Published by Institute of Physics Publishing for SISSA

Received: *February 8, 2006*
Accepted: *April 5, 2006*
Published: *April 19, 2006*

# Masses of fermions in supersymmetric models

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Abstract: We consider the mass generation for the usual quarks and leptons in some supersymmetric models. The masses of the top, the bottom, the charm, the tau and the muon are given at the tree level. All the other quarks and the electron get their masses at the one loop level in the Minimal Supersymmetric Standard Model (MSSM) and in two Supersymmetric Left-Right Models, one model uses triplets (SUSYLRT) to break $SU(2)_R$-symmetry and the other use doublets(SUSYLRD).

Keywords: Supersymmetry Phenomenology, Discrete and Finite Symmetries, Supersymmetric gauge theory.

## Contents





## 1. Introduction

One of important issues for particle physics is the small mixing angles in the charged fermion sector and the hierarchy of quark and lepton masses [1].This issue has brought into the focus due to the success of the Standard Model (SM) in describing the available experimental data except such mass spectrum and due to the recent experiments are showing it is likely neutrinos have such mass hierarchy but their mix pattern differs from that of the quarks, e.g., the 2-3 lepton mixing angle is close to the maximal value while the analogous quark mixing angle is small ($\theta^q_{23} \sim 2^\circ$). The fermions mass spectrum is an aspect of a problem named as fermion flavour structure [2] which includes the suppression of flavour change neutral current, strong CP-problem, etc.

Approaches based on Supersymmetry (SUSY) have been proposed in order to explain the values of these masses and the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements. They are guided by the pattern of hierarchy and one pattern used is the following *horizontal hierarchy*:

$$
\begin{aligned}
m_t : m_c : m_u &\sim 1 : \varepsilon_u : \varepsilon_u^2 \quad \varepsilon_u \simeq \frac{1}{500}, \\
m_b : m_s : m_d &\sim 1 : \varepsilon_d : \varepsilon_d^2 \quad \varepsilon_d \simeq \frac{1}{50}, \\
m_\tau : m_\mu : m_e &\sim 1 : \varepsilon_e : \varepsilon_e^2 \quad \varepsilon_e \simeq \frac{1}{50},
\end{aligned}
\tag{1.1}
$$

where $m_u$ and $m_d$ are current quark masses. This pattern has suggested, e.g., the masses of different families are generated in different stages of chiral symmetry breaking: at the first stage only $t$ and $b$ quarks acquire mass and there is no mixing, at the second stage $c$ and $s$ get mass and there is a mixing between the third and second family and in the end $u$ and $d$ quarks get their masses. This can be realized by the *radiative mass generation mechanism* where the lowest quarks are prevented to acquire mass at tree level [3−5]. However this mechanism in supersymmetric models gives rise to the flavour changing problem in the loop that generates the masses. In order to avoid this problem a horizontal flavour symmetry has been proposed within supersymmetric extensions of SM [6] and unified $SO(10)$ model [7]. The last one assumes a pattern where the first family instead of third family plays a unique role and named it as *inverse hierarchy pattern*. This is inspired by the fact at GUT scale running masses of electron, u and d quarks are not strongly split.

Another pattern shows us two different scales for the masses of quarks, one is at MeV scale

$$m_u \sim 1 - 5\; MeV, \qquad m_d \sim 3 - 9 MeV, \qquad m_s \sim 75 - 170 MeV, \tag{1.2}$$

while the other is at GeV scale:

$$m_c \sim 1,15 - 1,35\; GeV, \qquad m_t \sim 174\; GeV, \qquad m_b \sim 4,0 - 4,4\; GeV. \tag{1.3}$$

This point of view has implications for nuclear physics. Due to $u$, $d$ and $s$ quarks are lights one is allowed to build an effective field theory as an expansion on masses of light quarks of the underlying theory. The Chiral Perturbation Theory (ChPT) [8] is the prototype of this approach. It respects all principles of the underlying theory but with effective degrees of freedom instead of quarks degrees of freedom. A model independent description of dynamics [9] and structure of nucleons [10] above MeV scale is obtained.

We explore the implications of this picture in MSSM and Left-Right Supersymmetric Model (LRSM). In the framework of SUSY models the Higgs mechanism can be extended by increasing the number of scalar particles, as a consequence the number of vacuum expectation values also is increased and one has the possibility of two scale of masses for the case of two scalar particles. However there is no constrains to the size of masses and it is likely they could be at the same scale.

We also study the mechanism of radiative mass generation in the last pattern. In this case an additional symmetry [3, 4] suppresses the mass generation of ligth quarks ($u$, $d$ and $s$ ) at tree level while the the heavier ones acquire masses and there is mixing of $t$, $b$, and $c$ quarks masses. For the leptons the same description is applied and a low value for the masses of light quarks and light lepton is obtained.

The outline of this work is as follows. The section 2 describes how the additional discrete symmetry $\mathcal{Z}_2'$ is introduced into the framework of MSSM in order to prevent the light quarks and the electron to acquire mass at tree level. The radiative mechanism is described in section 3, and $u$, $d$ and $s$ quarks together with the electron acquire mass at 1-loop level. We also show that our results are still valid in two supersymmetric left-right models. Our conclusions are found in the last section. All the details of the models (conventions) and computations of mass matrices are in the appendices.



JHEP04(2006)031

## 2. MSSM and $\mathcal{Z}_2'$ symmetry

In the MSSM [11], which the gauge group is $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$, let $\hat{L}$ ($\hat{l}^c$) denotes left-handed (right-handed) leptons,[1] $\hat{Q}(\hat{u}^c, \hat{d}^c)$ left-handed (right-handed) quarks and $\hat{H}_1, \hat{H}_2$ are the Higgs doublets respectively (a summary is in appendix A).

The fermion mass comes from the following terms of the superpotential (eq. (A.8)):

$$W = -\left( y^l_{ab} L_a H_1 l^c_b + y^d_{ij} Q_i H_1 d^c_j + y^u_{ij} Q_i H_2 u^c_j + h.c. \right) + \cdots , \tag{2.1}$$

where $y^l_{ab}$, $y^d_{ij}$ and $y^u_{ij}$ are the yukawa couplings of Higgs with leptons families, "down" sector quarks and "up" sector quarks respectively and ...stands for other terms which we are not concerned here. The family indices $a$ and $i$ run over $e, \mu, \tau$ and $1, 2, 3$, respectively.

Based on eq. (A.10), we get the following non-diagonal mass matrices $M^{l,d,u}_{ij}$ :

$$\begin{aligned} M^u_{ij} &= \frac{y^u_{ij}}{\sqrt{2}} v_2 (u_i u^c_j + h.c.), \\ M^d_{ij} &= \frac{y^d_{ij}}{\sqrt{2}} v_1 (d_i d^c_j + h.c.), \\ M^l_{ab} &= \frac{y^l_{ab}}{\sqrt{2}} v_1 (l_a l^c_b + h.c.). \end{aligned} \tag{2.2}$$

Where all the fermions fields are still Weyl spinors. The fields in the parenthesis define the basis to get the mass matrix. We can also rewrite all the equations above as

$$M^\psi_{ij} = -\left( \bar{\psi}_{iL} m_{ij} \psi_{jR} + h.c. \right), \tag{2.3}$$

where $\psi_i$ [2] is the Dirac spinor.

Therefore, the "down" quark sector ($d, s$ and $b$ quarks) as well as the $e, \mu$ and $\tau$ will have masses proportional to the vacuum expectation value $v_1$, whereas the "up" sector will have masses proportional to $v_2$. Note that the neutrinos remains massless due to lepton-number conservation, but we know that neutrinos have masses. In order to give mass to neutrinos one has to introduce $R$ -parity violating term $W_{RV}$ of eq. (A.8). We will focus our attention to the quark and lepton sector and for the case of neutrinos the reader is invited to look at ref. [12].

Although the Higgs mechanism and SUSY allow two different scale of masses, there is no underlying principle to keep them different from each other. The fact that $m_u, m_d, m_s$ and $m_e$ are many orders of magnitude smaller than the masses of others fermions may well be indicative of a radiative mechanism at work for these masses as considered at [3, 4].

The key feature of this kind of mechanism is to allow only the quarks $c, b, t$, and the leptons $\mu$ and $\tau$ have Yukawa couplings to the Higgs bosons. It means to prevent $u, d, s$ and $e$ from picking up tree-level masses, all one needs to do at this stage is to impose the

[1] $c$ stands for charge conjugation

[2] $\psi_i$ indicate any charged fermion, the translation of two-component formalism into four-component formalism can be found in [11, 13, 18]

following $\mathcal{Z}'_2$ symmetry on the Lagrangian

$$\widehat{d}^c_2 \longrightarrow -\widehat{d}^c_2, \qquad \widehat{d}^c_3 \longrightarrow -\widehat{d}^c_3, \qquad \widehat{u}^c_3 \longrightarrow -\widehat{u}^c_3, \qquad \widehat{l}^c_3 \longrightarrow -\widehat{l}^c_3, \tag{2.4}$$

the others superfields are even under this symmetry. On ref. [4], only the electron and the first quark family don't pick up tree-level masses.

