PHYSICAL REVIEW D **80,** 113005 (2009)

Explaining solar neutrinos with heavy Higgs masses in partial split supersymmetry

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(Received 21 October 2009; published 14 December 2009)

Partial Split Supersymmetry with violation of R-parity as a model for neutrino masses is explored. It is shown that at the one-loop level the model can give predictions that are in agreement with all present experimental values for the neutrino sector. An analytical result is that the small solar neutrino mass difference can be naturally explained in the decoupling limit for the heavy Higgs mass eigenstates.

DOI: 10.1103/PhysRevD.80.113005 PACS numbers: 14.60.Pq, 12.60.Jv

I. INTRODUCTION

Supersymmetric models softly broken at low energies introduce new sources of flavor changing neutral currents and new CP-violating phases that can create unobserved phenomenological effects. If supersymmetric particles are of the order of the electroweak scale, it is not obvious why they do not contribute to processes with flavor changing neutral currents with rates higher than observed [1]. Large values of sparticle masses is one mechanism that explains the inconspicuous contributions from supersymmetry to flavor observables. Split supersymmetric models were introduced advertising precisely this feature [2].

R-parity violation in supersymmetric models [3] is an attractive feature because it provides a mechanism for neutrino mass generation. It is especially compelling in the case of bilinear R-parity violation (BRPV) [4] because an atmospheric mass difference is generated at tree level by a low energy seesaw mechanism triggered by a mixing between neutrinos and neutralinos. In addition, a solar mass difference is generated at one-loop level by contributions from all particles [5], explaining the hierarchy between atmospheric and solar mass scales. Original split supersymmetry (SS) conserves R-parity, nonetheless, the possibility has been explored, proving, for example, that a SS with R-parity violation cannot generate a solar mass [6–9].

In partial split supersymmetry (PSS) [7], all sfermions are very heavy and decouple from the low energy theory, alleviating the flavor constraints present in supersymmetric models. Nevertheless, as opposed to the original split supersymmetric model, both Higgs boson doublets remain light in comparison with the split supersymmetric scale $\tilde{m}$. Additionally we assume that R-parity is not conserved. As a consequence, a solar neutrino mass difference is generated at one-loop, while the atmospheric mass difference is generated at tree level.

In contrast to gravitationally inspired models [10], in PSS the solar neutrino mass difference is generated by loops involving the CP-odd Higgs A and the CP-even Higgs bosons h and H. We are particularly interested in the limit $m_A \gg m_Z$, called the decoupling limit [11], where the light Higgs h decouples from the heavy ones, acquiring SM-like couplings to fermions. Under this condition, the Higgs boson h does not contribute to the solar mass, and it is experimentally challenging to distinguish it from the standard model (SM) Higgs boson. In this article we study this scenario, and the ability of the mechanism to generate a solar mass from loops involving the heavy Higgs bosons H and A.

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

Partial split supersymmetry (SUSY) is an effective theory whose Lagrangian is valid at scales lower than $\tilde{m}$. The low energy Lagrangian includes two Higgs doublets H_u and H_d, and it is characterized by the following R-parity conserving terms,

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

In the first two lines we have the Higgs potential, including both mass and self-interaction terms. The electroweak symmetry is spontaneously broken as in the minimal supersymmetric standard model (MSSM) when the Higgs fields acquire nonvanishing vacuum expectation values. In the third line we have the Yukawa interactions, and in the fourth line we have the interactions between Higgs, gauginos, and Higgsinos. This Lagrangian is to be compared with the supersymmetric Lagrangian valid above the scale $\tilde{m}$, which includes the analogous terms,

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

These two models are connected through boundary conditions at the scale $\tilde{m}$. For the Higgs self-couplings they are,

$$\lambda_1 = \lambda_2 = \tfrac{1}{4}(g^2 + g'^2), \qquad \lambda_3 = \tfrac{1}{4}(g^2 - g'^2), \qquad \lambda_4 = -\tfrac{1}{2} g^2, \tag{3}$$

which are typical matching conditions between the MSSM and two Higgs doublet models. In an analogous way we have for the Yukawa couplings at $\tilde{m}$,

$$h_u = \lambda_u, \qquad h_d = \lambda_d, \qquad h_e = \lambda_e, \tag{4}$$

and for the Higgsino-gaugino Yukawa couplings at $\tilde{m}$,

$$\tilde{g}_u = \tilde{g}_d = g, \qquad \tilde{g}'_u = \tilde{g}'_d = g'. \tag{5}$$

