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NEUTRINO–NEUTRALINO MIXING AND ONE-LOOP CORRECTIONS IN THE R-PARITY VIOLATING SUPERSYMMETRIC MODEL

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We consider the supersymmetric extension of the Standard Model with neutrino Yukawa interactions and R-parity violation. We calculate one-loop corrections to physical neutrino mass in the case of neutrino–neutralino mixing and discuss the influence of this mass shift on parameter constraints.

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INTRODUCTION

The searches for the Higgs boson and new phenomena beyond the Standard Model (SM) of fundamental interactions are important tasks for the Large Hadron Collider. The most popular direction beyond the SM is low energy supersymmetry. However, it is not clear how supersymmetry is realized. The simplest case, the Minimal Supersymmetric Standard Model (MSSM) [1–4], is studied in detail; however, possible deviations from it are of great interest as well. There exists a wide class of models which contain the so-called R-parity breaking interactions leading to the violation of lepton and baryon numbers [5]. These models have a number of new coupling constants, some of them are badly constrained, for instance, by rare processes, other ones are less restricted.

In this paper we consider the model with the neutrino Yukawa interactions and the related R-parity violating term in the superpotential [6]. Among interesting consequences of including the R-parity violating term $\lambda_i \bar{\nu}_i H_u H_d$ is mixing between neutrino and neutralino. We study a constraint connected to small neutrino masses at tree level and evaluate one-loop corrections.

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1. SUPERSYMMETRIC STANDARD MODEL WITH THE RIGHT-HANDED NEUTRINO AND R-PARITY BREAKING

The superpotential of the Minimal Supersymmetric Standard Model

$$W_{\mathrm{RMSSM}} = y_u^{ij}\bar{u}_i Q_j \cdot H_u - y_d^{ij}\bar{d}_i Q_j \cdot H_d - y_e^{ij}\bar{e}_i L_j \cdot H_d + \mu H_u \cdot H_d \tag{1}$$

is constructed under the assumption that neutrinos are massless (there are no Yukawa interactions for the neutrinos) and the R-parity is conserved. However, it is believed nowadays that neutrinos have masses, even tiny, then the neutrino Yukawa term

$$y_\nu^{ij}\bar{\nu}_i L_j \cdot H_u \tag{2}$$

in the superpotential is possible and should be included ($\bar{\nu}_i$ here are $SU(2)$ singlet right-handed neutrino superfields, and y_ν^{ij} are neutrino Yukawa couplings). The latter implies that one also has to include a term

$$\lambda_\nu^i \bar{\nu}_i H_u \cdot H_d \tag{3}$$

in the R-violating part. Therefore, we consider a model with the following superpotential:

$$W = W_{\mathrm{RMSSM}} + y_\nu^{ij}\bar{\nu}_i L_j \cdot H_u + \lambda_\nu^i \bar{\nu}_i H_u \cdot H_d. \tag{4}$$

The soft supersymmetry-breaking Lagrangian also includes the following terms:

$$\mathcal{L}_{\mathrm{SSB}} = \ldots + A_\nu^{ij}\bar{\nu}_i L_j \cdot H_u + A_\lambda^i \bar{\nu}_i H_u \cdot H_d \tag{5}$$

($\bar{\nu}_i, L_j, H_u, H_d$ here are scalar components of the corresponding superfields).

The model with this kind of superpotential has been previously studied [7,8]. However, the authors were mainly interested in the solution of the μ problem rather than considering Higgs and neutrino mass predictions in the model. Also we suppose sneutrino has no v.e.v. to avoid spontaneous lepton number violation.

The new term (3) gives the F-type contribution to the Higgs self-coupling, and the Higgs scalar potential now reads

$$V = V_{\mathrm{MSSM}} + \left|\lambda_\nu^i \lambda_\nu^i\right| \left|H_u^+ H_d^- - H_u^0 H_d^0\right|^2. \tag{6}$$

Minimization conditions and the curvature at the minimum are modified as well, leading to an increase of the lightest Higgs boson mass [6].

