The “Right” Sneutrino as the LSP

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(Dated: March 25, 2022)

Abstract

It is shown that “right” sneutrino can be the lightest supersymmetric particle. Clearly, this possibility will drastically change decay chains of SUSY particles. The sneutrino production at next linear colliders has been analyzed in this scenario.

PACS numbers: 11.30.Pb, 12.60.Jv, 14.80.Ly
I. INTRODUCTION

R-parity conserving Minimal Supersymmetric Standard Model (MSSM) is the most studied scenario of physics beyond the SM [1]. Searching for supersymmetric particles is an essential part of experimental programs of future colliders. Obviously the search strategy is strongly dependent on mass and mixing patterns of the SUSY particles, especially on an assumption about Lightest Supersymmetric Particle (LSP). Usually, the lightest neutralino is taken as the LSP. The sneutrino can be considered as an another candidate for LSP. However, LEP1 data exclude this possibility for “left” sneutrino [2]. It should be emphasized that this statement is valid for superpartners of left-handed neutrinos. On the other hand, neutrino oscillation experiments show that neutrinos have non-zero masses and it turns out that right-handed neutrinos must exist. Thus superpartners of the right-handed neutrinos should be included into the MSSM. As noted in [3] the LEP1 data does not essentially constrain the masses of superpartners of the right-handed neutrinos if their mixings with “left” sneutrinos are sufficiently small.

In this paper we present a scenario which assumes the superpartner of right-handed neutrino (“Right” Sneutrino) to be the Lightest Supersymmetric Particle; which is hereafter called RS-LSP scenario.

In section 2 we illustrate that “right” sneutrino can be the LSP and analyze constraints on mass and mixing of “right” sneutrino coming from LEP1 data. Next in section 3 we briefly discuss decay chains of SUSY particles in this scenario. A search for sneutrino at future linear colliders is considered in section 4. Finally, we give some concluding remarks in section 5.

II. RS-LSP SCENARIO

The huge number of free parameters in the three family MSSM leads to consideration of some simplified versions, such as the constrained MSSM (see [6] and references therein). In general, these simplifications ignore interfamily mixings and possible existence of right-handed neutrinos, and consequently their super-partners. As a result one avoids possible conflicts with experimental data on flavor violating processes. But at the same time we miss very interesting possible phenomenology. In this study we deal with the three family MSSM
and do not consider possible R-parity violation, as well as GUT and SUGRA extensions.

A. MSSM with “right” sneutrinos.

There are a number of arguments favoring the existence of right-handed neutrinos. First of all, in the framework of the SM $\nu_R$’s are counterparts of the right-handed components of the up-type quarks according to the quark-lepton symmetry. Then, almost all extensions of the SM, with the SU(5) GUT as a possible exception, naturally contain right-handed neutrinos. Finally, as it was mentioned in introduction, observation of the neutrino oscillations provides the experimental confirmation for $\nu_R$. For these reasons we consider three family MSSM with right-handed neutrinos. Therefore we deal with the six species that constitute a family, rather than five species considered in [5] (in the argument given below, we adopt the notations used in this paper): $q$, $\bar{d}$, $\bar{u}$, $l$, $\bar{e}$, and $\bar{\nu}$, where $q$ and $l$ denote weak iso-doublets and the rest are iso-singlets. The masses of the SM fermions and their super-partners are generated due to

$$L = L_{\text{scalar}} + L_{\text{Yukawa}} + L_{\text{triscalar}}.$$

The first term has the form

$$L_{\text{scalar}} = \sum_{A,i,j} m^2_{Ai,j} \tilde{A}^*_i \tilde{A}_j,$$

where $A$ labels six species mentioned above, the tilde labels sparticle and $i, j=1,2,3$ are family labels. Therefore this part contains six $3 \times 3$ Hermitian mass matrices. Each matrix contains six real parameters and three phases. The Yukawa part of the lagrangian is derived from the superpotential

$$W_{\text{Yukawa}} = \sum_{i,j} (q_i \lambda_{uij} \bar{u}_j H_u + q_i \lambda_{dij} \bar{d}_j H_d + l_i \lambda_{\nu ij} \bar{\nu}_j H_u + l_i \lambda_{eij} \bar{e}_j H_d),$$

