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CP violation in decays of the lightest supersymmetric particle with bilinearly broken R parity

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Abstract: Supersymmetric models with broken R-parity induced by lepton number violating terms provide a calculable framework for neutrino masses and mixings. Within models with bilinear R-parity breaking six new physical phases appear which are potential sources of novel CP-violating phenomena compared to the minimal supersymmetric extension of the standard model. We consider CP-violating observables in the decays of the lightest supersymmetric particle in this class of models. We show that: (i) Neutrino physics requires a strong correlation between three different pairs of phases, thus reducing the effective number of new phases to three. (ii) CP-violating phenomena in decays of the lightest supersymmetric particle due to new R-parity breaking phases turn out to be small, once constraints from neutrino physics are taken into account. We demonstrate that this feature does not depend on the nature of the lightest supersymmetric particle.

Keywords: supersymmetry, R parity, violation, violation, CP.

Contents

1. Introduction

The standard model (SM) contains a source for CP violation, the complex phase of the Kobayashi-Maskawa matrix [1]. However, it has often been argued that the SM is not able to explain the observed baryon/antibaryon asymmetry of the universe (see for example [2]) and thus new physics might be expected to show up in CP violating phenomena.

In supersymmetric (SUSY) extensions of the SM, with their much larger particle content, a considerable number of parameters could in principle be complex and consequently CP violation phenomena differ considerably from SM expectations [3]. It is therefore not surprising that the study of CP violation in the minimal supersymmetric extension of the standard model (MSSM) [4] or its constrained version (sometimes called CMSSM or mSUGRA in the literature) has received quite some attention recently [5, 6].

The new phases present in SUSY models, however, can be restricted by existing upper limits on electric dipole moments [3]. The general consensus is that (at least) one of the following three conditions has to be realized: (i) The phases are severely suppressed [7]. (ii) Supersymmetric particles of the first two generations are rather heavy, with masses of the order of a few TeV at least [8]. (iii) There is a rather strong correlation between phases [9], leading to a cancellation of the different SUSY contribution to EDMs *.

Far less work, however, has been devoted up to now to CP violation in supersymmetry with broken R-parity ($\not{R}_p$) [10]. R-parity breaking implies a violation of either lepton or baryon number. It is phenomenologically unacceptable that both types of terms are present [11]. We will discuss only the lepton number violating terms in the following, because we focus on connections to neutrino physics.

R-parity can be broken by bilinear and by trilinear terms. If both types were present up to $(36+3)\times 2$ new parameters appear in the theory, all of which could be complex. *A priori* one could therefore expect that CP violating phenomena differ considerably from both the SM as well as the MSSM.

Electric dipole moments do not constrain in a significant way the phases of either trilinear or bilinear terms individually, essentially since the leading contributions to the EDMs occur only at 2-loop level, if either only trilinear or only bilinear terms are present [12]. † If both types of R-parity breaking terms were present contributions to EDMs appear at 1-loop level [13] and thus the imaginary part of a certain product of bilinear and trilinear terms is more tightly constrained. However, even in the latter case the limit is not especially strong, considering the typical size of R-parity violating parameters expected from neutrino physics. It is thus fair to say that very little is known currently about the phases of R-parity breaking parameters.

In this paper we study CP violation in SUSY with bilinear R-parity breaking, focusing mainly on aspects of those phases which are not present in the MSSM. The assumption of having only bilinear R-parity violation reduces the number of new phases to only six, as will be discussed in more detail later in the paper.

Studies of bilinear R-parity breaking SUSY at this moment are mainly motivated by the recent discoveries in neutrino physics. Observations of atmospheric neutrinos by the Super-K collaboration [14] have confirmed the deficit of atmospheric muon neutrinos, especially at small zenith angles, and thus strongly point to non-zero neutrino masses and mixings. The preferred range of oscillation parameters from atmospheric neutrino data is currently (at 3σ and 1 d.o.f.) [14, 15]

$$0.3 \leq \sin^2\theta_{Atm} \leq 0.7\,, \quad 1.2\times 10^{-3}\ \mathrm{eV}^2 \leq \Delta m^2_{Atm} \leq 4.8\times 10^{-3}\ \mathrm{eV}^2\,. \tag{1.1}$$

Also the long-standing solar neutrino problem now provides strong evidence for neutrino flavour conversion, especially considering the recent measurement of the neutral current rate for solar neutrinos by the SNO collaboration [16]. If interpreted in terms of neutrino

*The correlation "solution" is somewhat debated, see for example [7].

†1-loop contributions are proportional to the neutrino mass, see the first paper in [12].

oscillations, the data indicate a large mixing angle between ν_e and $\nu_\mu - \nu_\tau$, with a strong preference towards the large mixing angle MSW solution (LMA). At 3σ one has [15, 17]

$$0.25 \leq \tan^2\theta_\odot \leq 0.83 \tag{1.2}$$

for 1 d.o.f., the best-fit-parameters being

$$\tan^2\theta_\odot = 0.44 \quad \Delta m^2_\odot = 6.6 \times 10^{-5}\ \mathrm{eV}^2 \tag{1.3}$$

This nicely confirms earlier hints found in Ref. [18].

Calculated at the 1-loop level [19, 20], SUSY with bilinear R-parity breaking can explain the solar and atmospheric neutrino data [20]. However, [19, 20] considered only real parameters and thus were not concerned about possible CP-violating effects.

We will essentially follow [20], extending the calculation to the complex case. In the next section we will give mass matrices and the Higgs potential of the model. In Section 3 we will discuss the constraints on the various phases of the model implied by current neutrino data. Section 4 will discuss possible CP-violating observables, before we close with a short conclusion.

2. The Model

2.1 Superpotential and soft SUSY breaking

The supersymmetric Lagrangian is specified by the superpotential W

$$W = \varepsilon_{ab}\left[h_U^{ij}\widehat{Q}_i^a\widehat{U}_j\widehat{H}_u^b + h_D^{ij}\widehat{Q}_i^b\widehat{D}_j\widehat{H}_d^a + h_E^{ij}\widehat{L}_i^b\widehat{R}_j\widehat{H}_d^a - \mu\widehat{H}_d^a\widehat{H}_u^b + \epsilon_i\widehat{L}_i^a\widehat{H}_u^b\right], \tag{2.1}$$

where $i,j = 1,2,3$ are generation indices, $a,b = 1,2$ are $SU(2)$ indices, and ε is the completely antisymmetric 2×2 matrix, with $\varepsilon_{12} = 1$. The symbol "hat" over each letter indicates a superfield, with $\widehat{Q}_i$, $\widehat{L}_i$, $\widehat{H}_d$, and $\widehat{H}_u$ being $SU(2)$ doublets with hypercharges $\frac{1}{3}$, -1, -1, and 1, respectively, and $\widehat{U}$, $\widehat{D}$, and $\widehat{R}$ being $SU(2)$ singlets with hypercharges $-\frac{4}{3}$, $\frac{2}{3}$, and 2, respectively. The couplings h_U, h_D and h_E are 3×3 Yukawa matrices, and μ and ϵ_i are parameters with units of mass.

Supersymmetry breaking is parameterized with a set of soft supersymmetry breaking terms,

$$\begin{aligned} V_{soft} =\ & M_Q^{ij2}\widetilde{Q}_i^{a*}\widetilde{Q}_j^a + M_U^{ij2}\widetilde{U}_i\widetilde{U}_j^* + M_D^{ij2}\widetilde{D}_i\widetilde{D}_j^* + M_L^{ij2}\widetilde{L}_i^{a*}\widetilde{L}_j^a + M_R^{ij2}\widetilde{R}_i\widetilde{R}_j^* \\ & + m_{H_d}^2 H_d^{a*}H_d^a + m_{H_u}^2 H_u^{a*}H_u^a - \left[\tfrac{1}{2}M_3\lambda_3\lambda_3 + \tfrac{1}{2}M_2\lambda_2\lambda_2 + \tfrac{1}{2}M_1\lambda_1\lambda_1 + \text{h.c.}\right] \\ & + \varepsilon_{ab}\left[A_U^{ij}\widetilde{Q}_i^a\widetilde{U}_jH_u^b + A_D^{ij}\widetilde{Q}_i^b\widetilde{D}_jH_d^a + A_E^{ij}\widetilde{L}_i^b\widetilde{R}_jH_d^a - B\mu H_d^aH_u^b + B_i\epsilon_i\widetilde{L}_i^aH_u^b + \text{h.c.}\right] \end{aligned} \tag{2.2}$$

In the following we assume that there is no intergenerational mixing in the soft terms. Let us first list the parameters which may be complex in the model defined by Eq. (2.1) and Eq. (2.2). Decomposed in modulus and phase these are given by $\epsilon_i = |\epsilon_i|e^{i\varphi_{\epsilon_i}}$ and $\mu =$

$|\mu|e^{i\varphi_\mu}$ in Eq. (2.1), and $A_U = \mathrm{diag}\{|A_u|e^{i\varphi_{A_u}}, |A_c|e^{i\varphi_{A_c}}, |A_t|e^{i\varphi_{A_t}}\}$, $A_D = \mathrm{diag}\{|A_d|e^{i\varphi_{A_d}}, |A_s|e^{i\varphi_{A_s}}, |A_b|e^{i\varphi_{A_b}}\}$, $A_E = \mathrm{diag}\{|A_e|e^{i\varphi_{A_e}}, |A_\mu|e^{i\varphi_{A_\mu}}, |A_\tau|e^{i\varphi_{A_\tau}}\}$, $M_1 = |M_1|e^{i\phi_1}$, $M_2 = |M_2|e^{i\phi_2}$, $M_3 = |M_3|e^{i\phi_3}$, $B\mu = |B\mu|e^{i\varphi_B}$ and $B_i\epsilon_i = |B_i\epsilon_i|e^{i\varphi_{B_i}}$ in Eq. (2.2). This means, in addition to the MSSM parameters the following six parameters are in general complex.

$$\epsilon_i = |\epsilon_i|e^{i\varphi_{\epsilon_i}}, \qquad B_i\epsilon_i = |B_i\epsilon_i|e^{i\varphi_{B_i}}. \tag{2.3}$$

As mentioned no restrictions on the size of the phases in Eq. (2.3) exist; they can be $O(1)$.

