-parity symmetry. There is a second way to avoid these small couplings if the susy breaking scale is large, this is the case of Split-SUSY. Since in this case there is no need to impose any symmetry by hand we stick to this possibility and study the generation of neutrino masses in this context.

Let us discuss how to generate neutrino masses through this mechanism in three different scenarios:

- **MSSM with SPLIT SUSY**

In this supersymmetric scenario called Split SUSY all scalars are very heavy, except for one Higgs doublet. Integrating out the heavy scalars all possible R-parity conserving interactions in split supersymmetric scenarios are given by [2]:

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

where

$$H = \begin{pmatrix} H^+ \\ \frac{1}{\sqrt{2}}(v + \phi^0 + i\varphi^0) \end{pmatrix}, \quad (3)$$

is the SM Higgs. In the above equations we have the SM fields q_L , u_R , d_R , l_L , e_R and the superpartners of the Higgs and gauge bosons present in the MSSM. Following our notation \tilde{G} , \tilde{W} , and \tilde{B} are the gauginos associated to the $SU(3)$, $SU(2)$, and $U(1)$ gauge groups, respectively. While \tilde{H}_u and \tilde{H}_d correspond to the up and down higgsinos. The parameters in Eq. (2) are the following: m is the Higgs mass parameter, λ is the Higgs quartic self coupling; Y_u , Y_d , and Y_e are the Yukawa couplings; M_3 , M_2 , and M_1 are the gaugino masses, μ the higgsino mass, and \tilde{g}_u , \tilde{g}'_u , \tilde{g}_d , and \tilde{g}'_d are trilinear couplings between the Higgs boson, gauginos, and higgsinos.

The Higgs-gaugino-higgsino couplings in Eq. (2) satisfy matching conditions at the scale \tilde{m} . Above this scale, the theory is supersymmetric and the squarks, sleptons, and heavy Higgs doublet have a mass assumed to be nearly degenerate equal to \tilde{m} . The supersymmetric lagrangian includes the terms,

$$\mathcal{L}_{susy} \ni -\mu \tilde{H}_u^T i\sigma_2 \tilde{H}_d - \frac{H_u^\dagger}{\sqrt{2}} \left(g\sigma \tilde{W} + g' \tilde{B} \right) \tilde{H}_u - \frac{H_d^\dagger}{\sqrt{2}} \left(g\sigma \tilde{W} - g' \tilde{B} \right) \tilde{H}_d, \quad (4)$$

which implies the following boundary conditions at \tilde{m} :

$$\begin{aligned} \tilde{g}_u(\tilde{m}) &= g(\tilde{m}) \sin \beta(\tilde{m}) \quad , & \tilde{g}_d(\tilde{m}) &= g(\tilde{m}) \cos \beta(\tilde{m}), \\ \tilde{g}'_u(\tilde{m}) &= g'(\tilde{m}) \sin \beta(\tilde{m}) \quad , \text{ and } & \tilde{g}'_d(\tilde{m}) &= g'(\tilde{m}) \cos \beta(\tilde{m}), \end{aligned} \quad (5)$$

where $g(\tilde{m})$ and $g'(\tilde{m})$ are the gauge coupling constants evaluated at the scale \tilde{m} . At the same time the angle β is the mixing angle between the two Higgs doublets H_d and H_u of the supersymmetric model. In order to set our notation the two doublets are given by,

$$H_d = \begin{pmatrix} \frac{1}{\sqrt{2}}(v_d + \phi_d^0 + i\varphi_d^0) \\ H_d^- \end{pmatrix}, \quad H_u = \begin{pmatrix} H_u^+ \\ \frac{1}{\sqrt{2}}(v_u + \phi_u^0 + i\varphi_u^0) \end{pmatrix}, \quad (6)$$

and $\tan \beta = v_u/v_d$. In terms of these two Higgs doublets of the MSSM, the light fine-tuned Higgs doublet H in the low energy effective model is $H = -i\sigma_2 H_d^* \cos \beta(\tilde{m}) + H_u \sin \beta(\tilde{m})$.

As we mentioned before in SPLIT SUSY scenarios at low energy we have the SM fields, the charginos and neutralinos. Using the above notation the chargino mass matrix is given by:

$$\mathbf{M}_{\chi^\pm}^{SS} = \begin{bmatrix} M_2 & \frac{1}{\sqrt{2}}\tilde{g}_u v \\ \frac{1}{\sqrt{2}}\tilde{g}_d v & \mu \end{bmatrix}, \quad (7)$$

while the neutralino mass matrix reads as:

$$\mathbf{M}_{\chi^0}^{SS} = \begin{bmatrix} M_1 & 0 & -\frac{1}{2}\tilde{g}'_d v & \frac{1}{2}\tilde{g}'_u v \\ 0 & M_2 & \frac{1}{2}\tilde{g}_d v & -\frac{1}{2}\tilde{g}_u v \\ -\frac{1}{2}\tilde{g}'_d v & \frac{1}{2}\tilde{g}_d v & 0 & -\mu \\ \frac{1}{2}\tilde{g}'_u v & -\frac{1}{2}\tilde{g}_u v & -\mu & 0 \end{bmatrix}. \quad (8)$$

Now, since we are interested in the possibility to describe the neutrino masses in Split-SUSY, we write all relevant R-Parity violating interactions:

$$\mathcal{L}_{RpV}^{split} = \epsilon_i \tilde{H}_u^T i\sigma_2 L_i - \frac{1}{\sqrt{2}} a_i H^T i\sigma_2 (-\tilde{g}_d \sigma \tilde{W} + \tilde{g}'_d \tilde{B}) L_i + h.c., \quad (9)$$

where ϵ_i are the parameters that mix higgsinos with leptons, and a_i are dimensionless parameters that mix gauginos with leptons. Notice that the first term is the usual bilinear term, while the last

two terms are obtained once we integrate out the sleptons using the bilinear soft terms ($\tilde{L}_i H_u$) which break explicitly R-parity. As it is well-known we can also write the usual R-parity violating trilinear terms ($\hat{Q}\hat{D}^C\hat{L}$, $\hat{L}\hat{L}\hat{E}^C$). However, since the sfermions are very heavy in SPLIT SUSY and the contributions to the neutrino mass matrix coming from those terms are at one-loop level, those interactions cannot play any important role. Using Eq. (9), after the Higgs acquires a vev, we find the relevant terms for neutrino masses:

$$\mathcal{L}_{RpV}^{split} = - \left[\epsilon_i \tilde{H}_u^0 + \frac{1}{2} a_i v \left(\tilde{g} c_\beta \tilde{W}_3 - \tilde{g}' c'_\beta \tilde{B} \right) \right] \nu_i + h.c. + \dots, \quad (10)$$

where v is the vacuum expectation value of the SM-like Higgs field H . Knowing all R-parity violating interactions we can write the neutralino/neutrino mass matrix as:

$$\mathcal{M}_N^{SS} = \begin{bmatrix} M_{\chi_0^0}^{SS} & (m^{SS})^T \\ m^{SS} & 0 \end{bmatrix}, \quad (11)$$

where $M_{\chi_0^0}^{SS}$ is given by Eq. (7) and m^{SS} reads as:

$$m^{SS} = \begin{bmatrix} -\frac{1}{2} \tilde{g}'_d a_1 v & \frac{1}{2} \tilde{g}_d a_1 v & 0 & \epsilon_1 \\ -\frac{1}{2} \tilde{g}'_d a_2 v & \frac{1}{2} \tilde{g}_d a_2 v & 0 & \epsilon_2 \\ -\frac{1}{2} \tilde{g}'_d a_3 v & \frac{1}{2} \tilde{g}_d a_3 v & 0 & \epsilon_3 \end{bmatrix}. \quad (12)$$

We define the parameters $\lambda_i \equiv a_i \mu + \epsilon_i$, which are related to the traditional BRpV parameters Λ_i [9] by $\Lambda_i = \lambda_i v_d$. Integrating out the neutralinos, we find that the neutrino mass matrix is given by:

$$\mathbf{M}_\nu^{eff} = -m^{SS} (M_{\chi_0^0}^{SS})^{-1} (m^{SS})^T = \frac{v^2}{4 \det M_{\chi_0^0}^{SS}} (M_1 \tilde{g}_d^2 + M_2 \tilde{g}_d'^2) \begin{bmatrix} \lambda_1^2 & \lambda_1 \lambda_2 & \lambda_1 \lambda_3 \\ \lambda_2 \lambda_1 & \lambda_2^2 & \lambda_2 \lambda_3 \\ \lambda_3 \lambda_1 & \lambda_3 \lambda_2 & \lambda_3^2 \end{bmatrix}, \quad (13)$$

where the determinant of the neutralino mass matrix is:

$$\det M_{\chi_0^0}^{SS} = -\mu^2 M_1 M_2 + \frac{1}{2} v^2 \mu (M_1 \tilde{g}_u \tilde{g}_d + M_2 \tilde{g}_u' \tilde{g}_d') + \frac{1}{16} v^4 (\tilde{g}_u' \tilde{g}_d - \tilde{g}_u \tilde{g}_d')^2. \quad (14)$$

Notice that the effective neutrino mass matrix \mathbf{M}_ν^{eff} has only one eigenvalue different from zero. As in the case of R-parity violation in the MSSM with bilinear terms, at tree level only one neutrino is massive. Therefore, we have to investigate all possible one loop contributions to the neutrino mass matrix which help us to generate the atmospheric and solar neutrino masses. It has been argued in the literature [8] that using the bilinear terms it is not possible to explain the neutrino masses and mixing. We study this issue in detail and as we will show in the next section that once we include the one-loop contributions to the neutrino mass matrix it is not possible to generate all neutrino masses in agreement with the experiment.

• **MSSM with Partial SPLIT SUSY**

Let us study the same issue, how to generate the neutrino masses through the R-parity violating interactions, in Partial SPLIT SUSY where only the sfermions are very heavy while the Higgs can be light. Notice that in this case proton decay can be suppressed and the unification of the gauge interactions at the high scale is possible as well. We will show that in this scenario the contributions from the light Higgs bosons is enough to generate the neutrino masses at one-loop, and study the decoupling limit in order to have a better understanding of the results presented in the previous section.

We integrate out the heavy squarks and sleptons and find that the R-parity conserving (RpC) interactions below the scale \tilde{m} are given by

$$\begin{aligned} \mathcal{L}_{PSS}^{RpC} \ni & - \left[m_1^2 H_d^\dagger H_d + m_2^2 H_u^\dagger H_u - m_{12}^2 (H_d^T \epsilon H_u + h.c.) \right. \\ & + \frac{1}{2} \lambda_1 (H_d^\dagger H_d)^2 + \frac{1}{2} \lambda_2 (H_u^\dagger H_u)^2 + \lambda_3 (H_d^\dagger H_d) (H_u^\dagger H_u) + \lambda_4 |H_d^T \epsilon H_u|^2 \left. \right] \\ & + h_u \bar{u}_R H_u^T \epsilon q_L - h_d \bar{d}_R H_d^T \epsilon q_L - h_e \bar{e}_R H_d^T \epsilon l_L - \\ & - \frac{1}{\sqrt{2}} H_u^\dagger (\tilde{g}_u \sigma \tilde{W} + \tilde{g}'_u \tilde{B}) \tilde{H}_u - \frac{1}{\sqrt{2}} H_d^\dagger (\tilde{g}_d \sigma \tilde{W} - \tilde{g}'_d \tilde{B}) \tilde{H}_d + h.c. \end{aligned} \quad (15)$$

In the above equations, the two Higgs doublets that survive at the weak scale are H_d and H_u . The parameters in Eq. (15) not defined before are the following: m_1^2 , m_2^2 , and m_{12}^2 are the Higgs mass parameters, λ_i , $i = 1, 2, 3, 4$ are the Higgs quartic self couplings; and h_u , h_d , and h_e are the Yukawa couplings. The Higgs-gaugino-higgsino, gauge, and Yukawa couplings in Eq. (15) satisfy matching conditions at the scale \tilde{m} . Above this scale, the theory is supersymmetric and the squarks and sleptons have a mass assumed to be nearly degenerate to \tilde{m} . The supersymmetric lagrangian above \tilde{m} includes the terms,

$$\begin{aligned} \mathcal{L}_{susy}^{RpC} \ni & - \left[m_1^2 H_d^\dagger H_d + m_2^2 H_u^\dagger H_u - m_{12}^2 (H_d^T \epsilon H_u + h.c.) + \frac{1}{8} (g^2 + g'^2) (H_d^\dagger H_d)^2 \right. \\ & + \frac{1}{8} (g^2 + g'^2) (H_u^\dagger H_u)^2 + \frac{1}{4} (g^2 - g'^2) (H_d^\dagger H_d) (H_u^\dagger H_u) - \frac{1}{2} g^2 |H_d^T \epsilon H_u|^2 \left. \right] \\ & + \lambda_u \bar{u}_R H_u^T \epsilon q_L - \lambda_d \bar{d}_R H_d^T \epsilon q_L - \lambda_e \bar{e}_R H_d^T \epsilon l_L \\ & - \frac{1}{\sqrt{2}} H_u^\dagger (g \sigma \tilde{W} + g' \tilde{B}) \tilde{H}_u - \frac{1}{\sqrt{2}} H_d^\dagger (g \sigma \tilde{W} - g' \tilde{B}) \tilde{H}_d + h.c. \end{aligned} \quad (16)$$

Consequently, at the scale \tilde{m} we have the following boundary conditions for the Higgs couplings,

$$\lambda_1 = \lambda_2 = \frac{1}{4} (g^2 + g'^2), \quad \lambda_3 = \frac{1}{4} (g^2 - g'^2), \quad \lambda_4 = -\frac{1}{2} g^2, \quad (17)$$

for the Yukawa couplings $h_u = \lambda_u$, $h_d = \lambda_d$, $h_e = \lambda_e$, and for the higgsino-gaugino Yukawa couplings, $\tilde{g}_u = \tilde{g}_d = g$, $\tilde{g}'_u = \tilde{g}'_d = g'$. All of them evaluated at the scale \tilde{m} . Note the difference

between these boundary conditions and the corresponding ones in the original Split Supersymmetric model: the former do not involve the angle β . At the weak scale, the minimization of the Higgs potential leads to a vacuum expectation value for both Higgs doublets which satisfy $v_d^2 + v_u^2 = v^2$, such that $m_W^2 = \frac{1}{2}g^2v^2$ and $m_Z^2 = \frac{1}{2}(g^2 + g'^2)v^2$, as usual for a two Higgs doublet model (2HDM).

