Higgs Bosons in a minimal R-parity conserving left-right supersymmetric model

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Abstract

We revisit the Higgs sector of the left-right supersymmetric model. We study the scalar potential in a version of the model in which the minimum is the charge conserving vacuum state, without $R$–parity violation or additional non-renormalizable terms in the Lagrangian. We analyze the dependence of the potential and of the Higgs mass spectrum on the various parameters of the model, pinpointing the most sensitive ones. We also show that, contrary to previous expectations, the model can predict light neutral flavor-conserving Higgs bosons, while the flavor-violating ones are heavy, and within the limits from $K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$ and $B^0_{d,s} - \bar{B}^0_{d,s}$ mixings. We study variants of the model in which at least one pair of doubly-charged Higgs bosons is light, and show that the parameter space for such Higgs masses and mixings is very restrictive, thus making the model more predictive.

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I. INTRODUCTION

Within this decade, the LHC will play a significant role in probing the Standard Model (SM) of electroweak interactions and disentangling the models beyond it. The progress expected in experimental high energy physics will complement theoretical explorations of various scenarios of new physics. The experimental data could confirm any of the many theoretical models of new physics advanced over the last decades.

One of the first observations expected at the LHC is the Higgs boson. This is the one remaining piece of puzzle missing from the SM and on this finding rests our understanding of mass generation. However, most models beyond the SM also predict the existence of one or more Higgs bosons. Some of them might be heavy, but several are expected to be light. While the standard model contains one neutral Higgs boson, many models predict one or more Higgs doublets, thus at least one charged Higgs boson (such as the many variants of the two Higgs doublet models and supersymmetry). Finding a light charged Higgs boson would raise problems as to which fundamental gauge symmetry is responsible for its existence. The hope of a clearer signal rests on more exotic Higgs bosons, such as the ones predicted in left-right models [1]. Left-right symmetric models with seesaw neutrino mass generation [2] predict doubly-charged Higgs bosons [3], which, if light, would give distinctive and spectacular signals at the colliders.

Including supersymmetry adds several attractive features to the left-right model [4]. Softly-broken supersymmetry resolves some of the inconsistencies of the standard model: it provides a solution to the gauge hierarchy problem, a natural candidate for weakly-interacting dark matter, and allows for gauge coupling unification. In addition, the left-right supersymmetric model (LRSUSY) accounts for neutrino masses [1], parity violation, offers a solution to the strong and weak CP violation without introduction of the axion [5], and explains the absence of excessive SUSY CP violation. Left-right symmetry is favored by many extra-dimensional models, and many gauge unification scenarios, such as $SO(10)$ [6].

However the model seems to suffer from a serious shortcoming. Minimization of the Higgs potential requires either spontaneous $R$–parity breaking by the vacuum expectation value (VEV) of the right-chiral scalar neutrino [7]; or introduction of higher scale non-renormalizable operators [8, 9]. Since an attractive characteristic of the left-right supersymmetric model is that explicit $R$–parity breaking is forbidden by the symmetry of the
model, spontaneous breaking is not a desirable feature. Ditto for higher order operators at the Planck scale. The shortcoming comes from the fact that, in the simplest version of the model, the global minimum of the theory breaks electric charge, making the theory unacceptable. This can be remedied by allowing a VEV for the right sneutrino. The Higgs boson spectrum was previously analyzed in this variant of the model with $R-$parity violation where sneutrinos and sleptons mix with the Higgs bosons [10].

However, a new version of the theory suggested by Babu and Mohapatra [11], allows for both $R-$parity conservation and the absence of higher-dimensional operators by inclusion of the Yukawa coupling of the heavy Majorana neutrino in the effective Lagrangian. We study the Higgs sector of such a model and examine the masses of the doubly-charged, singly charged and neutral bosons (both scalar and pseudoscalar sectors). Although the model depends on many parameters, we show that the masses are sensitive to only a few, and thus the model is more predictive. Light doubly-charged Higgs bosons emerge naturally. The LRSUSY model predicts neutral scalar and pseudoscalar Higgs bosons that violate flavor at tree level. We impose conditions coming from phenomenology: $K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$ and $B^0_{d,s} - \bar{B}^0_{d,s}$ mixing. We show that, contrary to previous expectations, one can have light neutral and charged Higgs bosons that conserve flavor, while the flavor violating bosons are in the 600 GeV- 100 TeV scale, as required by meson mixing constraints. We pinpoint the parameters that the masses are most sensitive to, and show that they satisfy the constraints in a limited range of these parameters. We set up the structure of the Higgs potential, masses and mixing, including the constraints, while leaving the study of the characteristic signals at the LHC for a future study.

The paper is organized as follows. In Section II we summarize the particular LRSUSY model we use, with emphasis on the Higgs structure. In the following Section III we present analytic formulas for the mass matrices in the neutral, singly-charged and doubly-charged sectors. In Section IV we present the results of the constraints from $K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$ and $B^0_{d,s} - \bar{B}^0_{d,s}$ mixings on the Higgs masses and mixings. We illustrate our results by showing two numerical scenarios for desirable Higgs mass values for the model which satisfy the constraints in Section V as well as presenting plots for masses consistent with the constraints. We summarize our findings and conclude in Section VI.
II. R-PARITY CONSERVING LEFT-RIGHT SUPERSYMMETRIC MODEL

The supersymmetric left-right model incorporates supersymmetry in the left-right model based on the gauge symmetry $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. Including the $B-L$ (where $B$ and $L$ stand for baryon and lepton numbers) in a gauge symmetry, the only quantum number left ungauged in SM, is an additional attractive feature of the model. The model contains left and right fermion doublets, as well as triplet gauge bosons for $SU(2)_L$ and $SU(2)_R$, and a neutral gauge boson for $U(1)_{B-L}$. $R$-parity, defined as $R_P = (-1)^{3(B-L)+2s}$ (with $s$ the spin of the particle), is imposed in the Minimal Supersymmetric Standard Model (MSSM) to avoid dangerous baryon and lepton number violating operators, otherwise explicit Yukawa terms that violate $R$-parity can exist in the Lagrangian. This explicit $R$-parity breaking is forbidden in LRSUSY models by the symmetries of the model. In early left-right symmetric models $SU(2)_R$ doublets were used to break the gauge symmetry. Later $SU(2)_{L,R}$ triplets were introduced to provide the seesaw mechanism for neutrino masses \cite{2}, and both left- and right-handed triplet Higgs bosons are required by parity conservation. The model was described extensively in several previous works \cite{4}. However $R$-parity may not be conserved in this setup. The reason is that the minimum of the potential prefers a solution in which the right-chiral scalar neutrino gets a VEV, thus breaking $R$-parity spontaneously. Two scenarios have been proposed which remedy this situation. One is the model of Babu and Mohapatra \cite{11} where an extra singlet Higgs boson is added to the model and one-loop corrections to the potential show that an $R$-parity conserving minimum can be found. The second model is that of Aulakh et. al. \cite{12}, where the addition of two more triplets, $\Omega(1,3,1,0)$ and $\Omega_c(1,1,3,0)$, with zero lepton number, achieves left-right symmetry breaking with conserved $R$-parity at tree-level. In our work, we adopt the former, as it is a minimal model, and present a short description below.

