# NEUTRINO AND LIGHTEST SUPERSYMMETRIC PARTICLE PROPERTIES IN N=1 AND N=2 SUPERSYMMETRIC MODELS

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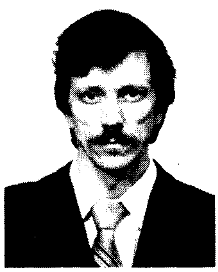

ABSTRACT

In the first Section the properties of neutrino (mass, decays, magnetic dipole moment) are studied in R-even and R-odd N=1 supersymmetric approach. It is shown that in R-odd SUSY models the mass of unstable $\tau$-neutrino can be in the following range: 1 MeV$< m_{\nu_\tau} <$ 35 MeV.

In the second Section the viability of the existence of light photino is studied in the N=2 SUSY model. The see saw mechanism is suggested for the light photino mass generation.

## 1. THE PROBLEMS OF NEUTRINO MASS AND DECAYS IN THE R-EVEN AND R-ODD MODELS OF N=1 SUPERSYMMETRY

From the viewpoint of the naturality of theory one might expect zero neutrino masses but, on the other hand, the masslessness a particle requires the existence of a rigorous local gauge symmetry. Since such a symmetry is unknown in the case with neutrino, one should not expect its mass to be zero either. According to the 1st experiment conducted in 1980, which gave an evidence for a nonzero mass of electron antineutrino, 20 eV$<m_\nu<$45 eV. All other experimental results[1] yield just the upper bounds on the values of masses for known neutrinos

$$m(\nu_e) < 18 \text{ eV} \qquad m(\nu_\mu) < 250 \text{ KeV},$$
$$m(\nu_\tau) < 35 \text{ MeV}, \qquad m(\nu_e) < 27 \text{ eV}. \qquad (1)$$

The data on the 2nd momentum of neutrino emission through the explosion of the Supernova 1987A from Magellanic cloud registrated with detectors Kamiokande II and IMB do not contradict the idea on the existence of massive neutrino states ($\nu_e$) in the $0 < m_\nu < 20$ eV[2] range.

Additional bounds on the masses of all sorts of neutrinos can be obtained by reffering to cosmological estimates. A universally recognized standard model giving the true description of the Universe evolution is considered to be the model of Big Bang (MBB). According to the MBB estimates, the mass sum of all sorts of neutrinos does not exceed 100 eV. One should note that these cosmological bounds are no longer valid if neutrino decays sufficiently fast. The MBB suggests that a stable neutrino is lighter than 100 eV or heavier 4 GeV, while the mass of an unstable one may range for 100 eV to 4 GeV providing its lifetime satisfies the conditions[3]

$$[m(\nu_i)]^2 \tau(\nu_i) \leqslant 2\cdot 10^{20} \text{ eV}^2 \text{ sec, if } m(\nu_i) \leqslant O(1) \text{ MeV},$$
$$(2)$$
$$[m(\nu_i)]^{-4} \tau(\nu_i) \leqslant 1.5\cdot 10^{-22} \text{ eV}^{-4} \text{ sec, if } m(\nu_i) \geqslant O(1) \text{ MeV},$$

In the standard SU(2)xU(1) theory of electroweak interactions, $\nu$ is massless because the simple higgs structure of the theory leads to the global symmetry corresponding to the conservation of leptonic number. A massive neutrino appears unstable as a rule and therefore the study of possible ranges of neutrino mass becomes expedient from the viewpoint of cosmological bounds (2). In this light, it is sufficient to study two and three-particle neutrino decay channels possible in the extended versions of SM:

$$\nu_i \to \nu_j + \nu_k + \nu_\ell ; \quad \nu_i \to \nu_j + e^+ + e^- ;$$
$$\nu_i \to \nu_j + \gamma ; \quad \nu_i \to \nu_j + \gamma + \gamma . \tag{3}$$

In the model of horizontal gauge interactions the decay $\nu_i \to 3\nu_j$ due to flavour-changing neutral currents occuring at the tree level is possible. The estimate of the width for the decays $\nu_\mu \to \nu_e \nu_e \bar{\nu}_e$ and $\nu_\tau \to \nu_j \nu_k \bar{\nu}_k$ ($\nu_\tau \to \nu_\mu \nu_\mu \bar{\nu}_\mu$, $\nu_\mu \nu_e \bar{\nu}_e$, $\nu_e \nu_\mu \bar{\nu}_\mu$, $\nu_e \nu_e \bar{\nu}_e$) together with experimental (1) and cosmological bounds (2) yields the following bounds on the neutrino mass[4)]: $m(\nu_\mu) > 90$ MeV, $m(\nu_\tau) > 70$ MeV (for $M_H \gtrsim 1$ TeV) the former ones excluding decay modes (3). On the contrary, in the L-R symmetric $SU(2)_L x SU(2)_R x U(1)$ gauge model the situation may vary, i.e., the decay $\nu_\tau \to 3\nu$ is possible and the mass of $\nu_\tau$ can be observed in the 1 MeV$< m(\nu_\tau) <$35 MeV range.