where four Yukawa matrices $\lambda$ are general $3 \times 3$ matrices. Each of these matrices contains nine real parameters and nine phases. Finally, the triscalar part is given by

$$L_{\text{triscalar}} = \sum_{i,j} (\tilde{q}_i a_{uij} \tilde{u}_j H_u + \tilde{q}_i a_{dij} \tilde{d}_j H_d + \tilde{l}_i a_{\nu ij} \tilde{\nu}_j H_u + \tilde{l}_i a_{eij} \tilde{e}_j H_d) \times M,$$
where $a$ are general $3 \times 3$ matrices and $M$ is some mass parameter. Here in eq. (4) in order to avoid confusion with eq. (3) we introduce the notations which are slightly different from that in [2]. As a result we have 108 real parameters (masses and mixing angles) and 90 phases. However, part of them are unobservable because of $U(3)^6$ symmetry of the gauge sector, so 18 angles and 34 phases can be rotated out and remaining two phases correspond to baryon number in quark sector and general lepton number. Finally, matter sector of the MSSM (fundamental fermions and their superpartners) contains 90 observable real parameters, namely 36 masses and 54 mixing angles, and 56 phases.

B. Flavor Democracy

Now let us recall the main assumptions of the flavor democracy i.e., democratic mass matrix hypothesis (in the framework of the n family SM):

i) Before the spontaneous symmetry breaking fermions with the same quantum numbers are indistinguishable. Therefore Yukawa couplings are equal within each type of fermions $\lambda_{uij} = \lambda_u$, $\lambda_{dij} = \lambda_d$, $\lambda_{lij} = \lambda_l$ and $\lambda_{vij} = \lambda_\nu$.

ii) There is only one Higgs doublet, which gives Dirac masses to all four types of fermions. Therefore Yukawa constants for different types of fermions should be nearly equal $\lambda_u \approx \lambda_d \approx \lambda_l \approx \lambda_\nu \approx \lambda$.

The first statement result in $n-1$ massless particles and one massive particle with $m = n\lambda_F$ ($F=\text{u, d, l, } \nu$) for each type of the SM fermions. The masses of the first $n-1$ families, as well as observable interfamily mixings, are generated due to a small deviation from the full flavor democracy. Taking into account the mass values for the third generation, the second statement leads to the assumption that the fourth SM family should exist. Alternatively, masses of up and down type fermions should be generated by the interaction with different Higgs doublets, as it takes place in the MSSM.
C. “Right” sneutrino as the LSP

It is straightforward to apply flavor democracy to the MSSM. For example according to the flavor democracy, sneutrino mass matrix has the form

\[
\begin{pmatrix}
    m_{LL}^2 & m_{LL}^2 & m_{LL}^2 & m_{LR}^2 & m_{LR}^2 & m_{LR}^2 \\
    m_{LL}^2 & m_{LL}^2 & m_{LL}^2 & m_{LR}^2 & m_{LR}^2 & m_{LR}^2 \\
    m_{LL}^2 & m_{LL}^2 & m_{LL}^2 & m_{LR}^2 & m_{LR}^2 & m_{LR}^2 \\
    m_{RL}^2 & m_{RL}^2 & m_{RL}^2 & m_{RR}^2 & m_{RR}^2 & m_{RR}^2 \\
    m_{RL}^2 & m_{RL}^2 & m_{RL}^2 & m_{RR}^2 & m_{RR}^2 & m_{RR}^2 \\
    m_{RL}^2 & m_{RL}^2 & m_{RL}^2 & m_{RR}^2 & m_{RR}^2 & m_{RR}^2 \\
\end{pmatrix}
\]

As a result we deal with four massless sneutrinos and two sneutrino s having the masses

\[
m_{3,6}^2 = \frac{3}{2} \left( m_{LL}^2 + m_{RR}^2 + \sqrt{(m_{LL}^2 - m_{RR}^2)^2 + 4m_{LR}^2m_{RL}^2} \right).
\]

