A Simple Phenomenological Parametrization of Supersymmetry without R-Parity

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Abstract

We present a parametrization of the supersymmetric standard model without R-parity that permits efficient phenomenological analyses of the full model without a priori assumptions. Under the parametrization, which is characterized by a single vacuum expectation value for the scalar components of the $Y = -1/2$ superfields, the expressions for tree-level mass matrices are quite simple. They do not involve the trilinear R-parity violating couplings; however, the bilinear $\mu_i$ terms do enter and cannot be set to zero without additional assumptions. We set up a framework for doing phenomenology and show some illustrative results for fermion mass matrices and related bounds on parameters. We find in particular that large values of $\tan\beta$ can suppress R-parity violating effects, substantially weakening experimental constraints.

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I. INTRODUCTION

A softly broken supersymmetry (SUSY) around or below the scale of a TeV is no doubt the most popular extension of the Standard Model. Most SUSY studies concentrate on a “minimal” version of such a model which contains two electroweak symmetry breaking Higgs doublets and an \textit{ad hoc} discrete symmetry, called R-parity, which essentially distinguishes particles from superparticles. The phenomenological role of R-parity is to forbid B- or L-number violating couplings for which there are important experimental bounds, for instance those from superparticle mediated proton decays. However, strict R-parity conservation is not required to satisfy these bounds. Furthermore, allowing R-parity to be broken, either spontaneously through sneutrino VEV’s or explicitly in the Lagrangian, gives rise to interesting phenomenology. This has become a subject of much interest recently (for some recent reviews and references, see [1,2]). Because of the large number of possible R-parity violating (RPV) couplings, most studies impose assumptions to restrict the analyses to a particular subset of RPV couplings. We wish to adopt a purely phenomenological point of view that allows an analysis of all the RPV couplings without \textit{a priori} assumptions. In this letter we present our perspective on the parametrization of the general RPV Lagrangian and show some initial results. Here we concentrate on tree level results for mass matrices in the fermion sector, especially in the large $\tan \beta$ regime; a more detailed study will appear elsewhere.

II. PARAMETRIZATION OF R-PARITY VIOLATION

The most general renormalizable superpotential for the supersymmetric standard model without R-parity can be written as

$$W = \varepsilon_{ab} [\mu_{\alpha} \hat{L}_{\alpha}^a \hat{H}_{u}^b + h_{ik} \hat{Q}_{i}^a \hat{R}_{u}^{b} \hat{U}_{k}^c + \lambda_{\alpha k}^i \hat{L}_{\alpha}^a \hat{D}_{k}^c + \lambda_{\alpha \beta k} \hat{L}_{\alpha}^a \hat{E}_{k}^c + \lambda_{\alpha \beta \gamma} \hat{D}_{i}^c \hat{D}_{j}^c \hat{U}_{k}^c],$$

(1)

where $(a, b)$ are $SU(2)$ indices, $(i, j, k)$ are family (flavor) indices, and $(\alpha, \beta)$ are (extended) flavor indices from 0 to 3 with $\hat{L}_{\alpha}$’s denoting the four doublet superfields with $Y = -1/2$. \lambda
and $\lambda''$ are antisymmetric in the first two indices as required by $SU(2)$ and $SU(3)$ product rules respectively. In the limit where $\lambda_{ijk}, \lambda'_{ijk}, \lambda''_{ijk}$ and $\mu_i$ all vanish, one recovers the expression for the R-parity preserving minimal supersymmetric standard model (MSSM), with $\hat{L}_0$ identified as $\hat{H}_d$. In that case, the two Higgses acquire vacuum expectation values (VEV’s) and the lepton and quark Yukawa couplings are given by $h^e_{ik} \equiv 2\lambda_{ik0} (= -2\lambda_{0ik})$, $h^d_{ik} \equiv \lambda'_{ik0}$, and $-h^u_{ik}$, respectively. In the general case, the full expression for $W$ together with all admissible soft SUSY breaking terms should be used to construct the scalar potential and solve for the vacuum. The solution is then expected to involve VEV’s for all five neutral scalars, i.e. for both the Higgses and sneutrinos in the usual terminology.