After the diagonalization procedure of mass matrices of fermions (see appendix B), we can write $M_{diag} = \mathrm{diag}(m_{f_1}, m_{f_2}, m_{f_3})$ where

$$m_{f_1} = \frac{1}{2}\left(t_f + r_f\right), \qquad m_{f_2} = \frac{1}{2}\left(t_f - r_f\right), \qquad m_{f_3} = 0, \tag{2.5}$$

with $t_f$, $r_f$ are given at eq. (B.4) and $f$ runs over fermions. Taking $M_{diag}$ into account we can do the following phenomenological identification:

$$m_{u_1} \equiv m_t, \qquad m_{d_1} \equiv m_b, \qquad m_{u_2} \equiv m_c, \qquad m_{l_1} \equiv m_\tau,\ m_{l_2} \equiv m_\mu\,. \tag{2.6}$$

In order to fit the experimental data we make the following choices into eq. (2.5):

$$\begin{aligned}
&t_l = m_\tau + m_\mu, \qquad t_u = m_t + m_c, \qquad t_d = m_b, \\
&(y^l_{13})^2 u_l - 2y^l_{12} y^l_{13} v_l + y_l^2 + (y^l_{12})^2 z_l = \frac{1}{4}\left[(m_\tau + m_\mu)^2 - (m_\tau - m_\mu)^2\right], \\
&(y^u_{13})^2 u_u - 2y^u_{12} y^u_{13} v_u + y_u^2 + (y^u_{12})^2 z_u = \frac{1}{4}\left[(m_t + m_c)^2 - (m_t - m_c)^2\right].
\end{aligned}$$

Thus the quarks $u, d, s$ and the electron come about be massless due to $\mathcal{Z}'_2$ symmetry. It means a discrete symmetry like $\mathcal{Z}'_2$ is protecting the Chiral symmetry to be broken in the $SU(3)_F$ sector.

Now, we want to show that a consistent picture with the experimental values of the Cabibbo-Kobayashi-Maskawa (CKM) matrix is obtained even in the presence of the $\mathcal{Z}'_2$ symmetry. The CKM matrix comes from the fact that the mass eigenstates of physical quarks are a mixture of different quarks eigenstates of symmetry and for three generation of quarks one has:

$$V_{CKM} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \tag{2.7}$$

where the matrix element $V_{ij}$ indicates the contribution of quark $(j\ )$ to quark $(i)$. The experimental values are [19]:

$$\begin{pmatrix} 0.9739 - 0.9751 & 0.221 - 0.227 & 0.0029 - 0.0045 \\ 0.221 - 0.227 & 0.9730 - 0.9744 & 0.039 - 0.044 \\ 0.0048 - 0.014 & 0.037 - 0.043 & 0.9990 - 0.9992 \end{pmatrix}. \tag{2.8}$$

As the quarks $t$ and $c$ get masses at tree-level their states can be mixed and we can write the eigenvector of "up" quark sector [3] as

$$E^u_L = \begin{pmatrix} \cos\theta & \sin\theta & 0 \\ -\sin\theta & \cos\theta & 0 \\ 0 & 0 & 1 \end{pmatrix}. \tag{2.9}$$

[3] $(t, c, u)^T = (u_1, u_2, u_3)^T E^{uT}_L$

For another hand, in the "down" quark sector only the quark $b$ get mass at tree-level and there is no mixing on this sector. Therefore we can write

$$E_L^d = I_{3\times 3}\,. \tag{2.10}$$

where $I_{3\times3}$ is the identity matrix $3\times 3$. Then, with eq. (2.9,2.10), we can get an expression to the CKM matrix as follows:

$$V_{CKM} = E_L^{u\dagger} E_L^d = \begin{pmatrix} \cos\theta & -\sin\theta & 0 \\ \sin\theta & \cos\theta & 0 \\ 0 & 0 & 1 \end{pmatrix}. \tag{2.11}$$

Comparing eqs. (2.8), (2.11), we can conclude that the $\mathcal{Z}_2'$ symmetry in the MSSM can explain the lower masses of the $u, d$ and $s$ quarks and also gives a hint about the mixing angles of quarks.

## 3. Radiative mechanism to the fermions masses

The discrete symmetry $\mathcal{Z}_2'$ has to be broken in order to allow the generation of fermions masses by radiative corrections and the most general soft supersymmetry breaking Lagrangian eq. (A.9) has already the following $\mathcal{Z}_2'$ breaking terms

$$\begin{aligned} \mathcal{L}_{soft} =& \left[ \sum_{i=1}^{3} A^d_{i3} H_1 \tilde{Q}_i \tilde{d}^c_{3L} + \sum_{i=1}^{3} A^d_{i2} H_1 \tilde{Q}_i \tilde{d}^c_{2L} + \sum_{i=1}^{3} A^u_{i3} H_2 \tilde{Q}_i \tilde{u}^c_{3L} \right. \\ &+ \left. \sum_{a=1}^{3} A^l_{a3} H_1 \tilde{L}_a \tilde{l}^c_{3L} + h.c. \right] + \cdots, \end{aligned} \tag{3.1}$$

where $\cdots$ stands for other terms. It means that the squarks $\tilde{q}$ and $\tilde{q}^c$ will mix, the same will happen with the sleptons $\tilde{l}$ and $\tilde{l}^c$ [4]. We can write

$$\begin{aligned} \mathcal{L}_{soft} =& M_Q^2 \left( \tilde{u}_3^* \tilde{u}_3 + \tilde{d}_3^* \tilde{d}_3 + \tilde{d}_2^* \tilde{d}_2 \right) + M_u^2 \tilde{u}_3^{c*} \tilde{u}_3^c + M_d^2 \left( \tilde{d}_3^{c*} \tilde{d}_3^c + \tilde{d}_2^{c*} \tilde{d}_2^c \right) + M_L^2 \tilde{l}_3^* \tilde{l}_3 \\ &+ M_l^2 \tilde{l}_3^{c*} \tilde{l}_3^c + \left[ A^l_{33} v_1 \tilde{l}_3 \tilde{l}_3^c + A^u_{33} v_2 \tilde{u}_3 \tilde{u}_3^c + A^d_{33} v_1 \tilde{d}_3 \tilde{d}_3^c + A^d_{22} v_1 \tilde{d}_2 \tilde{d}_2^c + h.c. \right] \\ &+ \cdots \end{aligned} \tag{3.2}$$

For the case of the physical u-squark states, that we will denotated as $\tilde{u}_1, \tilde{u}_2$, it gives rise to the following eigenstates of mass as functions of symmetry eigenstates:[5]

$$\begin{aligned} \tilde{u}_1 &= \cos\theta_{\tilde{u}} \tilde{u}_3 + \sin\theta_{\tilde{u}} \tilde{u}_3^{c*}, \\ \tilde{u}_2 &= -\sin\theta_{\tilde{u}} \tilde{u}_3 + \cos\theta_{\tilde{u}} \tilde{u}_3^{c*}\,. \end{aligned} \tag{3.3}$$

[4] In supersymmetric theories the sfermions masses come from the scalar potetial given by $V = V_F + V_D + V_{soft}$ [13, 18], here we are not showing all the details

[5] Mixing between squarks of different generations can cause severe problems due to too large loop contributions to flavour changing neutral currents (FCNC) process. Due this fact we are ignoring intergenerational mixing.

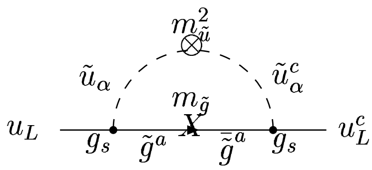


**Figure 1:** The diagram which gives mass to quark $u$ which does not apperar in the superpotential, $\tilde{g}$ is the gluino while $\tilde{u}$ is the squark.

Similar expressions to the d-squark ($\tilde{d}$), s-squark ($\tilde{s}$), selectron ($\tilde{e}$) and smuon ($\tilde{\mu}$) can be obtained. By another side, the interaction between quark-squark-gluino is given by:

$$\mathcal{L}_{q\tilde{q}\tilde{g}} = -i\sqrt{2}g_s\bar{T}^a(\tilde{u}^c_i\bar{u}^c_i\bar{\lambda}^a_C - \bar{\tilde{u}}^c_i u^c_i\lambda^a_C + \tilde{d}^c_i\bar{d}^c_i\bar{\lambda}^a_C - \bar{\tilde{d}}^c_i d^c_i\lambda^a_C) + \cdots . \tag{3.4}$$

We must remember that, mixing between sfermions of different generations is model dependent. Such mixing can cause severe phenomenological problems, by producing unacceptably large flavor changing neutral currents (FCNC) between ordinary quarks and leptons through 1–loop processes. There are three ways to suppress this FCNC [13]. The most popular way is assuming that the quark-squark mixing is flavor conserving. But on this case, $V_{td}, V_{ts}, V_{cb}$ and $V_{ub}$ (which are zero at tree level), remains zero after the radiative one loop correction for the quark mass matrices. In order to get a small values for these mixing angle we can use the mass insertion method [14]. Another possibility one can add higher-dimension (nonrenormalizable) operators at the superpotential, that arise from new physics at some scale $\Lambda$ [15−17]. This subject will be useful for further study.