At this point it is appropriate to note that not all of the phases quoted above have a physical meaning. Any two of these CP–odd phases may be eliminated by employing the Peccei–Quinn- and the R symmetry $U(1)_{PQ}$ and $U(1)_R$, respectively. We will use this phase freedom later on to remove two unphysical phases.

2.2 Scalar potential

Next we consider the scalar potential and derive the tadpole equations. The scalar potential at tree level is

$$V = \sum_i \left|\frac{\partial W}{\partial z_i}\right|^2 + V_D + V_{soft}. \tag{2.4}$$

where z_i is any one of the scalar fields in the superpotential in Eq. (2.1), V_D are the D-terms, and V_{soft} is given in Eq. (2.2). The vacuum expectation values (vevs) of the Higgs (H_u, H_d) and lepton $(\tilde{L}_i)$ fields are complex in general. In order to determine the ground state of the potential Eq. (2.4), we use the linear parametrization

$$H_d = e^{i\theta}\begin{pmatrix} \frac{1}{\sqrt{2}}[\sigma_d^0 + v_d + i\varphi_d^0] \\ H_d^- \end{pmatrix}, \qquad H_u = \begin{pmatrix} H_u^+ \\ \frac{1}{\sqrt{2}}[\sigma_u^0 + v_u + i\varphi_u^0] \end{pmatrix},$$
$$\widetilde{L}_i = e^{i\eta_i}\begin{pmatrix} \frac{1}{\sqrt{2}}[\tilde{\nu}_i^R + v_i + i\tilde{\nu}_i^I] \\ \tilde{\ell}_i^- \end{pmatrix}, \tag{2.5}$$

where v_d, v_u and v_i are positive, and we have set the phase of H_u to zero since only relative phases are meaningful. The stationary conditions then read

$$\begin{aligned}
\frac{\partial\langle V\rangle}{\partial v_d} &= \left(m_{H_d}^2 + |\mu|^2\right)v_d + v_d D - |B\mu|v_u\cos(\varphi_B+\theta) - v_i|\mu||\epsilon_i|\cos(\varphi_\mu+\theta-\varphi_{\epsilon_i}-\eta_i) = 0\\
\frac{\partial\langle V\rangle}{\partial v_u} &= \left(m_{H_u}^2 + |\mu|^2 + |\epsilon_i|^2\right)v_u - v_u D + v_i|B_i\epsilon_i|\cos(\varphi_{B_i}+\eta_i) - |B\mu|v_d\cos(\varphi_B+\theta) = 0\\
\frac{\partial\langle V\rangle}{\partial v_1} &= v_1 D + |\mu||\epsilon_1|v_d\cos(\varphi_\mu+\theta-\varphi_{\epsilon_1}-\eta_1) + v_1\left(|\epsilon_1|^2 + M_{L_1}^2\right) + |B_1\epsilon_1|v_u\cos(\varphi_{B_1}+\eta_1)\\
&\quad + v_2\epsilon_2\epsilon_1\cos(\varphi_{\epsilon_1}+\eta_1-\varphi_{\epsilon_2}-\eta_2) + v_3\epsilon_3\epsilon_1\cos(\varphi_{\epsilon_1}+\eta_1-\varphi_{\epsilon_3}-\eta_3) = 0\\
\frac{\partial\langle V\rangle}{\partial v_2} &= v_2 D + |\mu||\epsilon_2|v_d\cos(\varphi_\mu+\theta-\varphi_{\epsilon_2}-\eta_2) + v_2\left(|\epsilon_2|^2 + M_{L_2}^2\right) + |B_2\epsilon_2|v_u\cos(\varphi_{B_2}+\eta_2)\\
&\quad + v_1\epsilon_1\epsilon_2\cos(\varphi_{\epsilon_2}+\eta_2-\varphi_{\epsilon_1}-\eta_1) + v_3\epsilon_3\epsilon_2\cos(\varphi_{\epsilon_2}+\eta_2-\varphi_{\epsilon_3}-\eta_3) = 0\\
\frac{\partial\langle V\rangle}{\partial v_3} &= v_3 D + |\mu||\epsilon_3|v_d\cos(\varphi_\mu+\theta-\varphi_{\epsilon_3}-\eta_3) + v_3\left(|\epsilon_3|^2 + M_{L_3}^2\right) + |B_3\epsilon_3|v_u\cos(\varphi_{B_3}+\eta_3)\\
&\quad + v_1\epsilon_1\epsilon_3\cos(\varphi_{\epsilon_3}+\eta_3-\varphi_{\epsilon_1}-\eta_1) + v_2\epsilon_2\epsilon_3\cos(\varphi_{\epsilon_3}+\eta_3-\varphi_{\epsilon_2}-\eta_2) = 0
\end{aligned} \tag{2.6}$$

where $D = \frac{1}{8}(g^2 + g'^2)(v_1^2 + v_2^2 + v_3^2 + v_d^2 - v_u^2)$, and

$$\begin{aligned}
\frac{\partial\langle V\rangle}{\partial\theta} &= v_u v_d |B\mu| \sin(\varphi_B + \theta) + v_i v_d |\epsilon_i||\mu| \sin(\varphi_\mu + \theta - \varphi_{\epsilon_i} - \eta_i) = 0 \\
\frac{\partial\langle V\rangle}{\partial\eta_1} &= v_1 v_d |\mu||\epsilon_1| \sin(\varphi_\mu + \theta - \varphi_{\epsilon_1} - \eta_1) + v_1 v_u |B_1\epsilon_1| \sin(\varphi_{B_1} + \eta_1) \\
&\quad + v_1 v_2 \epsilon_1 \epsilon_2 \sin(\varphi_{\epsilon_1} + \eta_1 - \varphi_{\epsilon_2} - \eta_2) + v_1 v_3 \epsilon_1 \epsilon_3 \sin(\varphi_{\epsilon_1} + \eta_1 - \varphi_{\epsilon_3} - \eta_3) = 0 \\
\frac{\partial\langle V\rangle}{\partial\eta_2} &= v_2 v_d |\mu||\epsilon_2| \sin(\varphi_\mu + \theta - \varphi_{\epsilon_2} - \eta_2) + v_2 v_u |B_2\epsilon_2| \sin(\varphi_{B_2} + \eta_2) \\
&\quad - v_2 v_1 \epsilon_2 \epsilon_1 \sin(\varphi_{\epsilon_1} + \eta_1 - \varphi_{\epsilon_2} - \eta_2) + v_2 v_3 \epsilon_2 \epsilon_3 \sin(\varphi_{\epsilon_2} + \eta_2 - \varphi_{\epsilon_3} - \eta_3) = 0 \\
\frac{\partial\langle V\rangle}{\partial\eta_3} &= v_3 v_d |\mu||\epsilon_3| \sin(\varphi_\mu + \theta - \varphi_{\epsilon_3} - \eta_3) + v_3 v_u |B_3\epsilon_3| \sin(\varphi_{B_3} + \eta_3) \\
&\quad - v_3 v_1 \epsilon_3 \epsilon_1 \sin(\varphi_{\epsilon_1} + \eta_1 - \varphi_{\epsilon_3} - \eta_3) - v_3 v_2 \epsilon_2 \epsilon_3 \sin(\varphi_{\epsilon_2} + \eta_2 - \varphi_{\epsilon_3} - \eta_3) = 0
\end{aligned} \tag{2.7}$$

where the repeated index i implies a summation over $i = 1, 2, 3$.

It was shown in Ref. [21] that spontaneous CP violation does not occur in this model with $v_i \neq 0$, see also [22]. In the case of explicit CP violation, where $\{\varphi_\mu, \varphi_B, \varphi_{\epsilon_i}, \varphi_{B_i}\} \neq 0, \pi$, the vevs are complex in general, as can be seen from Eq. (2.7).

2.3 Scalar mass matrices

With the solutions to Eq. (2.6) and (2.7) we can give the mass matrices of the matter fields. The mass matrix for the neutral boson fields is a 10×10 symmetric matrix. The corresponding Lagrangian is

$$\mathcal{L}^S_{\text{mass}} = -\tfrac{1}{2} S^T \boldsymbol{M}^2_S S, \tag{2.8}$$

where we have defined $S^T \equiv [\sigma^0_d, \sigma^0_u, \tilde{\nu}^R_i, \varphi^0_d, \varphi^0_u, \tilde{\nu}^I_i]$. The mass matrix in Eq. (2.8) can be decomposed into three 5×5 blocks:

$$\boldsymbol{M}^2_S = \begin{bmatrix} \boldsymbol{M}^2_{SS} & \boldsymbol{M}^2_{SP} \\ \boldsymbol{M}^T_{SP} & \boldsymbol{M}^2_{PP} \end{bmatrix} \tag{2.9}$$

This mass matrix is diagonalized through

$$R \boldsymbol{M}^2_S R^T = \text{diag}(0, m^2_{S_1}, ..., m^2_{S_9}), \tag{2.10}$$

where R is a orthogonal matrix. We give here only the scalar–pseudoscalar mixing block $\boldsymbol{M}^2_{SP}$. The scalar–scalar block $\boldsymbol{M}^2_{SS}$ and the pseudoscalar–pseudoscalar block $\boldsymbol{M}^2_{PP}$ can be found in Ref. [20] for real parameters. They are valid also for the complex case after obvious substitutions for complex parameters.