It is not necessary, however, to retain all five VEV’s in parametrizing the model, because the freedom associated with the choice of flavor basis creates redundancy amongst the parameters. For example, in the MSSM, the two $3 \times 3$ complex mass matrices $\frac{\mu_i}{\sqrt{2}} h^u_{ik}$ and $\frac{\mu_i}{\sqrt{2}} h^d_{ik}$ correspond only to ten real parameters describing the six mass eigenvalues and the CKM-matrix. And although there are models attempting to construct the full high-energy mass matrices in particular flavor bases, so far as low-energy phenomenology is concerned the different bases cannot be distinguished. A fruitful strategy in the MSSM is therefore to choose a flavor basis that parametrizes the two Yukawa matrices $h^u$ and $h^d$ with exactly ten parameters, namely assuming one mass matrix to be given by the diagonal eigenvalues and the other a multiple of the CKM-matrix and the other diagonal eigenvalue matrix. Similarly, in parametrizing the general supersymmetric standard model without R-parity, $U(3)$ flavor rotations for $\hat{Q}_i, \hat{U}_i^c, \hat{D}_i^c$ and $\hat{E}_i^c$ as well as a $U(4)$ rotation for $\hat{L}_a$ can be exploited.

The above observation is not new. The popular parametrization exploiting the flavor rotations has all $\mu_0$’s except $\mu_0$ as well as two of the sneutrino VEV’s ($\tilde{\nu}_i$’s) set to zero (rotated away) \[3\]. (In this context sneutrinos refer to the scalar components of the three $\hat{L}_i$ superfields as defined in the basis where the $\mu_i$’s are zero.) Note that the lepton Yukawa matrix $h^e_{ik}$ cannot then be taken as diagonalized, since the full $\hat{L}_a$ basis is already fixed. This (single-$\mu$) parametrization was introduced originally in studies of RPV effects on the leptons.
under the assumption that the trilinear RPV couplings are zero \[4\]. Extending its usage to
the most general RPV scenario proves difficult. For instance, the chargino-charged-lepton
mass matrix in a generic basis is given by

$$
\mathcal{M}_c = \begin{pmatrix}
M_2 & \frac{g v_u}{\sqrt{2}} & 0 \\
\frac{g v_u}{\sqrt{2}} & \mu_0 & -h_{ik}^e \frac{\hat{v}_i}{\sqrt{2}} \\
\frac{g v_u}{\sqrt{2}} & \mu_i & h_{ik}^e \frac{\hat{v}_i}{\sqrt{2}} + 2\lambda_{ijk} \frac{\hat{v}_i}{\sqrt{2}}
\end{pmatrix}.
$$

(2)

It is easy to see that under the single-\(\mu\) parametrization a set of \(\lambda\)-couplings associated
with the nonzero sneutrino VEV still remain, making the analysis of the mass eigenstates
quite complicated. Analogously, \(\lambda'\)-couplings will enter the down-quark mass matrix along
with the nonzero sneutrino VEV. It has also been pointed out that rotating away the \(\mu_i\)'s
does not simplify the analysis of the scalar potential because the RPV soft mass terms also
contribute to the same quadratic terms \[5\].

We wish to find a relatively simple parametrization not requiring \textit{a priori} assumptions.
We propose here to use the \(U(4)\) rotation to set all sneutrino VEV's \((\frac{\hat{v}_i}{\sqrt{2}} \equiv \langle \hat{L}_i \rangle)\) to zero,
leaving a single VEV for \(\langle \hat{L}_0 \rangle\). We keep all of the \(\mu_0\) while the rest of the leptonic flavor
rotations are used to set \(h_{ik}^e\) diagonal. In our new basis, the (tree-level) mass matrices for
all the fermions \textit{do not} involve any trilinear RPV couplings\[5\]. In particular, \(\mathcal{M}_c\) is given by

$$
\mathcal{M}_c = \begin{pmatrix}
M_2 & \frac{g v_u}{\sqrt{2}} & 0 & 0 \\
\frac{g v_u}{\sqrt{2}} & \mu_0 & 0 & 0 \\
0 & \mu_1 & m_1 & 0 \\
0 & \mu_2 & 0 & m_2 \\
0 & \mu_3 & 0 & m_3
\end{pmatrix},
$$