The Higgs sector in this minimal left-right supersymmetric model under the gauge group, together with the Higgs VEVs, is given in Table I.

The superpotential of this model is given by

$$W = Y_u Q^T \tau_2 \Phi_1 \tau_2 Q^c + Y_d Q^T \tau_2 \Phi_2 \tau_2 Q^c + Y_\nu L^T \tau_2 \Phi_1 \tau_2 L^c + Y_\ell L^T \tau_2 \Phi_2 \tau_2 L^c + h.c. + i \left( f^* L^T \tau_2 \Delta L + f L^T \tau_2 \Delta^c L^c \right) + S \left[ \lambda \text{Tr} \left( \Delta^* \Delta^c + \Delta^c \Delta^* \right) + \lambda_{ij} \text{Tr} \left( \Phi_i^T \tau_2 \Phi_j \tau_2 \right) - M_R^2 \right] + W'$$

(2.1)
where

\[ W' = \left[ M_\Delta \text{Tr}(\Delta \tilde{\Delta}) + M_\Delta^* \text{Tr}(\Delta^c \tilde{\Delta}^c) \right] + \mu_{ij} \text{Tr} \left( \Phi_i^T \tau_2 \Phi_j \tau_2 \right) + M_S S^2 + \lambda_S S^3 . \] (2.2)

Here \( Y_{u,d} \) and \( Y_{\nu,\ell} \) in Eq. (2.1) are quark and lepton Yukawa coupling matrices, while \( f \) is the Majorana neutrino Yukawa coupling. We choose to work with \( W' = 0 \), which leads to an enhanced \( R^- \) symmetry and a natural interpretation of the supersymmetric \( \mu \) term, as explained below.

The model is minimal in the following sense: \( \Delta^c \) and \( \bar{\Delta}^c \) fields are needed for breaking \( SU(2)_R \otimes U(1)_{B-L} \) symmetry without \( R \)-parity violation, the \( \Delta \) and \( \bar{\Delta} \) fields are for parity invariance, and the two bidoublets \( \Phi_1 \) and \( \Phi_2 \) are needed to generate the quark and lepton masses and Cabibbo Kobayashi Maskawa (CKM) mixings. The singlet field \( S \) is introduced so that \( SU(2)_R \otimes U(1)_{B-L} \) symmetry breaking occurs in the supersymmetric limit. The charge is defined as

\[ Q = I_3^L + I_3^R + \frac{B - L}{2} . \]

The VEVs of the Higgs fields in this model needed to break the symmetries as described above, are given in Table 1. If we assume that the VEVs of the bidoublet Higgs are real, the fermion mass matrices become Hermitian. The VEVs of the left–handed triplet fields \( \Delta, \bar{\Delta} \), which determine the tree-level left-handed neutrino masses must be extremely small and are assumed to be zero.

In the supersymmetric limit, the VEV of the singlet \( S \) Higgs boson is zero, but after SUSY breaking, \( \langle S \rangle \sim m_{\text{SUSY}} \). Thus the \( \mu \) term for the bidoublet \( \Phi \) will arise from the coupling \( \lambda_{ij} \), with a magnitude of order \( m_{\text{SUSY}} \). In the SUSY limit,

\[ |v_R| = |\bar{v}_R|, \quad \lambda v_R \bar{v}_R = M^2_R, \quad \langle S \rangle = 0 . \] (2.3)

The VEV of \( S \) field, generated after SUSY breaking, arises from linear terms in SUSY breaking

\[ V_{\text{soft}} = A_\lambda \lambda S \text{Tr}(\Delta^c \tilde{\Delta}^c + \Delta^* \tilde{\Delta}^*) - C_\lambda M^2_R S + \text{h.c.} \] (2.4)

Minimization of the resulting potential yields \( \langle S^* \rangle = \frac{1}{2\lambda}(C_\lambda - A_\lambda) \), which is of order \( m_{\text{SUSY}} \). If the coupling \( \lambda \) is small, then \( \langle S \rangle \) can be above the SUSY breaking scale. This feature can be used to make one pair of Higgs doublet superfields heavier than the SUSY breaking scale. However, the masses of doubly charged fermionic fields, which are equal to \( \lambda \langle S \rangle \) must
| Higgs Field | Matrix Representation | Vacuum Expectation Values |
|-------------|-----------------------|--------------------------|
| $\Delta(1,3,1,2)$ | $\left( \begin{array}{cc} \delta^+ & \delta^{++} \\ \frac{\delta^0}{\sqrt{2}} & -\delta^+ \end{array} \right)$ | $\left( \begin{array}{cc} 0 & 0 \\ v_L & 0 \end{array} \right)$ |
| $\Delta(1,3,1,-2)$ | $\left( \begin{array}{cc} \delta^- & \delta^0 \\ \frac{-\delta^-}{\sqrt{2}} & \delta^0 \end{array} \right)$ | $\left( \begin{array}{cc} 0 & \bar{v}_L \\ 0 & 0 \end{array} \right)$ |
| $\Delta^c(1,1,3,-2)$ | $\left( \begin{array}{cc} \delta^{c^-} & \delta^c \\ \frac{-\delta^{c^-}}{\sqrt{2}} & \delta^c \end{array} \right)$ | $\left( \begin{array}{cc} 0 & v_R \\ 0 & 0 \end{array} \right)$ |
| $\Delta^c(1,1,3,2)$ | $\left( \begin{array}{cc} \delta^{c^+} & \delta^{c^{++}} \\ \frac{-\delta^{c^+}}{\sqrt{2}} & \delta^{c^+} \end{array} \right)$ | $\left( \begin{array}{cc} 0 & 0 \\ \bar{v}_R & 0 \end{array} \right)$ |
| $\Phi_1(1,2,2,0)$ | $\left( \begin{array}{cc} \phi_1^+ & \phi_2^0 \\ \phi_1^0 & \phi_2^- \end{array} \right)$ | $\left( \begin{array}{cc} 0 & \kappa_1' \\ \kappa_1 & 0 \end{array} \right)$ |
| $\Phi_2(1,2,2,0)$ | $\left( \begin{array}{cc} \chi_1^+ & \chi_2^0 \\ \chi_1^0 & \chi_2^- \end{array} \right)$ | $\left( \begin{array}{cc} 0 & \kappa_2 \\ \kappa_2' & 0 \end{array} \right)$ |
| $S(1,1,1,0)$ | $S_0$ | $\langle S \rangle$ |
remain below a TeV. Consistency of the model (non-vanishing CKM mixing angle) requires the asymmetry $\mu_{12} = \mu_{21}$.