As we mentioned before in SPLIT SUSY scenarios, the charginos and neutralinos survive at low energies. Using the above notation the chargino mass matrix is given by:

$$\mathbf{M}_{\chi^\pm}^{PSS} = \begin{bmatrix} M_2 & \frac{1}{\sqrt{2}}v\tilde{g}_u s_\beta \\ \frac{1}{\sqrt{2}}v\tilde{g}_d c_\beta & \mu \end{bmatrix}, \quad (18)$$

while the neutralino mass matrix reads as:

$$\mathbf{M}_{\chi^0}^{PSS} = \begin{bmatrix} M_1 & 0 & -\frac{1}{2}\tilde{g}'_d c_\beta v & \frac{1}{2}\tilde{g}'_u s_\beta v \\ 0 & M_2 & \frac{1}{2}\tilde{g}_d c_\beta v & -\frac{1}{2}\tilde{g}_u s_\beta v \\ -\frac{1}{2}\tilde{g}'_d c_\beta v & \frac{1}{2}\tilde{g}_d c_\beta v & 0 & -\mu \\ \frac{1}{2}\tilde{g}'_u s_\beta v & -\frac{1}{2}\tilde{g}_u s_\beta v & -\mu & 0 \end{bmatrix}. \quad (19)$$

The difference with the Split Supersymmetric case in Eqs. (7) and (8) is in the mixings between higgsinos and gauginos. Now, with the neutrino masses in mind, we write all relevant R-Parity violating interactions in Partial SPLIT SUSY:

$$\mathcal{L}_{PSS}^{RpV} = -\epsilon_i \tilde{H}_u^T \epsilon L_i - \frac{1}{\sqrt{2}} b_i H_u^T \epsilon (\tilde{g}_d \sigma \tilde{W} - \tilde{g}'_d \tilde{B}) L_i + h.c., \quad (20)$$

with b_i dimensionless parameters. Using Eq. (20), after the Higgs acquires a vev, we find the relevant terms for neutrino masses:

$$\mathcal{L}_{PSS}^{RpV} = - \left[\epsilon_i \tilde{H}_u^0 + \frac{1}{2} b_i v_u \left(\tilde{g}_d \tilde{W}_3 - \tilde{g}'_d \tilde{B} \right) \right] \nu_i + h.c. + \dots, \quad (21)$$

where $v_d = v c_\beta$ and $v_u = v s_\beta$ are the vev of the two Higgs doublets. The neutralino/neutrino mass matrix still has the form given in Eq. (11), but in this scenario the matrix m reads as,

$$m^{PSS} = \begin{bmatrix} -\frac{1}{2}\tilde{g}'_d b_1 v_u & \frac{1}{2}\tilde{g}_d b_1 v_u & 0 & \epsilon_1 \\ -\frac{1}{2}\tilde{g}'_d b_2 v_u & \frac{1}{2}\tilde{g}_d b_2 v_u & 0 & \epsilon_2 \\ -\frac{1}{2}\tilde{g}'_d b_3 v_u & \frac{1}{2}\tilde{g}_d b_3 v_u & 0 & \epsilon_3 \end{bmatrix}. \quad (22)$$

The effective neutrino mass matrix obtained after diagonalizing by blocks is,

$$\mathbf{M}_\nu^{eff} = -m^{PSS} (\mathbf{M}_{\chi^0}^{PSS})^{-1} (m^{PSS})^T = \frac{M_1 \tilde{g}_d^2 + M_2 \tilde{g}'_d^2}{4 \det M_{\chi^0}^{PSS}} \begin{bmatrix} \Lambda_1^2 & \Lambda_1 \Lambda_2 & \Lambda_1 \Lambda_3 \\ \Lambda_2 \Lambda_1 & \Lambda_2^2 & \Lambda_2 \Lambda_3 \\ \Lambda_3 \Lambda_1 & \Lambda_3 \Lambda_2 & \Lambda_3^2 \end{bmatrix}, \quad (23)$$

with $\Lambda_i = \mu b_i v_u + \epsilon_i v_d$, and with the determinant of the neutralino submatrix equal to,

$$\det M_{\chi^0}^{PSS} = -\mu^2 M_1 M_2 + \frac{1}{2} v_u v_d \mu (M_1 \tilde{g}_u \tilde{g}_d + M_2 \tilde{g}'_u \tilde{g}'_d), \quad (24)$$

which is analogous to Eq. (14).

• SUSY SU(5) with Split SUSY

Now, let us discuss how one can find the R-parity violating couplings in the context of the simplest UV completion of the MSSM, the minimal SUSY SU(5). In this context the relevant superpotential is given by

$$\mathcal{W}_{NR}^{SU(5)} = \eta_i \hat{5}_i \hat{5}_H + c_i \hat{5}_i \hat{24}_H \hat{5}_H + \Lambda_{ijk} \hat{10}_i \hat{5}_j \hat{5}_k, \quad (25)$$

where our notation is $\bar{5}^T = (\hat{D}^C, -\hat{L}^T i\sigma_2)$, $10 = (\hat{Q}, \hat{U}^C, \hat{E}^C)$, $5_H^T = (\hat{T}, \hat{H}_u)$, and $\hat{24}_H = (\hat{\Sigma}_8, \hat{\Sigma}_3, \hat{\Sigma}_{(3,2)}, \hat{\Sigma}_{(\bar{3},2)}, \hat{\Sigma}_{24})$. Since all trilinear terms are coming from the same term in SU(5) one finds

$$\alpha_{ijk}/2 = \beta_{ijk} = \gamma_{ijk} = \Lambda_{ijk} = -\Lambda_{ikj}, \quad (26)$$

and the relevant interactions for the generation of neutrinos masses are given by

$$\mathcal{L}_{RpV} = -a_i \nu_i \tilde{H}_u^0 + \frac{1}{2} c_i \nu_i \tilde{\Sigma}_3^0 H_u^0 + \frac{3c_i}{2\sqrt{15}} \nu_i \tilde{\Sigma}_{24} H_u^0 + \text{h.c.}, \quad (27)$$

where at the renormalizable level $M_{\Sigma_3} = 5M_{\Sigma_{24}} = M_\Sigma$. Therefore, in this case one has the usual contribution from the bilinear term plus an extra contribution for the neutrino masses once we integrate out the neutral component of Σ_3 and Σ_{24} . It is important to mention that $a_i = \eta_i - 3\langle \Sigma_{24} \rangle c_i / 2\sqrt{15}$. Now, integrating out the fields Σ_3 and Σ_{24} one finds that the mass matrix for neutrinos is given by

$$M_{ij}^{SU(5)} = M_{ij}^{SS} + \frac{v_u^2}{M_\Sigma} c_i c_j, \quad (28)$$

where one can have $M_\Sigma \approx 10^{15-16}$ GeV in agreement with the unification constraints [10].

III. ONE-LOOP CORRECTIONS TO THE NEUTRINO MASS MATRIX

The one loop corrections are crucial for the correct characterization of neutrino phenomena. In the MSSM usually the most important one-loop contributions to the neutrino mass matrix are the bottom squarks, charginos, and neutralinos contributions.

A. Split SUSY Case

In Split SUSY all scalars, except for one light Higgs boson, are superheavy. Therefore, in this case the only potentially important contributions are charginos and neutralinos together with W , Z , and light Higgs inside the loop. We show in Appendix A that Z and W loops are just a small renormalization of the tree-level contribution. The Higgs boson loop together with neutralinos has the same property in the decoupling limit. We discuss those contributions in detail in this section.

In general, the one loop contributions to the neutrino mass matrix can be written as [6]:

$$\Delta M_\nu^{ij} = \Pi_{ij}(0) = -\frac{1}{16\pi^2} \sum_{f,b} G_{ijfb} m_f B_0(0; m_f^2, m_b^2), \quad (29)$$

where the sum is over the fermions (f) and the bosons (b) inside the loop, m_f is the fermion mass, and G_{ijfb} is defined by the couplings between the neutrinos and the fermions and bosons inside the loop. Once the smallness of the ϵ_i and λ_i parameters is taken into account, each contribution can be expressed in the form

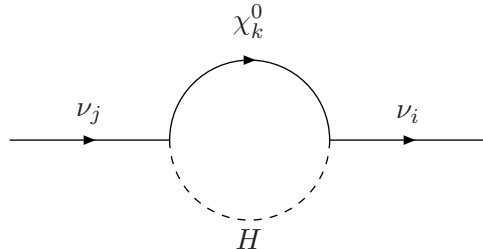
$$\Delta \Pi_{ij} = A^{(1)} \lambda_i \lambda_j + B^{(1)} (\epsilon_i \lambda_j + \epsilon_j \lambda_i) + C^{(1)} \epsilon_i \epsilon_j, \quad (30)$$

with $A^{(1)}$, $B^{(1)}$, and $C^{(1)}$ parameters independent of ϵ_i and λ_i , but dependent on the other SUSY parameters. The super-index (1) refers to the one-loop contribution. The tree-level neutrino mass matrix in Eq. (13) has the form $\mathbf{M}_\nu^{eff} = A^{(0)} \lambda_i \lambda_j$ with

$$A^{(0)} = \frac{v^2}{4 \det M_{\chi^0}^{SS}} (M_1 \tilde{g}_d^2 + M_2 \tilde{g}_d'^2), \quad (31)$$

and we define the one-loop corrected parameters $A = A^{(0)} + A^{(1)}$, $B = B^{(1)}$, and $C = C^{(1)}$.

In the MSSM with BRpV the neutral Higgs bosons mix with the sneutrinos forming two sets of 5 scalars and 5 pseudo-scalars. Nevertheless, in Split SUSY, all the sneutrinos are extremely heavy and decouple from the light Higgs boson H . In addition, the heavy Higgs boson also has a very large mass, leaving the light Higgs as the only neutral scalar able to contribute to the neutrino masses. This contribution is represented by the following Feynman graph,



which is proportional to the neutralino mass $m_{\chi_k^0}$. Here χ_k^0 and H are the neutralino and Higgs mass eigenstates, but the graph is calculated in the basis where ν_i are not mass eigenstates. The fields ν_i are the neutrino fields associated to the effective mass matrix given in Eq. (13). This contribution to Eq. (29) proceeds with the coupling [6]

$$G_{ijk}^h = \frac{1}{2}(O_{Ljk}^{nnh} O_{Lki}^{nnh} + O_{Rjk}^{nnh} O_{Rki}^{nnh}), \quad (32)$$

where the relevant vertex is:

$$= i \left[O_{Lij}^{nnh} \frac{(1-\gamma_5)}{2} + O_{Rij}^{nnh} \frac{(1+\gamma_5)}{2} \right]$$

Here F_i^0 are the seven eigenvectors linear combination of the higgsinos, gauginos, and neutrinos. The O_L and O_R couplings satisfy $O_{Lij}^{nnh} = (O_{Rji}^{nnh})^*$ and above the scale \tilde{m} we have,

$$O_{Rij}^{nnh} = \frac{1}{2} \left\{ \mathcal{N}_{i4} (g s_\beta \mathcal{N}_{j2} - g' s_\beta \mathcal{N}_{j1}) - \mathcal{N}_{i3} (g c_\beta \mathcal{N}_{j2} - g' c_\beta \mathcal{N}_{j1}) + \mathcal{N}_{i\ell+4} (g s_\ell \mathcal{N}_{j2} - g' s_\ell \mathcal{N}_{j1}) + (i \leftrightarrow j) \right\}, \quad (33)$$

where we have an implicit sum over $\ell = 1, 2, 3$. We allow the matrix elements of the matrix \mathcal{N} to be imaginary when one of the eigenvalues is negative, such that we do not need to include explicitly the sign called η_i in Ref. [6]. The difference with the MSSM couplings given in Ref. [11] lies in the fact that in our case \mathcal{N} is a 7×7 matrix, and the Higgs mixing angle has been replaced by $\alpha = \beta - \pi/2$, valid in the decoupling limit [12]. In addition, the third term is not present in the MSSM and comes from the second term in the supersymmetric lagrangian of Eq. (9).