The two Higgs fields H_u and H_d acquire a vacuum expectation value v_u and v_d, defining the usual mixing angle $\tan\beta = v_u/v_d$. The neutralino mass matrix in this scenario is written as

$$\mathbf{M}_{\chi^0}^{\rm PSS} = \begin{bmatrix} M_1 & 0 & -\frac{1}{2}\tilde{g}'_d v_d & \frac{1}{2}\tilde{g}'_u v_u \\ 0 & M_2 & \frac{1}{2}\tilde{g}_d v_d & -\frac{1}{2}\tilde{g}_u v_u \\ -\frac{1}{2}\tilde{g}'_d v_d & \frac{1}{2}\tilde{g}_d v_d & 0 & -\mu \\ \frac{1}{2}\tilde{g}'_u v_u & -\frac{1}{2}\tilde{g}_u v_u & -\mu & 0 \end{bmatrix}. \tag{6}$$

In partial split supersymmetry with nonconserved R-parity, neutrino masses are generated [7]. Trilinear R-parity violating (RPV) terms are irrelevant because high sfermion masses make their loop contributions negligible. Bilinear RPV terms do contribute via neutrino/neutralino mixing. The relevant terms are

$$\begin{aligned}\mathcal{L}_{\rm PSS}^{\rm RPV} = & -\epsilon_i \tilde{H}_u^T \epsilon L_i - \frac{1}{\sqrt{2}} b_i H_u^T \epsilon (\tilde{g}_d \sigma \tilde{W} - \tilde{g}'_d \tilde{B}) L_i \\ & + \text{h.c.}, \end{aligned} \tag{7}$$

where ϵ_i are the usual mass parameters that mix Higgs with lepton superfields, and b_i are effective couplings between Higgs, gauginos, and leptons. After the Higgs fields acquire vacuum expectation values, mixing terms are generated between neutrinos, on one hand, and Higgsinos and gauginos on the other hand,

$$\mathcal{L}_{\rm PSS}^{\rm RPV} = -[\epsilon_i \tilde{H}_u^0 + \tfrac{1}{2} b_i v_u (\tilde{g}_d \tilde{W}_3 - \tilde{g}'_d \tilde{B})] \nu_i + \text{h.c.} + \ldots \tag{8}$$

and they extend the 4×4 neutralino mass matrix, in Eq. (6), to a 7×7 mass matrix that includes the neutrinos. The off diagonal mixing block is

$$m^{\rm PSS} = \begin{bmatrix} -\frac{1}{2}\tilde{g}'_d b_1 v_u & \frac{1}{2}\tilde{g}_d b_1 v_u & 0 & \epsilon_1 \\ -\frac{1}{2}\tilde{g}'_d b_2 v_u & \frac{1}{2}\tilde{g}_d b_2 v_u & 0 & \epsilon_2 \\ -\frac{1}{2}\tilde{g}'_d b_3 v_u & \frac{1}{2}\tilde{g}_d b_3 v_u & 0 & \epsilon_3 \end{bmatrix}, \tag{9}$$

while the neutrino-neutrino block is a null 3×3 matrix. After a low energy seesaw mechanism, the effective neutrino mass matrix is

$$\begin{aligned}\mathbf{M}_\nu^{\rm eff} & = -m^{\rm PSS} (\mathbf{M}_{\chi^0}^{\rm PSS})^{-1} (m^{\rm PSS})^T \\ & = \frac{M_1 \tilde{g}_d^2 + M_2 \tilde{g}_d'^2}{4 \det M_{\chi^0}^{\rm PSS}} \begin{bmatrix} \Lambda_1^2 & \Lambda_1\Lambda_2 & \Lambda_1\Lambda_3 \\ \Lambda_2\Lambda_1 & \Lambda_2^2 & \Lambda_2\Lambda_3 \\ \Lambda_3\Lambda_1 & \Lambda_3\Lambda_2 & \Lambda_3^2 \end{bmatrix}, \end{aligned} \tag{10}$$

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

$$\begin{aligned}\det M_{\chi^0}^{\rm PSS} = & -\mu^2 M_1 M_2 + \tfrac{1}{2} v_u v_d \mu (M_1 \tilde{g}_u \tilde{g}_d + M_2 \tilde{g}'_u \tilde{g}'_d) \\ & + \tfrac{1}{16} v_u^2 v_d^2 (\tilde{g}'_u \tilde{g}_d - \tilde{g}_u \tilde{g}'_d)^2. \end{aligned} \tag{11}$$

As it is well known, the neutrino mass matrix in Eq. (10) has only one eigenvalue different from zero, and quantum corrections must be added in order to generate a solar mass.