2. NEUTRINO MASSES AND NEUTRINO–NEUTRALINO MIXING

As the electroweak symmetry is broken and the Higgs fields acquire v.e.v.'s v_u and v_d, the R-parity breaking term (3) mixes neutrino and neutralino and thus gives a contribution to their mass matrix. This matrix represents quadratic interactions of the following particles: four superpartners of Higgs and gauge bosons $\tilde{G}^0 = (\tilde{B}\tilde{W}^0\tilde{H}_d^0\tilde{H}_u^0)^T$, three right-handed neutrinos ($SU(2)$ singlets) $\bar{N} = (\bar{\nu}_e, \bar{\nu}_\mu, \bar{\nu}_\tau)^T$, and three left-handed neutrinos (components of $SU(2)$ doublets) $N = (\nu_e, \nu_\mu, \nu_\tau)^T$:

$$-\frac{1}{2}(\tilde{G}^{0T}\bar{N}^T N^T) \begin{pmatrix} \mathbf{M}_{\tilde{G}^0} & \mathbf{M}_{\tilde{G}^0\bar{N}} & 0 \\ \mathbf{M}^T_{\tilde{G}^0\bar{N}} & \mathbf{M}_{\bar{N}} & \mathbf{M}_D \\ 0 & \mathbf{M}_D^T & 0 \end{pmatrix} \begin{pmatrix} \tilde{G}^0 \\ \bar{N} \\ N \end{pmatrix} + \text{h.c.} \qquad (7)$$

The 4×4 matrix $\mathbf{M}_{\tilde{G}^0}$ is the usual neutralino mass matrix [4]. The left-handed neutrinos cannot have any mass terms due to the $SU(2)$ gauge symmetry of the Lagrangian. The new superpotential term (2) generates a Dirac neutrino mass $(\mathbf{M}_D)^{ij} = y_\nu^{ij} v_u$, and the new term (3) generates

$$\mathbf{M}_{\tilde{G}^0\bar{N}} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ -\lambda_\nu^1 v_u & -\lambda_\nu^2 v_u & -\lambda_\nu^3 v_u \\ -\lambda_\nu^1 v_d & -\lambda_\nu^2 v_d & -\lambda_\nu^3 v_d \end{pmatrix} = M_Z \sqrt{\frac{2}{g^2 + g'^2}} \begin{pmatrix} 0 \\ 0 \\ -\lambda_\nu^i \sin\beta \\ -\lambda_\nu^i \cos\beta \end{pmatrix},$$

where $\tan\beta = v_u/v_d$.

The Majorana mass terms $(\mathbf{M}_{\bar{N}})^{ij}\bar{\nu}_i\bar{\nu}_j$ are allowed by gauge invariance. The Majorana mass can origin from quadratic terms of a superpotential or from SUSY-breaking terms. Anyway $(\mathbf{M}_{\bar{N}})$ is arbitrary in supersymmetric models.

In order to explain the huge difference between the neutrino mass scale (eVs) and the mass scale of other fundamental fermions (e.g., GeVs for quarks), one may assume the elements of $\mathbf{M}_{\bar{N}}$ to be very large (about GUT scale). Then the mass matrix has three small eigenvalues («see-saw mechanism» [9]). In this scenario, neutrino is a Majorana particle.

Another possibility is $\mathbf{M}_{\bar{N}} = 0$ and y_ν^{ij} to be very small. However, these conditions are not sufficient to ensure the smallness of neutrino masses in the general case for significant values of $\varepsilon^2 = \lambda_\nu^i \lambda_\nu^i$. Nevertheless, the mass matrix (7) can still have three small eigenvalues for a special parameter set. This requirement gives a very strong constraint on the model parameters. Though it can be treated as not quite «natural» since it requires a fine-tuning, it is interesting to see if this constraint is in accordance with other constraints. Numerical calculations discussed bellow show that this is a feasible option.

In this scenario, neutrino is a «pseudo-Dirac» particle. It should be noted that all the experimental data on neutrino mixing can be successfully described in both scenarios. A possible discovery of neutrinoless double beta decay would be able to clarify the nature of neutrino (Dirac or Majorana particle) [10].

3. ONE-LOOP CORRECTIONS

One-loop corrections to neutralino masses have been considered in literature (e.g., [11]). We generalize those results for the extended neutrino–neutralino mass matrix (7). It can be diagonalized with an orthogonal mixing matrix Z: $ZMZ^T = \mathrm{diag}\,(m_i)$. Mass shift of a physical state i reads [12]

$$\delta m_i = m_i \Sigma^L_{ii}(m_i^2) \,+\, \frac{1}{2}\Big[\Sigma^M_{ii}(m_i^2) + \Sigma^{M*}_{ii}(m_i^2)\Big]. \tag{8}$$