The full potential of the model relevant for symmetry breaking includes $F$-term, $D$-term and soft SUSY breaking contributions. They are given by

$$
V_F = \left| \lambda \text{Tr}(\Delta^* \Delta^* + \Delta^c \Delta^c) + \lambda_{ij} \text{Tr}(\Phi_i^T \tau_2 \Phi_j \tau_2) - M_R^2 \right|^2 \\
+ \lambda^2 |S|^2 \left| \text{Tr}(\Delta^* \Delta^{*\dagger}) + \text{Tr}(\Delta^c \Delta^{c\dagger}) + \text{Tr}(\Delta^\dagger \Delta^c) \right|,
$$

$$
V_{\text{soft}} = M_1^2 \text{Tr}(\Delta^* \Delta^* + \Delta^c \Delta^c) + M_2^2 \text{Tr}(\Delta^* \Delta^* + \Delta^c \Delta^c) \\
+ M_3^2 \Phi_1^\dagger \Phi_1 + M_4^2 \Phi_2^\dagger \Phi_2 + M_5^2 |S|^2 \\
+ \{A_{\lambda} \lambda S \text{Tr}(\Delta^* \Delta^* + \Delta^c \Delta^c) - C_{\lambda} M_R^2 S + h.c.\},
$$

$$
V_D = \frac{g_2^2}{8} \sum_i \left| \text{Tr}(2 \Delta^* \tau_i \Delta^* + 2 \Delta^c \tau_i \Delta^c + \Phi_a \tau_i^T \Phi_b^\dagger) \right|^2 \\
+ \frac{g_R^2}{8} \sum_i \left| \text{Tr}(2 \Delta^{c\dagger} \tau_i \Delta^c + 2 \Delta^{\dagger} \tau_i \Delta^c + \Phi_a \tau_i^T \Phi_b^\dagger) \right|^2 \\
+ \frac{g_2^2}{2} \left| \text{Tr}(\Delta^* \Delta^* + \Delta^c \Delta^c - \Delta^{c\dagger} \Delta^c + \Delta^{\dagger} \Delta^c) \right|^2.
$$

All terms in the scalar potential are identical for the configurations in which VEVs are given to the neutral right-handed triplet Higgs, or the charged Higgs, except for the $D-$term, which is lower for the charge breaking configuration. Previous solutions suggested are breaking $R-$parity, which would have the attractive feature that $v_R \sim 1$ TeV, but which abandons the LSP as the candidate for dark matter [4], or introducing higher dimensional operators to lower the charge conserving vacuum, with $v_R \sim 10^{11}$ GeV, but loosing the solution to strong and weak CP violation [12]. More recently, a new version of the model [11] examined effective operators which generate terms of the form $\text{Tr}(\Delta^c \Delta^c) \text{Tr}(\Delta^{c\dagger} \Delta^{\dagger})$ at one loop level, induced by the Majorana neutrino $\nu_c$ couplings with $\Delta^c$. These operators mimic the effects of the higher dimensional operators in previous versions, without the need to introduce them explicitly, thus solving the problem of the global minimum (if their coefficient is positive). The advantage of such a formalism is that the masses are very predictive, as they do not depend on coefficients of ad-hoc higher order terms, or sneutrino VEVs. In the next section, we study explicitly the implications for the Higgs masses in this model.
III. HIGGS BOSON COMPOSITION AND MASSES

The Higgs boson spectrum was previously analyzed in a variant of the model \[10\] with $R$–parity violation. The new features of the present analysis are 1) we employ a version of the model that uses the right-chiral neutrino couplings to the triplet Higgs bosons to eliminate the need for $L$-number violation; and 2) we include constraints from FCNC processes to predict the range of Higgs masses and parameters in LRSUSY. Effectively, we are looking at a very different model and Higgs sector than in \[10\].

We proceed the usual way to find the masses and mixing matrices for the Higgs bosons in this model. We minimize the Higgs potential given in the previous section, taking into account corrections induced by the heavy Majorana neutrino Yukawa couplings. This insures that the minimum of the potential is charge conserving. We forgo providing explicit expressions for the equations obtained by taking

$$\frac{\partial V}{\partial \kappa_1} = \frac{\partial V}{\partial \kappa_2} = \frac{\partial V}{\partial v_L} = \frac{\partial V}{\partial \bar{v}_L} = \frac{\partial V}{\partial v_R} = \frac{\partial V}{\partial \bar{v}_R} = \frac{\partial V}{\partial \langle S \rangle} = 0,$$

instead we give the relevant mass matrices for the Higgs fields. For simplicity, we use the abbreviations

$$\kappa_{\text{dif}}^2 = \kappa_1^2 - \kappa_2^2, \quad \text{(3.1)}$$
$$\rho_{\text{dif}}^2 = v_R^2 - \bar{v}_R^2 + \frac{1}{2} (\kappa_1^2 - \kappa_2^2), \quad \text{(3.2)}$$
$$Y = A_\lambda \lambda S + \lambda (-M_R^2 - 2\lambda_21\kappa_1\kappa_2 + \lambda v_R\bar{v}_R), \quad \text{(3.3)}$$
$$M = 2\lambda_21 (-M_R^2 - 2\lambda_21\kappa_1\kappa_2 + \lambda v_R\bar{v}_R), \quad \text{(3.4)}$$
$$f(\epsilon) = \epsilon (\frac{M}{2\lambda_21} - 2\lambda_21\kappa_1\kappa_2 - \epsilon \kappa_1\kappa_2), \quad \text{(3.5)}$$
$$g(\epsilon) = \epsilon \lambda \kappa_1 \kappa_2, \quad \text{(3.6)}$$
$$h(\epsilon) = \epsilon \kappa_1 \kappa_2 (4\lambda_21 + \epsilon), \quad \text{(3.7)}$$

with $\epsilon = \mu_{21} - \mu_{12}$, small but non-zero after symmetry breaking.