If $\tau$-onic neutrino is heavier than 1 MeV, the decay $\nu_\tau \to \nu_e e\bar{e}$ is possible.

Applying cosmological bound (2) we obtain the following lower bound on the $\nu_\tau$ mass: $m(\nu_\tau) > 5$ MeV. However, there exists an additional astrophysical constraint[4)]

$$\tau(\nu) < 10^3 \div 10^4 \text{ sec} \tag{4}$$

on the decay $\nu_\tau \to \nu_j + e^+ + e^-$ and also on radiation decays, which brings the lower bound on the neutrino mass to $m(\nu_\tau) > 10$ MeV. It should be noted that condition (4) is a more stringent constraints than (2) on this and some other decays.

The combined analysis carried out in ref.[4)] in the light of cosmological and astrophysical constraints imposed on neutrino decays in the frames of the SM versions extended nonsupersymmetrically leads to the following results: muon and electron neutrinos should be light, $m(\nu_e)<18$ eV, $m(\nu_\mu)<100$ eV and $\tau$-neutrino may be relatively heavy, i.e., its mass should lie in the MeV-energy range.

Now, we study the neutrino properties, in the models of simple N=1 supergravity.

As in the models of simple SUSY and SUGRA the mass matrix of scalar SUSY-particles ($\tilde{\ell}$, $\tilde{q}$) has a nondiagonal form: not only particles mix in the generations but also do their left- and right handed states. Such mixings can enhance essentially processes taking place with SUSY particle exchange in a loop, e.g., radiative neutrino decays.

The mass matrix of charged scalar leptons and scalar quarks have in N=1 supergravity the forms

$$\begin{pmatrix} m_L^2 + m_\ell^+ m_\ell & A m_\ell^+ m_{3/2} \\ A m_\ell m_{3/2} & m_R^2 + m_\ell m_\ell^+ \end{pmatrix}; \quad \begin{pmatrix} m_L^{D^2} & \delta m_{3/2} T^+ m^D \\ \delta m_{3/2} m^D T & m_R^{D^2} \end{pmatrix}, \qquad (5)$$

where $m_\ell$ is the mass matrix of ordinary charged leptons, $m^D$ - ordinary quarks, $m_{L,R} \sim m_{3/2}$ ((5) is a 6x6 matrix). It is not a problem to diagonalize matrix (5). Note that here we come accross the mass degeneration of scalar leptons over generations because $m_\ell << m_{3/2}$. Besides, the flavour changing transition amplitudes are suppressed essentially. The situation can be remedied by introducing a right-handed neutrino in which case the superpotential will contain additional terms hLHN and MNN (N is the supermultiplet of right-hand neutrinos). The former coupling yields the Dirac neutrino mass $m_\nu^D = \bar{\nu}_L \nu_R h\langle H_o\rangle$ and the latter yields the Majorana one. As a result, following the mechanism of "see-saw" type one may construct a Majorana neutrino of mass $m_\nu \sim (m_\nu^D)^2/M$ and a heavy right-handed neutrino of mass M. In this case the mass matrix of left-handed scalar leptons will vary to take the form

$$\tilde{\ell}_L^+ (m_L^2 + m_\ell^+ m_\ell + c m_\nu^{D+} m_\nu^D)\tilde{\ell}_L \simeq \tilde{\ell}_L^+ (m_L^2 + c m_\nu^{D+} m_\nu^D)\tilde{\ell}_L, \qquad (6)$$

which provides a sufficiently large difference between the masses of left-handed scalar leptons of different generations. Parameter C can be calculated by solving the system of renormalization group equations for superpotential parameters. In our case, we follow ref. [4] and use the value of $C \simeq 10^{-1}$–$10^{-2}$. With a right-handed neutrino introduced into the theory, one may have an illusion about an essential increase of the radiative neutrino decay width (see fig.1a). However, the amplitudes of the processes are still proportional to the neutrino mass: the Yukawa coupling $\bar{\nu}_R \tilde{\ell}_L \tilde{H}$ is proportional to $m_\nu^D$. A physically light neutrino contains $\nu_R$ of weight $m^D/M$ and the transition amplitude $\nu \to \nu\gamma$ is proportional to $m_\nu^D m_\nu^D/M \simeq m_\nu$.