Comparing the lagrangian below the scale \tilde{m} in Eq. (9) with the relevant term of the supersymmetric lagrangian above \tilde{m} given by

$$\mathcal{L}_{SUSY} \ni -\frac{1}{\sqrt{2}} \tilde{L}_i^\dagger \left(g \sigma^a \tilde{W}^a - g' \tilde{B} \right) L_i, \quad (34)$$

and considering the mixing between sleptons and Higgs bosons above that scale, a correspondence is found when the replacement $\tilde{L}_i^* \rightarrow -s_i i \sigma_2 H$ is made. The relevant matching condition at \tilde{m} is

$$a_i(\tilde{m}) = \frac{s_i(\tilde{m})}{\cos \beta(\tilde{m})}, \quad (35)$$

where the parameters $s_i(\tilde{m})$ represent the amount of slepton \tilde{L}_i in the low energy Higgs H , and related to the sneutrino vev present above the scale \tilde{m} ,

$$\tilde{L}_i = \begin{pmatrix} \frac{1}{\sqrt{2}}(v_i + \tilde{\ell}_{si}^0 + i\tilde{\ell}_{pi}^0) \\ \tilde{\ell}_{Li}^- \end{pmatrix}, \quad (36)$$

as explained in the appendix. Using the approximation for the matrix \mathcal{N} from Appendix A, we obtain for the coupling below the scale \tilde{m} :

$$\begin{aligned} O_{Rik}^{\nu\chi h} = \frac{1}{2} \Big\{ & -(\tilde{g}s_\beta N_{k2} - \tilde{g}'s'_\beta N_{k1}) \xi_{i4} - N_{k4} (\tilde{g}s_\beta \xi_{i2} - \tilde{g}'s'_\beta \xi_{i1}) \\ & + (\tilde{g}c_\beta N_{k2} - \tilde{g}'c'_\beta N_{k1}) (\xi_{i3} - a_i) + N_{k3} (\tilde{g}c_\beta \xi_{i2} - \tilde{g}'c'_\beta \xi_{i1}) \Big\}. \end{aligned} \quad (37)$$

Notice that there is no term proportional to ϵ_i since there is a cancellation in $\xi_{i3} - a_i$. It can be checked using Eqs. (A4) and the definition of $\lambda_i = a_i\mu + \epsilon_i$. This implies that the contribution of the light Higgs boson has the form:

$$\Delta\Pi_{ij}^h = A^h \lambda_i \lambda_j, \quad (38)$$

which does not break the symmetry of the neutrino mass matrix at tree level. The detailed expression is given by,

$$\Delta\Pi_{ij}^h = -\frac{1}{16\pi^2} \sum_{k=1}^4 (\tilde{O}_k^{\nu\chi h})^2 \lambda_i \lambda_j m_{\chi_k^0} B_0(0; m_{\chi_k^0}^2, m_h^2), \quad (39)$$

with

$$\begin{aligned} \tilde{O}_k^{\nu\chi h} = \frac{1}{2} \Big\{ & -(\tilde{g}s_\beta N_{k2} - \tilde{g}'s'_\beta N_{k1}) \xi_4 - N_{k4} (\tilde{g}s_\beta \xi_2 - \tilde{g}'s'_\beta \xi_1) \\ & + (\tilde{g}c_\beta N_{k2} - \tilde{g}'c'_\beta N_{k1}) (\xi_3 - 1/\mu) + N_{k3} (\tilde{g}c_\beta \xi_2 - \tilde{g}'c'_\beta \xi_1) \Big\}. \end{aligned} \quad (40)$$

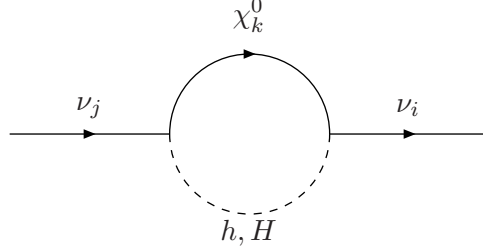
Since the gauge and Goldstone boson contribute to the neutrino mass matrix in the same form, as can be checked in the Appendix, we conclude that it is not possible to generate the neutrino masses in Split Supersymmetry with bilinear R-Parity violating interactions alone. This conclusion is in agreement with the results presented in Ref. [8], and in Ref. [13], where the contribution from the Higgs boson can be inferred taking the decoupling limit (see also [14]).

B. Partial Split SUSY Case

In this scenario the five physical Higgs states, h, H, A, H^\pm , are light and contribute to the neutrino mass matrix. In the following subsections we divide them in CP-even, CP-odd, and charged Higgs contributions.

1. CP-even neutral Higgs bosons

The two CP-even neutral Higgs bosons contribute to the neutrino mass matrix through the following graphs,



where the G factor in Eq. (29) is,

$$G_{ijkr}^s = \frac{1}{2}(O_{Ljkr}^{nns} O_{Lkir}^{nns} + O_{Rjkr}^{nns} O_{Rkir}^{nns}). \quad (41)$$

The relevant coupling above the scale \tilde{m} is the CP-even neutral scalar couplings to two neutral fermions, given by,

$$= i \left[O_{Lijk}^{nns} \frac{(1-\gamma_5)}{2} + O_{Rijk}^{nns} \frac{(1+\gamma_5)}{2} \right]$$

where,

$$O_{Lijk}^{nns} = \frac{1}{2} \left[(-R_{k1}^0 \mathcal{N}_{j3}^* + R_{k2}^0 \mathcal{N}_{j4}^* - R_{k\ell+2}^0 \mathcal{N}_{j\ell+4}^*) (g \mathcal{N}_{i2}^* - g' \mathcal{N}_{i1}^*) + (i \leftrightarrow j) \right], \quad (42)$$

and $O_{Rijk}^{nns} = (O_{Lijk}^{nns})^*$. The fields S_k^0 are linear combinations of CP-even Higgs and sneutrinos whose mass matrix in the basis $(\phi_d^0, \phi_u^0, \tilde{\ell}_{si}^0)$ is given in the Appendix B. In the PSSusy, the mass matrix can be diagonalized by,

$$\begin{pmatrix} h \\ H \\ \tilde{\nu}_s^i \end{pmatrix} = \begin{pmatrix} -s_\alpha & c_\alpha & -s_s^j \\ c_\alpha & s_\alpha & -t_s^j \\ -s_\alpha s_s^i + c_\alpha t_s^i & c_\alpha s_s^i + s_\alpha t_s^i & \delta_{ij} \end{pmatrix} \begin{pmatrix} \phi_d^0 \\ \phi_u^0 \\ \tilde{\ell}_{sj}^0 \end{pmatrix}, \quad (43)$$

where the angle α is analogous to the CP-even neutral Higgs bosons mixing angle of the MSSM. An expression for the mixing angles s_s^i and t_s^i above the scale \tilde{m} can be found in the Appendix B. Comparing

the supersymmetric lagrangian above the scale \tilde{m} in Eq. (34) with the terms of the PSSusy lagrangian in Eq. (20) we find the following matching conditions,

$$s_s^i(\tilde{m}) = -b_i(\tilde{m})c_\alpha; \quad t_s^i(\tilde{m}) = -b_i(\tilde{m})s_\alpha, \quad (44)$$

where $s_s^i(\tilde{m})$ represents the amount of slepton \tilde{L}_i present in the low energy light Higgs h , and analogously with $t_s^i(\tilde{m})$ for the low energy heavy Higgs H . In the limit where the sleptonic fields have a very large mass, they satisfy,

$$s_s^i \rightarrow -c_\alpha \frac{v_i}{v_u}, \quad t_s^i \rightarrow -s_\alpha \frac{v_i}{v_u}, \quad (45)$$

which tells us that the parameter b_i , defined below \tilde{m} , is directly proportional to the sneutrino vacuum expectation value v_i , defined above the scale \tilde{m} .

In the coupling in Eq. (42), we take the first neutral fermion as a neutrino and the second as a neutralino, obtaining the following couplings for both Higgs bosons h and H ,

$$\begin{aligned} O_{Lik}^{\nu\chi h} &= \frac{1}{2} \left[(s_\alpha N_{k3}^* + c_\alpha N_{k4}^*) (-g\xi_{i2} + g'\xi_{i1}) + (-s_\alpha \xi_{i3} - c_\alpha \xi_{i4} + s_s^i) (gN_{k2}^* - g'N_{k1}^*) \right], \\ O_{Lik}^{\nu\chi H} &= \frac{1}{2} \left[(-c_\alpha N_{k3}^* + s_\alpha N_{k4}^*) (-g\xi_{i2} + g'\xi_{i1}) + (c_\alpha \xi_{i3} - s_\alpha \xi_{i4} + t_s^i) (gN_{k2}^* - g'N_{k1}^*) \right]. \end{aligned} \quad (46)$$

After isolating the terms proportional to ϵ_i in the couplings, and using eq. (44), we find the following expressions valid below \tilde{m} ,

$$\begin{aligned} O_{Lik}^{\nu\chi h} &= \tilde{O}_{Lk}^{\nu\chi h} \Lambda_i + \frac{1}{2\mu s_\beta} \cos(\alpha - \beta) (gN_{k2}^* - g'N_{k1}^*) \epsilon_i, \\ O_{Lik}^{\nu\chi H} &= \tilde{O}_{Lk}^{\nu\chi H} \Lambda_i + \frac{1}{2\mu s_\beta} \sin(\alpha - \beta) (gN_{k2}^* - g'N_{k1}^*) \epsilon_i, \end{aligned} \quad (47)$$

with the term proportional to Λ_i given by,

$$\begin{aligned} \tilde{O}_{Lk}^{\nu\chi h} &= -\frac{1}{2} \left[(s_\alpha N_{k3}^* + c_\alpha N_{k4}^*) (g\xi_2 - g'\xi_1) + \left(s_\alpha \xi_3 + c_\alpha \xi_4 + \frac{c_\alpha}{\mu v_u} \right) (gN_{k2}^* - g'N_{k1}^*) \right], \\ \tilde{O}_{Lk}^{\nu\chi H} &= \frac{1}{2} \left[(c_\alpha N_{k3}^* - s_\alpha N_{k4}^*) (g\xi_2 - g'\xi_1) + \left(c_\alpha \xi_3 - s_\alpha \xi_4 - \frac{s_\alpha}{\mu v_u} \right) (gN_{k2}^* - g'N_{k1}^*) \right]. \end{aligned} \quad (48)$$

Notice that the presence of the term proportional to ϵ_i in Eq. (47) implies that the contribution of the CP-even Higgs bosons has the form:

$$\Delta\Pi_{ij} = A\Lambda_i\Lambda_j + B(\Lambda_i\epsilon_j + \Lambda_j\epsilon_i) + C\epsilon_i\epsilon_j, \quad (49)$$

breaking the symmetry of the neutrino mass matrix at tree level, and generating a solar mass. Explicitly, this contribution is:

$$\Delta\Pi_{ij}^{hH} = -\frac{1}{16\pi^2} \sum_{k=1}^4 \sum_{n=1}^2 (E_k^n \Lambda_i + F_k^n \epsilon_i) (E_k^n \Lambda_j + F_k^n \epsilon_j) m_{\chi_k^0} B_0(0; m_{\chi_k^0}^2, m_{H_n}^2), \quad (50)$$

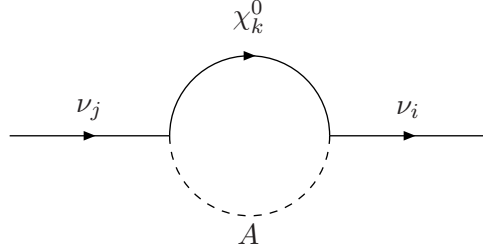
with

$$\begin{aligned} E_k^1 &= \tilde{O}_{Lk}^{\nu\chi^h}, & F_k^1 &= \frac{\cos(\alpha - \beta)}{2\mu s_\beta} (gN_{k2}^* - g'N_{k1}^*), \\ E_k^2 &= \tilde{O}_{Lk}^{\nu\chi^H}, & F_k^2 &= \frac{\sin(\alpha - \beta)}{2\mu s_\beta} (gN_{k2}^* - g'N_{k1}^*), \end{aligned} \quad (51)$$

where we work in the Feynman gauge.