$$\boldsymbol{M}_{SP}^2 = \begin{bmatrix} 0 & |B\mu| s_\psi & -|\mu||\epsilon_1| s_{\rho_1} & -|\mu||\epsilon_2| s_{\rho_2} & -|\mu||\epsilon_3| s_{\rho_3} \\ |B\mu| s_\psi & 0 & -|B_1\epsilon_1| s_{\sigma_1} & -|B_2\epsilon_2| s_{\sigma_2} & -|B_3\epsilon_3| s_{\sigma_3} \\ |\mu||\epsilon_1| s_{\rho_1} & -|B_1\epsilon_1| s_{\sigma_1} & 0 & |\epsilon_1||\epsilon_2| s_{\lambda_{12}} & |\epsilon_1||\epsilon_3| s_{\lambda_{13}} \\ |\mu||\epsilon_2| s_{\rho_2} & -|B_2\epsilon_2| s_{\sigma_2} & -|\epsilon_1||\epsilon_2| s_{\lambda_{12}} & 0 & |\epsilon_2||\epsilon_3| s_{\lambda_{23}} \\ |\mu||\epsilon_3| s_{\rho_3} & -|B_3\epsilon_3| s_{\sigma_3} & -|\epsilon_1||\epsilon_3| s_{\lambda_{13}} & -|\epsilon_2||\epsilon_3| s_{\lambda_{23}} & 0 \end{bmatrix}, \tag{2.11}$$

where we have used the shorthand notation: $s_\psi \equiv \sin(\theta + \varphi_B), s_{\rho_i} \equiv \sin(\varphi_\mu + \theta - \varphi_{\epsilon_i} - \eta_i), s_{\sigma_i} \equiv \sin(\eta_i + \varphi_{B_i})$ and $s_{\lambda_{ij}} \equiv \sin(\varphi_{\epsilon_i} + \eta_i - \varphi_{\epsilon_j} - \eta_j)$.

Eq. (2.7) together with Eq. (2.11) shows that scalar–pseudoscalar Higgs transitions appear already at tree–level, contrary to the MSSM where this is possible only at one loop level. However, considering the typical size of the $\not{R}_p$ violating parameters, $|\epsilon_i/\mu| \lesssim O(10^{-3})$ [20], such a transition has to be expected to be rather small.

2.4 Neutralino–neutrino mass matrix

Here we focus on the 7×7 Neutralino–neutrino mass matrix. In the basis $\Psi_0'^T = (-i\lambda_1, -i\lambda_2, \tilde{H}_d^1, \tilde{H}_u^2, \nu_e, \nu_\mu, \nu_\tau)$, it reads

$$\mathcal{M}^0 = \begin{pmatrix} \mathcal{M}_{\chi^0} & m^T \\ m & 0 \end{pmatrix}. \tag{2.12}$$

Here, the 3×4 sub-matrix m contains entries from the bilinear $\not{R}_p$ parameters,

$$m = \begin{pmatrix} -\frac{1}{2} g' v_1 e^{-i\eta_1} & \frac{1}{2} g v_1 e^{-i\eta_1} & 0 & |\epsilon_1| e^{i\varphi_{\epsilon_1}} \\ -\frac{1}{2} g' v_2 e^{-i\eta_2} & \frac{1}{2} g v_2 e^{-i\eta_2} & 0 & |\epsilon_2| e^{i\varphi_{\epsilon_2}} \\ -\frac{1}{2} g' v_3 e^{-i\eta_3} & \frac{1}{2} g v_3 e^{-i\eta_3} & 0 & |\epsilon_3| e^{i\varphi_{\epsilon_3}} \end{pmatrix}, \tag{2.13}$$

and the MSSM neutralino mass matrix is given by,

$$\mathcal{M}_{\chi^0} = \begin{pmatrix} |M_1| e^{i\phi_1} & 0 & -\frac{1}{2} g' v_d e^{-i\theta} & \frac{1}{2} g' v_u \\ 0 & |M_2| e^{i\phi_2} & \frac{1}{2} g v_d e^{-i\theta} & -\frac{1}{2} g v_u \\ -\frac{1}{2} g' v_d e^{-i\theta} & \frac{1}{2} g v_d e^{-i\theta} & 0 & -|\mu| e^{i\varphi_\mu} \\ \frac{1}{2} g' v_u & -\frac{1}{2} g v_u & -|\mu| e^{i\varphi_\mu} & 0 \end{pmatrix}. \tag{2.14}$$

The mass matrix $\mathcal{M}^0$ is diagonalized by

$$\mathcal{N}^{0*} \mathcal{M}^0 \mathcal{N}^{0^{-1}} = \mathrm{diag}(m_{\nu_i}, m_{\chi_j^0}). \tag{2.15}$$

where $(i = 1, ..., 3)$ for the neutrinos, and $(j = 1, ..., 4)$ for the neutralinos. The method of perturbative diagonalization of $\mathcal{M}^0$ presented in [23] holds also in the general case of a complex symmetric mass matrix $\mathcal{M}^0$.

Only one neutrino acquires mass at tree–level and its mass is approximately given by

$$m_{\nu_3} \simeq \left| \frac{M_1 g^2 + M_2 g'^2}{4 \det(\mathcal{M}_{\chi^0})} \right| |\vec{\Lambda}|^2 \tag{2.16}$$

where,

$$\Lambda_i = |\epsilon_i| v_d \, e^{i(\varphi_{\epsilon_i} - \theta)} + v_i |\mu| \, e^{i(\varphi_\mu - \eta_i)}. \tag{2.17}$$

As a result, for a realistic description of the neutrino spectrum one has to improve the calculation to 1-loop order.

A complete list of 1-loop contributions can be found, for example, in [20]. For our purposes it is sufficient to consider only the $\tilde{b} - b$ loop, which gives in a wide range of parameter space the most important contribution to the neutrino masses. The 1-loop corrections to the neutrino sector can be written as

$$(\Delta m_\nu)^{\tilde{b}b}_{ij} \simeq -\frac{3}{16\pi^2} m_b h_b^2 \sin 2\theta_{\tilde{b}} \, e^{i\varphi_{\tilde{b}}} \mathcal{N}^{0\,*}_{i3} \, \mathcal{N}^{0\,*}_{j3} \log\left(\frac{m^2_{\tilde{b}_2}}{m^2_{\tilde{b}_1}}\right), \tag{2.18}$$

where $h_b = e^{-i\theta}\sqrt{2} m_b / v_d$, $\theta_{\tilde{b}}$ is the mixing angle and $\varphi_{\tilde{b}}$ is the phase in the sbottom sector, respectively. $m_{\tilde{b}_i}$ are the two sbottom masses. The eigenvalues up to one loop are obtained by diagonalizing $\mathrm{diag}(m_{\nu_i}, m_{\chi^0_j}) + (\Delta m_\nu)^{\tilde{b}b}_{ij}$. The corresponding mixing matrix is $\mathcal{N}^1$, and the complete mixing matrix which relates weak basis and eigenbasis, is

$$\mathcal{N} = \mathcal{N}^1 \cdot \mathcal{N}^0. \tag{2.19}$$

The CP violating Dirac phase δ entering the oscillation formula can be extracted by using the following relation

$$|\delta| = \sin^{-1}\left(\left| \frac{8\, \Im m(\mathcal{N}_{25} \mathcal{N}^*_{26} \mathcal{N}_{16} \mathcal{N}^*_{15})}{\cos\theta_{13} \sin 2\theta_{13} \sin 2\theta_{12} \sin 2\theta_{23}} \right| \right), \tag{2.20}$$

where $\theta_{13} = \sin^{-1}(|\mathcal{N}_{35}|), \theta_{12} = \tan^{-1}(|\mathcal{N}_{25}/\mathcal{N}_{15}|)$ and $\theta_{23} = \tan^{-1}(|\mathcal{N}_{36}/\mathcal{N}_{37}|)$. Note that the quantity in the denominator of Eq. (2.20) is one of six equivalent representations of the Jarlskog invariant in the neutrino sector.

3. Neutrino data and CP violating phases

3.1 Analytical discussion

Data from neutrino oscillation experiments require that the bilinear R-parity breaking parameters must obey certain conditions [20]. As briefly discussed in the introduction, atmospheric neutrino data requires Δm^2_{Atm} to be of the order of $\Delta m^2_{Atm} \simeq \mathcal{O}(10^{-3})$ eV^2, while the solar neutrino problem can be solved for $\Delta m^2_\odot \sim$ (few) 10^{-5} eV^2 in the case of the LMA solution. ‡

‡The LOW solution $\Delta m^2_\odot \sim 10^{-7}$ eV^2 and the quasi vacuum solution $\Delta m^2_\odot \sim 10^{-9}$ eV^2 can not be excluded at present, but are strongly disfavoured after the SNO neutral current measurement, see for example [24].

In bilinear R-parity breaking the neutrino spectrum is hierarchical and therefore m_{ν_3} and m_{ν_2} are approximately given by $m_{\nu_3} \simeq \sqrt{\Delta m^2_{Atm}} \simeq 0.05$ eV and $m_{\nu_2} \simeq \sqrt{\Delta m^2_{\odot}}$. Since, as explained above, at tree level the model has only one non-zero mass eigenstate, it is straightforward to see that the ratio of $m_{\nu_2}/m_{\nu_3} \simeq \sqrt{\Delta m^2_{\odot}}/\sqrt{\Delta m^2_{Atm}}$ determines the relative importance of the one-loop mass with respect to the tree-level mass in our model.

The hierarchy of $m_{\nu_2}/m_{\nu_3} \simeq 0.1$ (in case of the LMA solution) implies that the terms in Eq. (2.17) must almost cancel. Thus, three pairs of phases have to obey the condition

$$\varphi_{\epsilon_i} + \eta_i - \theta - \varphi_\mu \pm \pi < O(10^{-2}) \ . \tag{3.1}$$

This observation reduces the effective number of free phases in the bilinear model from six to three, as we will now demonstrate.