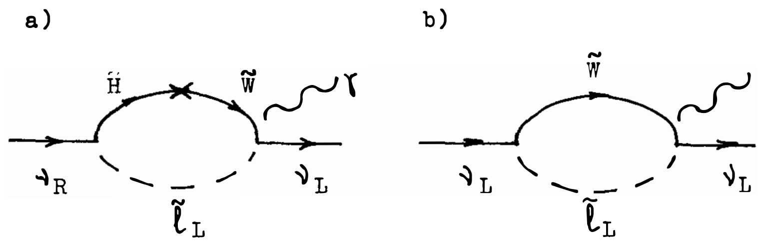


Fig. 1. One-photon neutrino decay (photon can be emitted by any virtual particle).

Mixing matrices of scalar leptons and scalar quarks have the forms[6]

$$S_{36}S_{35}S_{34}S_{26}S_{25}S_{24}S_{16}S_{15}S_{14}\begin{pmatrix} U & 0 \\ 0 & V \end{pmatrix};$$

$$B = P_{36}P_{35}P_{34}P_{26}P_{25}P_{24}P_{16}P_{15}P_{14}\begin{pmatrix} T & 0 \\ 0 & T \end{pmatrix}.$$

Here V is the lepton mixing matrix and U diagonalizes $m_\nu^D$. $S_{ik}$ and $P_{ik}$ diagonalizes (approximately) block-diagonal submatrices 1-i,2-i,3-i (i=4,5,6) of scalar lepton and scalar quarks mass matrices. The mixing angles $P_{ik}$ are given by the formulae

$$\mathrm{tg}\,2\alpha_{i,3+j} = \frac{2\delta m_{3/2} m_j^D a_{ji}}{m^2(\tilde{q}_i)_L - m^2(\tilde{q}_j)_R}, \quad i,j=1,2,3 \qquad (7)$$

$a_{ij}$ are the elements of the Kobayashi-Maskawa matrix.

Let us get down to considering radiative neutrino decays in the simple N=1 SUGRA model. The diagrams describing on one-photon neutrino decay are presented in fig.1. The decay width is expressed as:

$$\Gamma(\nu_i \to \nu_j \gamma) = \frac{\alpha G_F^2 [m(\nu_i)]^5}{64\pi^4} \left(\frac{m_W}{m_{\tilde{W}}}\right)^2 \left[\sum_a V_{ia} V_{ja} \left(F_1(x_a) + \frac{m_W}{m_{\tilde{W}}} F_2(x_a)\right)\right]^2, \quad (8)$$

where $m_{\tilde{W}}$ is the wino mass, $x_a = m^2(\tilde{\ell}_a)/m^2_{\tilde{W}}$

$$F_1(x) = \frac{1}{(1-x)^2}(1-x+x\ell nx), \quad F_2(x) = \frac{1}{2(1-x)^3}\left(\frac{1}{2} - \frac{1}{2}x^2 + x\ell nx\right). \quad (8')$$

Functions $F_1$ and $F_2$ correspond to the contributions from diagrams 1a and 1b, respectively. Applying limitations (2), (4) to expression (8) we obtain that, depending upon the value of $m_{\tilde{W}}$, $\nu_\tau$ neutrino may be heavier than 15÷20 MeV. As differed from the results of Section 1 of the reference[6], the N=1 SUGRA model allows a one-photon $\tau$-onic neutrino decay with the value of $m(\nu_\tau)$ ranging from 20 MeV$<m(\nu_\tau)<$35 MeV.