A. Doubly Charged Higgs Boson Masses

Mass matrices for the doubly charged Higgs fields are of block diagonal form of one two by two matrix for $(\delta^{++}, \bar{\delta}^{--})$ fields and one two by two matrix for $(\delta^{--}, \bar{\delta}^{++})$ fields.
respectively,

\[
M_{\delta^+ \bar{\delta}^-}^2 = \begin{pmatrix}
  \frac{1}{2} g_L^2 \kappa_1^2 - g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{M^2} Y' & Y' \\
  Y' & -\frac{1}{2} g_L^2 \kappa_1^2 + g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{M^2} Y'
\end{pmatrix},
\]

\[
M_{\delta^- \bar{\delta}^+}^2 = \begin{pmatrix}
  -2 g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{M^2} Y' & Y' \\
  Y' & 2 g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{M^2} Y'
\end{pmatrix},
\]

(3.8)

where \(Y' = Y - g(\epsilon)\). From these expressions we can find the exact analytic forms for the doubly charged Higgs masses. In the limit \(v_R \approx \bar{v}_R\), these are:

\[
M_{H^{++}}^2 \approx -Y' \pm \frac{1}{2} \sqrt{\left(\frac{1}{2} g_L^2 \kappa_1^2 - g_R^2 \rho_{\text{dif}}^2 \right)^2 + 2 Y'^2}
\]

\[
M_{H^{++}}^2 \approx -Y' \pm \frac{1}{2} \sqrt{4 g_R^4 \rho_{\text{dif}}^4 + 2 Y'^2}
\]

(3.9)

Thus, in all cases, the left-handed doubly charged Higgs fields are expected to be lighter than the right-handed ones.

B. Singly Charged Higgs Boson Masses

Mass matrices for the singly charged Higgs fields are of block diagonal form of one two by two matrix for \((\delta^+, \bar{\delta}^-)\) fields, one two by two matrix for \((\phi^+, \chi^-)\) fields and one four by four matrix for \((\delta^c-, \bar{\delta}^-c^+, \phi^c-, \chi^c^+)\) fields respectively,

\[
M_{\delta^+ \delta^-}^2 = \begin{pmatrix}
  -g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{M^2} Y' & Y' \\
  Y' & g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{M^2} Y'
\end{pmatrix},
\]

(3.10)

\[
M_{\phi^+ \chi^-}^2 = \begin{pmatrix}
  \frac{\kappa_1}{M'} & \frac{M'}{\kappa_1} M' \\
  M' & \frac{\kappa_2}{M'} M'
\end{pmatrix},
\]

(3.11)

where \(M' = M + f(\epsilon)\). The elements of the four by four matrix are

\[
M_{\delta^c- \delta^c}^2 = g_R^2 \bar{v}_R^2 = g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{\bar{v}_R} Y'
\]

(3.12)

\[
M_{\delta^c- \delta^c}^2 = -\frac{2}{v_R} v_R \bar{v}_R + Y'
\]

(3.13)

\[
M_{\delta^c- \phi^c}^2 = -\frac{1}{\sqrt{2}} g_R^2 \kappa_1 v_R
\]

(3.14)

\[
M_{\delta^c- \chi^c}^2 = -\frac{1}{\sqrt{2}} g_R^2 \kappa_2 v_R
\]

(3.15)

\[
M_{\delta^c+ \delta^c}^2 = \frac{g_R^2 v_R^2}{v_R} + g_R^2 \rho_{\text{dif}}^2 - \frac{v_R}{\bar{v}_R} Y'
\]

(3.16)
\[ M^2_{\delta^+ \phi_2^+} = \frac{1}{\sqrt{2}} g_R^2 \kappa_1 \bar{v}_R \] (3.17)
\[ M^2_{\delta^+ \chi_1^{++}} = \frac{1}{\sqrt{2}} g_R^2 \kappa_2 \bar{v}_R \] (3.18)
\[ M^2_{\phi_2^+ \phi_2^+} = \frac{1}{2} \kappa_1^2 (g_L^2 + g_R^2) - \frac{1}{\kappa_1} g_L^2 \kappa_1^2 - g_R^2 \rho_{dR}^2 + \frac{\kappa_2}{\kappa_1} M' \] (3.19)
\[ M^2_{\phi_2^+ \chi_1^{++}} = \frac{1}{2} \kappa_1 \kappa_2 (g_L^2 + g_R^2) + M' \] (3.20)
\[ M^2_{\chi_1^{++} \chi_1^{++}} = \frac{1}{2} \kappa_2^2 (g_L^2 + g_R^2) + \frac{1}{2} g_L^2 \kappa_2^2 + g_R^2 \rho_{dR}^2 + \frac{\kappa_1}{\kappa_2} M' \] (3.21)

C. Neutral Higgs Boson Masses

Mass matrices for the neutral scalar Higgs fields are of block diagonal form of one two by two matrix for \((\delta^0, \bar{\delta}^0)\) fields, one two by two matrix for \((\phi_1^0, \chi_1^0)\) fields and one five by five matrix for \((\delta^0, \bar{\delta}^0, \phi_1^0, \chi_2^0, S^0)\) fields respectively,

\[
M^2_{\delta^0,\bar{\delta}^0} = \begin{pmatrix}
-\frac{1}{2} g_L^2 \kappa_1^2 & -\frac{1}{2} g_R^2 \rho_{dR}^2 - \frac{\kappa_1}{\kappa_2} M' & Y' \\
Y' & \frac{1}{2} g_L^2 \kappa_1^2 + g_R^2 \rho_{dR}^2 \frac{\kappa_1}{\kappa_2} M' & \end{pmatrix}
\] (3.22)