The interaction, given at eq. (3.4), generate the radiative mechanism for the mass of the $u, d$ and $s$ quarks. On Fig.(1) we show the lowest order contribution. It was also shown in [3, 4] to current mass of up quark.

Notice that, all the mass insertion on this diagram came from the soft term, see eq. (3.2), while the two vertices come from eq. (3.4). Similar diagram can be drawn to the $d$ and $s$ quarks.

Following [4] we calculated their masses and we obtained:

$$\begin{aligned}
m_u &\propto \frac{\alpha_s \sin(2\theta_{\tilde{u}})}{\pi} m_{\tilde{g}} \left[ \frac{M^2_{\tilde{u}_1}}{M^2_{\tilde{u}_1} - m^2_{\tilde{g}}} \ln\left(\frac{M^2_{\tilde{u}_1}}{m^2_{\tilde{g}}}\right) - \frac{M^2_{\tilde{u}_2}}{M^2_{\tilde{u}_2} - m^2_{\tilde{g}}} \ln\left(\frac{M^2_{\tilde{u}_2}}{m^2_{\tilde{g}}}\right)\right], \\
m_d &\propto \frac{\alpha_s \sin(2\theta_{\tilde{d}})}{\pi} m_{\tilde{g}} \left[ \frac{M^2_{\tilde{d}_1}}{M^2_{\tilde{d}_1} - m^2_{\tilde{g}}} \ln\left(\frac{M^2_{\tilde{d}_1}}{m^2_{\tilde{g}}}\right) - \frac{M^2_{\tilde{d}_2}}{M^2_{\tilde{d}_2} - m^2_{\tilde{g}}} \ln\left(\frac{M^2_{\tilde{d}_2}}{m^2_{\tilde{g}}}\right)\right], \\
m_s &\propto \frac{\alpha_s \sin(2\theta_{\tilde{s}})}{\pi} m_{\tilde{g}} \left[ \frac{M^2_{\tilde{s}_1}}{M^2_{\tilde{s}_1} - m^2_{\tilde{g}}} \ln\left(\frac{M^2_{\tilde{s}_1}}{m^2_{\tilde{g}}}\right) - \frac{M^2_{\tilde{s}_2}}{M^2_{\tilde{s}_2} - m^2_{\tilde{g}}} \ln\left(\frac{M^2_{\tilde{s}_2}}{m^2_{\tilde{g}}}\right)\right],
\end{aligned} \tag{3.5}$$

where $m_{\tilde{g}}$ , $m_{\tilde{u}}$, $m_{\tilde{d}}$ and, $m^2_{\tilde{s}}$ are, respectively, the masses of the gluino, u-squark, d-squark and s-squark.

In order to obtain quark masses in agreement with the experimental limits [19] we set $m_{\tilde{g}} \approx 100$ GeV, $\sin(2\theta_{\tilde{u}}) \approx \sin(2\theta_{\tilde{d}}) \approx 10^{-3}$ and $\sin(2\theta_{\tilde{s}}) \approx 10^{-2}$.

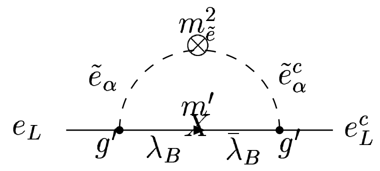


**Figure 2:** The diagram which gives mass to electron which does not apperar in the superpotential, $\tilde e$ is the selectron.

The electron couples with the gaugino $\lambda_B$ of $U(1)$ group as the following:

$$\mathcal{L}_{l\tilde l\tilde g} = -\frac{ig'}{\sqrt{2}}(2)\left(\tilde l^c_a \bar l^c_a \bar\lambda_B - \bar{\tilde l}^c_a l^c_a \lambda_B\right) + \cdots . \tag{3.6}$$

This allows the diagram of figure 2 to contribute to the electon mass. Therefore the electron mass is given by:

$$m_e \propto \frac{\alpha_{U(1)} \sin(2\theta_{\tilde e})}{\pi} m' \left[ \frac{M^2_{\tilde e_1}}{M^2_{\tilde e_1} - m'^2} \ln\left(\frac{M^2_{\tilde e_1}}{m'^2}\right) - \frac{M^2_{\tilde e_2}}{M^2_{\tilde e_2} - m'^2} \ln\left(\frac{M^2_{\tilde e_2}}{m'^2}\right)\right], \tag{3.7}$$

where $\alpha_{U(1)} = g'^2/(4\pi)$. Similar numerical analysis can be done as we performed in the quark sector above.

From the figures 1 and 2 one can see why quarks are heavier than leptons; they get color contribution while leptons not as was showed on ref. [3].

## 4. Supersymmetric left-right model (SUSYLR)

The supersymmetric extension of left-right models [20, 21] is based on the gauge group $SU(3)_C \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. Apart from its original motivation of providing a dynamic explanation for the parity violation observed in low-energy weak interactions, this model differs from the SM in another important aspect; it explains the observed lightness of neutrinos in a natural way and it can also solve the strong CP problem.

On the technical side, the left-right symmetric model has a problem similar to that in the SM: the masses of the fundamental Higgs scalars diverge quadratically. As in the SM, the SUSYLR can be used to stabilize the scalar masses and cure this hierarchy problem. SUSYLR models have the additional appealing characteristics of having automatic R-parity conservation.

On the literature there are two different SUSYLR models. They differ in their $SU(2)_R$ breaking fields: one uses $SU(2)_R$ triplets (SUSYLRT) and the other $SU(2)_R$ doublets (SUSYLRD). Theoretical consequences of these models can be found in various papers including [20] and [21] respectively. Some details of both models are described at appendix C and appendix D.

## 5. Masses of fermions in SUSYLR models

For SUSYLRT, the mass term to the quarks is (eqs. (C.4), (D.3)):

$$\mathcal{L}^{mass}_{quarks} = -\left[h^q_{ij} Q^T_i \Phi \imath\tau_2 Q^c_j + \tilde h^q_{ij} Q^T_i \Phi' \imath\tau_2 Q^c_j + h.c.\right]. \tag{5.1}$$

Using eq. (C.6) on the equation above, we get the following mass matrix in the non-diagonal form

$$
\begin{aligned}
M^u_{ij} &= \frac{1}{\sqrt{2}}\left[k_1 h^q_{ij} + k_2' \tilde{h}^q_{ij}\right](u_i u^c_j + hc), \\
M^d_{ij} &= \frac{1}{\sqrt{2}}\left[k_1' h^q_{ij} + k_2 \tilde{h}^q_{ij}\right](d_i d^c_j + hc).
\end{aligned} \tag{5.2}
$$

For the leptons (eq. (C.4)), the mass term is

$$
\begin{aligned}
\mathcal{L}^{mass}_{leptons} &= -\Big[f_{ab}(L^T_a \imath\tau_2 \Delta_L L_b) + f^c_{ab}(L^{cT}_a \imath\tau_2 \delta^c_L L^c_b) + h^l_{ab}(L^T_a \Phi \imath\tau_2 L^c_b) \\
&+ \tilde{h}^l_{ab}(L^T_a \Phi' \imath\tau_2 L^c_b) + h.c.\Big].
\end{aligned} \tag{5.3}
$$

Using eq. (C.6) on equation given above, we get the following mass matrix in the non-diagonal representation

$$
\begin{aligned}
M^l_{ab} &= \frac{1}{\sqrt{2}}\left[k_1' h^l_{ab} + k_2 \tilde{h}^l_{ab}\right](l_a l^c_b + hc), \\
M^\nu_{ab} &= \frac{1}{\sqrt{2}}\left[k_1 h^l_{ab} + k_2' \tilde{h}^l_{ab}\right](\nu_a \nu^c_b + hc) + \frac{v_R}{\sqrt{2}} f^c_{ab}(\nu^c_a \nu^c_b + hc) \\
&- \frac{v_L}{\sqrt{2}} f_{ab}(\nu_a \nu_b + hc).
\end{aligned} \tag{5.4}
$$

This result is in agreement with the presented in [22], if we take $v_L = 0$.