It would be quite interesting to consider the problem of neutrino mass origin in supersymmetric models with R-parity breaking $R=(-1)^{2S+3B+L}$, where S is spin, B is the baryonic number and L is the leptonic one. If R-number is conserved, lightest supersymmetric particle (LSP) is stable. In theories with R-parity violated, LSP decays into ordinary particles, for example, if $m_{\tilde{\gamma}} \gg m_\nu$, $\tilde{\gamma} \to \nu\gamma$. R parity can be violated spontaneously by providing scalar neutrinos with nonzero vacuum expectation values and/or/by adding terms violating the leptonic number into the superpotential. In these models, neutrino mixing with neutral Higgsino, gaugino ($\nu \leftrightarrow \tilde{\gamma} \leftrightarrow \tilde{H}_1 \leftrightarrow \tilde{H}_2 \leftrightarrow \tilde{Z}$) takes place and if photino is a sufficiently light particle $m_{\tilde{\gamma}} \sim m_\nu$, mixing $\nu \leftrightarrow \tilde{\gamma}$ may prove to be rather strong[7]. Therefore radiative neutrino decays (see figs.2,3) are enhanced in R-odd SUSY models. Though the mixing angles of neutrinos with neutral gauginos, left- and right-handed scalar particles (leptons, quarks) are small, radiative neutrino decays are enhanced essentially because the GIM mechanism is not valid and the transition amplitudes become proportional to the masses of virtual particles. In the R-odd theory, the width of an one-photon decay has the form[7] (Fig.2):

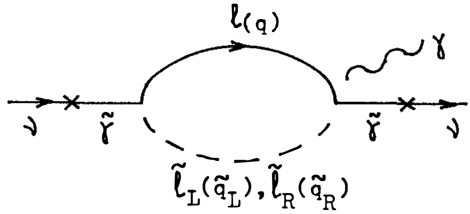


Fig. 2. One-photon neutrino decay in R-odd supersymmetry ($\tilde{\gamma}$ is a photino).

$$\Gamma(\nu_i \to \nu_j \gamma) = \frac{\alpha G_F^2 [m(\nu_i)]^3 m_W^2}{64\pi^4} \frac{m_W^2}{m_q^2} C^4_{\nu\tilde{\gamma}} (\sum_a B_{ia} B_{ja} F_1(x_a))^2. \tag{9}$$

With astrophysical limitation (4) applied, this yields the following lower bound on the neutrino mass: $m_\nu > 1$ MeV. ($C_{\nu\tilde{\gamma}} \sim 10^{-1}$).

If photino is lighter than neutrino, then R-odd supersymmetry allows the decay $\nu_i \to \gamma \tilde{\gamma}$ (see fig.3) having the width

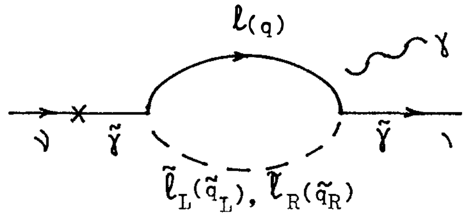


Fig. 3. Decay $\nu \to \gamma \tilde{\gamma}$.

$$\Gamma(\nu_i \to \gamma \tilde{\gamma}) = \frac{\alpha G_F^2 [m(\nu_i)]^3 m_W^2}{64\pi^4} \frac{m_W^2}{m_q^2} C^2_{\nu\tilde{\gamma}} (\sum_a B_{ia} B_{ja} F_1(x_a))^2. \tag{10}$$

Applying the same astrophysical limitation (4) to formula (10) we obtain the lower bound on the neutrino mass of about 1 MeV ($C_{\nu\tilde{\gamma}} \sim 10^{-2}$).

SUSY allows new possible neutrino decay channels (see figs. 1, 2, 3), in which 1 MeV $< m(\nu_\tau) <$ 35 MeV. The right-handed neutrino introduced into in N=1 SUGRA, removes degeneration of the masses of charged scalar leptons and therefore the width of radiative decays increases because the GIM mechanism does not work any longer. As a result, the $\nu_\tau$ mass can still lie in the MeV range, 20 MeV $< m(\nu_\tau) <$ 35 MeV and, correspondingly, the $\nu_\tau$ can decays into $\nu_{\mu,e}\gamma$. In R-odd SUSY models, the lower bound on the $\nu_\tau$ mass can be brought down to 1 MeV. Mixing between neutrinos and photinos, left- and right-handed scalar SUSY particles ($\tilde{\ell}$, $\tilde{q}$) enhances strongly neutrino radiative decays because the GIM does not work and the transition amplitudes become proportional to the masses of virtual particles.

Figure 2 presents the diagrams contributing into the neutrino magnetic dipole moment (mdm) in R-odd SUSY. Due to the mixing between left- and right-handed mass states of scalar SUSY-particles the value of the neutrino mdm is sufficiently large

$$\mu = \frac{G_F e m_q}{2\sqrt{2}\pi^2} \frac{m_W^2}{m_q^2} C^2_{\nu\tilde{\gamma}} (B_{ia} B_{ia} F_1(x_a)) \tag{11}$$

its maximum value being $\mu \sim 10^{-12} \div 10^{-13} \mu_B$.