III. HIGGS DECOUPLING LIMIT

The Higgs decoupling limit in the MSSM, or in the nonsupersymmetric two Higgs doublet model, is the regime where the lightest Higgs scalar mass is much smaller than all the other Higgs boson masses [11]. It has been studied in detail because if realized in nature, it will be experimentally difficult to distinguish the lightest Higgs boson properties from the ones of the SM Higgs boson.

If we neglect the running of the Higgs potential parameters, in first approximation the Higgs sector in PSS is analogous to the one in the MSSM. At tree level we have for the neutral CP-even Higgs bosons the following masses:

$$m_{h,H}^2 = \tfrac{1}{2}(m_A^2 + m_Z^2) \mp \tfrac{1}{2}\sqrt{(m_A^2 + m_Z^2)^2 - 4 m_A^2 m_Z^2 c_{2\beta}^2} \tag{12}$$

as a function of the neutral CP-odd Higgs mass m_A and $\tan\beta$. Although m_h receive large quantum corrections, they are negligible for the heavy Higgs H when m_A is large. Thus the tree-level formula for m_H is adequate in the decoupling limit. In this limit, when $m_A \gg m_Z$, we have

$$m_H^2 \approx m_A^2 + m_Z^2 \sin^2(2\beta) \tag{13}$$

which is a very good approximation already for $m_A > 200$ GeV. The CP-even Higgs mass matrix is diagonalized by a rotation with an angle α, which satisfies at tree level,

$$\cos^2(\alpha - \beta) = \frac{m_h^2(m_Z^2 - m_h^2)}{m_A^2(m_H^2 - m_h^2)}. \tag{14}$$

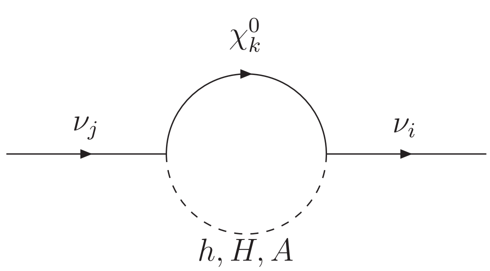
The quantity $\cos(\alpha - \beta)$ is important because it is equal to the ratio between the heavy Higgs coupling to two Z bosons (and two W bosons) and the same couplings for the SM Higgs boson, H_{SM}. Thus, it is a key parameter for the determination of the H production cross section in association with gauge bosons. Conversely, $\sin(\alpha - \beta)$ is proportional to the light Higgs boson couplings to two gauge bosons. In this limit we have

$$\cos^2(\alpha - \beta) \approx \frac{m_Z^4 \sin^2(4\beta)}{4m_A^4}. \tag{15}$$

This means that H decouples from the low energy theory. Thus the h production cross section becomes ever more similar to the H_{SM} one. The same happens to the h couplings to fermions: they become similar to the H_{SM} couplings in the decoupling limit, making it an experimental challenge to differentiate a H_{SM} from a h in this scenario.

IV. LOOP CORRECTIONS TO NEUTRINO MASSES IN THE DECOUPLING LIMIT

In PSS the only loops capable to contribute to the solar mass are loops involving neutralinos and neutral Higgs bosons.



Radiative corrections to the effective neutrino mass matrix are of the form

$$\Delta\Pi_{ij} = A\Lambda_i\Lambda_j + B(\Lambda_i\epsilon_j + \Lambda_j\epsilon_i) + C\epsilon_i\epsilon_j, \tag{16}$$

where the A-term is the only one that receives tree-level contributions, as indicated in Eq. (10). Charged and neutral gauge bosons contribute only to the A-term, i.e., to the atmospheric mass. The charged Higgs boson contributes to A and B-terms. However, the B-term is scale dependent and can be rendered zero with an appropriate choice for the arbitrary subtraction scale Q. Therefore, charged Higgs bosons do not contribute either to the solar mass [7]. This leaves only the neutral Higgs bosons, whose contributions to the C-term are

$$\Delta\Pi_{ij}^{h,H}|_{\epsilon\epsilon} = -\frac{1}{16\pi^2}\sum_{k=1}^{4}(F_k^{h,H})^2\epsilon_i\epsilon_j m_{\chi_k^0}B_0(0; m_{\chi_k^0}^2, m_{h,H}^2),$$
$$\Delta\Pi_{ij}^{A}|_{\epsilon\epsilon} = \frac{1}{16\pi^2}\sum_{k=1}^{4}(F_k^{A})^2\epsilon_i\epsilon_j m_{\chi_k^0}B_0(0; m_{\chi_k^0}^2, m_{A}^2), \tag{17}$$