The elements of the five by five matrix are

\[
M^2_{\delta^0,\bar{\delta}^0} = 2v_R^2 (g_{B-L}^2 + g_R^2) + \lambda^2 \bar{v}_R - \bar{v}_R Y' \] (3.23)
\[
M^2_{\delta^0,\bar{\delta}^0} = -2v_R \bar{v}_R (g_{B-L}^2 + g_R^2) + \lambda^2 \bar{v}_R \bar{v}_R + Y' \] (3.24)
\[
M^2_{\delta^0,\phi_1^0} = g_R^2 \kappa_1 v_R - 2\lambda \kappa_2 v_R \frac{\kappa_1}{\kappa_2} - \frac{2}{\kappa_1} \bar{v}_R g(\epsilon) \] (3.25)
\[
M^2_{\delta^0,\chi_2^0} = g_R^2 \kappa_2 v_R - 2\lambda \kappa_1 v_R \frac{\kappa_1}{\kappa_2} - \frac{2}{\kappa_1} \bar{v}_R g(\epsilon) \] (3.26)
\[
M^2_{\delta^0,\bar{S}^0} = 2\lambda^2 S v_R + A_\lambda \kappa_1 v_R \] (3.27)
\[
M^2_{\delta^0,\bar{S}^0} = 2(v_R^2 (g_{B-L}^2 + g_R^2) + \lambda^2 \bar{v}_R - \bar{v}_R Y' \] (3.28)
\[
M^2_{\delta^0,\phi_1^0} = -2g_R^2 \kappa_1 \bar{v}_R - 2\lambda \kappa_2 v_R \frac{\kappa_1}{\kappa_2} - \frac{2}{\kappa_1} \bar{v}_R g(\epsilon) \] (3.29)
\[
M^2_{\delta^0,\chi_2^0} = g_R^2 \kappa_2 \bar{v}_R - 2\lambda \kappa_1 v_R \frac{\kappa_1}{\kappa_2} - \frac{2}{\kappa_1} \bar{v}_R g(\epsilon) \] (3.30)
\[
M^2_{\delta^0,\bar{S}^0} = 2\lambda^2 S \bar{v}_R + A_\lambda \kappa_1 v_R \] (3.31)
\[
M^2_{\phi_1^0,\phi_1^0} = \frac{1}{2} \kappa_1^2 (g_L^2 + g_R^2) + 4\lambda^2 \kappa_2^2 + \frac{\kappa_2}{\kappa_1} [M' + h(\epsilon)] \] (3.32)
\[
M^2_{\phi_1^0,\chi_2^0} = -\frac{1}{2} \kappa_1 \kappa_2 (g_L^2 + g_R^2) + 4\lambda^2 \kappa_1 \kappa_2 - [M' - h(\epsilon)] \] (3.33)
The elements of the five by five matrix are

\[ M^2_{\phi^0_2S^0} = 0 \]  
\[ M^2_{\delta^{21}_{\bar{R}}} = \frac{1}{2} \kappa_2^2 (g_L^2 + g_R^2) + 4 \lambda^2_{21} \kappa_1^2 + \frac{\kappa_1}{\kappa_2} [M' + h(\epsilon)] \]  
\[ M^2_{\phi^0_1S^0} = 0 \]  
\[ M^2_{S^0S^0} = M^2_S + \lambda^2 (v_R^2 + \bar{v}_R^2) \]

Mass matrices for the neutral pseudoscalar Higgs fields are similar of block diagonal form of one two by two matrix for \((\delta^{0i}, \bar{\delta}^{0i})\) fields, one two by two matrix for \((\phi^0_2, \chi^0_2)\) fields and one five by five matrix for \((\delta^{0i}, \bar{\delta}^{0i}, \phi^0_1, \chi^0_2, S^0i)\) fields respectively,

\[
M^2_{\delta^{0i}\bar{\delta}^{0i}} = \begin{pmatrix} -\frac{1}{2} g^2_{2L} \kappa_{2f}^2 - g^2_{2R} \rho_{2f}^2 - \frac{\mu}{v_R} Y' & -Y' \\ -Y' & \frac{1}{2} g^2_{2L} \kappa_{2f}^2 + g^2_{2R} \rho_{2f}^2 - \frac{\mu}{v_R} Y' \end{pmatrix},
\]

\[
M^2_{\phi^0_1\chi^0_2} = \begin{pmatrix} -\frac{1}{2} g^2_{2L} \kappa_{2f}^2 - g^2_{2R} \rho_{2f}^2 + \frac{\sigma}{\alpha_1} M' & M' \\ M' & \frac{1}{2} g^2_{2L} \kappa_{2f}^2 + g^2_{2R} \rho_{2f}^2 + \frac{\sigma}{\alpha_2} M' \end{pmatrix}
\]

The elements of the five by five matrix are

\[
M^2_{\delta^{0i}\phi^0_1} = \lambda^2 \bar{v}_R^2 - \bar{v}_R^2 Y' \]  
\[ M^2_{\delta^{0i}\phi^0_1} = \lambda^2 v_R \bar{v}_R - Y' \]  
\[ M^2_{\delta^{0i}\phi^0_1} = -2 \lambda \lambda_2 \kappa_2 \bar{v}_R - \frac{\mu}{\kappa_1} g(\epsilon) \]  
\[ M^2_{\delta^{0i}\chi^0_2} = -2 \lambda \lambda_2 \kappa_1 \bar{v}_R - \frac{\mu}{\kappa_2} g(\epsilon) \]  
\[ M^2_{\delta^{0i}S^0i} = -A_\lambda \lambda \bar{v}_R \]  
\[ M^2_{\delta^{0i}\phi^0_1} = \lambda^2 v_R^2 - \frac{\mu}{v_R} Y' \]  
\[ M^2_{\delta^{0i}\phi^0_1} = -2 \lambda \lambda_2 \kappa_2 v_R - \frac{\mu}{\kappa_1} g(\epsilon) \]  
\[ M^2_{\delta^{0i}\chi^0_2} = -2 \lambda \lambda_2 \kappa_1 v_R - \frac{\mu}{\kappa_2} g(\epsilon) \]  
\[ M^2_{\delta^{0i}S^0i} = -A_\lambda \lambda v_R \]  
\[ M^2_{\phi^0_1\phi^0_1} = 4 \lambda^2_{21} \kappa_2^2 + \frac{\kappa_2}{\kappa_1} [M' + h(\epsilon)] \]  
\[ M^2_{\phi^0_1\chi^0_2} = 4 \lambda^2_{21} \kappa_2 \kappa_2 + [M' + h(\epsilon)] \]  
\[ M^2_{\phi^0_2S^0} = 0 \]  
\[ M^2_{\chi^0_2S^0} = 0 \]  
\[ M^2_{S^0S^0} = M^2_S + \lambda^2 (v_R^2 + \bar{v}_R^2) \]
IV. CONSTRAINTS ON THE HIGGS SECTOR

A. Flavor Changing Neutral Higgs Bosons

As any model with more than one Higgs doublet, the LRSUSY is plagued by tree-level FCNC-inducing Higgs bosons [13]. We proceed first by isolating the flavor-violating and flavor-conserving field combinations, then proceed to subject them to constraints coming from mixings in the kaon, B and D neutral meson states. We show more explicitly the expressions for the down-quark sector; the up-quark sector can be obtained simply by the same method. The Yukawa Lagrangian in the quark sector is given by

\[ L_Y = \bar{d}_L Y_u \phi^0 \phi^0 + \bar{d}_L Y_d \chi^0 \chi^0 + \bar{u}_L Y_u \phi^0 u_R + \bar{u}_L Y_d \chi^0 u_R + \text{h.c.}, \]

where \( Y_u \) and \( Y_d \) are \( 3 \times 3 \) Hermitian matrices in flavor space. When the bi-doublets acquire the VEV as in Table I, with \( \kappa_1, \kappa_2, \kappa'_1 \) and \( \kappa'_2 \) real, the up and the down type quark mass matrices are given by:

\[ M_u = Y_u \kappa_1 + Y_d \kappa'_2 \]
\[ M_d = Y_u \kappa'_1 + Y_d \kappa_2. \]