## 2. THE SYMMETRIES OF N=2 SUPERSYMMETRIC MODELS AND THE MASS OF LIGHTEST SUSY PARTICLE

The conclusions of previous chapter about $\nu_\tau$-mass in R-odd N=1 SUSY model was highly connected to the fact that photino (the LSP) was light enough, $m_{\tilde{\gamma}} \ll 1$ GeV, or at least $m_{\tilde{\gamma}} \sim 1$ GeV, because only in this case the big mixing of $\nu_\tau \leftrightarrow \tilde{\gamma}$ is possible. It is very difficult to obtain the low-mass photino in N=1 SUSY theories without fine tuning. In this chapter we will try to solve this problem by building the "see saw" mechanism for photino mass in N=2 SUSY theory.

Let us consider N=2 SUSY theories constructed from vector multiplets (gauge superfields) and hypermultiplets (matter and Higgs superfields). In terms of N=1 supersymmetry, the vector multiplet contains vector $V \supset (A_\mu, \lambda)$ and chiral $\Phi \supset (M, \psi)$ superfields transforming as adjoint representation of gauge group, the hypermultiplet is constructed from two chiral superfields $X \supset (x, \tilde{X})$ and $Y \supset (y, \tilde{Y})$ (bosonic components are given first in parenthesis)[8]. Let the $\tilde{R}$-charge of superfied $\Phi$ be 1/2. Than for X and Y $\tilde{R}$-charges one has the relation $\tilde{R}_X+\tilde{R}_Y=1/2$.

In the N=2 supersymmetric extension of Standard Model it is difficult to obtain the realistic spectrum of fermionic matter, and commonly, to make the mirror fermions heavy enough, not to contradict the experimental limitations in the models with N=2 superfield content given in present paper this difficulty is due to the fact, that fermionic matter fields are singlets under the transformations of internal SU(2)-group.

The extension of gauge group is a way of the solution of fermion mass problem. For example, the extension of electroweak group to SU(4)xU(1) makes it possible to construct the lagrangian, in which the ordinary and mirror matter is in one representation. This allows to give the tree-level masses to fermions spontaneously, by VEVs of scalar field M[9].

Let us consider how the masses of gauginos will arise in such a theory. In the terms of N=1 superfields given below

$V(15,0)$; $V'(1,0)$; $\Phi(15,0)$; $\Phi'(1,0)$ - gauge fields;

$X_L(4, -\frac{1}{2}) = \binom{L}{L'}$; $Y_L(\bar{4},\frac{1}{2}) = \binom{L'}{L}$ - matter fields,

$X_1(4, -\frac{1}{2})$; $Y_1(\bar{4},\frac{1}{2})$; $X_2(4,\frac{1}{2})$, $Y_2(\bar{4}, -\frac{1}{2})$ - Higgs fields

(L and $\bar{L}$ denote $SU(2)_L$ and $SU(2)_R$ doublets, while L' and $\bar{L}'$ - corresponding mirror doublets) the N=2 supersymmetric lagrangian will have the form:

$$L = \left[\frac{1}{8g^2}\mathrm{Tr}WW + i\sqrt{2}g\ Y\Phi X\right]_F + \mathrm{h.c.}$$

$$+ \left[2\mathrm{Tr}\Phi^+ e^{2gV}\Phi e^{-2gV} + X^+ e^{2gV}X + Y^+ e^{-2gV^T}Y\right]_D. \qquad (12)$$

The internal symmetry group of N=2 lagrangian (12) is larger, than the N=1 one, including the group $SU(2)xU(1)_{\tilde{R}}$ (where $U(1)_{\tilde{R}}$ is the $\tilde{R}$-invariance). In addition N=2 theories posess the following discrete symmetries: (i) R-parity, determined in chapter 1, (ii) the mirror interchange symmetry $X \to Y$, $\Phi \to \Phi^T$, $V \to -V^T$. (iii) Mirror symmetry. As ordinary and mirror matter in SU(4)xU(1) theory is in the same representation, the transformations under M-parity will have the following form: $\Phi \to e^{2\pi i t_{3R}}\Phi e^{-2\pi i t_{3R}}$

$$V \to e^{2\pi i t_{3R}} V e^{-2\pi i t_{3R}}; \quad X \to e^{2\pi i t_{3R}} X; \quad Y \to -e^{2\pi i t_{3R}} Y,$$

where $t_{3R} = \frac{1}{2}\mathrm{diag}(0,0,1,-1)$.