For another hand, in the case of SUSYLRD one extracts from eqs. (D.3) the mass term to the leptons

$$
\mathcal{L}^{mass}_{leptons} = -\left(h^l_{ab}(L^T_a \Phi \imath\tau_2 L^c_b) + \tilde{h}^l_{ab}(L^T_a \Phi' \imath\tau_2 L^c_b) + h.c.\right). \tag{5.5}
$$

Using eq. (C.6) above, we get the following mass matrix in the non diagonal representation

$$
\begin{aligned}
M^l_{ab} &= -\frac{1}{\sqrt{2}}\left[k_1' h^l_{ab} + k_2 \tilde{h}^l_{ab}\right](l_a l^c_b + hc), \\
M^\nu_{ab} &= -\frac{1}{\sqrt{2}}\left[k_1 h^l_{ab} + k_2' \tilde{h}^l_{ab}\right](\nu_a \nu^c_b + hc).
\end{aligned} \tag{5.6}
$$

From eqs. (5.4), (5.6) we see that the choice of the triplets is preferable to doublets because in the first case we can generate a large Majorana mass for the right-handed neutrinos [22].

The $d, s$ and $b$ quarks as well as the $e, \mu$ and $\tau$ leptons will have masses proportional to the vacuum expectation values $k_1', k_2$, whereas the $u, c$ and $t$ will have masses proportional to $k_1, k_2'$.

Now, we are deal with the charged fermions and we are going to present the results which are hold in both models. To avoid flavor-changing-neutral currents (see [23]), we can choose the vacuum expectations values of the bidoublets as

$$
\langle \Phi \rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} k_1 & 0 \\ 0 & 0 \end{pmatrix}; \qquad \langle \Phi' \rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & 0 \\ 0 & k_2 \end{pmatrix}. \tag{5.7}
$$

Where $k_1, k_2$ are of the order of the electroweak scale $10^2$GeV. Using this fact on eqs. (5.2), (5.4), we can rewrite the mass matrix of charged fermion as

$$
\begin{aligned}
M^u_{ij} &= \frac{h^q_{ij}}{\sqrt{2}} k_1 (u_i u^c_j + hc), \\
M^d_{ij} &= \frac{\tilde{h}^q_{ij}}{\sqrt{2}} k_2 (d_i d^c_j + hc), \\
M^l_{ab} &= -\frac{\tilde{h}^l_{ab}}{\sqrt{2}} k_2 (l_i l^c_j + hc).
\end{aligned} \tag{5.8}
$$

The equations above are very similar to ones we get on the MSSM case, see eq. (2.2). Following the references [4, 3] we can try to find a discrete symmetry in order to prevent the electron and the quarks $u$ and $d$ from acquire masses at tree level. We impose the following $\mathcal{Z}'_2$ symmetry

$$
\begin{aligned}
\hat{Q}^c_2 &\to \tau_3 \qquad \hat{Q}^c_2, \\
\hat{Q}^c_3 &\to -I \qquad \hat{Q}^c_3, \\
\hat{L}_3 &\to \tau_3 \qquad \hat{L}_3,
\end{aligned} \tag{5.9}
$$

and the others superfields are even under this symmetry (compare with eq. (2.4)). With eqs. (5.9), (5.10) at hand we can reproduce the results presented in the section 2. We want to emphasize this symmetry does not forbid a large Majorana mass for the right-handed neutrinos and the results presented in ref. [22] are still valid.

The mixing between the squarks and the sleptons is given by

$$
\begin{aligned}
\mathcal{L} &= m^2_{Q_L}(\tilde{u}^*_3 \tilde{u}_3 + \tilde{d}^*_3 \tilde{d}_3) + m^2_{Q_R}(\tilde{u}^{*c}_3 \tilde{u}^{*c}_3 + \tilde{d}^{*c}_3 \tilde{d}_3) + m^2_{L_L} \tilde{l}^*_3 \tilde{l}_3 + m^2_{L_R} \tilde{l}^c_3 \tilde{l}^{*c}_3 + \\
&+ \frac{1}{\sqrt{2}} \left[ A^{LR}_{33} k_1 \tilde{l}_3 \tilde{l}^c_3 + \tilde{A}^{LR}_{33} k'_2 \tilde{l}_3 \tilde{l}^c_3 + A^{QQ}_{33} k_1 \tilde{u}_3 \tilde{u}^c_3 + \tilde{A}^{QQ}_{33} k'_2 \tilde{d}_3 \tilde{d}^c_3 \right],
\end{aligned} \tag{5.10}
$$

which generates the diagrams shown in figures 1 and 2, and then we can reproduce the results presented in the section 3.

## 6. Conclusions

We show that we can introduce a discrete symmetry $\mathcal{Z}'_2$ in MSSM and in both SUSYLR in order to explain the lower masses of the quarks $u, d$ and $s$ and of the electron while a consistent picture with experimental data of CKM matrix is obtained. We have also shown that in the models studied in this work the heavy leptons ($\mu$ and $\tau$) acquire mass at tree level while the electron get their mass at 1-loop level.

A discrete symmetry like $\mathcal{Z}'_2$ protects the Chiral symmetry to be broken in $SU(3)$ sector. This allows Chiral symmetry to be broken at different scales and two scales of mass is obtained.

These results presented in this article is easily extended to the others supersymmetric models as of the ref. [29].

| Superfield | Usual Particle | Spin | Superpartner | Spin |
|---|---|---|---|---|
| $\hat{V}'$ (U(1)) | $V_m$ | 1 | $\lambda_B$ | $\frac{1}{2}$ |
| $\hat{V}^i$ (SU(2)) | $V^i_m$ | 1 | $\lambda^i_A$ | $\frac{1}{2}$ |
| $\hat{V}^a_c(SU(3))$ | $G^a_m$ | 1 | $\tilde{g}^a$ | $\frac{1}{2}$ |
| $\hat{Q}_i \sim (\mathbf{3}, \mathbf{2}, 1/3)$ | $(u_i,\, d_i)_L$ | $\frac{1}{2}$ | $(\tilde{u}_{iL},\, \tilde{d}_{iL})$ | 0 |
| $\hat{u}^c_i \sim (\mathbf{3}^*, \mathbf{1}, -4/3)$ | $\bar{u}^c_{iL}$ | $\frac{1}{2}$ | $\tilde{u}^c_{iL}$ | 0 |
| $\hat{d}^c_i \sim (\mathbf{3}^*, \mathbf{1}, 2/3))$ | $\bar{d}^c_{iL}$ | $\frac{1}{2}$ | $\tilde{d}^c_{iL}$ | 0 |
| $\hat{L}_a \sim (\mathbf{1}, \mathbf{2}, -1)$ | $(\nu_a,\, l_a)_L$ | $\frac{1}{2}$ | $(\tilde{\nu}_{aL},\, \tilde{l}_{aL})$ | 0 |
| $\hat{l}^c_a \sim (\mathbf{1}, \mathbf{1}, 2)$ | $\bar{l}^c_{aL}$ | $\frac{1}{2}$ | $\tilde{l}^c_{aL}$ | 0 |
| $\hat{H}_1 \sim (\mathbf{1}, \mathbf{2}, -1)$ | $(H^0_1,\, H^-_1)$ | 0 | $(\tilde{H}^0_1,\, \tilde{H}^-_1)$ | $\frac{1}{2}$ |
| $\hat{H}_2 \sim (\mathbf{1}, \mathbf{2}, 1)$ | $(H^+_2,\, H^0_2)$ | 0 | $(\tilde{H}^+_2,\, \tilde{H}^0_2)$ | $\frac{1}{2}$ |

**Table 1:** Particle content of MSSM.


## Acknowledgments

C.M. Maekawa was supported by FAPERGS PROADE-2 under contract number 02/1266-6.


## A. MSSM

This model contains the particle content given at table 1. The families index for leptons are $a, b$ and they run over $e, \mu, \tau$, while the families index for the quarks are $i, j = 1, 2, 3$. The parentheses in the first column are the transformation properties under the respective representation of $(SU(3)_C, SU(2)_L, U(1)_Y)$.

The superfields formalism is useful in writing the manifestly invariant supersymmetric Lagrangian [18]. The fermions and scalars are represented by chiral superfields while the gauge bosons by vector superfields. As usual the superfield of a field $\phi$ is denoted by $\hat{\phi}$ [11]. The chiral superfield of a multiplet $\phi$ is denoted by

$$\begin{aligned}\hat{\phi} \equiv \hat{\phi}(x, \theta, \bar{\theta}) &= \tilde{\phi}(x) + i\, \theta\sigma^m\bar{\theta}\, \partial_m\tilde{\phi}(x) + \frac{1}{4}\, \theta\theta\, \bar{\theta}\bar{\theta}\, \Box\tilde{\phi}(x) \\ &+ \sqrt{2}\, \theta\phi(x) + \frac{i}{\sqrt{2}}\, \theta\theta\, \bar{\theta}\bar{\sigma}^m\partial_m\phi(x) \\ &+ \theta\theta\, F_\phi(x), \end{aligned} \tag{A.1}$$

while the vector superfield is given by

$$\hat{V} \equiv \hat{V}(x, \theta, \bar{\theta}) = -\theta\sigma^m\bar{\theta}V_m(x) + i\theta\theta\bar{\theta}\bar{\tilde{V}}(x) - i\bar{\theta}\bar{\theta}\theta\tilde{V}(x) + \frac{1}{2}\theta\theta\bar{\theta}\bar{\theta}D_V(x). \tag{A.2}$$




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The fields $F$ and $D$ are auxiliary fields which are needed to close the supersymmetric algebra and eventually will be eliminated using their equations of motion.