The neutralino one-loop self-energies have the following form [11]:

$$\Sigma^L_{ij}(k^2) = \frac{-1}{(4\pi)^2} \sum_{\mathrm{gen}} N_C \sum_{f=u,d} \sum_{a=1,2} (a^{\tilde f}_{ai} a^{\tilde f}_{aj} + b^{\tilde f}_{ai} b^{\tilde f}_{aj})\, B_1(k^2, m_f^2, m^2_{\tilde f_a}), \tag{9}$$

$$\Sigma^M_{ij}(k^2) = \frac{1}{(4\pi)^2} \sum_{\mathrm{gen}} N_C \sum_{f=u,d} \sum_{a=1,2} (a^{\tilde f}_{ai} b^{\tilde f}_{aj} + a^{\tilde f}_{aj} b^{\tilde f}_{ai})\, m_f B_0(k^2, m_f^2, m^2_{\tilde f_a}), \tag{10}$$

where the neutralino–sfermion–fermion couplings are

$$a^{\tilde f}_{ak} = \sqrt{2}\, g \left[\left(e_f - I^{3L}_f\right) \tan\theta_W Z_{k1} + I^{3L}_f Z_{k2} \right] R^{\tilde f}_{a1} + y_f Z_{kx} R^{\tilde f}_{a2},$$
$$b^{\tilde f}_{ak} = -\sqrt{2}\, g e_f \tan\theta_W\, Z_{k1} R^{\tilde f}_{a2} + y_f Z_{kx} R^{\tilde f}_{a1},$$

with $x = 3$ for down-type and $x = 4$ for up-type fermions; e_f, I^{3L}_f and y_f stand for their electric charge, weak isospin and Yukawa couplings; Z and R are neutralino and sfermion mixing matrices; two-point functions B_0 and B_1 are given in [13].

We apply (8) for the light eigenstates corresponding to neutrinos and we expect m_i and δm_i to be small (in comparison to the top-quark mass; we verify this assumption below). Therefore, the first term in (8) proportional to m_i can be neglected (usually one neglects the second term, e.g., for heavy neutralino states). The major contribution to the second term comes from top quark and squarks. When the terms containing the Yukawa coupling y_t are taken into account and gauge constant g is neglected, a simple approximation for δm_i can be derived. In this approximation, divergences are cancelled and δm_i is finite:

$$\delta m_i = \frac{3}{4\pi^2}\, m_t y_t^2 \left(Z_{i4}\right)^2 R_{11} R_{21} \times$$
$$\times \left(\log\frac{m_{\tilde t_1}}{m_{\tilde t_2}} + \frac{m_t^2}{m^2_{\tilde t_1} - m_t^2} \log\frac{m_{\tilde t_1}}{m_t} - \frac{m_t^2}{m^2_{\tilde t_2} - m_t^2} \log\frac{m_{\tilde t_2}}{m_t} \right). \tag{11}$$

Note that it vanishes in the case of degenerate squark masses.

4. PARAMETER SPACE AND NUMERICAL RESULTS

We adopt the mSUGRA universality hypothesis and suppose the following parameters are free: the universal (GUT-scale) scalar mass m_0, the spinor mass

$m_{1/2}$, and $\tan\beta$. The trilinear coupling A_0 is chosen to be 0 and $\mathrm{sign}\,(\mu) = +1$. We investigate the case $\mathbf{M}_{\tilde{N}} = 0$.

We use one-loop renormalization group equations [14]. Also we take into account one- and two-loop leading corrections from heavy quarks to the Higgs potential [15].

The numerical results described below were obtained for a simplified model with only third generation being massive. This means that $\lambda_1 = \lambda_2 = 0$, $\varepsilon = \lambda_3^2$. Because of the tiny neutrino masses we assume $y_\nu^{ij} = 0$. Therefore, the mass matrix (7) has nontrivial 5×5 block. For m_0 = 1000 GeV, $m_{1/2}$ = 431.5 GeV (we explain this choice below) and $\varepsilon = 0.05$, the block reads

$$M = \begin{pmatrix} 195.4 & 0 & -0.89 & 44.6 & 0 \\ 0 & 364.7 & 1.59 & -79.5 & 0 \\ -0.89 & 1.59 & 0 & -668.9 & -28.8 \\ 44.6 & -79.5 & -668.9 & 0 & -0.58 \\ 0 & 0 & -28.8 & -0.58 & 0 \end{pmatrix} \text{GeV}. \tag{12}$$

The eigenvalues are 194.3, 357.3, 682.0, −673.5, −0.00147, and the eigenvectors form the mixing matrix

$$Z = \begin{pmatrix} 0.9968 & -0.0112 & 0.0744 & -0.0225 & -0.0110 \\ 0.0253 & 0.9824 & -0.1623 & 0.0885 & 0.0129 \\ -0.0654 & 0.1788 & 0.6880 & -0.6998 & -0.0285 \\ -0.0356 & 0.0531 & 0.7034 & 0.7072 & 0.0307 \\ 0.0099 & -0.0094 & 0.0009 & -0.0431 & 0.9990 \end{pmatrix}. \tag{13}$$

The computer code searches for the set of parameters where the neutralino mass determined by the small eigenvalue is 0. In the case of parameter values listed above, the estimation of the mass shift (11) is $\delta m_5 = 0.00146$ GeV and $m_5 + \delta m_5 \approx 0$ (the expression (11) can be used as a rough approximation since $|Z_{54}/Z_{51}| \approx |Z_{54}/Z_{52}| \approx 4.5$). Consequently, this point meets the constraint in question and that is why it has been considered.