2. CP-odd neutral Higgs bosons

Loops including the CP-odd Higgs boson A must be added through the graph,



where the G factor in Eq. (29) is,

$$G_{ijk}^p = -\frac{1}{2}(O_{Ljkr}^{np} O_{Lkir}^{np} + O_{Rjkr}^{np} O_{Rkir}^{np}). \quad (52)$$

The relevant coupling above the scale \tilde{m} is the CP-odd neutral scalar couplings to two neutral fermions, given by,

$$\begin{array}{c} \text{Diagram: A dashed line } P_k^0 \text{ splits into two solid lines } F_j^0 \text{ and } F_i^0. \end{array} = \left[O_{Lijk}^{np} \frac{(1-\gamma_5)}{2} + O_{Rijk}^{np} \frac{(1+\gamma_5)}{2} \right]$$

where,

$$O_{Lijk}^{np} = -\frac{1}{2} \left[\left(-R_{k1}^p \mathcal{N}_{j3}^* + R_{k2}^p \mathcal{N}_{j4}^* - R_{k\ell+2}^p \mathcal{N}_{j\ell+4}^* \right) (g\mathcal{N}_{i2}^* - g'\mathcal{N}_{i1}^*) + (i \leftrightarrow j) \right], \quad (53)$$

and $O_{Rijk}^{np} = -(O_{Ljik}^{np})^*$. The fields P_k^0 are linear combinations of CP-odd Higgs and sneutrinos whose mass matrix in the basis $(\varphi_d^0, \varphi_u^0, \tilde{\ell}_{pi}^0)$ is given in the Appendix B. In the PSSusy, the mass matrix can be

diagonalized by,

$$\begin{pmatrix} G \\ A \\ \tilde{\nu}_p^i \end{pmatrix} = \begin{pmatrix} -c_\beta & s_\beta & -s_p^j \\ s_\beta & c_\beta & -t_p^j \\ -c_\beta s_p^i + s_\beta t_p^i & s_\beta s_p^i + c_\beta t_p^i & \delta_{ij} \end{pmatrix} \begin{pmatrix} \varphi_d^0 \\ \varphi_u^0 \\ \tilde{\ell}_{pj}^0 \end{pmatrix}. \quad (54)$$

An expression for the mixing angles s_p^i and t_p^i above the scale \tilde{m} can be found in the Appendix B. Comparing the supersymmetric lagrangian above the scale \tilde{m} in Eq. (34) with the terms of the PSSusy lagrangian in Eq. (20) we find the following matching conditions,

$$s_p^i(\tilde{m}) = b_i(\tilde{m})s_\beta; \quad t_p^i(\tilde{m}) = b_i(\tilde{m})c_\beta, \quad (55)$$

where $s_p^i(\tilde{m})$ represents the amount of slepton \tilde{L}_i present in the Goldstone boson G , and analogously with $t_p^i(\tilde{m})$ for the low energy CP-odd Higgs A . In the limit where the sleptonic fields have a very large mass,

$$s_p^i \rightarrow s_\beta \frac{v_i}{v_u}, \quad t_p^i \rightarrow c_\beta \frac{v_i}{v_u}, \quad (56)$$

which indicates $b_i = v_i/v_u$ in agreement with the CP-even case.

If we take the coupling in Eq. (53) and expand on small R-Parity violating parameters we find for the CP-odd Higgs bosons couplings,

$$O_{Lik}^{\nu\chi^a} = -\frac{1}{2} \left[(-s_\beta N_{k3}^* + c_\beta N_{k4}^*) (-g\xi_{i2} + g'\xi_{i1}) + (s_\beta \xi_{i3} - c_\beta \xi_{i4} + t_p^i) (gN_{k2}^* - g'N_{k1}^*) \right]. \quad (57)$$

If we isolate the terms proportional to ϵ_i , using eq. (55), we find,

$$O_{Lik}^{\nu\chi^a} = \tilde{O}_{Lk}^{\nu\chi^a} \Lambda_i + \frac{1}{2\mu s_\beta} (gN_{k2}^* - g'N_{k1}^*) \epsilon_i. \quad (58)$$

It is shown in Appendix A that the Goldstone boson contribution completely cancels out when gauge dependent terms from gauge couplings and tadpoles are included. The \tilde{O} coupling is defined by,

$$\tilde{O}_{Lk}^{\nu\chi^a} = -\frac{1}{2} \left[(s_\beta N_{k3}^* - c_\beta N_{k4}^*) (g\xi_2 - g'\xi_1) + \left(s_\beta \xi_3 - c_\beta \xi_4 + \frac{c_\beta}{\mu v_u} \right) (gN_{k2}^* - g'N_{k1}^*) \right]. \quad (59)$$

In this way, the CP-odd contribution is

$$\Delta\Pi_{ij}^A = \frac{1}{16\pi^2} \sum_{k=1}^4 (E_k^3 \Lambda_i + F_k^3 \epsilon_i) (E_k^3 \Lambda_j + F_k^3 \epsilon_j) m_{\chi_k^0} B_0(0; m_{\chi_k^0}^2, m_A^2), \quad (60)$$

with

$$E_k^3 = \tilde{O}_{Lk}^{\nu\chi^a}, \quad F_k^3 = \frac{1}{2\mu s_\beta} (gN_{k2}^* - g'N_{k1}^*). \quad (61)$$

Note that the CP-odd contribution in Eq. (61) has the opposite sign of the CP-even contribution. In addition, the $\epsilon_i \epsilon_j$ terms in the limit of equal neutral Higgs masses. This is because the CP-even terms are proportional to $\cos^2(\alpha - \beta) B_0(0; m_{\chi_k^0}^2, m_h^2)$ and $\sin^2(\alpha - \beta) B_0(0; m_{\chi_k^0}^2, m_H^2)$, while the CP-odd term is proportional to $-B_0(0; m_{\chi_k^0}^2, m_A^2)$.

TABLE I: PSS and neutrino mass matrix parameters.

Parameter	Solution	Units
$\tan \beta$	10	-
μ	450	GeV
M_2	300	GeV
M_1	150	GeV
m_h	120	GeV
m_A	1000	GeV
Q	830	GeV
A	-2.7	eV/GeV ⁴
B	-0.0005	eV/GeV ³
C	0.315	eV/GeV ²

IV. NUMERICAL RESULTS

A. Partial Split SUSY

As seen in the previous chapters, Partial Split Supersymmetry is determined by the following supersymmetric parameters: the supersymmetric Higgs mass μ , the gaugino masses M_1 and M_2 , the mass of the lightest CP-even Higgs m_h , the CP-odd Higgs mass m_A , and the tangent of the CP-odd Higgs mixing angle $\tan \beta$. As a working scenario we choose the numerical values given in Table I. In this scenario the four neutralino masses are $m_\chi = 147, 282, 455, 476$ GeV, with the lightest neutralino the LSP. In the Higgs sector, the charged Higgs mass is $m_{H^\pm} = 1003.2$ GeV, the heavy neutral CP-even Higgs mass is $m_H = 1000.2$ GeV, and the CP-even Higgs mixing angle is given by $\sin \alpha = 0.101$.

The one-loop corrected parameters A , B , and C introduced in Eq. (49) are calculated with the results in Eq. (50) for the neutral CP-even Higgs bosons, in Eq. (60) for the neutral CP-odd Higgs boson, and in Eq. (A24) for the charged Higgs boson. These contributions give rise to a set of parameters A , B , and C given in Table I. The value of $A = -2.7 \text{ eV/GeV}^2$ is mainly due to the tree level contribution, and $C = 0.315 \text{ eV/GeV}^4$ is completely generated by radiative corrections.

The parameter C is subtraction scale independent, while the parameters A and B depend on the subtraction scale Q . As a way of fixing this scale, we have chosen Q such that it minimizes the parameter B , making the solar mass completely scale independent. For the scenario in Table I we find that $Q = 830$ GeV

TABLE II: BRpV parameters and neutrino observables.

Parameter	Solution	Units
ϵ_1	0.0346	GeV
ϵ_2	0.265	GeV
ϵ_3	0.322	GeV
Λ_1	-0.0269	GeV ²
Λ_2	-0.00113	GeV ²
Λ_3	0.0693	GeV ²
Δm_{atm}^2	2.34×10^{-3}	eV ²
Δm_{sol}^2	8.16×10^{-5}	eV ²
$\tan^2 \theta_{atm}$	1.04	-
$\tan^2 \theta_{sol}$	0.455	-
$\tan^2 \theta_{13}$	0.0247	-
m_{ee}	0.00394	eV

gives rise to $B = -0.0005 \text{ eV/GeV}^3$, which is already negligible.

We notice that in the decoupling limit scenario the light CP-even Higgs h contribution to the solar mass (or equivalently, to the parameter C) is negligible, since it is proportional to $\cos(\alpha - \beta) \rightarrow 0$. Therefore, it can be said properly that the solar mass comes exclusively from the contributions of the heavy Higgs bosons H and A . Further more, as indicated by Eqs. (51) and (61) the contributions from H and A have opposite signs and tend to cancel each other in the decoupling limit, where $\sin(\alpha - \beta) \rightarrow 1$ and $m_H \rightarrow m_A$. In our scenario, $\cos(\alpha - \beta) = 0.0016$ and $m_H - m_A = 0.2 \text{ GeV}$, and the cancellation between H and A contributions to C is at the 0.07%. Within the scenario in Table I we look for a solution to the neutrino observables varying $\vec{\epsilon}$ and $\vec{\Lambda}$. An example solution is given in Table II. This solution satisfy $\epsilon_1 \ll \epsilon_2, \epsilon_3$ and $|\Lambda_2| \ll |\Lambda_1|, \Lambda_3$. The sign of these parameters have a very small influence. Also in Table II we list the neutrino observables. The atmospheric mass $\Delta m_{atm}^2 = 2.34 \times 10^{-3} \text{ eV}^2$ and the solar mass $\Delta m_{sol}^2 = 8.16 \times 10^{-5} \text{ eV}^2$ are practically at the center of the experimentally allowed regions. The atmospheric angle $\tan^2 \theta_{atm} = 1.04$ is slightly deviated from maximal mixing, while the solar angle $\tan^2 \theta_{sol} = 0.455$ is non-maximal and with a value centered on the experimentally allowed region. The other two parameters, the reactor angle $\tan^2 \theta_{13} = 0.0247$ and the neutrinoless double beta decay mass $m_{ee} = 0.00394 \text{ eV}$ have not been experimentally measured and the predictions of our model are well below the experimental upper bounds.

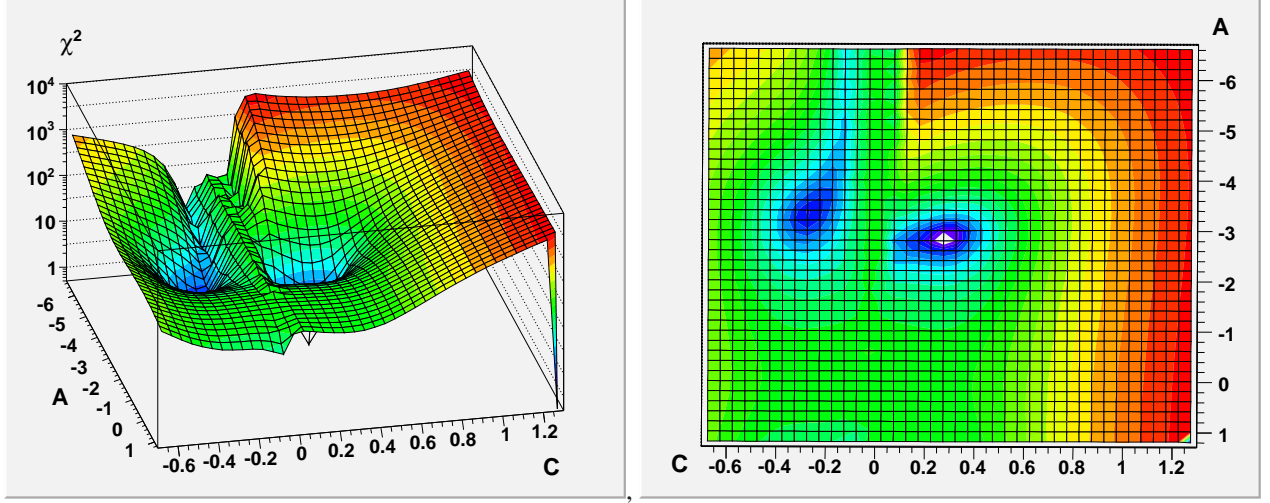


FIG. 1: Neutrino physics χ^2 as a function of the neutrino mass matrix parameters A and C , keeping $\vec{\epsilon}$ and $\vec{\Lambda}$ as indicated in Table I.