where the sum is over all four neutralinos, with masses $m_{\chi_k^0}$, and B_0 is the usual Veltman function for two-point Green functions. Note that h and H contributions have an overall minus sign, while the A contribution does not. The reason is that A has CP-odd couplings to fermions. The couplings F_k are equal to

$$F_k^h = \frac{\cos(\alpha - \beta)}{2\mu s_\beta}(gN_{k2}^* - g'N_{k1}^*),$$
$$F_k^H = \frac{\sin(\alpha - \beta)}{2\mu s_\beta}(gN_{k2}^* - g'N_{k1}^*), \tag{18}$$
$$F_k^A = \frac{1}{2\mu s_\beta}(gN_{k2}^* - g'N_{k1}^*).$$

The factor in the parenthesis indicates that it is the gaugino component of the neutralinos, the one that contributes to the solar neutrino mass. In addition, the B_0 Veltman's function for zero external momentum is

$$B_0(0; M^2, m^2) = \Delta + 1 - \frac{M^2 \ln\frac{M^2}{Q^2} - m^2 \ln\frac{m^2}{Q^2}}{M^2 - m^2} \tag{19}$$

where Q is the arbitrary subtraction scale and Δ is the regulator. In this way, the one-loop contribution to the C coefficient of the $\epsilon_i\epsilon_j$ term in Eq. (16) is

$$C^{AHh} = \frac{1}{64\pi^2\mu^2 s_\beta^2}\sum_{k=1}^{4} m_{\chi_k^0}(gN_{k2} - g'N_{k1})^2 \times [B_0(0; m_{\chi_k^0}^2, m_A^2) - \sin^2(\alpha - \beta)B_0(0; m_{\chi_k^0}^2, m_H^2) - \cos^2(\alpha - \beta)B_0(0; m_{\chi_k^0}^2, m_h^2)]. \tag{20}$$

The fact that the divergent term in B_0 is independent of all masses implies that the C coefficient is finite, which in turn leads to the finiteness of the solar mass.

The neutrino masses generated from Eq. (16) include a massless neutrino $m_{\nu_1} = 0$, and the two massive ones given by

$$m_{\nu_{3,2}} = \tfrac{1}{2}(A|\vec{\Lambda}|^2 + C|\vec{\epsilon}|^2) \pm \tfrac{1}{2}\sqrt{(A|\vec{\Lambda}|^2 + C|\vec{\epsilon}|^2)^2 - 4AC|\vec{\Lambda} \times \vec{\epsilon}|^2} \tag{21}$$

where the sign is chosen such that $|m_{\nu_2}| < |m_{\nu_3}|$. We are interested in the behavior of the solar mass in the Higgs decoupling limit, where $m_A \gg m_Z$. The B_0 Veltman function in Eq. (19) has the following expansion when one of the masses is much larger than the other, $M \gg m$,

$$B_0(0; M^2, m^2) \approx \Delta - \ln\frac{M^2}{Q^2} + 1 - \frac{m^2}{M^2}\ln\frac{M^2}{m^2} - \frac{m^4}{M^4}\ln\frac{M^2}{m^2}. \tag{22}$$

In the regime where the neutralinos are much lighter than the CP-odd Higgs mass, we find for the C coefficient

$$C^{AHh} \approx \frac{m_Z^2 \sin^2 2\beta}{64\pi^2\mu^2 s_\beta^2 m_A^2}\sum_{k=1}^{4} m_{\chi_k^0}(gN_{k2} - g'N_{k1})^2 \tag{23}$$

where the terms in parenthesis correspond to the zino component of each neutralino. This motivates us to define

$$\langle m_{\tilde{Z}}\rangle \equiv \sum_{k=1}^{4} m_{\chi_k^0}(c_W N_{k2} - s_W N_{k1})^2 \tag{24}$$

as the zino effective mass. In this way, we find that the CP-odd Higgs mass is related to the solar neutrino mass difference by the following simple relation:

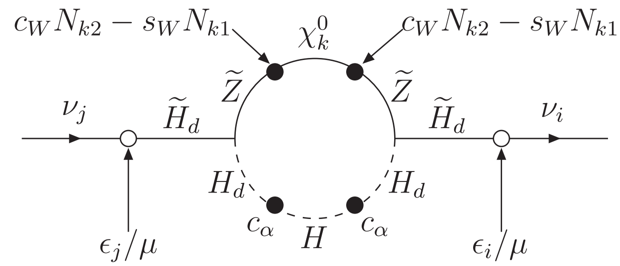
$$m_A^2 \approx \frac{g^4 m_Z^2 \cos^2\beta}{64\pi^2 c_W^4}\frac{\langle m_{\tilde{Z}}\rangle m_{\tilde{\gamma}}}{M_1 M_2}\sqrt{\frac{\delta}{1+\delta}}\frac{|\vec{\Lambda}\times\vec{\epsilon}|^2}{\mu^4\Delta m_{\mathrm{sol}}^2} \tag{25}$$

where $\delta = \Delta m_{\mathrm{sol}}^2/\Delta m_{\mathrm{atm}}^2 \approx 0.035$ and $m_{\tilde{\gamma}} = c_W^2 M_1 + s_W^2 M_2$ is the photino mass.

This formula is remarkable. First we notice that the CP-odd Higgs mass squared is inversely proportional to the solar mass difference. The reason behind this feature is as follows. In the Higgs decoupling limit the light Higgs has SM-like couplings and does not contribute to the solar mass. More precisely, h contribution to the neutrino mass matrix is proportional to $\cos^2(\alpha - \beta)$, which rapidly approaches zero as $m_A \gg m_Z$, as indicated by Eq. (15). The other two neutral Higgs bosons, H and A, have large contributions to the neutrino mass matrix, but with opposite signs, and as it can be seen already from Eq. (17) they tend to cancel each other in the decoupling limit. Therefore, the fact that supersymmetry forces $m_H \to m_A$ when $m_A \gg m_Z$, produces a fine cancellation that eventually generates a small solar mass. This fine cancellation is not a fine-tuning because it is a cancellation forced by symmetry. In addition, the CP-odd Higgs mass is dependent on the atmospheric mass through the ratio δ between solar atmospheric scales. Thus, the atmospheric mass also affects m_A^2 inversely although in an indirect way. Our model explains the smallness of δ because the atmospheric mass is generated at tree-level while the solar mass is generated at one-loop.

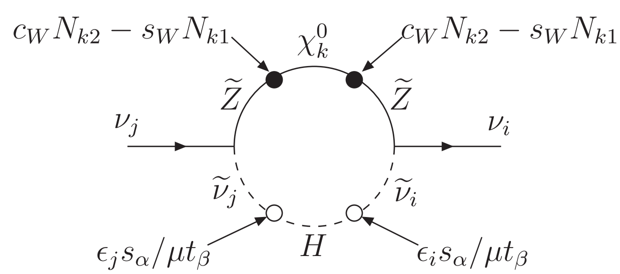
The CP-odd squared mass is proportional to $|\vec{\Lambda}\times\vec{\epsilon}|$, with an extra term μ^4 in the denominator. This cross product is there because if $\vec{\epsilon}$ and $\vec{\Lambda}$ are parallel, the symmetry of the neutrino mass matrix observed at tree-level in Eq. (11) is not removed, and no solar mass is generated. Furthermore, m_A^2 is proportional to the Z boson mass, indicating that a Majorana neutrino mass needs not only R-parity violation but a broken $SU(2)_L$ symmetry as well. Additionally, m_A^2 is proportional to $\cos^2\beta$ because the relevant vacuum expectation value is $\langle H_d\rangle$.

We also notice that the CP-odd Higgs mass squared is proportional to the effective zino mass and to the photino mass, normalized by both relevant gaugino masses. The appearance of the photino mass is due to the tree-level contribution to the effective neutrino mass matrix. The origin of the zino mass is that it is the zino component of each neutralino, the one that contributes to the neutrino mass matrix, and it multiplies the neutralino masses from the fermionic propagator. In addition, for this mechanism to work, R-parity must be broken, and the neutrino mixing with down Higgsinos in the one hand and Higgs bosons mixing with sneutrinos in the other hand, are crucial. The following schematic diagram, corresponding to the heavy Higgs boson loops, may help us to understand the origin of the one-loop contributions to the neutrino mass matrix,



The supersymmetric vertex behind this diagram is the zino coupling to down Higgs and down Higgsino. For this reason, the zino component from each neutralino is selected, $c_W N_{k2} - s_W N_{k1}$, which weights the corresponding neutralino mass picked up from the propagator. In addition, the down Higgs component from the heavy Higgs H, given by c_α, is selected. These mixings are represented in the diagram by full circles, as oppose to open circles which violate R-parity. Indeed, R-parity is violated at the mixing between neutrinos and down Higgsinos. These two mixings at the external legs also violate the lepton number by two units, as it should be for a Majorana neutrino mass.