(3)

where \(m_i = h_{ii}^e \frac{\hat{v}_i}{\sqrt{2}}\). The quark mass matrix for each sector involves only one VEV and
therefore assumes the same form as in the MSSM. The neutralino-neutrino mass matrix is
then given by

\[1\]The scalar mass matrices are also much simplified as there are only two non-zero VEV's, \(v_\chi\) and
\(v_u\). Further details will be discussed in \[\text{\[5\]}\].
\[ M_{\nu} = \begin{pmatrix}
M_1 & 0 & -\frac{g_1 v_u}{2} & -\frac{g_1 v_d}{2} & 0 & 0 & 0 \\
0 & M_2 & -\frac{g_2 v_u}{2} & -\frac{g_2 v_d}{2} & 0 & 0 & 0 \\
-\frac{g_1 v_u}{2} & -\frac{g_2 v_u}{2} & 0 & -\mu_0 & -\mu_1 & -\mu_2 & -\mu_3 \\
-\frac{g_1 v_d}{2} & -\frac{g_2 v_d}{2} & -\mu_0 & 0 & 0 & 0 & 0 \\
0 & 0 & -\mu_1 & 0 & 0 & 0 & 0 \\
0 & 0 & -\mu_2 & 0 & 0 & 0 & 0 \\
0 & 0 & -\mu_3 & 0 & 0 & 0 & 0
\end{pmatrix}, \quad (4) \]

The simplicity of the mass matrix expressions is obvious. We would like to re-emphasize that this is achieved without any a priori assumptions; we have simply chosen to parametrize the model in a specific flavor basis \[7\]. In our basis, we identify the Higgs \( \tilde{H}_1(\equiv \tilde{L}_0) \) as the \( Y = -1/2 \) doublet that bears the full VEV among the \( \tilde{L}_\alpha \)'s, and in the interesting region of relatively small \( \mu_i \)'s our three \( L_i \)'s align well with the charged-lepton mass eigenstates.

Our single-VEV parametrization helps to simplify analysis of the model in both the fermion and scalar sectors. We will take advantage of this to illustrate below some novel features of the model in the large \( \tan \beta \) regime. The parametrization also provides a framework that easily allows a full phenomenological analysis of the model, with both bilinear and trilinear RPV-terms admitted. Finally, all parameters are assumed to be real — potentially rich CP violating features of the model are not considered here.

Before going on to our analysis, it is worthwhile to further clarify some issues about the parametrization of R-parity violation. The single-VEV parametrization, with explicit bilinear RPV-terms, does not manifestly exhibit sneutrino VEV’s. Nevertheless, it is possible to use this framework to describe a spontaneously broken R-parity scenario, which is defined by the existence of sneutrino VEV’s in some basis where all explicit RPV terms vanish. A rotation of the \( \tilde{L}_\alpha \) would connect such a basis to ours. Note that if we perform such a rotation to our single-VEV basis, in general not only bilinear \( (\mu_i) \) terms but also trilinear and soft SUSY breaking RPV-terms will be introduced as well. A similar perspective holds for models with R-parity broken spontaneously via VEV(s) of extra singlet superfield(s) at
some higher scale. In that case, the extra superfield(s) have to be integrated out to recover the supersymmetric standard model. Such models are naturally formulated in a mixed parametrization \[2\], *i.e.* with both \(\mu_i\)’s and \(\langle \hat{L}_i \rangle\)’s.

There are also discussions in the literature which consider the MSSM with only a few RPV-terms added while the “sneutrino VEV’s” are assumed to be zero. One should be particularly careful in interpreting the meaning of the assumptions in such cases. As implied in the above discussion, imposing such assumptions at the Lagrangian level is not a (flavor) basis or parametrization independent procedure. Combining such assumptions with a specific choice of basis is a potential source of confusion and sometimes inconsistency. In our opinion, a clear interpretation is provided by beginning with the Lagrangian in the single-VEV parametrization we proposed above. A given model is then identified by specifying which RPV-terms are not admitted. An important related point is that so long as the \(\mu_i\)’s are not all assumed to be zero, the \(L_i\)’s are not exactly indentifiable as the physical charged leptons, nor are the charginos and neutralinos, for instance, the same states as in the R-parity conserving MSSM. The vanishing of the \(\mu_i\)’s, if taken, would be an assumption. All in all, while specific RPV models may be more naturally formulated in a particular parametrization, the single-VEV parametrization, we believe, provides a particularly efficient framework for a model independent study of constraints on and phenomenology of R-parity violation.