The (small) masses of the remaining four species can be generated due to the violation of flavor democracy. Including F- and D-term contributions, the elements of the matrix (5) have the following form [6]:

\[
\begin{align*}
    m_{LL}^2 &= m_i^2 + (\lambda_\nu v_u)^2 + \frac{1}{2} m_Z^2 \cos 2\beta, \\
    m_{LR}^2 &= m_{RL}^2 = a_\nu (M v_u - \mu v_d), \\
    m_{RR}^2 &= m_\nu^2 + (\lambda_\nu v_u)^2,
\end{align*}
\]

where \(v_u\) and \(v_d\) are vacuum expectation values of the Higgs fields \(H_u\) and \(H_d\), \(\tan \beta = v_u/v_d\) and \(\mu\) is the supersymmetry-conserving Higgs mass parameter.

In [2] it is shown that LEP1 data leads to a lower bound 44.6 GeV on the sneutrino masses and, consequently in the framework of the constrained MSSM [4], “left” sneutrino cannot be the LSP. But LEP1 data does not essentially constrain the masses of superpartners of the right-handed neutrinos if the LR mixings are sufficiently small. Indeed, the contribution of the “right” sneutrino to the invisible Z width is given by

\[
\Delta \Gamma_{inv} = 0.5 \times |\delta|^2 \times \left( 1 - \left( \frac{2 \bar{m}_\nu}{m_Z} \right)^2 \right)^{3/2} \times \Gamma_\nu
\]
where \( \delta \) denotes the “left” sneutrino fraction due to corresponding mixings and \( \Gamma_\nu = 167 \) MeV. The experimental value \( \Delta \Gamma_{\text{inv}} \leq 2.0 \) MeV leads to \( |\delta| \leq 0.155 \) for sufficiently light “right” sneutrino. If there are two light species one obtains \( |\delta_1|^2 + |\delta_2|^2 \leq 0.024 \). Therefore “right” sneutrino still can be considered as the LSP both in constrained and unconstrained MSSM.

III. SNEUTRINO DECAPS IN RS-LSP SCENARIO

Decay chains of SUSY particles depend on their masses and mixing patterns. In general one has to deal with a \( 6 \times 6 \) mass matrix for up-type and down-type squarks, charged sleptons and sneutrinos. Following [10], we ignore the interfamily mixings so that in the sneutrino sector we are left with

\[
\tilde{\nu}_1^l = \cos \phi \tilde{\nu}_L^l + \sin \phi \tilde{\nu}_R^l \\
\tilde{\nu}_2^l = -\sin \phi \tilde{\nu}_L^l + \cos \phi \tilde{\nu}_R^l
\]

(9)

for each family. From now on \( \tilde{\nu}_2^e \) is assumed to be the LSP, so that it is stable due to R-parity conservation. In this case the parameter \( \delta \) in eq. (8) is going to be \( \sin \phi_e \). The mass patterns of sleptons and squarks are assumed to satisfy \( \tilde{m}_2 \ll \tilde{m}_1 < \tilde{m}_l < \tilde{m}_q \), where \( \tilde{m}_2 \) and \( \tilde{m}_1 \) denote the masses of \( \tilde{\nu}_2^e \) and \( \tilde{\nu}_1^e \), respectively. If \( m_Z < \tilde{m}_1 \ll \tilde{m}_w, \tilde{m}_Z \) the sole two body decay mode is \( \tilde{\nu}_1^e \to Z\tilde{\nu}_2^e \). For heavier \( \tilde{\nu}_1^e \), two additional decay modes \( \tilde{\nu}_1^e \to \tilde{\nu}_1^e \) and \( \tilde{\nu}_1^e \to \tilde{\nu}_2^e \) do appear. In Fig. 1 decay width \( \Gamma(\tilde{\nu}_1^e \to Z\tilde{\nu}_2^e) \) is plotted as a function of \( \tilde{m}_2 \) for the following chosen values of \( \tilde{m}_1 = 150 \) GeV and 200 GeV. Decay widths \( \Gamma(\tilde{\nu}_1^e \to \tilde{\nu}_1^e) \) and \( \Gamma(\tilde{\nu}_1^e \to \tilde{\nu}_2^e) \) are plotted in Fig. 2 as functions of \( \tilde{m}_w \) and \( \tilde{m}_Z \) with \( \tilde{m}_1 = 200 \) GeV. Dependences of these decay widths on \( \sin \phi \) are shown in Fig. 3 by taking \( \tilde{m}_1 = 200 \) GeV and \( \tilde{m}_w = \tilde{m}_Z = 150 \) GeV. As it is seen from Figs. 1, 2, \( \Gamma(\tilde{\nu}_1^e \to Z\tilde{\nu}_2^e) \ll \Gamma(\tilde{\nu}_1^e \to \tilde{\nu}_1^e) < \Gamma(\tilde{\nu}_1^e \to \tilde{\nu}_2^e) \) when \( \sin \phi \) is less than 0.2 and \( \tilde{m}_1 - \tilde{m}_{w,Z} > 10 \) GeV.