The Lagrangian of this model is written as

$$\mathcal{L}_{MSSM} = \mathcal{L}_{SUSY} + \mathcal{L}_{soft}\,, \tag{A.3}$$

where $\mathcal{L}_{SUSY}$ is the supersymmetric piece and can be divided as follows

$$\mathcal{L}_{SUSY} = \mathcal{L}_{lepton} + \mathcal{L}_{Quarks} + \mathcal{L}_{Gauge} + \mathcal{L}_{Higgs}, \tag{A.4}$$

where each term is given by

$$\begin{aligned}
\mathcal{L}_{lepton} &= \int d^4\theta\ \left[\bar{\hat{L}}_a e^{2g\hat{V}+g'\left(-\frac{1}{2}\right)\hat{V}'}\hat{L}_a + \bar{\hat{l}}^c{}_a e^{g'\hat{V}'}\hat{l}^c_a\right],\\
\mathcal{L}_{Quarks} &= \int d^4\theta\ \left[\bar{\hat{Q}}_i e^{2g_s\hat{V}^a_c+2g\hat{V}+g'\left(\frac{1}{6}\right)\hat{V}'}\hat{Q}_i + \bar{\hat{u}}^c{}_i e^{2g_s\hat{V}^a_c+g'\left(-\frac{2}{3}\right)\hat{V}'}\hat{u}^c_i\right.\\
&\quad\left. + \bar{\hat{d}}^c{}_i e^{2g_s\hat{V}^a_c+g'\left(\frac{1}{3}\right)\hat{V}'}\hat{d}^c_i\right],\\
\mathcal{L}_{Gauge} &= \frac{1}{4}\left\{\int d^2\theta\ \left[\sum_{a=1}^{8} W^{a\alpha}_s W^a_{s\alpha} + \sum_{i=1}^{3} W^{i\alpha}W^i_\alpha + W'^{\alpha}W'_\alpha\right] + h.c.\right\},
\end{aligned} \tag{A.5}$$

the subscripts $a$ and $i$ are family index, summed over $e, \mu, \tau$ and $1, 2, 3$, respectively, on repetition. The last piece of our Lagrangian is written as

$$\begin{aligned}
\mathcal{L}_{Higgs} &= \int d^4\theta\ \left[\bar{\hat{H}}_1 e^{2g\hat{V}+g'\left(-\frac{1}{2}\right)\hat{V}'}\hat{H}_1 + \bar{\hat{H}}_2 e^{2g\hat{V}+g'\left(\frac{1}{2}\right)\hat{V}'}\hat{H}_2\right]\\
&\quad + \int d^2\theta \quad W + \int d^2\bar{\theta} \quad \bar{W}.
\end{aligned} \tag{A.6}$$

The field strength are given by [18]

$$\begin{aligned}
W^a_{s\alpha} &= -\frac{1}{8g_s}\,\bar{D}\bar{D}e^{-2g_s\hat{V}^a_c}D_\alpha e^{2g_s\hat{V}^a_c} \qquad \alpha = 1, 2\,,\\
W^i_\alpha &= -\frac{1}{8g}\,\bar{D}\bar{D}e^{-2g\hat{V}^i}D_\alpha e^{2g\hat{V}^i}\,,\\
W'_\alpha &= -\frac{1}{4}\,DD\bar{D}_\alpha\hat{V}'\,.
\end{aligned} \tag{A.7}$$

The superpotential is given by

$$\begin{aligned}
W &= W_{RC} + \bar{W}_{RC} + W_{RV} + \bar{W}_{RV},\\
W_{2RC} &= \mu\epsilon\hat{H}_1\hat{H}_2 + y^l_{ab}\epsilon\hat{L}_a\hat{H}_1\hat{l}^c_b + y^u_{ij}\epsilon\hat{Q}_i\hat{H}_2\hat{u}^c_j + y^d_{ij}\epsilon\hat{Q}_i\hat{H}_1\hat{d}^c_j,\\
W_{RV} &= \mu_{1a}\epsilon\hat{L}_a\hat{H}_2 + \lambda_{abc}\epsilon\hat{L}_a\hat{L}_b\hat{l}^c_c + \lambda^1_{aij}\epsilon\hat{L}_a\hat{Q}_i\hat{d}^c_j + \lambda^2_{ijk}\hat{u}^c_i\hat{d}^c_j\hat{d}^c_k\,.
\end{aligned} \tag{A.8}$$

Where $W_{RC}$ ($W_{RV}$) conserves (violates) $R$-parity. The indices are summed on repetition.

The terms that break Supersymmetry softly and do not induce quadratic divergence [24] are

$$\mathcal{L}_{\text{soft}} = -\frac{1}{2}\left(\sum_{i=1}^{8} m_{\tilde{g}}\tilde{g}^i\tilde{g}^i + \sum_{p=1}^{3} m_\lambda\lambda^p_A\lambda^p_A + m'\lambda_B\lambda_B + h.c.\right) - M^2_L\tilde{L}^\dagger\tilde{L}$$



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$$- M_l^2\tilde{l^c}^\dagger\tilde{l^c} - M_Q^2\tilde{Q}^\dagger\tilde{Q} - M_u^2\tilde{u^c}^\dagger\tilde{u^c} - M_d^2\tilde{d^c}^\dagger\tilde{d^c} - M_1^2\tilde{H_1}^\dagger\tilde{H}_1 - M_2^2\tilde{H_2}^\dagger\tilde{H}_2$$

$$- \left[A^l H_1\tilde{L}\tilde{l^c} + A^u H_2\tilde{Q}\tilde{u^c} + A^d H_1\tilde{Q}\tilde{d^c} + M_{12}^2 H_1 H_2 + h.c.\right]. \tag{A.9}$$

We have omitted generation indices and we do the same to all the soft terms that we will write on this article. The parameters $m_{\tilde{g}}, m_\lambda$, and $m'$ are the $SU(3), SU(2)$, and $U(1)$ gaugino masses, respectively. $M_1^2, M_2^2$ and $M_{12}^2$ are mass terms for the Higgs fields. The scalar mass terms $M_L^2, M_l^2, M_Q^2, M_u^2$, and $M_d^2$ are in general Hermitian $3\times 3$ matrices.

The vacuum expectation value of this model is given by

$$\langle H_1\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \qquad \langle H_2\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_2 \end{pmatrix}. \tag{A.10}$$

## B. Mass diagonalization

The mass matrix of the "up" quark sector ($Y_u$) (quark with charge $+2/3$) and of the charged leptons ($Y_l$) are written as the following:

$$Y_u = \begin{pmatrix} y_{11}^u & y_{12}^u & 0 \\ y_{21}^u & y_{22}^u & 0 \\ y_{31}^u & y_{32}^u & 0 \end{pmatrix}\cdot v_2\,, \qquad Y_l = \begin{pmatrix} y_{11}^l & y_{12}^l & 0 \\ y_{21}^l & y_{22}^l & 0 \\ y_{31}^l & y_{32}^l & 0 \end{pmatrix}\cdot v_1\,, \tag{B.1}$$

while to the case of "down " quark sector the mass matrix is given by

$$Y_d = \begin{pmatrix} 0 & 0 & y_{13}^d \\ 0 & 0 & y_{23}^d \\ 0 & 0 & y_{33}^d \end{pmatrix}\cdot v_1\,, \tag{B.2}$$

where $v_1$ and $v_2$ are VEVs of $H_1$ and $H_2$ respectively.

The fermion's mass matrix is diagonalized using two unitary matrices, $D$ and $E$. Then we can write the diagonal mass matrix as

$$M_{diag}^2 = DY_f^T\cdot Y_f D^{-1} = E^* Y_f\cdot Y_f^T (E^*)^{-1}, \tag{B.3}$$

where $f$ can represent any "up", "down" quarks or any charged lepton.