In table 1 we list the M-quantum numbers and in table 2 the R-numbers of superfields under consideration.

Table 1. The M quantum numbers of superfields

| | L | $\bar{L}$ | L' | $\bar{L}'$ | $\varphi_A$ | $\varphi_B$ | $\varphi_C$ | $\varphi_D$ | $\varphi'$ | $V_A$ | $V_B$ | $V_C$ | $V_D$ | V' |
|---|---|---|---|---|---|---|---|---|---|---|---|---|---|---|
| M | + | + | − | − | − | + | + | − | − | + | − | − | + | + |

Table 2. R-parities of component fields

| | | | | | | R |
|---|---|---|---|---|---|---|
| $A_\mu$ | M | X | Y | $\tilde{X}_{1;2}$ | $\tilde{\tilde{Y}}_{1;2}$ | + |
| $\lambda$ | $\Phi$ | $\tilde{\tilde{X}}$ | $\tilde{Y}$ | $X_{1;2}$ | $Y_{1;2}$ | − |

We have used the following notation $\Phi = \begin{pmatrix}\varphi_A & \varphi_B \\ \varphi_C & \varphi_D\end{pmatrix}$ and similarly for V.

And finally the attractive property of N=2 SUSY theory is the fact, that it can be made finite by choosing matter field representations so, that the condition $C_2(G) = \sum_\sigma T(t^\sigma)$ be satisfied. Here G is the gauge group with generators t,

$\sigma$ is the index of representation, $T(t^{\sigma})\delta_{ij} = Tr(t_i^{\sigma}t_j^{\sigma})$, and $c_2(G)\delta_{i\ell} = f_{ijk}f_{\ell jk}$ ($f_{igk}$ are the group structure constants). Note, that the consideration of quark sector (and gauge group $SU(3)_c \times SU(4) \times U(1)$) makes the theory under consideration finite.

There are five types of soft-SUSY breaking operators the adding of which to lagrangian (12) does not destroy the finiteness of theory:

1) SUSY mass terms: $m_4 Tr[\varphi^2]_F$; $[X_iY_j]_F$; $[X_iX_j]_F$;
2) scalar masses $[X_iX_j]_A$+h.c.; $Tr[\varphi^2]_A$+ h.c.;
3) the set of scalar mass terms $\alpha M^+M + \gamma_i \tilde{X}_i^+\tilde{X}_i + \beta_i \tilde{Y}_i \tilde{Y}_i$ such that $-\alpha = \beta_i + \gamma_i$;
4) following operators $-\frac{1}{2}m\lambda\lambda - \sqrt{2}\, img \sum_i \tilde{Y}_i M X_i$+h.c.$-|m|^2 M^+M$;
5) trilinear scalar couplings

$$k^{ijk}(\tilde{X}_i\tilde{X}_j\tilde{X}_k + \tilde{Y}_i^+\tilde{Y}_j^+\tilde{Y}_k^+) + h.c.; \quad k'^{ijk}(\tilde{X}_i\tilde{X}_j\tilde{Y}_k + \tilde{Y}_i^+\tilde{Y}_j^+\tilde{X}_k^+) + h.c.$$

with restrictions given in references[9,10].

In order to obtain the realistic fermion mass spectrum let us add to lagrangian (12) the soft SUSY-breaking mass terms of type 1, breaking N=2 SUSY, and of type 3, breaking supersymmetry completely. This terms will give the following contributions to the scalar potential of model

$$V_{SOFT} = -m_i^2|\tilde{X}_i|^2 - \tilde{m}_i^2|\tilde{Y}_i|^2 + m^2|\tilde{X}_{L_Q}|^2 + \tilde{m}^2|\tilde{Y}_{L_Q}|^2 + m_o^2 M_a^* M_a +$$

$$+ i\sqrt{2} g m_4 Y_j M^+ X_j + h.c. \qquad (i=1,2) \qquad (13)$$

This potential will induce the VEVs of Higgs scalars and field M. Let the non-zero VEVs be $(\tilde{X}_1)_1 = v_1$; $(\tilde{X}_2)_2 = v_2$; $(\tilde{Y}_1)_3 = V_1$; $(\tilde{Y}_2)_4 = V_2$ $\langle M_{13}\rangle_o = \tilde{v}_1$; $\langle M_{24}\rangle_o = \tilde{v}_2$. We will use further the notation $gV_2 = M_R$. As the value of v determines the masses of W and Z bosons, while V determines the scale of $SU(2)_R$ breaking, it is natural set up the following hierarchy $|v| \ll |\tilde{v}| \ll |V|$.