IV. SNEUTRINO PRODUCTION AT FUTURE LINEAR COLLIDERS

We have four processes \( e^+ e^- \to \tilde{\nu}_1^e \tilde{\nu}_1^e \) ( \( \tilde{\nu}_2^e \tilde{\nu}_2^e \), \( \tilde{\nu}_1^e \tilde{\nu}_1^e \), \( \tilde{\nu}_2^e \tilde{\nu}_1^e \) ) in RS-LSP scenario instead of one process \( e^+ e^- \to \tilde{\nu}_e \tilde{\nu}_e \) in MSSM without right neutrino (we should remember that we ignore
interfamily mixings and consider only the first lepton family). These processes proceed via t-channel $\tilde{w}$ exchange and s-channel Z boson exchange. The differential cross-sections for the processes $e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_2^c$, $\tilde{\nu}_1^c\tilde{\nu}_1^c$ are obtained as follows:

$$\frac{d\sigma(e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_2^c,\tilde{\nu}_2^c\tilde{\nu}_1^c)}{dt} = \frac{C_{12}g_W^4[(t-\tilde{m}_1^2)(t-\tilde{m}_2^2)+st]}{16\pi s^2}\Delta$$

where $C_{12}=\sin^2\varphi \cos^2\varphi$ and

$$\Delta = \frac{1}{4(t-\tilde{m}_w^2)^2} + \frac{1}{16\cos^4\theta_W} \left[ \frac{1-4\sin\theta_W+8\sin^2\theta_W}{(s-M_Z^2)^2+\Gamma_Z^2M_Z^2} \right]$$

$$- \frac{(2\sin\theta_W-1)(s-M_Z^2)}{2\cos^2\theta_W[(s-M_Z^2)^2+\Gamma_Z^2M_Z^2](t-\tilde{m}_w^2)^2}.$$ 

The differential cross-section for $e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_1^c$ is obtained from the general expression eq. (10) with the replacements $C_{12}\to C_{11} = \cos^4\varphi$ and $\tilde{m}_2\to\tilde{m}_1$ and for $e^+e^-\to\tilde{\nu}_2^c\tilde{\nu}_2^c$ with the replacements $C_{12}\to C_{22} = \sin^4\varphi$ and $\tilde{m}_1\to\tilde{m}_2$. In the case of $\sin\varphi = 0$ we obtain the well-known formula for $e^+e^-\to\tilde{\nu}_e\tilde{\nu}_e$. For numerical evaluations center of mass energy is taken 500 GeV, and the values $\tilde{m}_w = 150$, $M_Z = 92$ GeV, $\sin\varphi = 0.1$, $\sin^2\theta_W = 0.22$, $g_W = 0.66$ are used. In Figs. 4 and 5 we plot total cross sections of the processes $e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_1^c$ and $e^+e^-\to\tilde{\nu}_2^c\tilde{\nu}_2^c$ as functions of $\tilde{m}_1$ and $\tilde{m}_2$ respectively. The cross section of the process $e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_1^c$ is shown in Fig. 6 as a function of $\tilde{m}_1$ by setting $\tilde{m}_2 = 0$. Finally, for illustration we give the $\sin\varphi$ dependence of $\sigma(e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_1^c)$ in Fig. 7 by taking $\tilde{m}_1 = 200$ GeV and $\tilde{m}_2 = 0$.