In order to study the dependence of the neutrino physics solutions on different parameters we have implemented the following χ^2

$$\chi^2 = \left(\frac{10^3 \Delta m_{atm}^2 - 2.35}{0.95} \right)^2 + \left(\frac{10^5 \Delta m_{sol}^2 - 8.15}{0.95} \right)^2 + \left(\frac{\sin^2 \theta_{atm} - 0.51}{0.17} \right)^2 + \left(\frac{\sin^2 \theta_{sol} - 0.305}{0.075} \right)^2. \quad (62)$$

In each of these terms we evaluated how many standard deviation the prediction is from the measured experimental central values [15]. In Fig. 1 we have χ^2 in the vertical axis as a function of A and C , in perspective in the left frame and level contours in the right frame. The preferred solution of Table I appears at the center of the graphs. Neutrino observables are very sensitive to the parameters A and C as shown by contours, where the darkest ellipsoid (blue) corresponds to $\chi^2 \lesssim 10$, while the white center corresponds to $\chi^2 \lesssim 1$. There is a second minima, but it does not reach values near unity.

A good approximation for the neutrino masses in this scenario is the following,

$$\begin{aligned} m_3 &= C|\vec{\epsilon}|^2 + A \frac{(\vec{\epsilon} \cdot \vec{\Lambda})^2}{|\vec{\epsilon}|^2}, \\ m_2 &= A \frac{|\vec{\epsilon} \times (\vec{\Lambda} \times \vec{\epsilon})|^2}{|\vec{\epsilon}|^4}, \end{aligned} \quad (63)$$

with the third neutrino massless [16]. Despite the fact that C is one-loop generated and A receives contributions at tree level, the first term in m_3 is dominant, and thus more important for the atmospheric mass scale. The A term is the only one contributing to the solar mass, as indicated in Eq. (63).

In Fig. 2 we plot χ^2 as a function of ϵ_2 and ϵ_3 in two frames as described for the previous figure. The rest of the BRpV parameters are fixed to the values in Table I, while the values of A and C are calculated

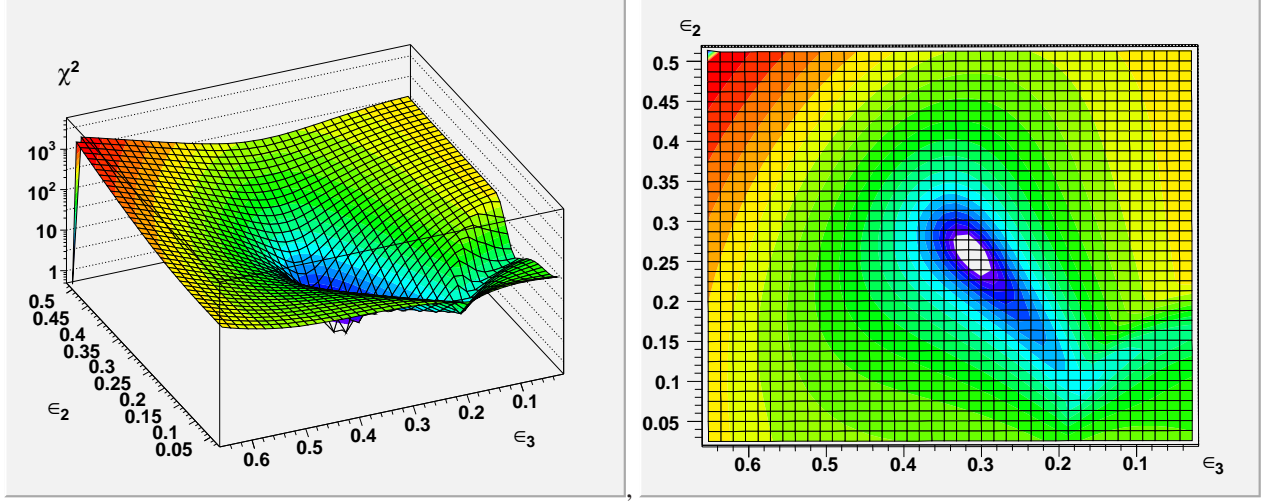


FIG. 2: Neutrino physics χ^2 as a function of the BRpV parameters ϵ_2 and ϵ_3 , keeping the rest of the parameters as indicated in Table I.

from the loop contributions. In our scenario approximated expression can be found when ϵ_1 and Λ_2 are neglected. It turns out that the atmospheric angle and mass squared difference depend strongly on ϵ_2 and ϵ_3 . They are given by,

$$\begin{aligned}\Delta m_{\text{atm}}^2 &\approx C^2(\epsilon_2^2 + \epsilon_3^2)^2, \\ \tan^2 \theta_{\text{atm}} &\approx \left(\frac{\epsilon_2}{\epsilon_3}\right)^2.\end{aligned}\quad (64)$$

Notice that it is the atmospheric mass who receives the main contribution from loop corrections, with C generated entirely at one-loop. Equal values for the atmospheric mass correspond to circles around the origin in the ϵ_2 - ϵ_3 plane, while equal values for the atmospheric angle are represented by straight lines passing through the origin. This geometry can be visualized in Fig. 2.

In Fig. 3 we plot χ^2 as a function of Λ_1 and Λ_3 with the other parameters as indicated in Table I. The solar mass squared difference and angle depend strongly on Λ_1 and Λ_3 as indicated by the following approximations,

$$\begin{aligned}\Delta m_{\text{sol}}^2 &\approx A^2 \left[\Lambda_1^2 + \frac{\Lambda_3^2}{1 + (\epsilon_3/\epsilon_2)^2} \right]^2, \\ \tan^2 \theta_{\text{sol}} &\approx \frac{\Lambda_1^2}{\Lambda_3^2} \left[1 + \left(\frac{\epsilon_3}{\epsilon_2}\right)^2 \right].\end{aligned}\quad (65)$$

When the ϵ parameters are kept constant, equal values for the solar mass are represented by ellipses, while constant values for the solar angle are represented by straight lines passing through the origin. As with the previous figure, this geometry can be visualized also in Fig. 3.

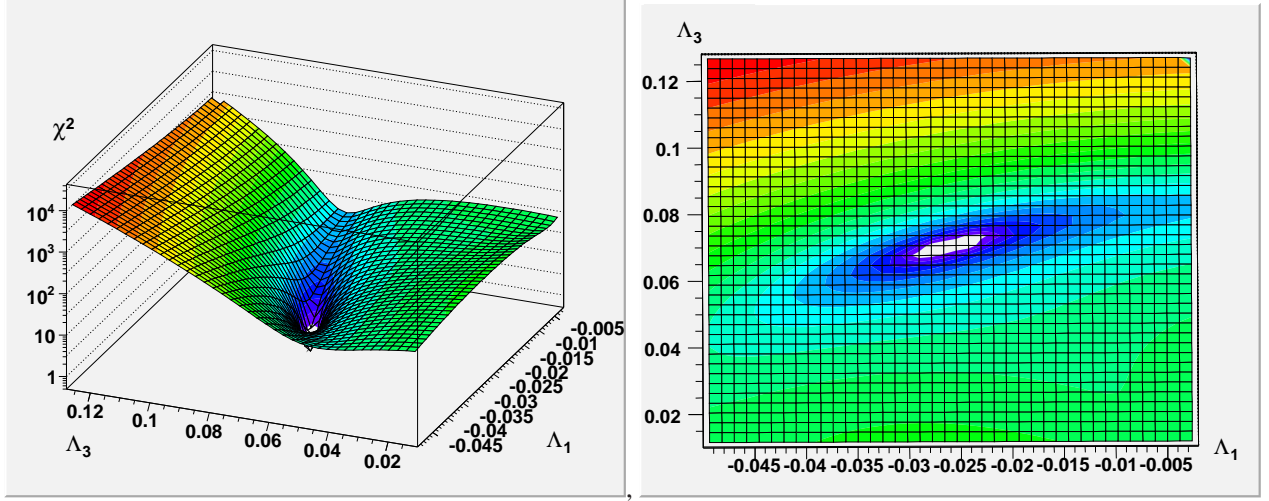


FIG. 3: Neutrino physics χ^2 as a function of the BRpV parameters Λ_1 and Λ_3 , keeping the rest of the parameters as indicated in Table I.

TABLE III: SU(5) SS and neutrino mass matrix parameters.

Parameter	Solution I	Solution II	Units
$\tan \beta$	10	10	-
μ	450	450	GeV
M_2	300	300	GeV
M_1	150	150	GeV
M_Σ	9×10^{15}	5×10^{15}	GeV
A	-1.7×10^3	-1.7×10^3	eV/GeV ²
C	6.7×10^{-3}	1.2×10^{-2}	eV

B. SUSY SU(5) Scenario

As expressed in Eq. (28) the tree level contribution from BRpV to the neutrino mass matrix is complemented in our $SU(5)$ supersymmetric model by a contribution suppressed by one power of the M_Σ mass scale. The dimensionless coefficients c_i are expected to be of order unity, but different from each other due to RGE effects. Despite that in Split Supersymmetric scenarios the light Higgs cannot contribute at one-loop to the neutrino mass matrix, this extra $SU(5)$ term is capable to generate a solar mass.

Keeping the low energy supersymmetric parameters equal to their values in the examples shown for Partial Split Supersymmetry in the previous section, we look for solutions in the case of SU(5) Split Susy

TABLE IV: SU(5) BRpV parameters and neutrino observables.

Parameter	Solution I	Solution II	Units
c_1	0.62	0.51	-
c_2	-0.52	-1.49	-
c_3	0.85	1.38	-
λ_1	0.0008	0.0015	GeV
λ_2	-0.0037	-0.0016	GeV
λ_3	-0.0038	-0.0011	GeV
Δm_{atm}^2	2.4×10^{-3}	2.6×10^{-3}	eV^2
Δm_{sol}^2	8.2×10^{-5}	8.3×10^{-5}	eV^2
$\tan^2 \theta_{atm}$	1.02	1.00	-
$\tan^2 \theta_{sol}$	0.45	0.50	-
$\tan^2 \theta_{13}$	0.026	0.049	-
m_{ee}	0.004	0.005	eV

with BRpV. In Table III we show two solutions for two different values of the scale M_Σ . The resulting neutrino mass coefficients A and C are also shown in the same Table. The A coefficient is independent of the mass scale M_Σ , but C is inversely proportional to it.

In Solution I, with a high value for $M_\Sigma = 9 \times 10^{15}$ GeV, the A term in the neutrino mass matrix dominates over the C term, such that the atmospheric mass comes mainly from $A\lambda_i\lambda_j$, and the smallness of the reactor angle is achieved with a small value for λ_1 . The solar mass is generated with the Cc_ic_j term, with the c_i of order unity.

In Solution II we lower the value for $M_\Sigma = 5 \times 10^{15}$ GeV, reversing the situation. Now the Cc_ic_j term dominates generating the atmospheric mass. Since we look for solutions with c_i of order one (we accept $0.5 < c_i < 1.5$), the value of $\tan^2 \theta_{13}$ grows to values close to its experimental upper bound. In this way, lower values of M_Σ are severely restricted. In this solution the solar mass is generated by the $A\lambda_i\lambda_j$ term.

V. SUMMARY

We have studied in detail the possibility to describe the neutrino masses and mixing angles in the context of split supersymmetric scenarios where the sfermions and/or Higgses are very heavy. We have considered all relevant contributions to the neutrino mass matrix up to one-loop level coming from the R-parity violating

interactions, showing the importance of the Higgs one-loop corrections in the case of Partial Split SUSY, where only the sfermions are very heavy. We have found new contributions in the context of the minimal supersymmetric SU(5) which can help us to generate the neutrino masses in agreement with the experiments in the SPLIT SUSY scenario.