After the diagonalization we have defined the followings parameters

$$\begin{aligned}
t_u &= \left((y_{11}^u)^2 + (y_{12}^u)^2 + (y_{21}^u)^2 + (y_{22}^u)^2 + (y_{31}^u)^2 + (y_{32}^u)^2\right)\cdot v_2, \\
t_l &= \left((y_{11}^l)^2 + (y_{12}^l)^2 + (y_{21}^l)^2 + (y_{22}^l)^2 + (y_{31}^l)^2 + (y_{32}^l)^2\right)\cdot v_1, \\
r_u &= \sqrt{(t_u)^2 - 4((y_{12}^u)^2 u_u - 2y_{11}^u y_{12}^u v_u + x_u^2 + (y_{11}^u)^2 z_u)}\cdot v_2, \\
u_u &= (y_{21}^u)^2 + (y_{31}^u)^2, \qquad v_u = y_{31}^u y_{32}^u + y_{21}^u y_{22}^u, \\
x_u &= y_{21}^u y_{22}^u - y_{31}^u y_{32}^u, \qquad z_u = (y_{22}^u)^2 + (y_{32}^u)^2, \\
r_l &= \sqrt{(t_l)^2 - 4((y_{12}^l)^2 u_l - 2y_{11}^l y_{12}^l v_l + x_l^2 + (y_{11}^l)^2 z_l)}\cdot v_1, \\
u_l &= (y_{21}^l)^2 + (y_{31}^l)^2, \qquad v_l = y_{31}^l y_{32}^l + y_{21}^l y_{22}^l, \\
x_l &= y_{21}^l y_{22}^l - y_{31}^l y_{32}^l, \qquad z_l = (y_{22}^l)^2 + (y_{32}^l)^2, \\
t_d &= \left((y_{13}^d)^2 + (y_{23}^d)^2 + (y_{33}^d)^2\right)\cdot v_1, \qquad r_d = t_d.
\end{aligned} \tag{B.4}$$



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| Superfield | Usual Particle | Spin | Superpartner | Spin |
|---|---|---|---|---|
| $\hat{V}'$ (U(1)) | $B_m$ | 1 | $\tilde{B}$ | $\frac{1}{2}$ |
| $\hat{V}^i_L$ $(SU(2)_L)$ | $W^i_{mL}$ | 1 | $\tilde{W}^i_L$ | $\frac{1}{2}$ |
| $\hat{V}^i_R$ $(SU(2)_R)$ | $W^i_{mR}$ | 1 | $\tilde{W}^i_R$ | $\frac{1}{2}$ |
| $\hat{V}^a_c(SU(3))$ | $g^a_m$ | 1 | $\tilde{g}^a$ | $\frac{1}{2}$ |
| $\hat{Q}_i \sim (\mathbf{3}, \mathbf{2}, \mathbf{1}, 1/3)$ | $(u_i,\, d_i)_{iL}$ | $\frac{1}{2}$ | $(\tilde{u}_{iL},\, \tilde{d}_{iL})$ | 0 |
| $\hat{Q}^c_i \sim (\mathbf{3}^*, \mathbf{1}, \mathbf{2}, -1/3)$ | $(d^c_i,\, -u^c_i)_{iL}$ | $\frac{1}{2}$ | $(\tilde{d}^c_{iL},\, -\tilde{u}^c_{iL})$ | 0 |
| $\hat{L}_a \sim (\mathbf{1}, \mathbf{2}, \mathbf{1}, -1)$ | $(\nu_a,\, l_a)_{aL}$ | $\frac{1}{2}$ | $(\tilde{\nu}_{aL},\, \tilde{l}_{aL})$ | 0 |
| $\hat{L}^c_a \sim (\mathbf{1}, \mathbf{1}, \mathbf{2}, 1)$ | $(l^c_a,\, -\nu^c_a)_{aL}$ | $\frac{1}{2}$ | $(\tilde{l}^c_{aL},\, -\tilde{\nu}^c_{aL})$ | 0 |
| $\hat{\Delta}_L \sim (\mathbf{1}, \mathbf{3}, \mathbf{1}, 2)$ | $\begin{pmatrix} \frac{\delta^+_L}{\sqrt{2}} & \delta^{++}_L \\ \delta^0_L & \frac{-\delta^+_L}{\sqrt{2}} \end{pmatrix}$ | 0 | $\begin{pmatrix} \frac{\tilde{\delta}^+_L}{\sqrt{2}} & \tilde{\delta}^{++}_L \\ \tilde{\delta}^0_L & \frac{-\tilde{\delta}^+_L}{\sqrt{2}} \end{pmatrix}$ | $\frac{1}{2}$ |
| $\hat{\Delta}'_L \sim (\mathbf{1}, \mathbf{3}, \mathbf{1}, -2)$ | $\begin{pmatrix} \frac{\delta'^-_L}{\sqrt{2}} & \delta'^0_L \\ \delta'^{--}_L & \frac{-\delta'^-_L}{\sqrt{2}} \end{pmatrix}$ | 0 | $\begin{pmatrix} \frac{\tilde{\delta}'^-_L}{\sqrt{2}} & \tilde{\delta}'^0_L \\ \tilde{\delta}'^{--}_L & \frac{-\tilde{\delta}'^-_L}{\sqrt{2}} \end{pmatrix}$ | $\frac{1}{2}$ |
| $\hat{\delta}^c_L \sim (\mathbf{1}, \mathbf{1}, \mathbf{3}, -2)$ | $\begin{pmatrix} \frac{\lambda^-_L}{\sqrt{2}} & \lambda^0_L \\ \lambda^{--}_L & \frac{-\lambda^-_L}{\sqrt{2}} \end{pmatrix}$ | 0 | $\begin{pmatrix} \frac{\tilde{\lambda}^-_L}{\sqrt{2}} & \tilde{\lambda}^0_L \\ \tilde{\lambda}^{--}_L & \frac{-\tilde{\lambda}^-_L}{\sqrt{2}} \end{pmatrix}$ | $\frac{1}{2}$ |
| $\hat{\delta}'^c_L \sim (\mathbf{1}, \mathbf{1}, \mathbf{3}, 2)$ | $\begin{pmatrix} \frac{\lambda'^+_L}{\sqrt{2}} & \lambda'^{++}_L \\ \lambda'^0_L & \frac{-\lambda'^+_L}{\sqrt{2}} \end{pmatrix}$ | 0 | $\begin{pmatrix} \frac{\tilde{\lambda}'^+_L}{\sqrt{2}} & \tilde{\lambda}'^{++}_L \\ \tilde{\lambda}'^0_L & \frac{-\tilde{\lambda}'^+_L}{\sqrt{2}} \end{pmatrix}$ | $\frac{1}{2}$ |
| $\hat{\Phi} \sim (\mathbf{1}, \mathbf{2}, \mathbf{2}, 0)$ | $\begin{pmatrix} \phi^0_1 & \phi^+_1 \\ \phi^-_2 & \phi^0_2 \end{pmatrix}$ | 0 | $\begin{pmatrix} \tilde{\phi}^0_1 & \tilde{\phi}^+_1 \\ \tilde{\phi}^-_2 & \tilde{\phi}^0_2 \end{pmatrix}$ | $\frac{1}{2}$ |
| $\hat{\Phi}' \sim (\mathbf{1}, \mathbf{2}, \mathbf{2}, 0)$ | $\begin{pmatrix} \chi^0_1 & \chi^+_1 \\ \chi^-_2 & \chi^0_2 \end{pmatrix}$ | 0 | $\begin{pmatrix} \tilde{\chi}^0_1 & \tilde{\chi}^+_1 \\ \tilde{\chi}^-_2 & \tilde{\chi}^0_2 \end{pmatrix}$ | $\frac{1}{2}$ |

**Table 2:** Particle content of SUSYLRT.

## C. Triplet model (SUSYLRT)

The particle content of the model is given at Tab.(2) (for recent work see for example [25] and references therein). In parentheses it appears the transformation properties under the respective $(SU(3)_C, SU(2)_L, SU(2)_R, U(1)_{B-L})$. The Lagrangian is given by:

$$\mathcal{L}_{SUSYLRT} = \mathcal{L}_{Lepton} + \mathcal{L}_{Quarks} + \mathcal{L}_{Gauge} + \mathcal{L}_{Higgs}, \tag{C.1}$$

where

$$
\begin{aligned}
\mathcal{L}_{Lepton} &= \int d^4\theta \left[ \bar{\hat{L}}_{aL} e^{2gT^i\hat{V}_L^i + g'\left(\frac{-1}{2}\right)\hat{V}'} \hat{L}_{aL} + \bar{\hat{L}}^c_{aL} e^{2gT^i\hat{V}_R^i + g'\left(\frac{1}{2}\right)\hat{V}'} \hat{L}^c_{aL} \right], \\
\mathcal{L}_{Quarks} &= \int d^4\theta \left[ \bar{\hat{Q}}_{iL} e^{2g_sT^a\hat{V}_c^a + 2gT^i\hat{V}_L^i + g'\left(\frac{1}{6}\right)\hat{V}'} \hat{Q}_{iL} \right. \\
&+ \left. \bar{\hat{Q}}^c_{iL} e^{2g_s\bar{T}^a\hat{V}_c^a + 2gT^i\hat{V}_R^i + g'\left(\frac{-1}{6}\right)\hat{V}'} \hat{Q}^c_{iL} \right], \\
\mathcal{L}_{Gauge} &= \frac{1}{4} \left\{ \int d^2\theta \left[ \sum_{a=1}^{8} W_s^{a\alpha} W^a_{s\alpha} + \sum_{i=1}^{3} W_L^{i\alpha} W^i_{L\alpha} + \sum_{i=1}^{3} W_R^{i\alpha} W^i_{R\alpha} \right.\right. \\
&+ \left.\left. W'^{\alpha} W'_{\alpha} \right] + h.c. \right\}, \qquad \text{(C.2)}
\end{aligned}
$$

with $T^i = \tau^i/2$ is the generator of $SU(2)$ group while $T^a = \lambda^a/2$ is the generator of triplets of $SU(3)$ while $\bar{T}^a = \bar{\lambda}^a/2$ is the generator of the anti-triplet of $SU(3)$. We make the usual assumption that the left and right couplings are equal, $g_L = g_R = g$. The terms $W_s^{a\alpha}, W_L^{i\alpha}, W_R^{i\alpha}$ and $W'^{\alpha}$ are calculated using expressions analogous to that at eq. (A.7).