In RS-LSP scenario $\tilde{\nu}_2^c$ is stable and hence the process $e^+e^-\to\tilde{\nu}_2^c\tilde{\nu}_2^c$ is not observable. In principle the process $e^+e^-\to\tilde{\nu}_2^c\tilde{\nu}_2^c\gamma$ is observable, but it has too small cross-section for $\sin\varphi < 0.2$.

For the process $e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_1^c$ the most spectacular manifestation is two Z bosons which are not coplanar, with missing energy due to $\tilde{\nu}_1^c\to Z\tilde{\nu}_2^c$ and $\tilde{\nu}_1^c\to Z\tilde{\nu}_2^c$. Assuming these decays are dominant and considering leptonic decays $Z\to e^+e^-$ and $Z\to \mu^+\mu^-$ we expect 700 events per year for an integrated luminosity at 100 fb$^{-1}$. For the process $e^+e^-\to\tilde{\nu}_1^c\tilde{\nu}_2^c$ we consider two possibilities. In the first $\tilde{\nu}_1^c\to Z\tilde{\nu}_2^c$ is dominant decay mode, and in the second $\tilde{\nu}_1^c\to \tilde{\nu}_1^c$ is dominant. In the first case we expect 90 mono-Z decaying into $e^+e^-$ and $\mu^+\mu^-$, with large missing energy. In the second case we expect 1500 $e^+e^-$ pairs which are not coplanar, with large missing energy.
V. CONCLUSION

The RS-LSP scenario should be seriously considered as an alternative to the neutralino-LSP scenario. Obviously in the first case decay chains of the supersymmetric particles drastically differ from those of the second case. Therefore the RS-LSP scenario should be taken into account in the SUSY search programs of future colliders. In this paper the process $e^+e^- \rightarrow \tilde{\nu}\tilde{\nu}$ has been analyzed. The associated sneutrino-squark production at future lepton-hadron colliders was considered in [10].

Acknowledgments

This work is partially supported by Turkish State Planning Organization under the grant no DPT2002K120250 and Abant Izzet Baysal University Research Found. Authors are grateful to A. Senol and A.T. Tasci for useful discussions.

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FIG. 1: Partial decay width for $\bar{\nu}_1^c \to Z \bar{\nu}_2^c$ versus $\tilde{m}_2$ with $\tilde{m}_1 = 200$ GeV (upper line) and $\tilde{m}_1 = 150$ GeV (lower line).

FIG. 2: Partial decay width for $\bar{\nu}_1^c \to \bar{w} e$ versus $\tilde{m}_w$ (upper line) and for $\bar{\nu}_1^c \to \tilde{Z} \nu_e$ versus $\tilde{m}_Z$ (lower line) with $\tilde{m}_1 = 200$ GeV in both cases.
FIG. 3: Partial decay width for $\tilde{\nu}_1^\tau \rightarrow \tilde{\tau}_e$ (upper line) and $\tilde{\nu}_1^\tau \rightarrow \tilde{\tau}_\nu$ (lower line) versus $\sin \varphi$ with $\tilde{m}_1=200$ GeV and $\tilde{m}_\nu=\tilde{m}_Z=150$ GeV.

FIG. 4: Production cross section for $e^+e^- \rightarrow \tilde{\nu}_1^\tau \tilde{\nu}_1^\tau$ versus $\tilde{m}_1$ at $\sqrt{s}=500$ GeV.
FIG. 5: Production cross section for $e^+e^- \rightarrow \tilde{\nu}_2^c \tilde{\nu}_2^c$ versus $\tilde{m}_2$ at $\sqrt{s}=500$ GeV.

FIG. 6: Production cross section for $e^+e^- \rightarrow \tilde{\nu}_1^c \tilde{\nu}_2^c$ versus $\tilde{m}_1$ with $\tilde{m}_2 = 0$ at $\sqrt{s}=500$ GeV.

FIG. 7: Production cross section for $e^+e^- \rightarrow \tilde{\nu}_1^c \tilde{\nu}_2^c$ versus $\sin \varphi$ with $\tilde{m}_1=200$ GeV and $\tilde{m}_2=0$. 