Acknowledgments

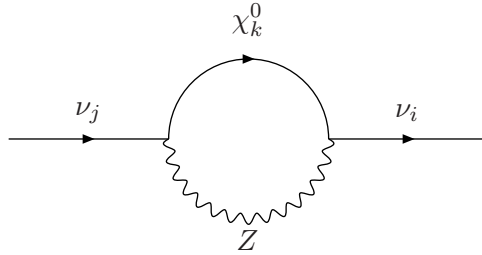
The work of M.A.D. was partly founded by Conicyt grant 1060629 and Conicyt and Banco Mundial grant “Anillo Centro de Estudios Subatómicos”. The work of P.F.P. has been supported by *Fundação para a Ciência e a Tecnologia* (FCT, Portugal) through the project CFTP, POCTI-SFA-2-777 and a fellowship under project POCTI/FNU/44409/2002. Also P.F.P. has been supported in part by the U.S. Department of Energy contract No. DE-FG02-08ER41531 and in part by the Wisconsin Alumni Research Foundation. We would like to thank E. J. Chun and G. Senjanović for a discussion. P. F. P thanks the Instituto de Física Corpuscular (IFIC) in Valencia for hospitality. The work by C. M. was partly supported by Conicyt fellowship, and by ALFA-EC in the framework of HELEN program. C. M. would like to thank CERN for their hospitality.

APPENDIX A: GAUGE AND GOLDSTONE BOSON LOOPS IN SPLIT SUSY

In this appendix we show the properties of the gauge boson one-loop contributions to the neutrino mass matrix.

1. Z and Neutral Goldstone Boson Loops

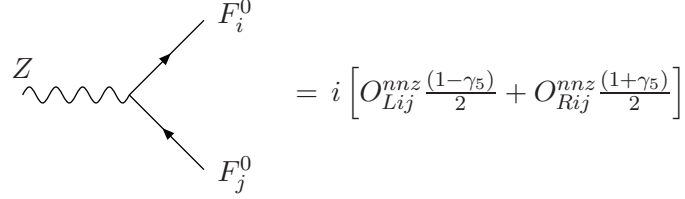
In Z loops the fermionic sum in Eq. (29) is over neutral fermions F_k^0 , of which only the neutralinos are relevant. There is no bosonic sum since only Z contributes.



The coupling G_{ijk}^Z is equal to

$$G_{ijk}^Z = -2(O_{Ljk}^{nnz} O_{Rki}^{nnz} + O_{Rjk}^{nnz} O_{Lki}^{nnz}), \quad (\text{A1})$$

where the coupling of a Z boson to two neutral fermions is [17],



$$= i \left[O_{Lij}^{nnz} \frac{(1-\gamma_5)}{2} + O_{Rij}^{nnz} \frac{(1+\gamma_5)}{2} \right]$$

with

$$O_{Lij}^{nnz} = -(O_{Rij}^{nnz})^*, \quad O_{Rij}^{nnz} = -\frac{g}{2c_W} \left(\mathcal{N}_{i4}^* \mathcal{N}_{j4} - \mathcal{N}_{i3}^* \mathcal{N}_{j3} - \sum_{a=1}^3 \mathcal{N}_{ia+4}^* \mathcal{N}_{ja+4} \right). \quad (\text{A2})$$

The matrix \mathcal{N} diagonalizes the 7×7 neutrino/neutralino mass matrix, giving non-negative eigenvalues. Without including the final rotation on the neutrino sector, it can be approximated in the following way [6]:

$$\mathcal{N} \approx \begin{bmatrix} N & N\xi^T \\ -\xi & 1 \end{bmatrix}, \quad (\text{A3})$$

where N diagonalizes the 4×4 neutralino mass sub-matrix. The parameters ξ are defined by

$$\begin{aligned} \xi_{i1} &= \frac{\tilde{g}'_d \mu M_2}{2 \det M_{\chi^0}} \Lambda_i, \\ \xi_{i2} &= -\frac{\tilde{g}_d \mu M_1}{2 \det M_{\chi^0}} \Lambda_i, \\ \xi_{i3} &= \frac{v_u}{4 \det M_{\chi^0}} (M_1 \tilde{g}_u \tilde{g}_d + M_2 \tilde{g}'_u \tilde{g}'_d) \Lambda_i - \frac{\epsilon_i}{\mu}, \\ \xi_{i4} &= -\frac{v_d}{4 \det M_{\chi^0}} (M_1 \tilde{g}_d^2 + M_2 \tilde{g}_d'^2) \Lambda_i. \end{aligned} \quad (\text{A4})$$

For notational brevity we define the ξ_i parameters as: $\lambda_i \xi_1 = \xi_{i1}$, $\lambda_i \xi_2 = \xi_{i2}$, $\lambda_i \xi_3 - \epsilon_i/\mu = \xi_{i3}$, and $\lambda_i \xi_4 = \xi_{i4}$. The couplings in Eq. (A2) can be approximated with the help of Eq. (A3) to

$$O_{Rik}^{\nu\chi z} \approx \frac{g}{2c_W} (2N_{k4}\xi_{i4} + N_{k1}\xi_{i1} + N_{k2}\xi_{i2}), \quad (\text{A5})$$

where i labels the three neutrinos and k labels the four neutralinos. Considering Eq. (A4) we conclude,

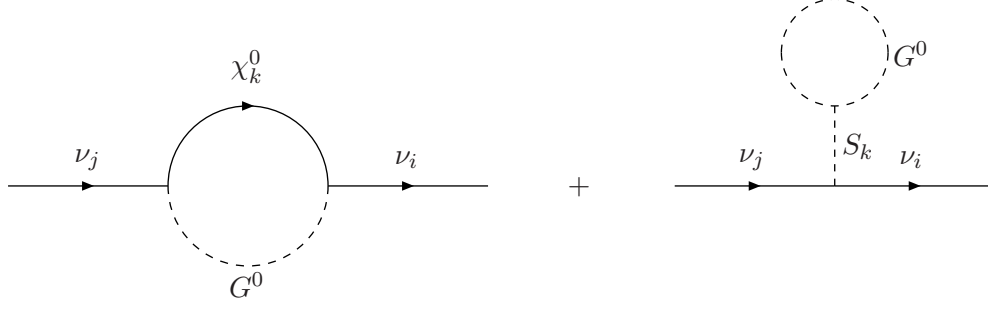
$$\Delta \Pi_{ij}^Z = A^Z \lambda_i \lambda_j, \quad (\text{A6})$$

with

$$A^Z = -\frac{g^2}{16\pi^2 c_W^2} \sum_{k=1}^4 (2N_{k4}\xi_4 + N_{k1}\xi_1 + N_{k2}\xi_2)^2 m_{\chi_k^0} B_0(0; m_{\chi_k^0}^2, m_Z^2). \quad (\text{A7})$$

This contribution is only a renormalization of the tree level mass matrix which it does not break its symmetry, i.e., it does not generate mass to all neutrinos.

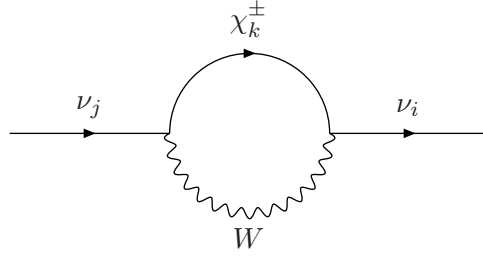
There is an extra contribution to A^Z dependent on the gauge parameter ξ . This is canceled by the following loops involving the neutral Goldstone boson,



as demonstrated in Ref. [6].

2. W and Charged Goldstone Boson Loops

In W loops the fermionic sum in Eq. (29) is over charged fermions F_k^+ , of which only the charginos are relevant. There is no bosonic sum since only W contributes.



The coupling G_{ijk}^W is equal to

$$G_{ijk}^W = -4(O_{Ljk}^{ncw} O_{Rik}^{ncw} + O_{Rjk}^{ncw} O_{Lik}^{ncw}), \quad (\text{A8})$$

where the coupling of a W boson to two fermions is

$$\begin{array}{c}
 F_i^0 \\
 \nearrow \\
 W \\
 \searrow \\
 F_j^+
 \end{array}
 = i \left[O_{Lij}^{ncw} \frac{(1-\gamma_5)}{2} + O_{Rij}^{ncw} \frac{(1+\gamma_5)}{2} \right]$$

with

$$\begin{aligned}
 O_{Lij}^{ncw} &= -g \left(\mathcal{N}_{i2}^* \mathcal{U}_{j1} + \frac{1}{\sqrt{2}} \mathcal{N}_{i3}^* \mathcal{U}_{j2} + \frac{1}{\sqrt{2}} \sum_{a=1}^3 \mathcal{N}_{ia+4}^* \mathcal{U}_{ja+2} \right), \\
 O_{Rij}^{ncw} &= -g \left(\mathcal{N}_{i2} \mathcal{V}_{j1}^* - \frac{1}{\sqrt{2}} \mathcal{N}_{i4} \mathcal{V}_{j2}^* \right).
 \end{aligned} \quad (\text{A9})$$

The \mathcal{U} and \mathcal{V} matrices diagonalize the 5×5 chargino/charged lepton mass matrix, and can be approximated to [6]

$$\mathcal{U} \approx \begin{bmatrix} U & U\xi_L^T \\ -\xi_L & 1 \end{bmatrix}, \quad \mathcal{V} \approx \begin{bmatrix} V & 0 \\ 0 & 1 \end{bmatrix}, \quad (\text{A10})$$

where U and V diagonalize the 2×2 chargino sub-matrix. The parameters ξ_L are

$$\xi_L^{i1} = \frac{\tilde{g}_d}{\sqrt{2} \det M_{\chi^+}} \Lambda_i, \quad \xi_L^{i2} = -\frac{\tilde{g}_u \tilde{g}_d v_u}{2\mu \det M_{\chi^+}} \Lambda_i - \frac{\epsilon_i}{\mu}, \quad (\text{A11})$$

with

$$\det M_{\chi^+} = \mu M_2 - \frac{1}{2} \tilde{g}_u \tilde{g}_d v_u v_d, \quad (\text{A12})$$

and similarly to what we did in the previous subsection, we define the parameters ξ_j^L , $j = 1, 2$, with the relations: $\xi_L^{i1} = \xi_1^L \lambda_i$ and $\xi_L^{i2} = \xi_2^L \lambda_i - \epsilon_i/\mu$. The couplings in Eq. (A9) can be approximated to

$$\begin{aligned} O_{Rij}^{\nu\chi^w} &\approx g \left(V_{j1}^* \xi_{i2} - \frac{1}{\sqrt{2}} V_{j2}^* \xi_{i4} \right), \\ O_{Lij}^{\nu\chi^w} &\approx g \left(U_{j1} \xi_{i2} - \frac{1}{\sqrt{2}} U_{j2} [\xi_L^{i2} - \xi_{i3}] - \frac{1}{\sqrt{2}} U_{j1} \xi_L^{i1} \right), \end{aligned} \quad (\text{A13})$$

where i labels the three neutrinos and j labels the two charginos. Similarly to what happened with the Z contributions, the W contribution depends only on the λ_i :

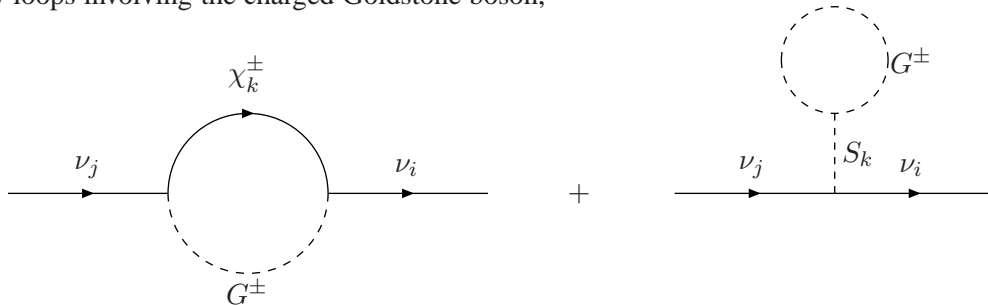
$$\Delta\Pi_{ij}^W = A^W \lambda_i \lambda_j, \quad (\text{A14})$$

with

$$A^W = \frac{g^2}{2\pi^2} \sum_{k=1}^2 \left[U_{k1} \xi_2 - \frac{U_{k2}}{\sqrt{2}} (\xi_2^L - \xi_3) + \frac{U_{k1}}{\sqrt{2}} \xi_1^L \right] \left(V_{k1} \xi_2 - \frac{V_{k2}}{\sqrt{2}} \xi_4 \right) m_{\chi_k^+} B_0(0; m_{\chi_k^+}^2, m_W^2). \quad (\text{A15})$$

Adding to the tree level contribution without changing the symmetry. Therefore the W and Z loops do not help us to generate mass to all neutrinos.