(4.2)

Inserting the expressions obtained for \( Y_u \) and \( Y_d \) in terms of masses, the Yukawa Lagrangian in the down type quark sector reads

\[ L^N_Y (d) = \frac{[d^s L M^u_{ij} d^j R (\kappa_2 \phi^0_2 - \kappa'_2 \phi^0_2) + d^s L M^d_{ij} d^j R (\kappa_1 \chi^0_2 - \kappa'_1 \phi^0_2)]}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} + \frac{[d^s R M^s u_{ij} d^j L (\kappa_2 \phi^0_2 - \kappa'_2 \phi^0_2) + d^s R M^s d_{ij} d^j L (\kappa_1 \chi^0_2 - \kappa'_1 \phi^0_2)]}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2}. \]

(4.3)

To obtain the physical states we diagonalize the mass matrices by the unitary transformations

\[ M^u_{ij} = U^i_k \hat{M}^u_{kl} W^j m_s \delta^{km}, \]
\[ M^d_{ij} = U^i_k \hat{M}^d_{kl} W^j m_s \delta^{km}, \]

(4.4)

where \( \hat{M}_u \) and \( \hat{M}_d \) are diagonal up and down type quark mass matrices. Since \( d_L \) and \( d_R \) are weak eigenstates, unitary transformations convert them into mass eigenstates

\[ d^j L \rightarrow U^{ij}_d d^j L, \]
\[ d^j R \rightarrow W^{ij}_d d^j R. \]

(4.5)
We define $U_d^{ij*} U_u^{ik} = V_L^{jk}$ and $W_u^{ij*} W_d^{jm} = V_R^{lm}$ where $V_L$ and $V_R$ are the components of the left-handed and right-handed CKM matrices. Then the Yukawa Lagrangian for down type quark fields is given by

$$\mathcal{L}_Y^N(d) = \frac{d_{\text{L}}^{n*} V_{\text{L}}^{kn*} \tilde{M}_u \delta_{\text{L}}^{km} (\kappa_2 \phi_0^0 - \kappa_2^{i*} \phi_0^0) + i d_{\text{L}}^{n*} \delta_{\text{L}}^{mk} \tilde{M}_d \delta_{\text{L}}^{kl} (\kappa_1 \chi_0^0 - \kappa_1^{i*} \phi_0^0)}{\kappa_2 \kappa_2' - \kappa_1^{i*} \kappa_1'} + \frac{d_{\text{R}}^{mn*} V_{\text{R}}^{km*} \tilde{M}_u \delta_{\text{R}}^{kl} (\kappa_2 \phi_0^0 - \kappa_2^{i*} \phi_0^0)}{\kappa_1^{i*} \kappa_1^2 - \kappa_1^{i*} \kappa_1'} + \frac{i d_{\text{R}}^{mn} \delta_{\text{R}}^{km} \tilde{M}_d \delta_{\text{R}}^{kl} (\kappa_1 \chi_0^0 - \kappa_1^{i*} \phi_0^0)}{\kappa_1^{i*} \kappa_1^2 - \kappa_1^{i*} \kappa_1'},$$

(4.6)

where the up and down mass matrices are Hermitian since the VEVs of bi-doublets are taken to be real. For simplicity, we assume $V_L = V_R = V$. The fields $\phi_0^0$ and $\chi_0^0$ are complex. Thus we can isolate two terms in the Lagrangian, one flavor violating, and one FCNC-conserving.

Writing the neutral and imaginary parts separately, the FCNC Lagrangian reads

$$\mathcal{L}_{\text{FCNC}}(d) = \frac{d_{\text{L}}^{n*} V_{\text{L}}^{kn*} \tilde{M}_u \delta_{\text{L}}^{km} (\kappa_2 \phi_0^{0i} - \kappa_2^{i*} \phi_0^{0i}) + i d_{\text{L}}^{n*} \delta_{\text{L}}^{mk} \tilde{M}_d \delta_{\text{L}}^{kl} (\kappa_1 \phi_0^{0i} - \kappa_1^{i*} \chi_0^{0i})}{\kappa_2 \kappa_2' - \kappa_1^{i*} \kappa_1'} + \frac{d_{\text{R}}^{mn*} V_{\text{R}}^{km*} \tilde{M}_u \delta_{\text{R}}^{kl} (\kappa_2 \phi_0^{0i} - \kappa_2^{i*} \phi_0^{0i}) - i d_{\text{R}}^{mn} \delta_{\text{R}}^{km} \tilde{M}_d \delta_{\text{R}}^{kl} (\kappa_1 \phi_0^{0i} - \kappa_1^{i*} \chi_0^{0i})}{\kappa_1^{i*} \kappa_1^2 - \kappa_1^{i*} \kappa_1'},$$

(4.7)

where $\phi_2^{0i}$ and $\chi_2^{0i}$ are the two of the nine bare scalar fields and $\phi_2^{0i}$ and $\chi_2^{0i}$ are the two of the nine bare pseudo-scalar fields appearing in LRSUSY Lagrangian. The $d - s$ coupling in Eq. (4.7) allows a $\Delta S = 2$ transition at tree level. To evaluate explicitly, we use the Wolfenstein parametrization with every parameter expanded as a power series in the parameter $\lambda = |V_{us}| = 0.2246 \pm 0.0012$ [14].

$$V = \begin{pmatrix}
1 - \frac{\lambda^2}{2} & \lambda & A \lambda^3 (\rho - i \eta) \\
-\lambda & 1 - \frac{\lambda^2}{2} & A \lambda^2 \\
A \lambda^3 (1 - \rho - i \eta) & -A \lambda^2 & 1
\end{pmatrix} + \mathcal{O}(\lambda^4).$$

(4.8)

For $\lambda = 0.2246, A = 0.832, \rho = 0.130$, and $\eta = 0.350$ [15, 16],

$$V^{kd} \tilde{M}_u^{kk} V^{ks} = (m_u - m_c)(\lambda - \frac{\lambda^3}{2}) - m_t A^2 \lambda^5 (1 - \rho + i \eta).$$