The allowed values of the parameters form a line $m_\nu = 0$. Its position depends on $\tan\beta$. For large $\tan\beta = 50$ the line constraint may be compatible with other constraints (Fig. 1).

As we can see from the plots, the influence of loop corrections is noticeable but moderate. The loop corrections are important for precise calculations but do not change qualitative predictions.

As a reference point we have shown the experimental LEP limit on the lightest neutralino mass $m_{\chi_1^0} > 46$ GeV [16] (it excludes a region corresponding to $m_{1/2} \lesssim 100$ GeV) and the constraint from the Higgs boson mass between 125 and 126 GeV [17, 18].

For small $\tan\beta = 3$ this line corresponds to small values of $m_{1/2}$ and contradicts other constraints, in particular, the neutralino constraint (Fig. 2). The

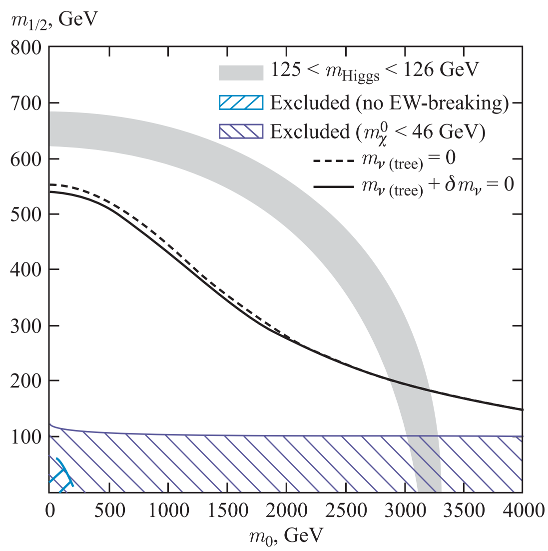

Fig. 1. Parameter space for $\tan\beta = 50$ and $\varepsilon^2 = 0.05$

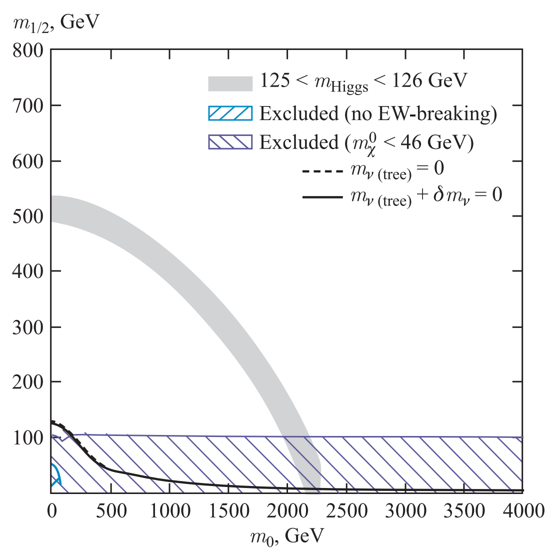

Fig. 2. Parameter space for $\tan\beta = 3$ and $\varepsilon^2 = 0.4$

constraints shown there were obtained for relatively large $\varepsilon = 0.4$. Unlike MSSM ($\varepsilon = 0$), this choice can be interesting because the Higgs boson mass constraint does not completely exclude the small $\tan\beta$ scenario.

CONCLUSIONS

We have shown that introducing the Yukawa interactions of neutrinos $y_\nu^{ij}\bar{\nu}_i L_j H_u$ leads to their Dirac mass terms, and the possible R-parity violating term $\lambda_\nu^i \bar{\nu}_i H_u H_d$ leads to neutrino–neutralino mixing.

There are two possible scenarios to ensure the smallness of the neutrino masses. In the see-saw mechanism scenario the Majorana mass is huge (close to the GUT scale). In the «pseudo-Dirac» neutrino scenario the Majorana mass is 0 and the Dirac mass is small. Moreover, a special parameter set is required to obtain three small eigenvalues of the neutrino–neutralino mass matrix. This constraint is compatible with other constraints for large $\tan\beta$, and the loop corrections from higgsino mixing in the physical neutrino state, though being noticeable, do not change this conclusion. However, the constraint may not be considered as quite natural, because a special choice of parameters can be treated as a fine-tuning.

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