The finiteness of theory allows one to give masses to LSP radiatively, without adition of soft mass terms of type 4.

In addition, in N=2 SUSY theories there is a possibility of spontaneous breaking of $\tilde{R}$-symmetry, by nonzero VEVs of mirror scalar gauge bosons M.

We will try to obtain the low-mass photino by "see saw" mechanism in the N=2 SUSY model, assuming that it is LSP. In such theories the new energy scale can arise, which is connected with the breaking of internal global SU(2) symmetry.

In the model under consideration we call "photino" the following combination of gauginos:

$$\lambda_\gamma = \left[\left(-\frac{1}{\sqrt{3}}\lambda_8 + \frac{2}{\sqrt{6}}\lambda_{15}\right)\sin\varphi + \lambda'\cos\varphi\right]\cos\theta + \lambda_3\sin\theta$$

where:

$$\sin\varphi = \frac{g'}{\sqrt{g^2+g'^2}}; \qquad \sin\theta = \frac{g'}{\sqrt{g^2+2g'^2}}.$$

The mass terms for photino will arise due to the one-loop diagrams.

Let us consider different cases:

(i) Let the main contributions to photino mass matrix come from diagrams with lepton exchange. In this case the mass matrix will have the form

$$L_m = -(\bar{\lambda}_\gamma, \bar{\Psi}, \bar{x}_2)\begin{bmatrix} a & b & 0 \\ b & \frac{m_4}{2} & -\frac{gV_2}{2} \\ 0 & -\frac{gV_2}{2} & 0 \end{bmatrix} \cdot \begin{bmatrix} \lambda_\gamma \\ \Psi \\ x_2 \end{bmatrix}, \qquad (14)$$

where $\Psi = \frac{1}{\sqrt{2}}\begin{bmatrix} \varphi_{12} - i\varphi_{11} \\ \varphi_{12} + i\varphi_{11} \end{bmatrix}$, $a = a_1 + a_2$,

$$a_1 = \frac{e^2 m_{l_m}\sin 2\vartheta_{l_m l}}{2\pi^2}\ln\frac{\tilde{M}_2}{\tilde{M}_1};$$

$$a_2 = \frac{e^2 m_{\ell_m} \sin 2\vartheta_{\tilde{\ell}_m}}{2\pi^2} \ln \frac{\tilde{M}_4}{\tilde{M}_3} ; \qquad a \approx \frac{e^2 m_\ell m_{\ell_m} m_4}{\pi^2 (m^2 - \tilde{m}^2)} \ln \frac{\tilde{M}_2 \tilde{M}_4}{\tilde{M}_1 \tilde{M}_3} ;$$

$$b = \frac{e g m_e}{8\sqrt{2}\pi^2} \ln \frac{\tilde{M}_4}{\tilde{M}_3} ; \qquad \sin 2\vartheta_{\tilde{\ell}_m} \approx \frac{2 m_4 m_b}{m^2 - \tilde{m}^2} \tag{15}$$

$\sin 2\vartheta_\ell$ can be obtained from $\sin 2\vartheta_{\tilde{\ell}_m}$ by replacing $\ell_m \to \ell$

In eq. (15) $\tilde{M}_{1,2,3,4}$ are the mass eigenvalues obtained after diagonalization of mass matrix of ordinary and mirror scalar leptons. The elements of mass matrix in (14) were calculated in the approximation $m, \tilde{m}, \tilde{M} \gg m_{\ell_m} \gg m_\ell$. In eq. (15) a is the contribution from diagram given on fig.4 (the mixing of left- and right-handed scalar leptons is due to the VEV of field M see eq. (13)), while b is the contribution from diagram on fig. 5.

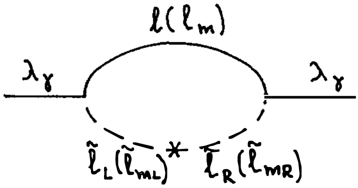


Fig. 4. Diagram for photino mass.

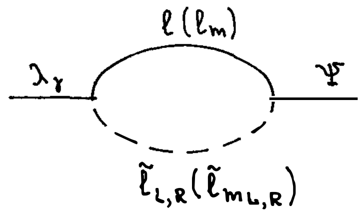


Fig. 5. Diagram for photino-gaugino mixing.