The phases η_i have to obey the tadpole equations Eq. (2.7) for an arbitrary set of input parameters. In the limit $\varphi_{\epsilon_i} + \eta_i - \theta - \varphi_\mu \pm \pi = 0$, the tadpole equations Eq. (2.7) reduce to

$$\begin{aligned} v_u v_d |B\mu| \sin(\varphi_B + \theta) &= 0 \ , \\ v_i v_u |B_i \epsilon_i| \sin(\varphi_{B_i} + \eta_i) &= 0 \ , \qquad i = 1, 2, 3 \end{aligned} \tag{3.2}$$

where the equation for θ is just the well-known MSSM relation. To find the correct minimum of the potential, we get from Eq. (3.2) $\theta = -\varphi_B$ and $\eta_i = \pm\pi - \varphi_{B_i}$. Eq. (3.1) now reads $\varphi_{\epsilon_i} = \varphi_{B_i} - \varphi_B + \varphi_\mu$ modulo 2π, which means that the number of independent phases is reduced to three. This leads immediately to the important result that the scalar–pseudoscalar mixing vanishes. This can be read of directly from Eq. (2.11) since every single term is zero in this limit.

Atmospheric neutrino measurements provide an additional constraint on the bilinear model, due to the large angle in $\nu_\mu \leftrightarrow \nu_\tau$ oscillations, which has to be $\tan^2\theta_{Atm} \simeq \frac{|\Lambda_2|^2}{|\Lambda_3|^2} \simeq 1$. For non-zero phases this condition could only be met if

$$\varphi_{\epsilon_2} + \eta_2 \simeq \pm(\varphi_{\epsilon_3} + \eta_3) \ . \tag{3.3}$$

We will investigate below to which extent Eq. (3.3) has to be satisfied numerically.

3.2 Numerical discussion

Before we describe our numerical analysis we specify the parameters. We remove two unphysical phases by applying the Peccei–Quinn symmetry $U(1)_{PQ}$ and the R symmetry $U(1)_R$. We choose $\varphi_B = 0$ and $\varphi_\mu = \pi$, and set all other phases of the MSSM part to zero. As input phases we take φ_{ϵ_i} and η_i randomly in the range $[-\pi, \pi]$, and use the stationary conditions Eq. (2.7) to solve for the phases φ_{B_i}. In order to reduce the number of parameters, the numerical calculations were performed in a constrained version of the MSSM. We have scanned the parameters in the following ranges: $M_2 \in [0, 1.2]$ TeV,

$|\mu| \in [0, 2.5]$ TeV, $m_0 \in [0, 1.2]$ TeV, A_0/m_0 and $B_0/m_0 \in [-3, 3]$ and $\tan\beta \in [2.5, 10]$. All randomly generated points were subsequently tested for consistency with the stationary conditions Eq. (2.6) and (2.7).

Next we consider to which extent the correlations (3.1) have to be obeyed. With our choice of the phases these correlations read: $\varphi_{\epsilon_i} + \eta_i = 0$ modulo 2π. We study the allowed region of the two sums of phases $\varphi_{\epsilon_2} + \eta_2$ and $\varphi_{\epsilon_3} + \eta_3$. From Fig. 1 we can see $|\varphi_{\epsilon_{2,3}} + \eta_{2,3}| \lesssim 5 \times 10^{-3}$ for most of the sums $\varphi_{\epsilon_{2,3}} + \eta_{2,3}$. If $\varphi_{\epsilon_2} + \eta_2 \simeq \pm(\varphi_{\epsilon_3} + \eta_3)$, then $|\varphi_{\epsilon_{2,3}} + \eta_{2,3}| \lesssim 10^{-2}$ is possible. This feature can be traced to the constraint $\tan^2\theta_{Atm} \simeq 1$, see Eq. (3.3). Fig. 2 confirms the additional phase correlation given in Eq. (3.3), where we have plotted $\tan^2\theta_{Atm}$ versus $(\varphi_{\epsilon_2}+\eta_2)/(\varphi_{\epsilon_3}+\eta_3)$ for points in which $\varphi_{\epsilon_{2,3}}+\eta_{2,3} \geq 3\times 10^{-3}$.

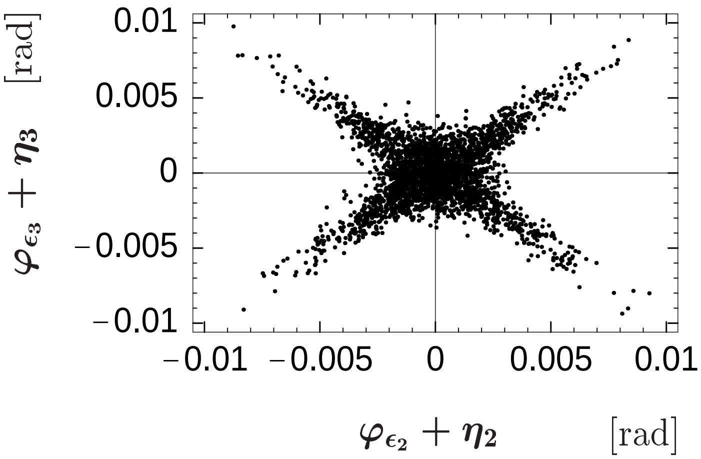

Figure 1: Phase correlation between $\varphi_{\epsilon_2} + \eta_2$ and $\varphi_{\epsilon_3} + \eta_3$ required by the constraints from atmospheric neutrino data.

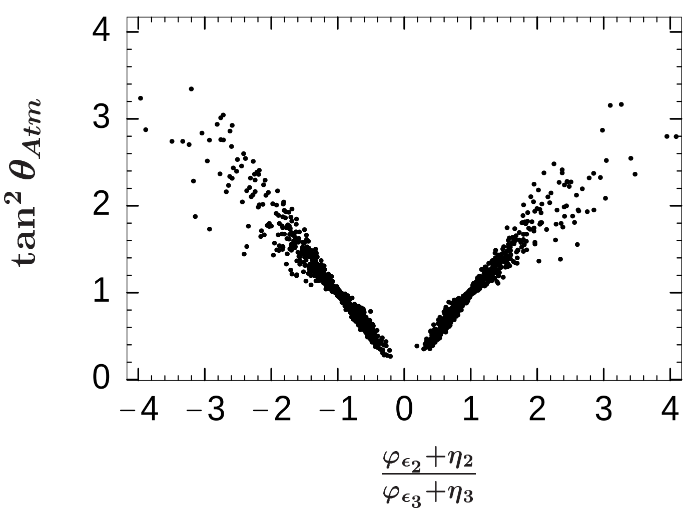

Figure 2: Correlation between $(\varphi_{\epsilon_2} + \eta_2)/(\varphi_{\epsilon_3} + \eta_3)$ and $\tan^2\theta_{Atm}$.

As has been pointed out previously, $\boldsymbol{M}_{SP}$=0 in the limit $\varphi_{\epsilon_i} + \eta_i = 0$. Even if this condition is not exactly fulfilled scalar–pseudoscalar Higgs mixing can be neglected in general provided that $|\varphi_{\epsilon_i} + \eta_i| \lesssim 10^{-2}$. In the sneutrino sector the situation may be sligthly different. Recall that in the MSSM the sneutrino is a complex scalar field. $\not{R}_p$ splits $\Re e(\tilde\nu)$ and $\Im m(\tilde\nu)$ by a small amount in each generation. Let us denote the associated mass eigenstates as $\tilde\nu^{1,2}$. Since $\tilde\nu^{1,2}$ are nearly degenerate large scalar–pseudoscalar mixing is possible in this sector also for $|\varphi_{\epsilon_i} + \eta_i| \lesssim 10^{-2}$. To demonstrate this feature, let us consider the third generation $\tilde\nu_\tau^{1,2}$. As an example we consider the $\not{R}_p$ coupling $C(\tilde\nu_\tau^1 bb)$, where $\tilde\nu_\tau^1$ does not contain imaginary parts in the limit $\varphi_{\epsilon_i} + \eta_i = 0$. The interaction is described by the Lagrangian

$$\mathcal{L} = \tilde\nu_\tau^1 \bar b \left[C_{\tilde\nu_\tau^1 bb} P_L + C^*_{\tilde\nu_\tau^1 bb} P_R\right] b, \tag{3.4}$$

where $C_{\tilde\nu_\tau^1 bb} \equiv -\frac{h_b}{\sqrt{2}}(R_{\tilde\nu_\tau^1 1} + iR_{\tilde\nu_\tau^1 6})$ and R is defined in Eq. (2.10). CP invariance means that $R_{\tilde\nu_\tau^1 6} = 0$. The extent to which this limit is saturated can be seen in Fig. 3 where we plot $(R^2_{\tilde\nu_\tau^1 1} - R^2_{\tilde\nu_\tau^1 6})/(R^2_{\tilde\nu_\tau^1 1} + R^2_{\tilde\nu_\tau^1 6})$ as a function of $\varphi_{\epsilon_3} + \eta_3$. For the light (dark) points the mass spliting is 10^{-3} eV $\leq |m_{\tilde\nu_\tau^1} - m_{\tilde\nu_\tau^2}| \leq 10^{-2}$ eV (0.1 eV $\leq |m_{\tilde\nu_\tau^1} - m_{\tilde\nu_\tau^2}|$). As can be seen, scalar–pseudoscalar mixing in the sneutrino sector vanishes only for very small $\varphi_{\epsilon_i} + \eta_i$ in this sector.

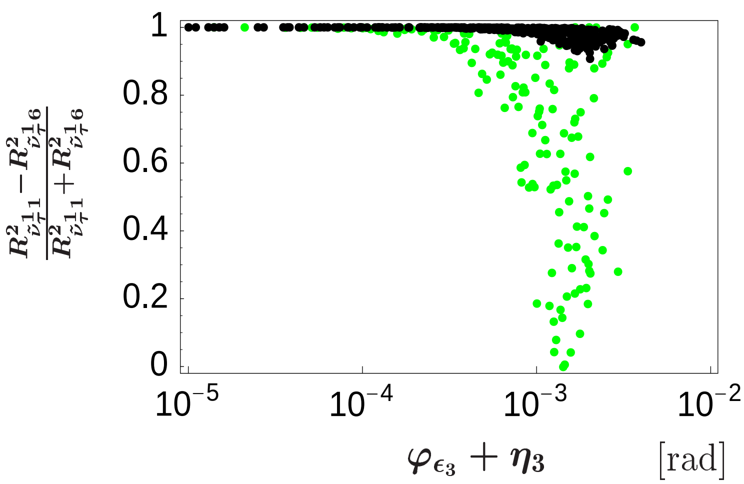

Figure 3: Scalar–pseudoscalar mixing in the nearly degenerate $\tilde\nu^R - \tilde\nu^I$ system. For a discussion see text.