(4.9)

We express the bare scalar $\psi_0^{0T} = (\phi_0^{0r} \phi_0^{0r} \phi_0^{0r} \phi_0^{0r} \phi_0^{0r} \chi_0^{0r} \chi_0^{0r} \chi_0^{0r} \chi_0^{0r})$ and pseudoscalar Higgs fields $\psi_0^{iT} = (\phi_0^{0r} \phi_0^{0r} \phi_0^{0r} \phi_0^{0r} \phi_0^{0r} \chi_0^{0r} \chi_0^{0r} \chi_0^{0r} \chi_0^{0r})$ as physical CP even Higgs fields $H_0^{0T} = (H_1^{0r} H_2^{0r} H_3^{0r} H_4^{0r} H_5^{0r} H_6^{0r} H_7^{0r} H_8^{0r} H_9^{0r})$ and physical CP odd Higgs fields $H_0^{0T} =$
\[ (H_1^{0i} H_2^{0i} H_3^{0i} H_4^{0i} H_5^{0i} H_6^{0i} H_7^{0i} H_8^{0i} H_9^{0i}) \]. Call \( A_{ij} \) the transformation matrix which transforms the bare scalar fields into the physical CP even ones, and \( B_{ij} \) matrix which transforms the bare pseudo-scalar fields into the physical CP odd ones: \( H_i^{0r} = A_{ij} \psi_j^{0r} \), \( H_i^{0i} = B_{ij} \psi_j^{0i} \) and substituting these into the Eq. \( (4.7) \), we obtain the explicit Lagrangian responsible for FCNC in the down-sector

\[
\mathcal{L}_{\Delta S=2}^{\text{FCNC}}(d) = \frac{m_t \lambda}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} \left( \left[ \left( \frac{m_u}{m_t} - \frac{m_c}{m_t} \right) \left( 1 - \frac{\lambda^2}{2} \right) - A^2 \lambda^4 \left( 1 - \rho \right) \right] (\kappa_2 A_{i6}^* - \kappa'_2 A_{i8}^*) H_i^{0r} \right)
\times (d_P R s + d_P L s) + A^2 \lambda^4 \eta (\kappa_2 B_{i6}^* - \kappa'_2 B_{i8}^*) H_i^{0i} (d_P R s - d_P L s) \right) + \frac{-i m_t \lambda}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} \left( \left[ \left( \frac{m_u}{m_t} - \frac{m_c}{m_t} \right) \left( 1 - \frac{\lambda^2}{2} \right) - A^2 \lambda^4 \left( 1 - \rho \right) \right] (\kappa_2 B_{i6}^* - \kappa'_2 B_{i8}^*) H_i^{0i} \right)
\times (d_P R s - d_P L s) - A^2 \lambda^4 \eta (\kappa_2 A_{i6}^* - \kappa'_2 A_{i6}^*) H_i^{0r} (d_P R s + d_P L s) . \tag{4.10}
\]

We proceed in similar fashion to evaluate the flavor-conserving and flavor-violating Higgs contributions to the up sector. The Yukawa Lagrangian for the up quark sector is

\[
\mathcal{L}_Y^N(u) = u_L^i Y_{ij}^u \phi_1^0 u_R^j + u_L^i Y_{ij}^d \lambda_1^0 u_R^j + u_L^i Y_{ij}^{0s} \phi_1^{0s} y_{ij}^s u_R^j + u_L^i Y_{ij}^{0v} \phi_1^{0v} y_{ij}^v u_R^j . \tag{4.11}
\]

We use the same substitutions as before and express the Lagrangian in terms of the complex fields \( \phi_2^0 \) and \( \chi_2^0 \). The first and third terms in the Lagrangian above are flavor-conserving.

Writing the neutral and imaginary parts separately, the FCNC Lagrangian reads

\[
\mathcal{L}_{\Delta C=2}^{\text{FCNC}}(u) = \frac{u_L^{ns} V^{nk} \tilde{M}_d^{kk} V^{lk*} Y_{ij}^u (\kappa_2 \phi_1^{0r} - \kappa'_2 \chi_1^{0r})}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} + \frac{i u_L^{ns} V^{nk} \tilde{M}_d^{kk} V^{lk*} Y_{ij}^d (\kappa_2 \phi_1^{0i} - \kappa'_2 \chi_1^{0i})}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} + \frac{i u_R^{ns} V^{nk} \tilde{M}_d^{kk} V^{lk*} Y_{ij}^{0s} u_R^j (\kappa_2 \phi_1^{0r} - \kappa'_2 \chi_1^{0r})}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} - \frac{i u_R^{ns} V^{nk} \tilde{M}_d^{kk} V^{lk*} Y_{ij}^{0v} u_R^j (\kappa_2 \phi_1^{0i} - \kappa'_2 \chi_1^{0i})}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} , \tag{4.12}
\]

where \( \phi_1^{0r} \) and \( \chi_1^{0r} \) are the two of the nine bare scalar fields and \( \phi_1^{0i} \) and \( \chi_1^{0i} \) are the two of the nine bare pseudo-scalar fields appearing in LRSUSY Lagrangian. The \( u - c \) coupling in Eq. \( (4.12) \) allows a \( \Delta C = 2 \) transition at tree level. Inserting \( V^{uk} \tilde{M}_u^{kk} V^{ck*} \) in terms of Wolfenstein parameters,

\[ V^{uk} \tilde{M}_u^{kk} V^{ck*} = (m_s - m_c) (\lambda - \frac{\lambda^3}{2}) - m_b A^2 \lambda^5 (-\rho + i \eta) , \tag{4.13} \]

and using physical states instead of \( \phi_1^{0r} \) and \( \chi_1^{0r} \) we obtain the explicit form of the Lagrangian responsible for FCNC in the up-sector

\[
\mathcal{L}_{\Delta C=2}^{\text{FCNC}}(u) = \frac{m_b \lambda}{\kappa_1 \kappa_2 - \kappa'_1 \kappa'_2} \left( \left[ \left( \frac{m_s}{m_b} - \frac{m_c}{m_b} \right) \left( 1 - \frac{\lambda^2}{2} \right) + A^2 \lambda^4 \rho \right] (\kappa_2 A_{i5}^* - \kappa'_2 A_{i7}^*) H_i^{0r} \right) .
\]
These expressions will be used to calculate the real and imaginary parts of the $K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$ and $B^0 - \bar{B}^0$ mixing.