The last part of our Lagrangian reads:

$$
\begin{aligned}
\mathcal{L}_{Higgs} &= \int d^4\theta \, Tr \left[ \bar{\hat{\Delta}}_L e^{2gT^i\hat{V}_L^i + g'(1)\hat{V}'} \hat{\Delta}_L + \bar{\hat{\Delta}}'_L e^{2gT^i\hat{V}_L^i + g'(-1)\hat{V}'} \hat{\Delta}'_L \right. \\
&+ \bar{\hat{\delta}}^c_L e^{2gT^i\hat{V}_R^i + g'(-1)\hat{V}'} \hat{\delta}^c_L + \bar{\hat{\delta}}'^c_L e^{2gT^i\hat{V}_R^i + g'(1)\hat{V}'} \hat{\delta}'^c_L + \bar{\hat{\Phi}} e^{2gT^i\hat{V}_L^i + 2gT^i\hat{V}_R^i} \hat{\Phi} \\
&+ \left. \bar{\hat{\Phi}}' e^{2gT^i\hat{V}_L^i + 2gT^i\hat{V}_R^i} \hat{\Phi}' \right] + \int d^2\theta W + \int d^2\bar{\theta} \overline{W}. \qquad \text{(C.3)}
\end{aligned}
$$

The most general superpotential $W$ [20] is given by

$$
\begin{aligned}
W &= M_\Delta Tr(\hat{\Delta}_L \hat{\Delta}'_L) + M_{\delta^c} Tr(\hat{\delta}^c_L \hat{\delta}'^c_L) + \mu_1 Tr(\imath\tau_2 \hat{\Phi} \imath\tau_2 \hat{\Phi}) + \mu_2 Tr(\imath\tau_2 \hat{\Phi}' \imath\tau_2 \hat{\Phi}') \\
&+ \mu_3 Tr(\imath\tau_2 \hat{\Phi} \imath\tau_2 \hat{\Phi}') + f_{ab} Tr(\hat{L}_a \imath\tau_2 \hat{\Delta}_L \hat{L}_b) + f^c_{ab} Tr(\hat{L}^c_a \imath\tau_2 \hat{\delta}^c_L \hat{L}^c_b) \\
&+ h^l_{ab} Tr(\hat{L}_a \hat{\Phi} \imath\tau_2 \hat{L}^c_b) + \tilde{h}^l_{ab} Tr(\hat{L}_a \hat{\Phi}' \imath\tau_2 \hat{L}^c_b) + h^q_{ij} Tr(\hat{Q}_i \hat{\Phi} \imath\tau_2 \hat{Q}^c_j) \\
&+ \tilde{h}^q_{ij} Tr(\hat{Q}_i \hat{\Phi}' \imath\tau_2 \hat{Q}^c_j) + W_{NR}. \qquad \text{(C.4)}
\end{aligned}
$$

Where $h^l, \tilde{h}^l, h^q$ and $\tilde{h}^q$ are the Yukawa couplings for the leptons and quarks, respectively, and $f$ and $f^c$ are the couplings for the triplets scalar bosons. We must emphasize that due to the conservation of $B - L$ symmetry, $\Delta'_L$ and $\delta'^c_L$ do not couple with the leptons and quarks. Here $W_{NR}$ denotes (possible) non-renormalizable terms arising from higher scale physics or Planck scale effects [26]. This model can be embedded in a supersymmetric grand unified theory as $SO(10)$ [27].

In addition, we have also to include soft supersymmetry breaking terms, they are:

$$
\begin{aligned}
\mathcal{L}_{soft} &= \left[ m^2_{L_L} \tilde{L}^\dagger_L \tilde{L}_L + m^2_{L_R} \tilde{L}^{c\dagger}_L \tilde{L}^c_L + m^2_{Q_L} \tilde{Q}^\dagger_L \tilde{Q}_L + m^2_{Q_R} \tilde{Q}^{c\dagger}_L \tilde{Q}^c_L + m^2_{\Phi\Phi} \Phi^\dagger \Phi \right. \\
&+ \left. m^2_{\Phi\Phi'} \Phi^\dagger \Phi' + m^2_{\Phi'\Phi'} \Phi'^\dagger \Phi' \right] - \left[ M_1^2 Tr(\Delta_L \Delta'_L) + M_2^2 Tr(\delta^c_L \delta'^c_L) + M_3^2 \Phi\Phi \right. \\
&+ \left. M_4^2 \Phi\Phi' + M_5^2 \Phi'\Phi' + h.c. \right] - \left[ A^{LL} Tr(\tilde{L}\tau_2 \Delta_L \tilde{L}) + \tilde{A}^{LL} Tr(\tilde{L}^c \tau_2 \delta^c_L \tilde{L}^c) \right. \\
&+ A^{LR} Tr(\tilde{L} \Phi \imath\tau_2 \tilde{L}^c) + \tilde{A}^{LR} Tr(\tilde{L} \Phi' \imath\tau_2 \tilde{L}^c) + A^{QQ} Tr(\tilde{Q} \Phi \imath\tau_2 \tilde{Q}^c)
\end{aligned}
$$

$$+ \ \tilde{A}^{QQ} Tr(\tilde{Q}\Phi' \imath\tau_2 \tilde{Q}^c) + h.c.\Big] - \frac{1}{2}\left(\sum_{i=1}^{8} m_{\tilde{g}} \tilde{g}^i \tilde{g}^i + \sum_{i=1}^{3} m_L \tilde{W}_L^i \tilde{W}_L^i \right.$$
$$\left. + \sum_{i=1}^{3} m_R \tilde{W}_R^i \tilde{W}_R^i + m' \tilde{B}\tilde{B} + h.c.\right). \tag{C.5}$$

The vacuum expectations values are given by [28]

$$\begin{aligned}
\langle \Phi \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} k_1 & 0 \\ 0 & k_1' \end{pmatrix}; \qquad \langle \Phi' \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} k_2' & 0 \\ 0 & k_2 \end{pmatrix}; \\
\langle \Delta_L \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ \upsilon_L & 0 \end{pmatrix}; \qquad ; \qquad \langle \Delta_L' \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & \upsilon_L' \\ 0 & 0 \end{pmatrix}; \\
\langle \delta_L^c \rangle &= \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & \upsilon_R \\ 0 & 0 \end{pmatrix}; \qquad \langle \delta_L'^c \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ \upsilon_R' & 0 \end{pmatrix}.
\end{aligned} \tag{C.6}$$

## D. Doublet model (SUSYLRD)

This model contains the particle content given at table 3.