4. Observables

In this section we discuss possible CP–odd observables arising due to the new phases in bilinear R-parity violating couplings. Recall that these $\not{R}_p$ couplings are typically two to three orders of magnitude smaller than the R parity conserving couplings if neutrino data are to be explained by R-parity violation. This already implies that CP violating effects induced by R-parity violating parameters can at most be seen in LSP decays, because in all other cases either the R-parity violating branching ratios are very small or the loop contributions due to the $\not{R}_p$ couplings to R-parity conserving decay modes are tiny.

In the following we will discuss various possibilities for the LSP: sfermions, neutralino and chargino. Here we will focus on observables arising in two–body decays of the LSP, which are either rate asymmetries or helicity asymmetries. Note that the assumption of two–body decays is well motivated by the fact that LEP has not found any supersymmetric particle which gives a lower bound on SUSY particle masses of order 100 GeV.

Let us first discuss a necessary condition whether a CP asymmetry is observable or not before considering the different LSP classes. The relevant quantity to decide whether an asymmetry is observable (at 1σ), is given by $(A_{\rm CP}^2 \times B)^{-1}$, where $A_{\rm CP}$ is the CP asymmetry and B is the branching ratio of the decay considered. If we assume rather optimistically that order 10^6 LSPs can be produced at a future collider experiment, this requires that $(A_{\rm CP}^2 \times B)^{-1} \lesssim O(10^6)$.

We want to recapitulate here, that in order to construct a rate asymmetry one needs final–state interactions, otherwise partial decay rates are equal due to CPT invariance even if CP is violated. The rate asymmetry is then built through an interference of tree–level amplitudes and one–loop amplitudes where a pair of intermediate particles in the loop is on–shell. Subsequently we will discuss various one–loop diagrams giving rise to CP asymmetries.

4.1 Sfermion LSP

Consider first the case were a squark is the LSP because here the discussion is rather simple. The possible final states are either $q' \, \ell_i$ or $q \, \nu_j$. The relevant diagrams contributing to possible CP asymmetries are shown in Fig. 4: a) self energy diagrams and b) vertex diagrams. It is obvious from this figure that in both cases all three couplings involved violate R–parity and, thus, the corresponding diagram gives only a tiny contribution leading to a negligible asymmetry.

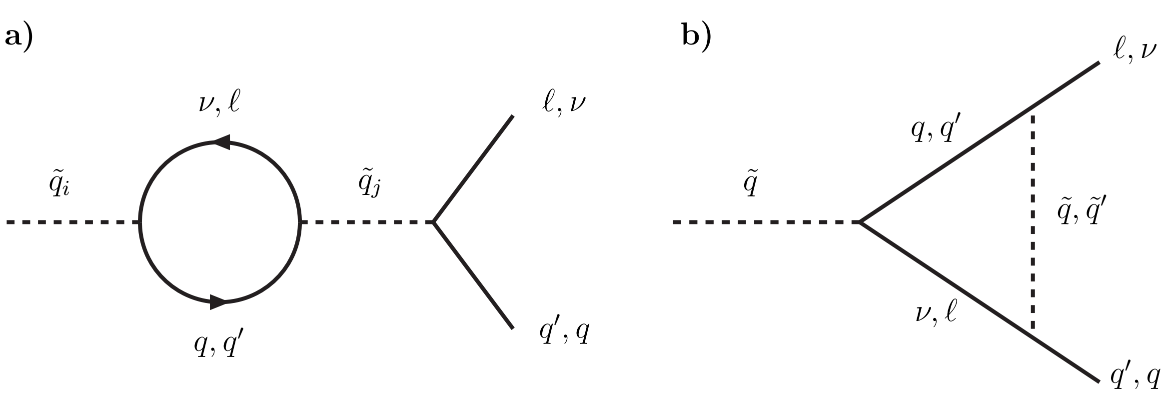

Figure 4: a) Self energy diagrams and b) vertex diagrams contributing to CP asymmetries of squark decays.

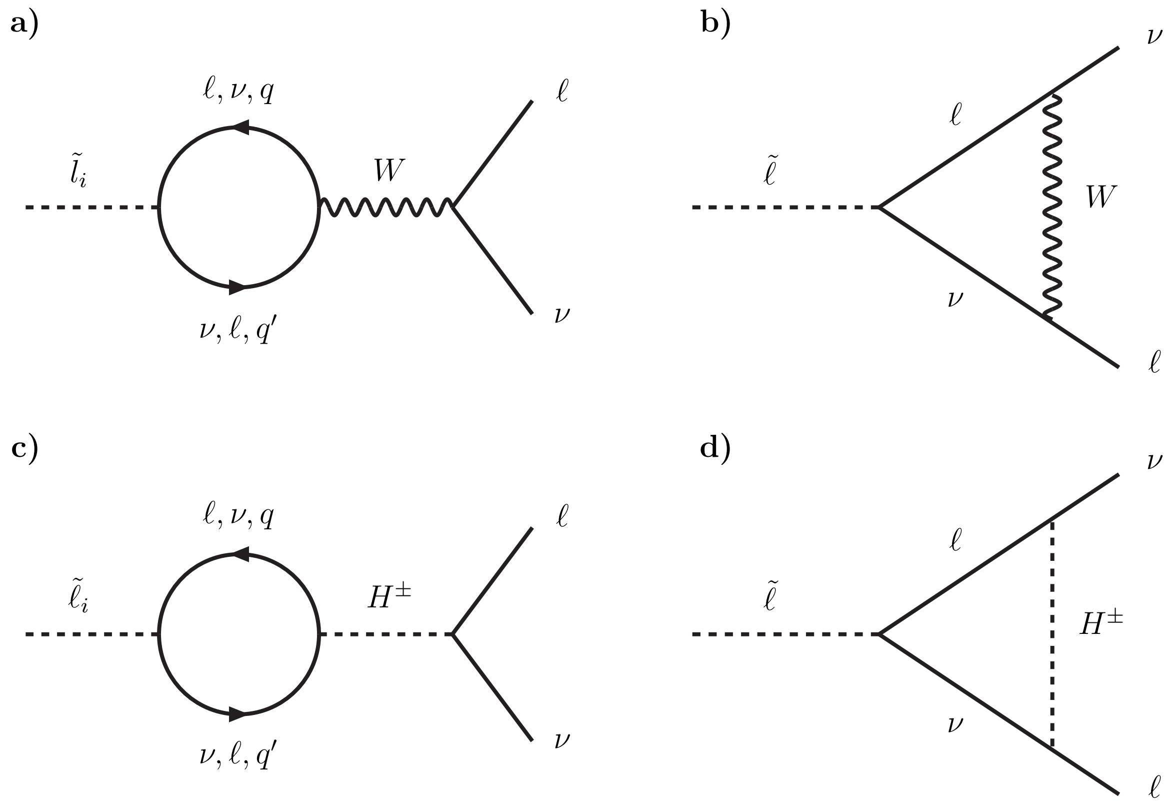

Figure 5: a) and c) Self energy diagrams and b) and d) vertex diagrams contributing to CP asymmetries of slepton decays.

For a charged slepton ($\tilde{\ell}$) LSP the possible final states are qq' and $\ell\nu$. We consider the following rate asymmetry,

$$\mathcal{A}_{\rm CP} = \frac{\Gamma(\tilde{\ell} \to \ell^- \nu) - \Gamma(\bar{\tilde{\ell}} \to \ell^+ \nu)}{\Gamma(\tilde{\ell} \to \ell^- \nu) + \Gamma(\bar{\tilde{\ell}} \to \ell^+ \nu)} \ . \tag{4.1}$$

In this case diagrams similar to those shown in Fig. 4 with appropriate replacements contribute. For the same reasoning as given above for the squark case their contribution is small. In addition to these diagrams one finds diagrams involving a W–boson and the corresponding ones where the W–boson is replaced by a charged Higgs boson as shown in Fig. 5. As can be easily shown the diagrams involving a W–boson are suppressed due to the small fermion masses involved. In case of a self–energy contribution with a top quark and a bottom quark in the loop, Fig. 5a, the CP asymmetry is suppressed by the factor $\frac{m_t m_\ell}{m_W^2}$. For the charged Higgs boson the situation is a little bit more subtle, because in the corresponding diagram of Fig. 5c the contribution with top/bottom quarks can be enhanced for large $\tan\beta$. In this case the CP asymmetry is given by

$$\mathcal{A}_{\rm CP} \sim \frac{3}{16\pi} \frac{\Im m\{h_\ell \sin\beta (g_1^L h_b \sin\beta + g_1^R h_t \cos\beta) g_4^*\}}{|g_4|^2} \ , \tag{4.2}$$

where $g_1^{L,R}$ (g_4) are the $R\!\!\!/_p$ and CP violating left and right (left) coupling of the $\tilde{\ell}$ to tb ($\nu\ell$). The relevant Lagrangian is given by

$$\mathcal{L} = \tilde{\ell}\,\bar{b}(g_1^L P_L + g_1^R P_R)\,t + g_4\tilde{\ell}\,\bar{\ell}P_L\,\nu + \text{h.c.} \tag{4.3}$$

where $P_{L,R} = 1/2(1 \mp \gamma_5)$. The couplings can be found in [25]. In principle large $\tan\beta$ may lead to a large CP asymmetry if $\ell = \tau$. However, for large values of $\tan\beta$ the 1-loop contributions to the neutrino masses tend to be too large. For this reason we have not found any points with $(A_{\text{CP}}^2 \times B)^{-1} \lesssim O(10^6)$, satisfying at the same time the constraints from neutrino physics.