The second diagram is



where the supersymmetric vertex supporting the diagram is the zino coupling to a neutrino and a sneutrino. The presence of the zino component of each neutralino is explained by the same argument as before. But in this case, R-parity and lepton number violation appear in the Higgs boson mixing with sneutrinos. The magnitude of this mixing, indicated in the diagram, is explained in Ref. [7].

Diagrams with one of each supersymmetric vertex also contribute, but are not shown.

V. NUMERICAL RESULTS

In our numerical analysis the contributions to the neutrino mass matrix in Eq. (16) were calculated and their influence on the neutrino observables such as mass differences and mixing angles was studied. The agreement with the experimental boundaries (3σ) [12] was quantified by calculating

$$\chi^2 = \left(\frac{10^3\Delta m^2_{\rm atm} - 2.4}{0.4}\right)^2 + \left(\frac{10^5\Delta m^2_{\rm sol} - 7.7}{0.6}\right)^2 + \left(\frac{\sin^2\theta_{\rm atm} - 0.505}{0.165}\right)^2 + \left(\frac{\sin^2\theta_{\rm sol} - 0.33}{0.07}\right)^2. \quad (26)$$

Additionally it was demanded that the upper bounds $\sin^2\theta_{\rm reac} < 0.05$ and $m_{\beta\beta} < 0.84$ eV have to be fulfilled. Thus, the model is in agreement with the experimental values if $\chi^2 < 4$. As a working scenario, we select a typical point in the PSS parameter space, denoted as P_t. This scenario consists of fixed values for the gaugino and Higgsino mass parameters, $\tan\beta$, the light CP-even and CP-odd Higgs masses, and the BRPV parameters, all given in Table I. This scenario satisfies the neutrino experimental constraints, predicting atmospheric and solar mass differences and mixing angles well within the 3σ regions indicated in Eq. (26). In addition, the predicted reactor angle and the neutrinoless double-beta decay mass parameter are below the upper bound. All these predictions of our P_t scenario are given in Table I.

TABLE I. PSS and RPV parameters and neutrino observables for the working scenario P_t.

SUSY parameter	P_t	Scanned range	Units
$\tan\beta$	10	[2, 50]	$\cdots$
$\lvert\mu\rvert$	450	[0, 1000]	GeV
M_2	300	[80, 1000]	GeV
M_1	150	$M_2/2$	GeV
m_h	120	[114, 140]	GeV
m_A	1000	[50, 6000]	GeV
Q	951.7	$\cdots$	GeV
RPV parameter			
ϵ_1	0.0346	$\cdots$	GeV
ϵ_2	0.2516	$[-1, 1]$	GeV
ϵ_3	0.3504	$[-1, 1]$	GeV
Λ_1	-0.0259	$[-1, 1]$	GeV^2
Λ_2	-0.0011	$\cdots$	GeV^2
Λ_3	0.0709	$[-1, 1]$	GeV^2

Observable	Solution	Units
$\Delta m^2_{\rm atm}$	2.45×10^{-3}	eV^2
$\Delta m^2_{\rm sol}$	7.9×10^{-5}	eV^2
$\tan^2\theta_{\rm atm}$	0.824	$\cdots$
$\tan^2\theta_{\rm sol}$	0.487	$\cdots$
$\tan^2\theta_{13}$	0.027	$\cdots$
m_{ee}	0.0016	eV

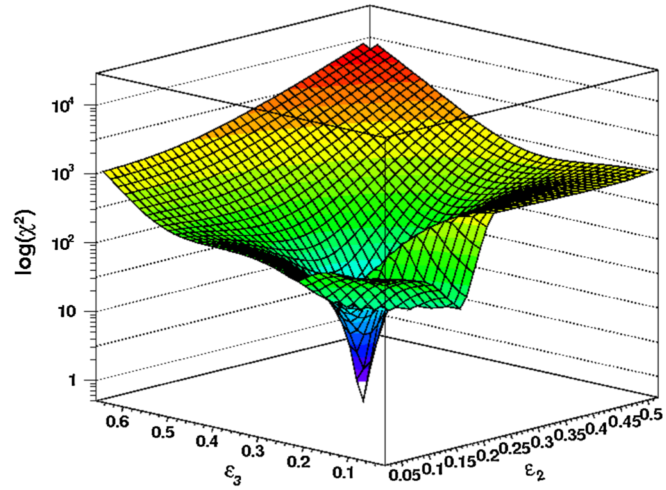


FIG. 1 (color online). χ^2 in dependence of ϵ_2 and ϵ_3, while the other parameters are fixed around the central value from Table I.