The diagonalization of mass matrix (14) gives the following results: the light fermion which is almost photino is the combination

$$\lambda_o = \frac{1}{N}\left[\lambda_\gamma + 2bm_4 \frac{\sin^2\varphi_1(m_4^2 + M_R^2) - \cos^2\varphi_1(2am_4 + M_R^2)}{(m_4^2 + M_R^2)(2am_4 + M_R^2)}\Psi - \right.$$

$$\left. - \frac{bm_4 \sin 2\varphi_1 (m_4^2 + 2M_R^2)}{(m_4^2 + M_R^2)(2am_4 + M_R^2)} X_2 \right] \qquad (16)$$

with mass eigenvalue

$$\mu_o = a - \frac{2\cos^2\varphi_1 b^2 m_4}{m_4^2 + M_R^2} + \frac{2\sin^2\varphi_1 b^2 m_4}{2am_4 + M_R^2} . \qquad (17)$$

1/N is the normalization coefficient and $\sin 2\varphi_1 = -\frac{2M_R}{\sqrt{m_4^2 + 4M_R^2}}$.

Let us perform some numerical estimates for illustration. When choosing mass parameters, we must take into account, that they must satisfy the equations:

$$M_R^3 + 2(m_{\ell_m}^2 - 2\tilde{m}_{1;2}^2)M_R - 4m_{\ell_m} m_4 m = 0 \qquad (18)$$

$$m_{\ell_m}^3 + m_o^2 m_{\ell_m} - M_R m_u m = 0 \qquad (19)$$

comming from the minimization of scalar potential of model.

If the main contributions to matrix (14) come from diagrams with electron exchange, with the following values for mass parameters: m =0.5 MeV, $m_{\ell_m}$ =100 GeV, $m_4 = M_R = 1$ TeV, $\tilde{M}_{2,4}$=3,4 TeV, $\tilde{M}_{1,3}$=3 TeV, than the eq. (17) transforms into a

$$\mu_o \approx \frac{e^2 g^2}{736\pi^4} \frac{m_{\ell_m}^2}{M_R} \ln^2 \frac{m}{\tilde{m}} \qquad (20)$$

giving $\mu_o$=100 eV.

If the dominant diagrams for photino mass are the ones with $\tau$-lepton exchange, than assuming that $m_\tau$ =1800 MeV and $M_{\tau_m}$ =100 GeV (other parameters are as in previous case) one gets $\mu_0$=0.2 MeV;

(ii) If the main contributions to photino mass are from diagrams with quark exchanges, than photino will mix with gaugino $\varphi_{5-i4}$ and higgsino $X_1$ as well. So the resulting mass matrix will be of dimension 5x5. But if the t-quark diagrams are the dominant ones, neglecting b-quark contributions, the mass matrix will be reduced to the old 3x3 matrix (14) replacing $\varphi_{12-i11} \to \varphi_{5-i4}$, $X_2 \to X_1$ and $V_2 \to V_1$. If $m_t$=60 GeV and $m_{t_m}$=200 GeV (other parameters are the same) the photino mass will be $\mu_0$=3 MeV.

(iii) If it would be possible to take the masses of mirror fermions heavy enough ($\sim O$(1 TeV)), than due to the radiative mechanism considered in present paper, the photino mass may turn out to be in GeV range. For example, if $m_t$= =60 GeV, $m_{t_m}$=2 TeV, $m_4 \sim M_R$ = 20 TeV, $\widetilde{M}_{1,2,3,4} \sim 3$ TeV and assuming that the main contribution to photino mass comes from t-quark loop, one obtains $\mu_0$=2 GeV. But in the SU(4)xU(1) model it is impossible to obtain such a high scale mirror fermion masses. As the VEVs of field M giving masses to mirror fermions, at the same time break the gauge group SU(4)xU(1) down to $U(1)_{EM}$, the masses of these mirror fermions are related to W-boson masses[9]: $m^2_{u_m}+m^2_{d_m} \approx 2\xi M^2_W$, where $\xi$=1 when M-parity is unbroken and $\xi$=1.5 for models with broken M-parity. Therefore, if one wants to obtain heavy mirror fermions, the considered SU(4)xU(1) model must be generalized, say, in the frames of gauge group with bigger rank, for example, $E_7$ coming from D=10 heterotic string compactifications on Calabi-Yau manifold (or orbifolds) with SU(2)-group of holonomy (or discrete subgroup of SU(2)).

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