Let us now turn to the case where the sneutrino is the LSP. It can be shown on general grounds that the mass splitting between $\Re e(\tilde{\nu})$ and $\Im m(\tilde{\nu})$ is very small if the Majorana mass of the corresponding neutrino is tiny [26]. Taking the known neutrino data into account this implies that in our case the CP-even/CP-odd mass splitting between $\Re e(\tilde{\nu})$ and $\Im m(\tilde{\nu})$ is typically of order eV and thus negligibly small. In the case of CP violation the mass eigenstates are a superposition of $\tilde{\nu}_\ell^R$ and $\tilde{\nu}_\ell^I$. As a possible CP sensitive observable we consider the following helicity asymmetry,

$$\mathcal{A}_{\text{hel}} = \frac{\sum_i[\Gamma(\tilde{\nu}_\ell^i \to \tau_L^+\tau_L^-) - \Gamma(\tilde{\nu}_\ell^i \to \tau_R^+\tau_R^-)]}{\sum_i[\Gamma(\tilde{\nu}_\ell^i \to \tau_L^+\tau_L^-) + \Gamma(\tilde{\nu}_\ell^i \to \tau_R^+\tau_R^-)]}, \tag{4.4}$$

where $\ell = \tau, \mu, e$, and we sum over the two (nearly) degenerate sneutrino states, because they can not be resolved experimentally.

The same classes of diagrams appear as in the case of the charged slepton decays discussed previously. The dominant diagrams are now the self–energy diagrams with a bottom quark in the loop, see Fig. 6. The contributions where a Z is exchanged in the s–channel (Fig. 6b) is suppressed by a factor $\frac{m_\tau m_b}{m_Z^2}$. The contribution in Fig. 6a can be substantial for large values of $\tan\beta$. However, as was mentioned, this in turn tends to drive the 1-loop contribution to the neutrino masses to be too large. Moreover, there is a cancellation between the contributions of the two degenerate sneutrinos for $\tilde{\nu}_{\mu,e}^i$. This cancellation can be seen from Fig. 7 for $\tilde{\nu}_\mu^i$ where the dark points represent the result for $i = 1$ in Eq. (4.4), whereas for the light points the two contributions are summed up. For $\tilde{\nu}_e$ the result is the same. For $\tilde{\nu}_\tau^i$ we find that $\sum_i B(\tilde{\nu}_\tau^i \to \tau^+\tau^-) \sim O(10^{-5} - 10^{-4})$. With the expected number of events, it is practically impossible to observe such a helicity asymmetry.

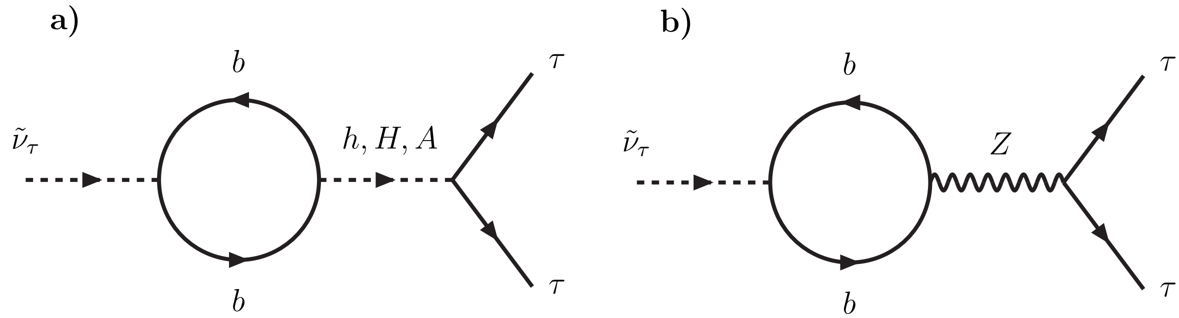

Figure 6: Dominant Feynman diagrams for the absorptive part of the amplitude $\tilde{\nu}^i \to \tau\tau$.

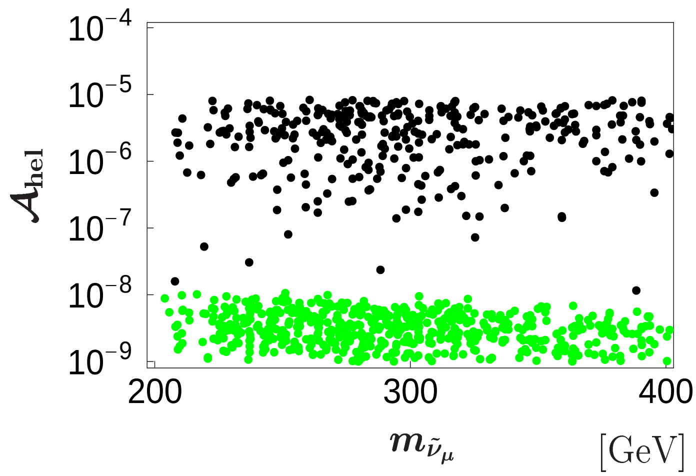

Figure 7: $\mathcal{A}_{\rm hel}$ as a function of $m_{\tilde{\nu}_\mu}$ where the dark (light) points represent the result without (with) sum of the contribution of the two degenerate $\tilde{\nu}^i_\mu$.

4.2 Neutralino LSP

If the $\tilde{\chi}^0_1$ is the LSP the decay modes may be $\tilde{\chi}^0_1 \to \{W\ell, Z\nu, h\nu\}$. As an example we consider rate asymmetries such as

$$\mathcal{A}_{\rm CP} = \frac{\Gamma(\tilde{\chi}^0_1 \to W^+\ell^-) - \Gamma(\tilde{\chi}^0_1 \to W^-\ell^+)}{\Gamma(\tilde{\chi}^0_1 \to W^+\ell^-) + \Gamma(\tilde{\chi}^0_1 \to W^-\ell^+)} \ . \tag{4.5}$$

The relevant part of the Lagrangian reads

$$\mathcal{L} = \bar{\nu} O_R^{\nu\chi h} P_R \ \tilde{\chi}^0_1 h + \bar{\ell}\gamma^\mu (O_R^{\ell\chi \rm w} P_R + O_L^{\ell\chi \rm w} P_L)\tilde{\chi}^0_1 W^-_\mu + \bar{\nu}\gamma^\mu O_L^{\nu\chi \rm z} P_L \ \tilde{\chi}^0_1 Z_\mu + \text{h.c.} \tag{4.6}$$

The relevant couplings in Eq. (4.6) can be found in [25]. Typically these couplings obey: $|O_R^{\ell\chi \rm w}| \sim |O_R^{\nu\chi h}| >> |O_L^{\ell\chi \rm w}|, |O_L^{\nu\chi \rm z}|$. This implies that final state interactions between $W\ell$ and $Z\nu$ will not contribute significantly to the CP asymmetry in Eq. (4.5). As the next possibility we consider a final–state interaction between $W\ell$ and $h\nu$, see Fig. 8. The

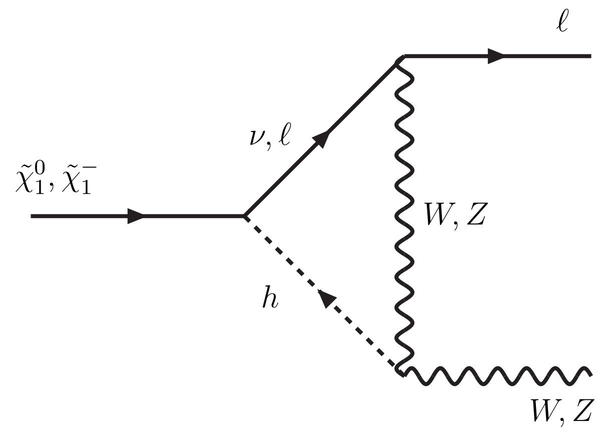

Figure 8: Dominant Feynman diagram for the absorptive part of the amplitude $\tilde{\chi}_1^0 \to \ell\, W$ and $\tilde{\chi}_1^- \to \ell\, Z$.

contribution where the W^+ is replaced by the charged Higgs H^+, is suppressed due to the Yukawa coupling and therefore not relevant. Setting small lepton masses to zero, a straightforward calculation gives an asymmetry proportional to $\Im m(O_R^{\nu\chi h} O_L^{\ell\chi \mathrm{w}*})$ (see Appendix). Since only the left coupling $O_L^{\ell\chi\mathrm{w}}$ (and not the much larger right coupling $O_R^{\ell\chi\mathrm{w}}$) appears in the denominator of Eq. (4.5) the resulting $(A_{\mathrm{CP}}^2 \times B)^{-1}$ is always above 10^6.