The Lagrangian of this model is given by:

$$\mathcal{L}_{SUSYLRD} = \mathcal{L}_{Lepton} + \mathcal{L}_{Quarks} + \mathcal{L}_{Gauge} + \mathcal{L}_{Higgs}, \tag{D.1}$$

where $\mathcal{L}_{Lepton}, \mathcal{L}_{Quarks}, \mathcal{L}_{Gauge}$ are given by eq. (C.2). The last part of our Lagrangian is given by:

$$\begin{aligned}
\mathcal{L}_{Higgs} =& \int d^4\theta \left[ \hat{\bar{\chi}}_1 e^{2gT^i \hat{V}_L^i + g'\left(\frac{1}{2}\right)\hat{V}'} \hat{\chi}_1 + \hat{\bar{\chi}}_2 e^{2gT^i \hat{V}_L^i + g'\left(\frac{-1}{2}\right)\hat{V}'} \hat{\chi}_2 + \hat{\bar{\chi}}_3^c e^{2gT^i \hat{V}_R^i + g'\left(\frac{-1}{2}\right)\hat{V}'} \hat{\chi}_3^c \right. \\
&+ \left. \hat{\bar{\chi}}_4^c e^{2gT^i \hat{V}_R^i + g'\left(\frac{1}{2}\right)\hat{V}'} \hat{\chi}_4^c + \hat{\bar{\Phi}} e^{2gT^i \hat{V}_L^i + 2gT^i \hat{V}_R^i} \hat{\Phi} + \hat{\bar{\Phi}}' e^{2gT^i \hat{V}_L^i + 2gT^i \hat{V}_R^i} \hat{\Phi}' \right] \\
&+ \int d^2\theta W + \int d^2\overline{\theta W}.
\end{aligned} \tag{D.2}$$

The most general superpotential and soft supersymmetry breaking Lagrangian for this model are:

$$\begin{aligned}
W =& M_\chi \hat{\chi}_1 \hat{\chi}_2 + M_{\chi^c} \hat{\chi}_3^c \hat{\chi}_4^c + \mu_1 Tr(\tau_2 \hat{\Phi} \tau_2 \hat{\Phi}) + \mu_2 Tr(\tau_2 \hat{\Phi}' \tau_2 \hat{\Phi}') + \mu_3 Tr(\tau_2 \hat{\Phi} \tau_2 \hat{\Phi}') \\
&+ h_{ab}^l Tr(\hat{L}_a \hat{\Phi} \imath\tau_2 \hat{L}_b^c) + \tilde{h}_{ab}^l Tr(\hat{L}_a \hat{\Phi}' \imath\tau_2 \hat{L}_b^c) + h_{ij}^q Tr(\hat{Q}_i \hat{\Phi} \imath\tau_2 \hat{Q}_j^c) \\
&+ \tilde{h}_{ij}^q Tr(\hat{Q}_i \hat{\Phi}' \imath\tau_2 \hat{Q}_j^c) + W_{NR}.
\end{aligned} \tag{D.3}$$

and the soft terms reads

$$\begin{aligned}
\mathcal{L}_{soft} =& -\Big[ m_{L_L}^2 \tilde{L}_L^\dagger \tilde{L}_L + m_{L_R}^2 \tilde{L}_L^{c\dagger} \tilde{L}_L^c + m_{Q_L}^2 \tilde{Q}_L^\dagger \tilde{Q}_L + m_{Q_R}^2 \tilde{Q}_L^{c\dagger} \tilde{Q}_L^c + m_{\Phi\Phi}^2 \Phi^\dagger \Phi \\
&+ \ m_{\Phi\Phi'}^2 \Phi^\dagger \Phi' + m_{\Phi'\Phi'}^2 \Phi'^\dagger \Phi' \Big] - \big[ M_1^2 \chi_1 \chi_2 + M_2^2 \chi_3^c \chi_4^c + M_3^2 \Phi\Phi + M_4^2 \Phi\Phi'
\end{aligned}$$

| Superfield | Usual Particle | Spin | Superpartner | Spin |
|---|---|---|---|---|
| $\hat{V}'$ (U(1)) | $B_m$ | 1 | $\tilde{B}$ | $\frac{1}{2}$ |
| $\hat{V}_L^i$ $(SU(2)_L)$ | $W_{mL}^i$ | 1 | $\tilde{W}_L^i$ | $\frac{1}{2}$ |
| $\hat{V}_R^i$ $(SU(2)_R)$ | $W_{mR}^i$ | 1 | $\tilde{W}_R^i$ | $\frac{1}{2}$ |
| $\hat{V}_c^a(SU(3))$ | $g_m^a$ | 1 | $\tilde{g}^a$ | $\frac{1}{2}$ |
| $\hat{Q}_i \sim (\mathbf{3}, \mathbf{2}, \mathbf{1}, 1/3)$ | $(u_i,\, d_i)_{iL}$ | $\frac{1}{2}$ | $(\tilde{u}_{iL},\, \tilde{d}_{iL})$ | 0 |
| $\hat{Q}_i^c \sim (\mathbf{3}^*, \mathbf{1}, \mathbf{2}, -1/3)$ | $(d_i^c,\, -u_i^c)_{iL}$ | $\frac{1}{2}$ | $(\tilde{d}_{iL}^c,\, -\tilde{u}_{iL}^c)$ | 0 |
| $\hat{L}_a \sim (\mathbf{1}, \mathbf{2}, \mathbf{1}, -1)$ | $(\nu_a,\, l_a)_{aL}$ | $\frac{1}{2}$ | $(\tilde{\nu}_{aL},\, \tilde{l}_{aL})$ | 0 |
| $\hat{L}_a^c \sim (\mathbf{1}, \mathbf{1}, \mathbf{2}, 1)$ | $(l_a^c,\, -\nu_a^c)_{aL}$ | $\frac{1}{2}$ | $(\tilde{l}_{aL}^c,\, -\tilde{\nu}_{aL}^c)$ | 0 |
| $\hat{\chi}_{1L} \sim (\mathbf{1}, \mathbf{2}, \mathbf{1}, 1)$ | $(\chi_{1L}^+,\quad \chi_{1L}^0)$ | 0 | $(\tilde{\chi}_{1L}^+,\quad \tilde{\chi}_{1L}^0)$ | $\frac{1}{2}$ |
| $\hat{\chi}_{2L} \sim (\mathbf{1}, \mathbf{2}, \mathbf{1}, -1)$ | $(\chi_{2L}^0,\quad \chi_{2L}^-)$ | 0 | $(\tilde{\chi}_{2L}^0,\quad \tilde{\chi}_{2L}^-)$ | $\frac{1}{2}$ |
| $\hat{\chi}_{3L}^c \sim (\mathbf{1}, \mathbf{1}, \mathbf{2}, -1)$ | $(\chi_{3L}^0,\quad \chi_{3L}^-)$ | 0 | $(\tilde{\chi}_{3L}^0,\quad \tilde{\chi}_{3L}^-)$ | $\frac{1}{2}$ |
| $\hat{\chi}_{4L}^c \sim (\mathbf{1}, \mathbf{1}, \mathbf{2}, 1)$ | $(\chi_{4L}^+,\quad \chi_{4L}^0)$ | 0 | $(\tilde{\chi}_{4L}^+,\quad \tilde{\chi}_{4L}^0)$ | $\frac{1}{2}$ |
| $\hat{\Phi} \sim (\mathbf{1}, \mathbf{2}, \mathbf{2}, 0)$ | $\begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}$ | 0 | $\begin{pmatrix} \tilde{\phi}_1^0 & \tilde{\phi}_1^+ \\ \tilde{\phi}_2^- & \tilde{\phi}_2^0 \end{pmatrix}$ | $\frac{1}{2}$ |
| $\hat{\Phi}' \sim (\mathbf{1}, \mathbf{2}, \mathbf{2}, 0)$ | $\begin{pmatrix} \chi_1^0 & \chi_1^+ \\ \chi_2^- & \chi_2^0 \end{pmatrix}$ | 0 | $\begin{pmatrix} \tilde{\chi}_1^0 & \tilde{\chi}_1^+ \\ \tilde{\chi}_2^- & \tilde{\chi}_2^0 \end{pmatrix}$ | $\frac{1}{2}$ |

**Table 3:** Particle content of SUSYLRD.

$$
\begin{aligned}
&+\ M_5^2\Phi'\Phi' + h.c.\Big] - \Big[A^{LR}Tr(\tilde{L}\Phi\imath\tau_2\tilde{L}^c) + \tilde{A}^{LR}Tr(\tilde{L}\Phi'\imath\tau_2\tilde{L}^c) \\
&+\ A^{QQ}Tr(\tilde{Q}\Phi\imath\tau_2\tilde{Q}^c) + \tilde{A}^{QQ}Tr(\tilde{Q}\Phi'\imath\tau_2\tilde{Q}^c) + h.c.\Big] \\
&-\ \frac{1}{2}\left(\sum_{i=1}^{8} m_{\tilde{g}}\tilde{g}^i\tilde{g}^i + \sum_{i=1}^{3} m_L\tilde{W}_L^i\tilde{W}_L^i + \sum_{i=1}^{3} m_R\tilde{W}_R^i\tilde{W}_R^i + m'\tilde{B}\tilde{B} + h.c.\right)
\end{aligned}
\tag{D.4}
$$

The vacuum expectation values of the new scalars are

$$
\begin{aligned}
\langle\chi_{1L}\rangle &= \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ \upsilon_L \end{pmatrix}, \qquad \langle\chi_{2L}\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} \upsilon'_L \\ 0 \end{pmatrix}, \\
\langle\chi_{3L}^c\rangle &= \frac{1}{\sqrt{2}}\begin{pmatrix} \upsilon_R \\ 0 \end{pmatrix}, \qquad \langle\chi_{4L}^c\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ \upsilon'_R \end{pmatrix}.
\end{aligned}
\tag{D.5}
$$

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