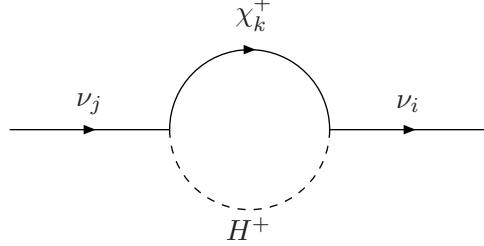
As for the case of A^Z , there is an extra contribution to A^W dependent on the gauge parameter ξ . This is canceled by loops involving the charged Goldstone boson,



The rest of the tadpoles form a gauge invariant set, and renormalize the vacuum expectation values [6].

3. Charged Higgs Boson Loops

The last loops we consider are the ones which include a charged scalar and a charged fermion. The loop is represented by the following graph,



where the G factor in Eq. (29) is,

$$G_{ijk}^{s+} = (O_{Ljkr}^{ncs} O_{Lkir}^{cns} + O_{Rjkr}^{ncs} O_{Rkir}^{cns}). \quad (\text{A16})$$

The relevant coupling above the scale \tilde{m} is the charged scalar couplings to a charged and a neutral fermion. It is given by,

$$S_k^+ \rightarrow F_i^+ F_j^0 = \left[O_{Lijk}^{cns} \frac{(1-\gamma_5)}{2} + O_{Rijk}^{cns} \frac{(1+\gamma_5)}{2} \right]$$

where the O_L^{cns} and O_R^{cns} couplings are,

$$\begin{aligned} O_{Lijk}^{cns} &= h_\tau R_{k1}^{S+} \mathcal{N}_{j7}^* \mathcal{V}_{i5}^* - R_{k2}^{S+} \left(\frac{g}{\sqrt{2}} \mathcal{N}_{j2}^* \mathcal{V}_{i2}^* + \frac{g'}{\sqrt{2}} \mathcal{N}_{j1}^* \mathcal{V}_{i2}^* + g \mathcal{N}_{j4}^* \mathcal{V}_{i1}^* \right) \\ &\quad - h_\tau R_{k5}^{S+} \mathcal{N}_{j3}^* \mathcal{V}_{i5}^* - \sqrt{2} g' R_{k\ell+5}^{S+} \mathcal{N}_{j1}^* \mathcal{V}_{i\ell+2}^*, \\ O_{Rijk}^{cns} &= R_{k1}^{S+} \left(\frac{g}{\sqrt{2}} \mathcal{N}_{j2} \mathcal{U}_{i2} + \frac{g'}{\sqrt{2}} \mathcal{N}_{j1} \mathcal{U}_{i2} - g \mathcal{N}_{j3} \mathcal{U}_{i1} \right) + h_\tau R_{k8}^{S+} (\mathcal{N}_{j7} \mathcal{U}_{i2} - \mathcal{N}_{j3} \mathcal{U}_{i5}) \\ &\quad + R_{k\ell+2}^{S+} \left(\frac{g}{\sqrt{2}} \mathcal{N}_{j2} \mathcal{U}_{i\ell+2} + \frac{g'}{\sqrt{2}} \mathcal{N}_{j1} \mathcal{U}_{i\ell+2} - g \mathcal{N}_{j\ell+4} \mathcal{U}_{i1} \right), \end{aligned} \quad (\text{A17})$$

with $O_{Lijk}^{cns} = O_{Rjik}^{ncs*}$ and $O_{Rijk}^{cns} = O_{Ljik}^{ncs*}$. The fields S_k^+ are eight linear combinations of charged Higgs bosons and charged sleptons, whose mass matrix in the $(H_d^+, H_u^+, \tilde{\ell}_{Lj}^+, \tilde{\ell}_{Rj}^+)$ basis is in Appendix B. This mass matrix is diagonalized in PSSusy by the rotation,

$$\begin{pmatrix} G^+ \\ H^+ \\ \tilde{\ell}_{Li}^+ \\ \tilde{\ell}_{Ri}^+ \end{pmatrix} = \begin{pmatrix} c_\beta & s_\beta & -s_L^j & 0 \\ -s_\beta & c_\beta & -t_L^j & 0 \\ c_\beta s_L^i - s_\beta t_L^i & s_\beta s_L^i + c_\beta t_L^i & \delta_{ij} & 0 \\ 0 & 0 & 0 & \delta_{ij} \end{pmatrix} \begin{pmatrix} H_d^+ \\ H_u^+ \\ \tilde{\ell}_{Lj}^+ \\ \tilde{\ell}_{Rj}^+ \end{pmatrix}. \quad (\text{A18})$$

An expression for the mixing angles s_L^i and t_L^i above the scale \tilde{m} can be found in the Appendix B. Comparing the supersymmetric lagrangian above the scale \tilde{m} in Eq. (34) with the terms of the PSSusy lagrangian in Eq. (20) we find the following matching conditions,

$$s_L^i(\tilde{m}) = b_i(\tilde{m})s_\beta; \quad t_L^i(\tilde{m}) = b_i(\tilde{m})c_\beta, \quad (\text{A19})$$

where $s_L^i(\tilde{m})$ represents the amount of slepton \tilde{L}_i present in the charged Goldstone boson G^+ , and analogously with $t_L^i(\tilde{m})$ for the low energy charged Higgs H^+ . In the limit where the sleptonic fields have a very large mass,

$$s_L^i \rightarrow s_\beta \frac{v_i}{v_u}, \quad t_L^i \rightarrow c_\beta \frac{v_i}{v_u}, \quad (\text{A20})$$

indicating that $b_i = v_i/v_u$ in agreement with the CP-even and CP-odd cases. Now we make an expansion of the couplings in Eq. (A17) and we find

$$\begin{aligned} O_{Lk}^{\nu\chi h+} &= c_\beta \left(\frac{g}{\sqrt{2}} \xi_{i2} V_{k2}^* + \frac{g'}{\sqrt{2}} \xi_{i1} V_{k2}^* + g \xi_{i4} V_{k1}^* \right), \\ O_{Rik}^{\nu\chi h+} &= s_\beta \left(\frac{g}{\sqrt{2}} \xi_{i2} U_{k2} + \frac{g'}{\sqrt{2}} \xi_{i1} U_{k2} + g \xi_{i3} U_{k1} \right) + g t_L^i U_{k1}, \end{aligned} \quad (\text{A21})$$

and isolating the terms proportional to ϵ_i , using Eq. (A19), we write,

$$O_{Lk}^{\nu\chi h+} = \tilde{O}_{Lk}^{\nu\chi h+} \Lambda_i, \quad O_{Rik}^{\nu\chi h+} = \tilde{O}_{Rk}^{\nu\chi h+} \Lambda_i - \frac{1}{\mu s_\beta} g U_{k1} \epsilon_i, \quad (\text{A22})$$

where we have defined,

$$\begin{aligned} \tilde{O}_{Lk}^{\nu\chi h+} &= c_\beta \left(\frac{g}{\sqrt{2}} \xi_{i2} V_{k2}^* + \frac{g'}{\sqrt{2}} \xi_{i1} V_{k2}^* + g \xi_{i4} V_{k1}^* \right), \\ \tilde{O}_{Rk}^{\nu\chi h+} &= s_\beta \left(\frac{g}{\sqrt{2}} \xi_{i2} U_{k2} + \frac{g'}{\sqrt{2}} \xi_{i1} U_{k2} + g \xi_{i3} U_{k1} \right) + g U_{k1} \frac{c_\beta}{\mu v_u}. \end{aligned} \quad (\text{A23})$$

Finally, the charged Higgs contribution to the neutrino mass matrix is,

$$\Delta \Pi_{ij}^{h+} = -\frac{1}{16\pi^2} \sum_{k=1}^2 \tilde{O}_{Lk}^{\nu\chi h+} \left[2\tilde{O}_{Rk}^{\nu\chi h+} \Lambda_i \Lambda_j - \frac{g U_{k1}}{\mu s_\beta} (\Lambda_i \epsilon_j + \Lambda_j \epsilon_i) \right] m_{\chi_k^+} B_0(0; m_{\chi_k^+}^2, m_{H^+}^2). \quad (\text{A24})$$

Note that there is no $\epsilon_i \epsilon_j$ term.

APPENDIX B: HIGGS SLEPTON SECTOR

Here we give details on the Higgs Slepton mass matrices and approximations in the case when the slepton masses are much heavier than the Higgs masses.

1. CP-even Higgs Sneutrino Mixing

The CP-even Higgs and sneutrino fields mix to form a set of five neutral mass eigenstates S_i^0 . We organize the mass terms in the lagrangian in the following way,

$$\mathcal{L} \ni -\frac{1}{2} \begin{bmatrix} \phi_d^0, \phi_u^0, \tilde{\ell}_{sj}^0 \end{bmatrix} \mathbf{M}_{S^0}^2 \begin{bmatrix} \phi_d^0 \\ \phi_u^0 \\ \tilde{\ell}_{sj}^0 \end{bmatrix}. \quad (\text{B1})$$

The mass matrix is divided into blocks [6],

$$\mathbf{M}_{S^0}^2 = \begin{bmatrix} \mathbf{M}_{S^0 hh}^2 & \mathbf{M}_{S^0 h\tilde{\nu}}^2 \\ \mathbf{M}_{S^0 h\tilde{\nu}}^{2T} & \mathbf{M}_{S^0 \tilde{\nu}\tilde{\nu}}^2 \end{bmatrix}. \quad (\text{B2})$$

The Higgs 2×2 sub-matrix is equal to,

$$\mathbf{M}_{S^0 hh}^2 = \begin{bmatrix} B_0 \mu \frac{v_u}{v_d} + \frac{1}{4} g_Z^2 v_d^2 + \mu \vec{\epsilon} \cdot \frac{\vec{v}}{v_d} + \frac{T_d}{v_d} & -B_0 \mu - \frac{1}{4} g_Z^2 v_d v_u \\ -B_0 \mu - \frac{1}{4} g_Z^2 v_d v_u & B_0 \mu \frac{v_d}{v_u} + \frac{1}{4} g_Z^2 v_u^2 - \vec{B}_\epsilon \cdot \frac{\vec{v}}{v_u} + \frac{T_u}{v_u} \end{bmatrix}, \quad (\text{B3})$$

where we call $g_Z^2 = g^2 + g'^2$, and in supergravity models we have $B_\epsilon^i = B_i \epsilon_i$. In this matrix we have eliminated the Higgs soft masses using the minimization conditions of the scalar potential (or tadpole equations) [6]. These Higgs tadpole equations at tree level are,

$$\begin{aligned} T_d &= (m_{H_d}^2 + \mu^2) v_d + v_d D - \mu (B_0 v_u + \vec{v} \cdot \vec{\epsilon}), \\ T_u &= -B_0 \mu v_d + (m_{H_u}^2 + \mu^2) v_u - v_u D + \vec{v} \cdot \vec{B}_\epsilon + v_u \vec{\epsilon}^2, \end{aligned} \quad (\text{B4})$$

with $D = \frac{1}{8}(g^2 + g'^2)(\vec{v}^2 + v_d^2 - v_u^2)$. At tree level, it is safe to set $T_u = T_d = 0$, and if we take the R-Parity conserving limit $\epsilon_i, v_i \rightarrow 0$, we can recognize the CP-even Higgs mass matrix of the MSSM. The 2×3 mixing sub-matrix is given by,

$$\mathbf{M}_{S^0 h\tilde{\nu}}^2 = \begin{bmatrix} M_{S^0 h_d \tilde{\nu}_i}^2 \\ M_{S^0 h_u \tilde{\nu}_i}^2 \end{bmatrix} = \begin{bmatrix} -\mu \epsilon_i + \frac{1}{4} g_Z^2 v_d v_i \\ B_\epsilon^i - \frac{1}{4} g_Z^2 v_u v_i \end{bmatrix}, \quad (\text{B5})$$

which vanishes in the R-Parity conserving limit. Finally, the sneutrino sub-matrix is given by,

$$(\mathbf{M}_{S^0 \tilde{\nu}\tilde{\nu}}^2)_{ij} = (M_{Li}^2 + D) \delta_{ij} + \frac{1}{4} g_Z^2 v_i v_j + \epsilon_i \epsilon_j, \quad (\text{B6})$$

where we have not yet used the corresponding tadpole equations, and we have assumed that the sneutrino soft mass matrix is diagonal. The sneutrino tadpole equations are given by,

$$T_i = v_i D + \epsilon_i (-\mu v_d + \vec{v} \cdot \vec{\epsilon}) + v_u B_\epsilon^i + v_i M_{Li}^2. \quad (\text{B7})$$

It is clear from this equation that if the sneutrino vev's are zero, $\mu\epsilon_i = B_\epsilon^i v_u/v_d$, and therefore, the mixing between the up and down Higgs fields with the sneutrino fields are related by $M_{S^0 h_d \tilde{\nu}}^2 = -\tan\beta M_{S^0 h_u \tilde{\nu}}^2$. Of course, this last relation is not valid if the sneutrino vev's are not zero.