4.3 Chargino LSP

If the χ_1^- is the LSP the possible final states are $Z\ell, W\nu$ and $h\,\ell$. We consider the CP asymmetry:

$$\mathcal{A}_{\mathrm{CP}} = \frac{\Gamma(\tilde{\chi}_1^- \to Z\ell^-) - \Gamma(\tilde{\chi}_1^+ \to Z\ell^+)}{\Gamma(\tilde{\chi}_1^- \to Z\ell^-) + \Gamma(\tilde{\chi}_1^+ \to Z\ell^+)} \tag{4.7}$$

The relevant part of the Lagrangian reads

$$\mathcal{L} = \bar{\ell}(O_L^{\ell\chi h} P_L + O_R^{\ell\chi h} P_R)\tilde{\chi}_1^- h + \bar{\nu}\gamma^\mu O_L^{\nu\chi\mathrm{w}} P_L \tilde{\chi}_1^- W_\mu + \bar{\ell}\gamma^\mu(O_L^{\ell\chi\mathrm{z}} P_L + O_R^{\ell\chi\mathrm{z}} P_R)\tilde{\chi}_1^- Z_\mu + \mathrm{h.c.} \tag{4.8}$$

The full form of the left and right couplings can be found in [20, 25]. Scanning the parameters over the parameter ranges as described in the previous section we find that the left and right couplings in Eq. (4.8) typically are: $|O_L^{\nu\chi\mathrm{w}}|, |O_L^{\ell\chi\mathrm{h}}|, |O_R^{\ell\chi\mathrm{z}}| << |O_R^{\ell\chi h}| \sim |O_L^{\ell\chi\mathrm{z}}|$. Therefore, similar to the neutralino we will consider only a final–state interaction between $Z\ell$ and $h\ell$. The dominant Feynman diagram is displayed in Fig. 8, while the calculation is given in the Appendix. The result of a numerical scan is shown Fig. 9a and b where we

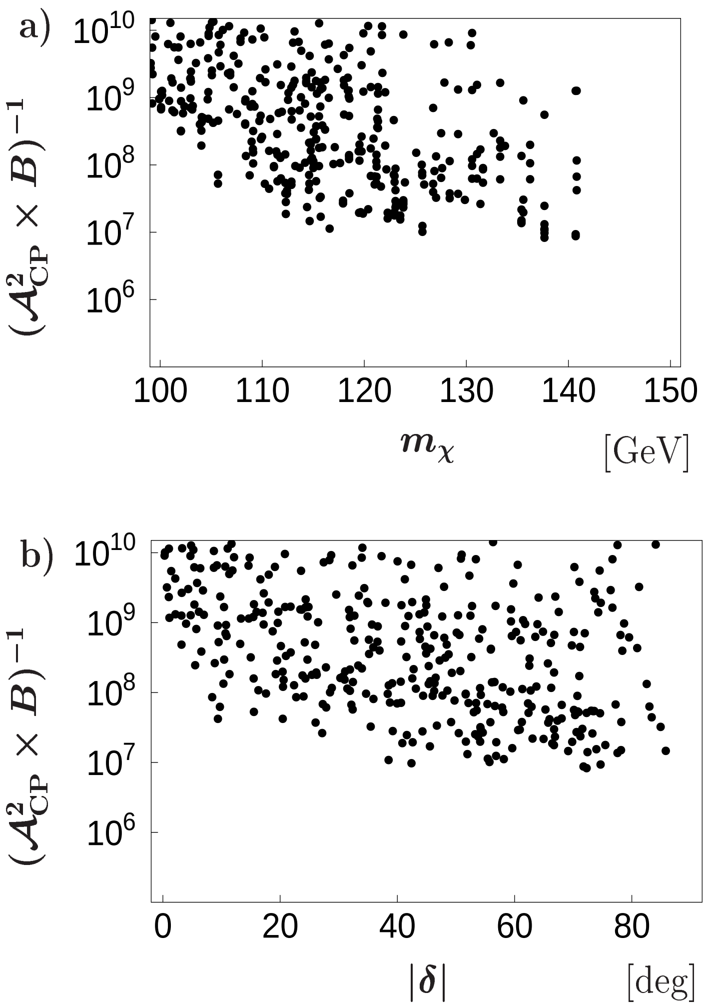

Figure 9: $(\mathcal{A}^2_{\rm CP} \times B)^{-1}$ for $\chi^\pm \to Z e^\pm$ as a function of a) m_χ and b) $|\delta|$, respectively.

plot $(\mathcal{A}^2_{\rm CP} \times B)^{-1}$ as a function of m_χ and of the modulus of the Dirac phase $|\delta|$ defined in Eq. (2.20), respectively. We show $(\mathcal{A}^2_{\rm CP} \times B)^{-1}$ for $\ell = e$; for $\ell = \mu, \tau$ the result is similar. As we can see in Fig. 9a and b, $(\mathcal{A}^2_{\rm CP} \times B)^{-1}$ is always above 10^6.

5. Conclusion

Supersymmetric models with bilinear R-parity breaking contain six new phases compared to the MSSM. These phases are currently only constrained by neutrino data. We have found that neutrino physics requires that the phases in the $\not{R}_p$ sector have to fulfill the relation $\varphi_{\epsilon_i} \simeq \varphi_{B_i}$, but are otherwise not necessarily small. This in turn implies that scalar–pseudoscalar mixing is vanishingly small even though the phases may be maximal CP violating. The only exception is if a CP–even and a CP–odd state are nearly degenerate, in which case the mixing between these two states can be large.

We have discussed CP asymmetries in the decays of various possible LSPs and concluded that all of them are unmeasurably small. Although we have considered only one class of CP violating observables, namely rate asymmetries, we conjecture that also other CP odd observables are to small to be measurable.

On the other hand, the neutrino oscillation Dirac phase δ is not necessarily small if bilinear $R\!\!\!/_p$ parameters are complex. This implies that possibly the only way to determine whether the $R\!\!\!/_p$ parameters are complex is to measure CP violation in neutrino oscillations themselves.

Acknowledgments

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6. Appendix

Here we outline the calculation of the rate asymmetry for neutralino and chargino decay. The calculation is carried out in the Feynman–'t Hooft gauge. In the limit where $m_\ell = 0$ and $|O_R^{\ell\chi h}| >> |O_L^{\ell\chi h}|$ the dominant one–loop amplitude, corresponding to the Feynman diagram in Fig. 8, has the same generic structure for chargino and neutralino decays. For the decay $\chi \to \ell^- V$, with a negative charged lepton in the final state, the amplitude reads:

$$\mathcal{M}^{(1)} = \frac{i}{(4\pi)^2}\bar{u}(p_\ell)[2p_{\ell\rho}(C_0 + C_1 + C_2)P_R - m_\chi C_2 \gamma_\rho P_L]u(p_\chi)\varepsilon^{\rho *}(-p_V) g_1^L g_0 \frac{2m_V^2}{v} \ , \quad (6.1)$$

C_i being the Passarino–Veltman functions, $g_0 \equiv O_R^{\ell\chi h}$ $[O_R^{\nu\chi h}]$ and $g_1^L \equiv -\frac{g}{\cos\theta_W}(-\frac{1}{2} + \sin^2\theta_W)[-\frac{g}{\sqrt{2}}\ \mathcal{N}^*_{i(\ell+4)}]$ for chargino [neutralino] decay. The subscript V stands for the appropriate vector boson and χ for the appropriate SUSY particle, respectively. The tree–level amplitude is given by:

$$\mathcal{M}^{(0)} = i\bar{u}(p_\ell)\gamma_\rho[g_3^L P_L + g_3^R P_R]u(p_\chi)\varepsilon^{\rho *}(-p_V) \ , \quad (6.2)$$

where $g_3^L \equiv O_L^{\ell\chi \mathrm{z}}$ $[O_L^{\ell\chi \mathrm{w}}]$ and $g_3^R \equiv O_R^{\ell\chi \mathrm{z}}$ $[O_R^{\ell\chi \mathrm{w}}]$ for chargino [neutralino] decay. The CP asymmetry is now

$$\mathcal{A}_{\mathrm{CP}} = \frac{\Gamma(\chi \to V\ell^-) - \Gamma(\bar{\chi} \to V\ell^+)}{\Gamma(\chi \to V\ell^-) + \Gamma(\bar{\chi} \to V\ell^+)} \simeq \frac{2\Re e(\mathcal{M}^{(0)\dagger}\mathcal{M}^{(1)})}{\mathcal{M}^{(0)\dagger}\mathcal{M}^{(0)}}$$

$$= \frac{4m_\chi\{\Im m B_0(2m_V^2 + 3m_\chi^2) + \Im m C_0(2m_V^4 + 2m_V^2 m_\chi^2 + m_\chi^4 + m_h^2(m_V^2 - m_\chi^2))\}}{2m_\chi^4 - 3m_V^4 + m_\chi^2 m_V^2}$$

$$\times \frac{1}{(4\pi)^2}\frac{2m_V^2}{v}\ \frac{\Im m(g_3^{L*}\ g_0\ g_1^L)}{(|g_3^L|^2 + |g_3^R|^2)}\ , \tag{6.3}$$

where

$$\Im m B_0 = \pi\frac{m_\chi^2 - m_h^2}{m_\chi^2}\ \Theta(m_\chi^2 - m_h^2)\ ,$$

$$\Im m C_0 = -\pi\frac{1}{m_\chi^2 - m_h^2}\log\left(1 + \frac{(m_\chi^2 - m_h^2)^2}{m_\chi^2 m_V^2}\right)\Theta(m_\chi^2 - m_h^2)\ , \tag{6.4}$$

where $v = 2m_W/g$ and Θ denotes the step–function. From the last line in Eq. (6.3) we can see that the CP asymmetry in the case of the neutralino decay is suppressed, since we have $|g_3^L| << |g_3^R|$. For the chargino decay this relation is reversed, as was mentioned above.

References

[1] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49** (1973) 652.

[2] G. R. Farrar and M. E. Shaposhnikov, Phys. Rev. Lett. **70** (1993) 2833 [Erratum-ibid. **71** (1993) 210] [arXiv:hep-ph/9305274]; P. Huet and E. Sather, Phys. Rev. D **51** (1995) 379 [arXiv:hep-ph/9404302]; M. Carena, M. Quiros and C. E. Wagner, Phys. Lett. B **380** (1996) 81 [arXiv:hep-ph/9603420].

[3] For a recent review see, for example A. Masiero and O. Vives, New J. Phys. **4** (2002) 4.

[4] H. E. Haber and G. L. Kane, Phys. Rept. **117** (1985) 75.