III. ANALYSIS OF NEUTRAL FERMION MASS MATRIX

The generation of non-zero neutrino mass(es) is one of the most prominent features of R-parity violation. As a result the experimental neutrino mass bound has been used to put constraints on various RPV-couplings \[8\]. Two neutrino eigenstates are left massless at the tree level, while the third one gains a mass through the RPV-couplings (\(\mu_i\)’s) to the higgsino, as can be seen from Eq. (4). Note that the massive eigenstate is in general a mixture of all three neutrino states. In fact one can use a simple rotation to decouple the massless states. The remaining \(5 \times 5\) mass matrix is then given by
\[ \mathcal{M}_N^{(5)} = \begin{pmatrix} M_1 & 0 & \frac{g_1 v d}{2} & -\frac{g_1 v u}{2} & 0 \\ 0 & M_2 & -\frac{g_2 v d}{2} & \frac{g_2 v u}{2} & 0 \\ \frac{g_1 v d}{2} & -\frac{g_2 v d}{2} & 0 & -\mu_0 & -\mu_5 \\ \frac{g_1 v u}{2} & \frac{g_2 v u}{2} & -\mu_0 & 0 & 0 \\ 0 & 0 & -\mu_5 & 0 & 0 \end{pmatrix}, \]  

where

\[ \mu_5 = \sqrt{\mu_1^2 + \mu_2^2 + \mu_3^2}; \]

and the corresponding massive neutrino state is given by

\[ |\nu_5\rangle = \frac{\mu_1}{\mu_5} |\nu_1\rangle + \frac{\mu_2}{\mu_5} |\nu_2\rangle + \frac{\mu_3}{\mu_5} |\nu_3\rangle. \]  

The common strategy to obtain the neutrino mass corresponds to assuming a small $\mu_5$ where $\mathcal{M}_N^{(5)}$ adopts a “seesaw” structure. This gives

\[ m_{\nu_5} = \frac{\det \mathcal{M}_N^{(5)}}{\det \mathcal{M}_{4\times4}} = -\frac{1}{2\mu_0} \frac{\mu_5^2 v^2 \cos^2 \beta (x g_2^2 + g_1^2)}{[2x M_2 \mu_0 - (x g_2^2 + g_1^2) v^2 \sin \beta \cos \beta]}, \]

where we have substituted $v_d = v \cos \beta$, $v_u = v \sin \beta$, and $M_1 = x M_2$. Note that for large $\tan \beta$, $\cos \beta$ is a suppression factor. For example, at $\tan \beta = 45$, saturation of the machine bound of 24 MeV [4] for $m_{\nu_5}$ allows a $\mu_5$ value as large as the chargino mass scale $M_2$ and the higgsino mass mixing parameter, $\mu_0$. Now, large $\mu_5$ values are beyond the validity of the “seesaw” analysis; however, an alternative perturbative analysis can be performed treating the EW-symmetry breaking terms in $\mathcal{M}_N^{(5)}$ as a perturbation. The first order part can then be diagonalized exactly without any assumptions about the magnitude of $\mu_5$. The resulting zero eigenvalue is lifted by the perturbation to give

\[ m_{\nu_5} = \frac{\mu_5^2 v^2 \cos^2 \beta (x g_2^2 + g_1^2)}{4 \left( \mu_0^2 + \mu_5^2 \right) x M_2}; \]

with the eigenvector in the original basis given by

\[ \left( \begin{array}{c} \frac{g_5 g_1 v \cos \beta}{2x M_2} \\ -\frac{g_5 g_2 v \cos \beta}{2x M_2} \\ 0 \\ -\mu_5 \\ \mu_0 \end{array} \right). \]
One can rewrite Eq.(9) to obtain a bound on \( \mu_5 \):

\[
\mu_5^2 < \frac{4x\mu_0^2 M_2 m_{\nu_5(\text{bound})}}{v^2 \cos^2 \beta (xg_1^2 + g_1^2) - 4xM_2 m_{\nu_5(\text{bound})}}.
\] (11)