Next we scan the parameter space varying the PSS parameters according to the intervals indicated in Table I. In Fig. 1, we vary ϵ_2 and ϵ_3 keeping all the other parameters as in P_t, and plot the logarithm of χ^2 associated to each point in parameter space. Solutions with $\chi^2 < 4$ are clearly visible; they are compatible with experiments, and P_t is inside this region. The value for χ^2 grows fast as we deviate from the experimentally accepted region around P_t. In Fig. 2 we see the dependence on Λ_1 and Λ_3 while keeping the rest of the PSS parameters as indicated by P_t. Again there is a steep growth on χ^2 when we deviate from the experimentally accepted region. Of course, these ex-

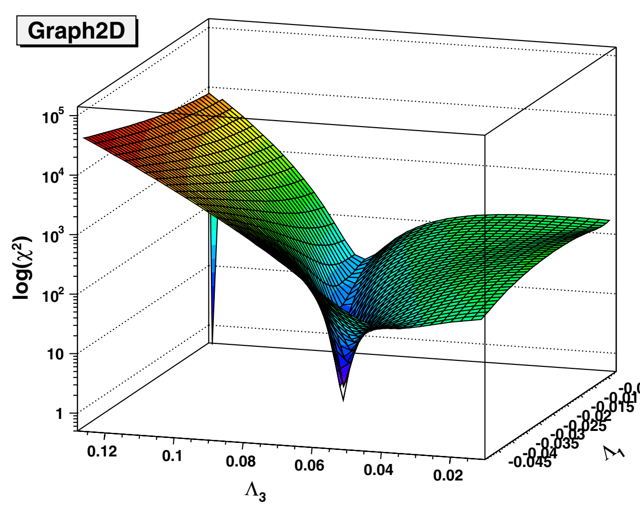


FIG. 2 (color online). χ^2 in dependence of Λ_1 and Λ_3, while the other parameters are fixed around the central value from Table I.

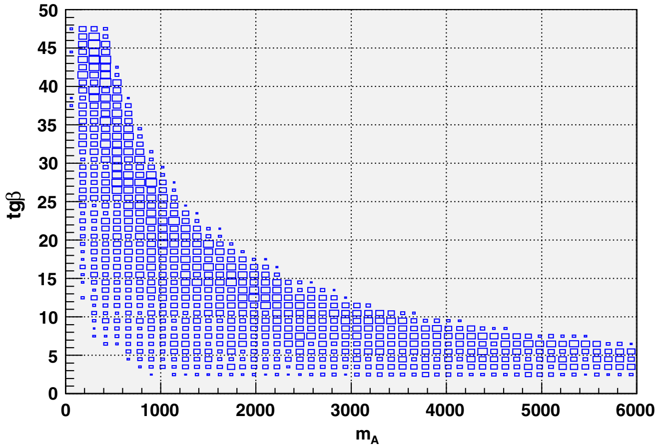


FIG. 3 (color online). Relation between $\tan\beta$ and m_A, while varying the other parameters as indicated in Table I.

perimentally compatible regions move around in the ϵ_2-ϵ_3 and Λ_1-Λ_3 planes when we change the fixed values of the other PSS parameters.

For the scan shown in Figs. 3 and 4 we keep the BRPV parameters ϵ_i and Λ_i fixed to their P_t values, and vary the rest of the PSS parameters as indicated in Table I. In Fig. 3 we have the relation between $\tan\beta$ and the CP-odd Higgs mass m_A. The scan shows that an agreement with the neutrino observables is disfavored if both SUSY parameters $\tan\beta$ and m_A take simultaneously large values. Because of this correlation a measurement of the value of one of the two parameters might give important information on the value of the other, even if the rest of the SUSY parameters are not known. A detailed study shows that the excluded region in the ($\tan\beta - m_A$) plane comes due to the observed constraint on $\tan^2\theta_{\rm atm}$. In Fig. 4 we show the relation between the Higgsino mass parameter μ and m_A from the same scan as before. This reveals that for large $m_A \approx 6000$ GeV good results compatible with neutrino experiments are only obtained for $200 < |\mu| < 350$ GeV, whereas for smaller $m_A \approx 1000$ GeV the preferred region