In the case of large slepton masses, the mass matrix in Eq. (B2) is diagonalized in two steps by the rotation matrix,

$$R_{S^0} = \begin{pmatrix} 1 & 0 & -s_s^j \\ 0 & 1 & -t_s^j \\ s_s^i & t_s^i & \delta_{ij} \end{pmatrix} \begin{pmatrix} -s_\alpha & c_\alpha & 0 \\ c_\alpha & s_\alpha & 0 \\ 0 & 0 & \delta_{ij} \end{pmatrix}, \quad (\text{B8})$$

with the mixing angles at the scale \tilde{m} satisfying,

$$s_s^i = \frac{-s_\alpha M_{S^0 h_d \tilde{\nu}_i}^2 + c_\alpha M_{S^0 h_u \tilde{\nu}_i}^2}{M_{L_i}^2 - m_h^2}, \quad t_s^i = \frac{c_\alpha M_{S^0 h_d \tilde{\nu}_i}^2 + s_\alpha M_{S^0 h_u \tilde{\nu}_i}^2}{M_{L_i}^2 - m_H^2}, \quad (\text{B9})$$

where the Higgs masses can be neglected in front of the slepton masses in this approximation. From Eq. (B5) we find the following limits for large slepton masses,

$$s_s^i \longrightarrow -c_\alpha \frac{v_i}{v_u}, \quad t_s^i \longrightarrow -s_\alpha \frac{v_i}{v_u}, \quad (\text{B10})$$

which links the smallness of the Higgs-sneutrino mixing needed for neutrino physics, with the smallness of the sneutrino vevs.

2. CP-odd Higgs-Sneutrino Mixing

The CP-odd Higgs bosons and sneutrinos mix to form a set of five CP-odd scalars, whose mass terms in the Lagrangian are,

$$\mathcal{L} \ni -\frac{1}{2} \begin{bmatrix} \varphi_d^0, \varphi_u^0, \tilde{\ell}_{pi}^0 \end{bmatrix} \mathbf{M}_{P^0}^2 \begin{bmatrix} \varphi_d^0 \\ \varphi_u^0 \\ \tilde{\ell}_{pj}^0 \end{bmatrix}, \quad (\text{B11})$$

where the 5×5 mass matrix we decompose in the following blocks,

$$\mathbf{M}_{P^0}^2 = \begin{bmatrix} \mathbf{M}_{P^0 hh}^2 & \mathbf{M}_{P^0 h \tilde{\nu}}^2 \\ \mathbf{M}_{P^0 h \tilde{\nu}}^{2T} & \mathbf{M}_{P^0 \tilde{\nu} \tilde{\nu}}^2 \end{bmatrix}. \quad (\text{B12})$$

The Higgs sector is given by the 2×2 mass sub-matrix,

$$\mathbf{M}_{P^0 hh}^2 = \begin{bmatrix} B_0 \mu \frac{v_u}{v_d} + \mu \vec{\epsilon} \cdot \frac{\vec{v}}{v_d} + \frac{T_d}{v_d} & B_0 \mu \\ B_0 \mu & B_0 \mu \frac{v_d}{v_u} - \vec{B}_\epsilon \cdot \frac{\vec{v}}{v_u} + \frac{T_u}{v_u} \end{bmatrix}, \quad (\text{B13})$$

where the tadpoles T_u and T_d are defined in Eq. (B4). In the R-Parity conserving limit we reproduce the CP-odd mass matrix in the MSSM. The higgs-sneutrino mixing is given by the 2×3 matrix,

$$\mathbf{M}_{Ph\tilde{\nu}}^2 = \begin{bmatrix} M_{Ph_d\tilde{\nu}_i}^2 \\ M_{Ph_u\tilde{\nu}_i}^2 \end{bmatrix} = \begin{bmatrix} -\mu\epsilon_i \\ -B_\epsilon^i \end{bmatrix}, \quad (\text{B14})$$

which vanishes in the R-Parity conserving limit. Finally, the sneutrino 3×3 mass matrix is,

$$(\mathbf{M}_{P\tilde{\nu}\tilde{\nu}}^2)_{ij} = (M_{Li}^2 + D) \delta_{ij} + \epsilon_i \epsilon_j, \quad (\text{B15})$$

where we have assumed diagonal soft slepton mass parameters.

If slepton masses are very large, the 5×5 mass matrix can be diagonalized with the following rotations,

$$R_{P0} = \begin{pmatrix} 1 & 0 & -s_p^j \\ 0 & 1 & -t_p^j \\ s_p^i & t_p^i & \delta_{ij} \end{pmatrix} \begin{pmatrix} -c_\beta & s_\beta & 0 \\ s_\beta & c_\beta & 0 \\ 0 & 0 & \delta_{ij} \end{pmatrix}, \quad (\text{B16})$$

with the mixing angles s_p^i and t_p^i satisfying at the scale \tilde{m} ,

$$s_p^i = \frac{-c_\beta M_{Ph_d\tilde{\nu}_i}^2 + s_\beta M_{Ph_u\tilde{\nu}_i}^2}{M_{Li}^2 - m_G^2}, \quad t_p^i = \frac{s_\beta M_{Ph_d\tilde{\nu}_i}^2 + c_\beta M_{Ph_u\tilde{\nu}_i}^2}{M_{Li}^2 - m_A^2}, \quad (\text{B17})$$

and the Higgs masses m_G^2 and m_A^2 negligible in front of the slepton masses. Using Eqs. (B14) and (B17) we find the following mixing angles in the limit of large slepton masses,

$$s_p^i \longrightarrow s_\beta \frac{v_i}{v_u}, \quad t_p^i \longrightarrow c_\beta \frac{v_i}{v_u}, \quad (\text{B18})$$

also proportional to the sneutrino vacuum expectation values.

3. Charged Higgs Slepton Mixing

The charged Higgs boson and slepton fields mix to form a set of eight charged eigenstates S_i^+ , whose mass terms in the lagrangian are organized according to,

$$\mathcal{L} \ni - \left[H_d^-, H_u^-, \tilde{\ell}_{Li}^-, \tilde{\ell}_{Ri}^- \right] \mathbf{M}_{S^+}^2 \begin{bmatrix} H_d^+ \\ H_u^+ \\ \tilde{\ell}_{Li}^+ \\ \tilde{\ell}_{Ri}^+ \end{bmatrix}. \quad (\text{B19})$$

The 8×8 mass matrix is written as,

$$\mathbf{M}_{S^+}^2 = \begin{bmatrix} \mathbf{M}_{S^+hh}^2 & \mathbf{M}_{S^+h\tilde{\ell}}^2 \\ \mathbf{M}_{S^+h\tilde{\ell}}^{2T} & \mathbf{M}_{S^+\tilde{\ell}\tilde{\ell}}^2 \end{bmatrix}, \quad (\text{B20})$$

with the following charged Higgs boson 2×2 block,

$$\mathbf{M}_{S^{+}hh}^2 = \begin{bmatrix} B_0\mu\frac{v_u}{v_d} + \mu\vec{\epsilon} \cdot \frac{\vec{v}}{v_d} + \frac{1}{4}g^2(v_u^2 - \vec{v}^2) + \frac{1}{2}h_{\ell k}^2 v_k^2 + \frac{T_d}{v_d} & B_0\mu + \frac{1}{4}g^2 v_d v_u \\ B_0\mu + \frac{1}{4}g^2 v_d v_u & B_0\mu\frac{v_d}{v_u} - \vec{B}_\epsilon \cdot \frac{\vec{v}}{v_u} + \frac{1}{4}g^2(v_d^2 + \vec{v}^2) + \frac{T_u}{v_u} \end{bmatrix}. \quad (\text{B21})$$

This mass matrix reduces to the charged Higgs mass matrix of the MSSM when the BRpV parameters are taken equal to zero. Mixing between charged Higgs bosons and left and right charged sleptons appear through terms in the following 2×6 block,

$$\mathbf{M}_{S^{+}h\tilde{\ell}}^2 = \begin{bmatrix} M_{S^{+}h_d\tilde{\ell}_{Li}}^2 & M_{S^{+}h_d\tilde{\ell}_{Ri}}^2 \\ M_{S^{+}h_u\tilde{\ell}_{Li}}^2 & M_{S^{+}h_u\tilde{\ell}_{Ri}}^2 \end{bmatrix} = \begin{bmatrix} -\mu\epsilon_i + \frac{1}{4}g^2 v_d v_i - \frac{1}{2}h_{\ell i}^2 v_d v_i & -\frac{1}{\sqrt{2}}h_{\ell i}^2 v_u \epsilon_i - \frac{1}{\sqrt{2}}A_{\ell i} v_i \\ -B_\epsilon^i + \frac{1}{4}g^2 v_u v_i & -\frac{1}{\sqrt{2}}h_{\ell i}(\mu v_i + \epsilon_i v_d) \end{bmatrix}, \quad (\text{B22})$$

which as expected vanishes in the R-Parity conserving limit. The charged slepton sub-matrix is further divided into left and right slepton sectors,

$$\mathbf{M}_{S^{+}\tilde{\ell}\tilde{\ell}}^2 = \begin{bmatrix} \mathbf{M}_{LL}^2 & \mathbf{M}_{LR}^2 \\ \mathbf{M}_{LR}^{2T} & \mathbf{M}_{RR}^2 \end{bmatrix}, \quad (\text{B23})$$

which are given by the following expressions,

$$\begin{aligned} M_{LL}^2 &= \left[M_{L_i}^2 + \frac{1}{8}(g^2 - g'^2)(v_u^2 - v_d^2 - \vec{v}^2) + \frac{1}{2}h_{\ell i}^2 v_d^2 \right] \delta_{ij} + \frac{1}{4}g^2 v_i v_j + \epsilon_i \epsilon_j, \\ M_{LR}^2 &= \frac{1}{\sqrt{2}}(v_d A_{\ell i} - \mu v_u h_{\ell i}) \delta_{ij}, \\ M_{RR}^2 &= \left[M_{R_i}^2 + \frac{1}{4}g'^2(v_u^2 - v_d^2 - \vec{v}^2) + \frac{1}{2}h_{\ell i}^2(v_d^2 + \vec{v}^2) \right] \delta_{ij}. \end{aligned} \quad (\text{B24})$$

Slepton soft mass parameters are taken diagonal, and the MSSM expressions are recovered when we make $\epsilon_i = v_i = 0$. As before, if slepton soft masses are large, a diagonalization can be accomplished by the rotations,

$$R_{S^{+}} = \begin{pmatrix} 1 & 0 & -s_L^j & -s_R^j \\ 0 & 1 & -t_L^j & -t_R^j \\ s_L^i & t_L^i & \delta_{ij} & 0 \\ s_R^i & t_R^i & 0 & \delta_{ij} \end{pmatrix} \begin{pmatrix} c_\beta & s_\beta & 0 & 0 \\ -s_\beta & c_\beta & 0 & 0 \\ 0 & 0 & \delta_{ij} & 0 \\ 0 & 0 & 0 & \delta_{ij} \end{pmatrix}, \quad (\text{B25})$$

with the following mixing angles at the scale \tilde{m} ,

$$\begin{aligned} s_L^i &= \frac{c_\beta M_{S^{+}h_d\tilde{\ell}_{Li}}^2 + s_\beta M_{S^{+}h_u\tilde{\ell}_{Li}}^2}{M_{L_i}^2 - m_{H^{+}}^2}, & t_L^i &= \frac{-s_\beta M_{S^{+}h_d\tilde{\ell}_{Li}}^2 + c_\beta M_{S^{+}h_u\tilde{\ell}_{Li}}^2}{M_{L_i}^2 - m_{G^{+}}^2}, \\ s_R^i &= \frac{c_\beta M_{S^{+}h_d\tilde{\ell}_{Ri}}^2 + s_\beta M_{S^{+}h_u\tilde{\ell}_{Ri}}^2}{M_{R_i}^2 - m_{H^{+}}^2}, & t_R^i &= \frac{-s_\beta M_{S^{+}h_d\tilde{\ell}_{Ri}}^2 + c_\beta M_{S^{+}h_u\tilde{\ell}_{Ri}}^2}{M_{R_i}^2 - m_{G^{+}}^2}. \end{aligned} \quad (\text{B26})$$

When slepton masses are very large, the right mixing angles vanish while the left mixing angles are proportional to the slepton vevs,

$$s_L^i \longrightarrow s_\beta \frac{v_i}{v_u}, \quad t_L^i \longrightarrow c_\beta \frac{v_i}{v_u}, \quad s_R^i \longrightarrow 0, \quad t_R^i \longrightarrow 0, \quad (\text{B27})$$

in a similar way as the previous two cases.

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