As \( M_2 \) increases, the denominator above drops to zero, beyond which there is no bound on \( \mu_5 \). For \( \tan\beta = 45 \), this happens at \( M_2 \sim 210 \text{ GeV} \). We note that in order to use the \( \nu_r \) machine mass bound of 24 MeV for \( m_{\nu_5(\text{bound})} \), we have assumed \( \mu_1 = \mu_2 = 0 \), i.e. \( \mu_5 = \mu_3 \) (\( \mu_1 : \mu_2 : \mu_3 = 0 : 0 : 1 \)).

The perturbative result in Eq.(11) is borne out by exact numerical results from diagonalizing the neutral fermion mass matrix, as illustrated in Figure 1, for \( \tan\beta = 45 \); the corresponding result for \( \tan\beta = 2 \) is also shown for comparison. The difference between \( \mu_5 \) bounds for the two cases is striking. In both cases we see that the neutrino mass bound on \( \mu_5 \) (= \( \mu_3 \) here) is tighter for low values of \( |\mu_0| \) and weakens as \( |\mu_0| \) is increased. For \( \tan\beta = 45 \), values of \( \mu_5 \) in the hundreds of GeV are completely consistent with the 24 MeV bound for viable regions of the \( (M_2, \mu_0) \) parameter space. And again, for \( M_2 \) large enough, we get no limit on \( \mu_5 \) at all.

There are potentially much stronger bounds on neutrino masses from cosmological considerations which however depend on the decay modes and other assumptions so that a neutrino mass above an MeV is not definitely ruled out [10]. There are also other experimental constraints on neutrino masses and mixings. In this first paper, our purpose is simply to illustrate the advantage of performing the analysis in the single-VEV parametrization as well as to point out the interesting suppression of RPV effects in the large \( \tan\beta \) regime.

### IV. CONSTRAINTS FROM CHARGINO-CHARGED LEPTON MASS MATRIX

How are the \( \mu_i \)'s otherwise constrained particularly in the large \( \tan\beta \) regime? Couplings of the charged leptons to the \( Z^0 \) are well measured and can constrain the \( \mu_i \)'s through the chargino-charged lepton mass matrix. In Ref. [11], various \( Z^0 \)-couplings constraints on R-parity violation are studied assuming the trilinear RPV couplings vanish. We follow basically
the same strategy (but without the latter assumption) and our parametrization allows the constraints to be cast explicitly in terms of the magnitudes of the $\mu_i$. For the large tan$\beta$ regime, we find a weakening of the constraints as a result of $\cos \beta$ suppression factor(s) even stronger than that illustrated above for the neutrino case.

We begin with a perturbative approach as in the neutral fermion sector. After we diagonalize its upper $2 \times 2$ (chargino) block, $\mathcal{M}_c$ (Eq. (3)) consists of two diagonal blocks of fully diagonal sub-matrices and a lower $3 \times 2$ off-diagonal block containing the $\mu_i$ parameters. The latter is taken as a perturbation to the diagonal matrix. It is then a simple exercise to obtain the matrix elements (12):

\[
U_L^{\dagger}(i + 2, 1) = -\frac{\mu_i \sqrt{2} M_W \cos \beta}{\mu_0 M_2 - 2 M_W^2 \sin \beta \cos \beta},
\]

\[
U_L^{\dagger}(i + 2, 2) = -\frac{\mu_i M_2}{\mu_0 M_2 - 2 M_W^2 \sin \beta \cos \beta},
\]

\[
U_R^{\dagger}(i + 2, 1) = -\frac{m_i \mu_i \sqrt{2} M_W (M_2 \sin \beta + \mu_0 \cos \beta)}{(\mu_0 M_2 - 2 M_W^2 \sin \beta \cos \beta)^2},
\]

\[
U_R^{\dagger}(i + 2, 2) = -\frac{m_i \mu_i (M_2^2 + 2 M_W^2 \cos^2 \beta)}{(\mu_0 M_2 - 2 M_W^2 \sin \beta \cos \beta)^2},
\]

with

\[
U_L^{\dagger} \mathcal{M}_c U_R = \text{diag}\{\bar{M}_{e1}, \bar{M}_{e2}, \bar{m}_1, \bar{m}_2, \bar{m}_3\}, \tag{14}
\]

where index $i$ refers to one of the three leptonic states. These matrix elements are the ones needed for studying leptonic physics. They characterize the gaugino and higgsino contents of each leptonic mass eigenstate.