[5] W. Bernreuther and P. Overmann, Z. Phys. C **72** (1996) 461 [arXiv:hep-ph/9511256]; A. Bartl, E. Christova and W. Majerotto, Nucl. Phys. B **460** (1996) 235 [Erratum-ibid. B **465** (1996) 365] [arXiv:hep-ph/9507445]; S. Bar-Shalom, D. Atwood and A. Soni, Phys. Rev. D **57** (1998) 1495 [arXiv:hep-ph/9708357]; M. Aoki and N. Oshimo, Nucl. Phys. B **531** (1998) 49 [arXiv:hep-ph/9801294]; A. Bartl, E. Christova, T. Gajdosik and W. Majerotto, Phys. Rev. D **58** (1998) 074007 [arXiv:hep-ph/9802352]; A. Bartl, E. Christova, T. Gajdosik and W. Majerotto, Phys. Rev. D **59** (1999) 077503 [arXiv:hep-ph/9803426]; V. D. Barger, T. Han, T. J. Li and T. Plehn, Phys. Lett. B **475** (2000) 342 [arXiv:hep-ph/9907425]; S. Y. Choi, H. S. Song and W. Y. Song, Phys. Rev. D **61** (2000) 075004 [arXiv:hep-ph/9907474]; E. Christova, S. Fichtinger, S. Kraml and W. Majerotto, Phys. Rev. D **65** (2002) 094002 [arXiv:hep-ph/0108076]; A. Bartl, T. Kernreiter and W. Porod, Phys. Lett. B **538** (2002) 59 [arXiv:hep-ph/0202198]; T. Ibrahim and P. Nath, Talk at the 10th International Conference on Supersymmetry and Unification of Fundamental Interactions (SUSY02), DESY, Hamburg, Germany, June 17-22, 2002 [arXiv:hep-ph/0210251].

[6] S. Dimopoulos and L. J. Hall, Phys. Lett. B **344** (1995) 185 [arXiv:hep-ph/9411273]; R. Barbieri, A. Romanino and A. Strumia, Phys. Lett. B **369** (1996) 283

[arXiv:hep-ph/9511305]; R. Garisto and J. D. Wells, Phys. Rev. D **55** (1997) 1611 [arXiv:hep-ph/9609511]; D. Chang, W. Y. Keung and A. Pilaftsis, Phys. Rev. Lett. **82** (1999) 900 [arXiv:hep-ph/9811202]; [Erratum-ibid. **83** (1999) 3972]; T. Falk, K. A. Olive, M. Pospelov and R. Roiban, Nucl. Phys. B **560** (1999) 3 [arXiv:hep-ph/9904393]; A. Pilaftsis, Phys. Lett. B **471** (1999) 174 [arXiv:hep-ph/9909485]; M. Brhlik, L. L. Everett, G. L. Kane and J. Lykken, Phys. Rev. Lett. **83** (1999) 2124 [arXiv:hep-ph/9905215]; M. Brhlik, L. L. Everett, G. L. Kane and J. Lykken, Phys. Rev. D **62** (2000) 035005 [arXiv:hep-ph/9908326]; E. Lunghi and D. Wyler, Phys. Lett. B **521** (2001) 320 [arXiv:hep-ph/0109149].

[7] T. Falk and K. A. Olive, Phys. Lett. B **375** (1996) 196 [arXiv:hep-ph/9602299]; V. D. Barger, T. Falk, T. Han, J. Jiang, T. Li and T. Plehn, Phys. Rev. D **64** (2001) 056007 [arXiv:hep-ph/0101106].

[8] M. Dine, A. Kagan and S. Samuel, Phys. Lett. B **243** (1990) 250; S. Dimopoulos and G. F. Giudice, Phys. Lett. B **357** (1995) 573 [arXiv:hep-ph/9507282]; A. Pomarol and D. Tommasini, Nucl. Phys. B **466** (1996) 3 [arXiv:hep-ph/9507462]; A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Lett. B **388** (1996) 588 [arXiv:hep-ph/9607394]; J. Hisano, K. Kurosawa and Y. Nomura, Phys. Lett. B **445** (1999) 316 [arXiv:hep-ph/9810411]; Nucl. Phys. B **584** (2000) 3 [arXiv:hep-ph/0002286]; J. A. Bagger, J. L. Feng, N. Polonsky and R. J. Zhang, Phys. Lett. B **473** (2000) 264 [arXiv:hep-ph/9911255]; J. Bagger, J. L. Feng and N. Polonsky, Nucl. Phys. B **563** (1999) 3 [arXiv:hep-ph/9905292]; K. Agashe and M. Graesser, Phys. Rev. D **59** (1999) 015007 [arXiv:hep-ph/9801446]; J. L. Feng, C. F. Kolda and N. Polonsky, Nucl. Phys. B **546** (1999) 3 [arXiv:hep-ph/9810500].

[9] T. Ibrahim and P. Nath, Phys. Rev. D **57** (1998) 478; T. Ibrahim and P. Nath, Phys. Rev. D **57** (1998) 478 [Erratum-ibid. D **58** (1998 ERRAT,D60,079903.1999 ERRAT,D60,119901.1999) 019901] [arXiv:hep-ph/9708456]; T. Ibrahim and P. Nath, Phys. Rev. D **58** (1998) 111301 [Erratum-ibid. D **60** (1999) 099902] [arXiv:hep-ph/9807501]; Phys. Rev. D **58** (1998) 111301; A. Bartl, T. Gajdosik, W. Porod, P. Stockinger and H. Stremnitzer, Phys. Rev. D **60** (1999) 073003; M. Brhlik, G. J. Good and G. L. Kane, Phys. Rev. D **59** (1999) 115004 [arXiv:hep-ph/9810457]; A. Bartl, T. Gajdosik, E. Lunghi, A. Masiero, W. Porod, H. Stremnitzer and O. Vives, Phys. Rev. D **64** (2001) 076009 [arXiv:hep-ph/0103324].

[10] S. A. Abel, Phys. Lett. B **410** (1997) 173 [arXiv:hep-ph/9612272]; D. Guetta, Phys. Rev. D **58** (1998) 116008 [arXiv:hep-ph/9805274]; G. Bhattacharyya and A. Datta, Phys. Rev. Lett. **83** (1999) 2300 [arXiv:hep-ph/9903490]; R. Mohanta, Phys. Rev. D **63** (2001) 056006 [arXiv:hep-ph/0005240]; G. Bhattacharyya, A. Datta and A. Kundu, Phys. Lett. B **514** (2001) 47 [arXiv:hep-ph/0105057].

[11] For an introduction to explicit R-parity violation see, H. K. Dreiner, arXiv:hep-ph/9707435.

[12] R. M. Godbole, S. Pakvasa, S. D. Rindani and X. Tata, Phys. Rev. D **61** (2000) 113003 [arXiv:hep-ph/9912315]; S. A. Abel, A. Dedes and H. K. Dreiner, JHEP **0005** (2000) 013 [arXiv:hep-ph/9912429].

[13] K. Choi, E. J. Chun and K. Hwang, Phys. Rev. D **63**, 013002 (2001) [arXiv:hep-ph/0004101]; Y. Y. Keum and O. C. Kong, Phys. Rev. Lett. **86** (2001) 393 [arXiv:hep-ph/0004110]; Y. Y. Keum and O. C. Kong, Phys. Rev. D **63** (2001) 113012 [arXiv:hep-ph/0101113]; D. Chang, W. F. Chang, M. Frank and W. Y. Keung, Phys. Rev. D **62** (2000) 095002 [arXiv:hep-ph/0004170].

[14] Y. Fukuda *et al.* [Super-Kamiokande Collaboration], Phys. Rev. Lett. **81**, 1562 (1998) hep-ex/9807003; see also hep-ex/9803006 and hep-ex/9805006; for an update see the talk by M. Shiozawa in [15].

[15] Talks at XX. Int. Conf. on Neutrino Physics & Astrophysics (Neutrino 2002), http://neutrino2002.ph.tum.de/

[16] Q. R. Ahmad *et al.* [SNO Collaboration], Phys. Rev. Lett. **89** (2002) 011301 [arXiv:nucl-ex/0204008]; Q. R. Ahmad *et al.* [SNO Collaboration], Phys. Rev. Lett. **89** (2002) 011302 [arXiv:nucl-ex/0204009].

[17] See, for example, M. Maltoni, T. Schwetz, M. A. Tortola and J. W. Valle, arXiv:hep-ph/0207227, and references therein.

[18] M. C. Gonzalez-Garcia, P. C. de Holanda, C. Pena-Garay and J. W. Valle, Nucl. Phys. B **573** (2000) 3 [arXiv:hep-ph/9906469].

[19] R. Hempfling, *Nucl. Phys.* **B478**, 3 (1996).

[20] M. Hirsch, M. A. Diaz, W. Porod, J. C. Romao and J. W. Valle, Phys. Rev. D **62**, 113008 (2000) [Erratum-ibid. D **65**, 119901 (2002)] [arXiv:hep-ph/0004115]; J. C. Romao, M. A. Diaz, M. Hirsch, W. Porod and J. W. Valle, Phys. Rev. D **61** (2000) 071703 [arXiv:hep-ph/9907499].

[21] M. Masip and A. Rasin, Phys. Rev. D **58**, 035007 (1998).

[22] A. Joshipura and M. Nowakowski, Phys. Rev. D **51**, 5271–5275 (1995).

[23] J. Schechter and J. W. F. Valle, Phys. Rev. **D25**, 774 (1982); M. Nowakowski and A. Pilaftsis, Nucl. Phys. **B461**, 19 (1996) [hep-ph/9508271]; M. Hirsch and J. W. F. Valle, Nucl. Phys. **B557**, 60 (1999).

[24] J. N. Bahcall, M. C. Gonzalez-Garcia and C. Pena-Garay, JHEP **0207** (2002) 054 [arXiv:hep-ph/0204314].

[25] W. Porod, M. Hirsch, J. C. Romao and J. W. Valle, Phys. Rev. D **63** (2001) 115004 [arXiv:hep-ph/0011248].

[26] M. Hirsch, H.V. Klapdor-Kleingrothaus and S.G. Kovalenko, Phys. Lett. **B398** (1997) 311.