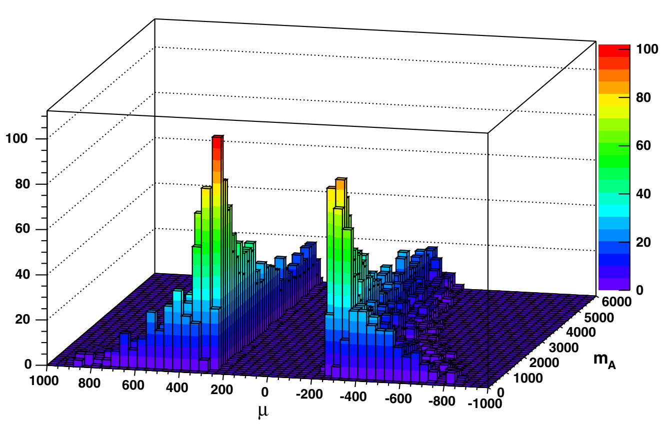


FIG. 4 (color online). Frequency plot in the μ-m_A plane, for points in parameter space that are consistent with neutrino experiments.

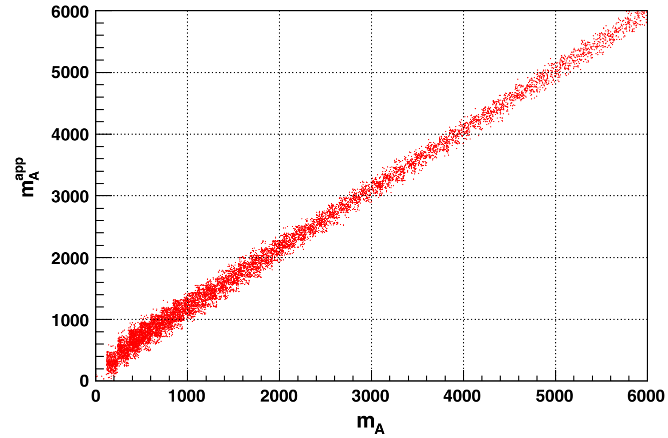


FIG. 5 (color online). Approximated $m_A^{\rm app}$ as a function of its exact value m_A, for a scan with 10^4 random points.

for μ widens up to $200 < |\mu| < 600$ GeV. A detailed study shows that the observed values for $\Delta m^2_{\rm sol}$ and $\Delta m^2_{\rm atm}$ forbid solutions with $|\mu| < 200$ GeV. Finally, in order to test the accuracy of the approximated formula given in Eq. (25), we perform a scan with 10^4 random points. In Fig. 5 we include the points that satisfy the neutrino experimental constraints, and we see that the approximated formula works in a very large range of m_A values, with larger percentage errors for smaller m_A, which is expected.

VI. SUMMARY

This paper explored partial split supersymmetry with RPV as a model for neutrino masses. It was shown that at the one-loop level the model can give predictions that are in very good agreement with all present experimental values for the neutrino sector. In contrast to this good agreement in PSS, it is not possible to generate all the neutrino mass parameters correctly in standard split supersymmetry with RPV. The difference between both models, lies solely in the fact that PSS allows for a larger Higgs sector, which contains the mass eigenstates A and H in addition the standard model-like state h. A continuous transition from PSS to SS can be achieved by raising the values for the heavy Higgs masses m_A and m_H (Higgs decoupling limit). An analytical study of this limit ($m_A \gg m_h$) reveals an approximate formula for PSS in which a large value for m_A is directly connected to a small solar neutrino mass difference $\Delta m^2_{\rm sol}$. Therefore, the small observed value for $\Delta m^2_{\rm sol}$ favors large values of m_A up to 6 TeV. Such large values for m_A would make PSS virtually indistinguishable from SS by using any observable other than neutrino masses.

ACKNOWLEDGMENTS

The work of M. A. D. was partly funded by Conicyt-PBCT Grant No. ACT028 (Anillo Centro de Estudios

Subatómicos), and by Conicyt-PBCT grant ACI35. B. K. was funded by Conicyt-PBCT grant PSD73.

Note Added.—While this article was being written, we read the paper ‘‘SUSY Splits, But Then Returns’’ by Professor Raman Sundrum [13], where models with two light Higgs doublets are also referred to as partial split supersymmetry. In the context of string theory, the expression ‘‘partial split supersymmetry’’ was used for a different kind of effective model [14].

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