The $Z^0$-boson coupling to the mass eigenstates is given by

\[
\mathcal{L}_{\text{int}}^{\chi^+ \chi^-} = \frac{g_2}{2 \cos \theta_w} Z^\mu \bar{\chi}_i^+ \gamma_\mu \left( \bar{A}_L^{ij} \frac{1 - \gamma_5}{2} + \bar{A}_R^{ij} \frac{1 + \gamma_5}{2} \right) \chi_j^-.
\]

Using the results above, together with unitarity of the $U_L^{\dagger}$ and $U_R^{\dagger}$ matrices, we have, for large tan$\beta$ where formulæ simplify,

\[
\bar{A}_L^{ij} = \frac{2 \mu_i \mu_j M_W^2 \cos^2 \beta}{\mu_0^2 M_2^2} + \delta_{ij} (1 - 2 \sin^2 \theta_w),
\]

\[
\bar{A}_R^{ij} = \frac{\mu_i \mu_j m_i m_j (M_2^2 + 4 M_W^2)}{\mu_0^4 M_2^2} - \delta_{ij} 2 \sin^2 \theta_w.
\]

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Note that while Eq. (13) is valid for all five charged fermion states, the above given formulæ for the couplings $\tilde{A}_{L,R}^{ij}$'s are only for the three leptonic states, to which we limit the present discussion. In accordance with the perturbational approach, the results contain "small" ratios such as $\frac{\mu_i}{\mu_0}$. The deviations from standard universality in $\tilde{A}_{ij}^R$ are typically ignored, being suppressed by two factors of the $\frac{m_i}{\mu_0}$ mass ratios. But the corresponding deviations in $\tilde{A}_{ij}^L$ have a $\cos^2\beta$ suppression, which for $\tan\beta = 45$, for example, gives a factor of $10^{-3}$. The important point to note here, in the large $\tan\beta$ regime, is that $\tilde{A}_{33}^R$ or $\tilde{A}_{23}^R$ could well be significant in comparison with the $\tilde{A}_{ij}^L$'s.

With the $\tilde{A}_{L,R}^{ij}$ formulæ, it is then straightforward to check the constraints different processes involving universality violation or FCNC impose on the $\frac{\mu_i}{\mu_0}$ ratios. Specifically, we check coupling universality and left-right asymmetry, as well as tree level $Z^0\ell_\ell k$ couplings through branching ratios of various $Z^0 \to 2\ell$, $\mu \to 3\ell$ and $\tau \to 3\ell$ processes. When a constraint allows the ratios to be larger than unity, it essentially goes away, from the present perturbative perspective. This happens for all of these constraints for a sufficiently large $M_2$ except the one from $\mu^- \to e^-e^+e^-$, which has a much stronger experimental bound. For $\tan\beta = 45$, $M_2 \gtrsim 15-35$ GeV eliminates all but the latter process from which we obtain

$$\frac{|\mu_1\mu_2|}{\mu_0^2} \lesssim 4.7 \times 10^{-7} M_2^2. \quad (17)$$

With an admissible gaugino mass $M_2$ of the order 100 GeV, this only surviving constraint gives a numerical bound on $\frac{|\mu_1\mu_2|}{\mu_0^2} \sim 10^{-3}$ only (the signs of the $\mu_i$ do not affect any of the results presented here). Furthermore, the same behavior appears when the $Z^0 \to \nu\nu$ width constraint is considered.

In the perturbative calculation, the $m_i$'s are approximated by the $\bar{m}_i$'s, the physical charged lepton masses. In the exact computations, we numerically integrate from $\mu_i = 0$ (for which the $m_i$ are exactly the $\bar{m}_i$) to the final $\mu_i$ values. This is necessary to find an acceptable set of $m_i$'s that yield the correct physical charged lepton masses for a given set of $\mu_i$'s. We also then find the chargino masses, which now depend on the $\mu_i$'s. For example, the minimum $\mu_i$ values required to give both chargino masses above 90 GeV for $\tan\beta = 45$. 

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and $\mu_1 : \mu_2 : \mu_3 = 0 : 1 : 1$ are shown in Figure 2.

The numerical results also bear out the perturbative analysis of the couplings (in lepton and $Z^0$ decays) mentioned above. Bounds from B.R.($\mu^- \rightarrow e^- e^+ e^-$) $< 1.0^{-12}$ \cite{9} are shown in Figure 3. Here we plot contours of $\mu_5$ assuming a ratio $\mu_1 : \mu_2 : \mu_3 = 1 : 1 : 0$. As with the neutrino mass bound, the limit on $\mu_5$ is more strict for low values of $|\mu_0|$. The substantial weakening of the constraint for large $\tan \beta$ is well illustrated in the figure. Numerical studies of the other processes mentioned above give no restriction on the maximum allowed values of $\mu_5$ for $|\mu_0|, M_2 \gtrsim 10$ GeV (again for $\tan \beta = 45$), as indicated by the perturbational results.

**V. SUMMARY**

In summary, we have presented and illustrated with examples the merits of the single-VEV parametrization of supersymmetry without R-parity (see also \cite{7}), which can be used in phenomenological studies without requiring specific model-dependent assumptions. Our analysis as outlined above also indicates a strong suppression of RPV effects from the bilinear $\mu_i$ terms in the large $\tan \beta$ regime. We note that the trilinear RPV couplings play no role in the analysis discussed here, though they are expected to have an important role in the other aspects of the model. This is, however, exactly what makes a comprehensive analysis of the full model feasible under the parametrization. Further details of such an analysis will be reported in future work.

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Figure Captions

Figure 1:
Maximum allowed values of $\mu_5$ (in GeV) consistent with $m_{\nu\tau} < 24$ MeV ($\mu_1 : \mu_2 : \mu_3 = 0 : 0 : 1$). $M_1 = x M_2$, with $x = \frac{5}{3} \tan^2 \theta_w$ assumed from gaugino unification ($M_Z = 91.19$ GeV, $\sin^2 \theta_w = 0.23$). The region above or outside of a given contour is excluded for $\mu_5$’s above the indicated value.

Figure 2:
Minimum values of $\mu_5$ (in GeV) required to give both chargino masses above 90 GeV ($\mu_1 : \mu_2 : \mu_3 = 0 : 1 : 1$). The area above or outside of a given contour has both chargino masses $> 90$ GeV for $\mu_5$’s above the indicated value.

Figure 3:
Maximum allowed values of $\mu_5$ (in GeV) consistent with B.R.($\mu^- \rightarrow e^- e^+ e^-$) $< 1.0^{-12}$ ($\mu_1 : \mu_2 : \mu_3 = 1 : 1 : 0$). The region above or outside of a given contour is excluded for $\mu_5$’s above the indicated value.
FIG. 1. Maximum allowed values of $\mu_5$ (in GeV) consistent with $m_{\nu_e} < 24$ MeV

($\mu_1 : \mu_2 : \mu_3 = 0 : 0 : 1$). $M_1 = xM_2$, with $x = \frac{5}{3}\tan^2\theta_w$ assumed from gaugino unification

($M_Z = 91.19$ GeV, $\sin^2\theta_w = 0.23$). The region above or outside of a given contour is excluded for $\mu_5$'s above the indicated value.
FIG. 2. Minimum values of $\mu_5$ (in GeV) required to give both chargino masses above 90 GeV ($\mu_1 : \mu_2 : \mu_3 = 0 : 1 : 1$). The area above or outside of a given contour has both chargino masses $> 90$ GeV for $\mu_5$'s above the indicated value.
FIG. 3. Maximum allowed values of $\mu_5$ (in GeV) consistent with $\text{B.R.}(\mu^- \rightarrow e^- e^+ e^-) < 10^{-12}$ ($\mu_1 : \mu_2 : \mu_3 = 1 : 1 : 0$). The region above or outside of a given contour is excluded for $\mu_5$'s above the indicated value.