B. $\epsilon_K$ and $K^0 - \bar{K}^0$ Mixing

We evaluate the real and imaginary parts of the $K^0 - \bar{K}^0$ transition. We assume a common mass for scalar and pseudoscalar Higgs fields.

\[
\Re\langle K^0|H_{\text{eff}}|K^0\rangle = \frac{m_i^2 \lambda^2}{4 M_1^2(\kappa_1 \kappa_2 - \kappa_1' \kappa_2')} \left\{ \left[ \frac{m_u}{m_t} - \frac{m_c}{m_t} \right] (2 - \lambda^2) - 2 A^2 \lambda^4 (1 - \rho) \right\} 
\times \left( \left[ (\kappa_2 A_{i6}^* - \kappa_2' A_{i8}^*)^2 - (\kappa_2 B_{i6}^* - \kappa_2' B_{i8}^*)^2 \right] (\langle \tilde{Q}_1(\mu) \rangle + \langle Q_1(\mu) \rangle) 
+ \left[ (\kappa_2 A_{i6}^* - \kappa_2' A_{i8}^*)^2 + (\kappa_2 B_{i6}^* - \kappa_2' B_{i8}^*)^2 \right] (\langle \tilde{Q}_2(\mu) \rangle + \langle Q_2(\mu) \rangle) \right) + 4 A^4 \lambda^8 \eta^2 \left( \left[ (\kappa_2 A_{i6}^* - \kappa_2' A_{i8}^*)^2 - (\kappa_2 B_{i6}^* - \kappa_2' B_{i8}^*)^2 \right] (\langle \tilde{Q}_1(\mu) \rangle + \langle Q_1(\mu) \rangle) 
+ \left[ (\kappa_2 A_{i6}^* - \kappa_2' A_{i8}^*)^2 + (\kappa_2 B_{i6}^* - \kappa_2' B_{i8}^*)^2 \right] (\langle \tilde{Q}_2(\mu) \rangle + \langle Q_2(\mu) \rangle) \right),
\]

(4.15)

and

\[
\Im\langle K^0|H_{\text{eff}}|\bar{K}^0\rangle = \frac{im_i^2 \lambda^2}{4 M_1^2(\kappa_1 \kappa_2 - \kappa_1' \kappa_2')} \left[ \frac{m_u}{m_t} - \frac{m_c}{m_t} \right] (2 - \lambda^2) A^2 \lambda^4 \eta - 2 A^4 \lambda^8 (1 - \rho) \eta 
\times \left( \left[ (\kappa_2 B_{i6}^* - \kappa_2' B_{i8}^*)^2 - (\kappa_2 A_{i6}^* - \kappa_2' A_{i8}^*)^2 \right] (\langle \tilde{Q}_1(\mu) \rangle + \langle Q_1(\mu) \rangle) 
- \left[ (\kappa_2 B_{i6}^* - \kappa_2' B_{i8}^*)^2 + (\kappa_2 A_{i6}^* - \kappa_2' A_{i8}^*)^2 \right] (\langle \tilde{Q}_2(\mu) \rangle + \langle Q_2(\mu) \rangle) \right),
\]

(4.16)

The quantities $Q_1, Q_2, \tilde{Q}_1$, and $\tilde{Q}_2$ are four quark operators and are given by

\[
Q_1 = (q_1^a P_L q_2^a) \otimes (q_1^\beta P_L q_2^\beta), \quad \tilde{Q}_1 = (q_1^a P_R q_2^a) \otimes (q_1^\beta P_R q_2^\beta), \\
Q_2 = (q_1^a P_L q_2^a) \otimes (q_2^\beta P_R q_2^\beta), \quad \tilde{Q}_2 = (q_1^a P_R q_2^a) \otimes (q_1^\beta P_L q_2^\beta),
\]

(4.17)

where $\alpha$ and $\beta$ are the color indices. The matrix elements are,

\[
\langle Q_1(\mu) \rangle = -\frac{5}{24} \left( \frac{m_a}{m_{q_1}(\mu) + m_{q_2}(\mu)} \right)^2 m_a F_a^2 B_1(\mu),
\]

where $m_a$ and $F_a$ are the mass and decay constant of the $a$th flavor quark.
\[ \langle Q_2(\mu) \rangle = \frac{1}{4} \left( \frac{m_a}{m_{q_1}(\mu) + m_{q_2}(\mu)} \right)^2 m_a F_a^2 B_2(\mu) \]  

(4.18)

where \( a = K, B_d, B_s, D \) mesons, and no summation is assumed. \( F_a \) is the decay constant of the corresponding meson and \( B_1(\mu) \) and \( B_2(\mu) \) are the bag parameters calculated in NDR scheme for an energy scale \( \mu \). The numerical values for all the parameters involved in the calculation of \( K^0 - \bar{K}^0 \), \( D^0 - \bar{D}^0 \) and \( B^0_{d,s} - \bar{B}^0_{d,s} \) mixings are summarized in Table II and the quark mass values in Table III. Same expressions for the operators \( Q_1 \) and \( Q_2 \) are valid for the operators \( \bar{Q}_1 \) and \( \bar{Q}_2 \). Substituting \( \mu = 2 \) GeV in the expressions for \( \Delta M_K \) and CP violating parameter \( \epsilon_K \) given below

\[ \Delta M_K = 2 \text{Re} \langle \bar{K}^0 | H_{eff} | K^0 \rangle, \quad \Delta \epsilon_K = \frac{1}{\sqrt{2} \Delta M_K} \text{Im} \langle \bar{K}^0 | H_{eff} | K^0 \rangle, \]  

(4.19)

| \( \mu \) | \( m_\mu \) | \( m_d \) | \( m_s \) | \( m_b \) |
|---|---|---|---|---|
| \( 2 \) GeV | \( 2 \) GeV | \( 2 \) GeV | | |
| \( q_1 \) | \( s \) | \( b \) | \( b \) | \( u \) |
| \( q_2 \) | \( d \) | \( d \) | \( s \) | \( c \) |
| \( m_a \) | \( 498 \) MeV | \( 5.28 \) GeV | \( 5.37 \) GeV | \( 1.86 \) GeV |
| \( F_a \) | \( 160 \) MeV | \( 0.21 \) GeV | \( 0.25 \) GeV | \( 232 \) MeV |
| \( B_1(\mu) \) | \( 0.76 \) | \( 0.82 \) | \( 0.83 \) | \( 1 \) |
| \( B_2(\mu) \) | \( 1.30 \) | \( 1.16 \) | \( 1.17 \) | \( 1 \) |

**TABLE II.** QCD parameters used for meson mixing

| \( m_u \) (2 GeV) | \( m_d \) (2 GeV) | \( m_s \) (2 GeV) |
|---|---|---|
| \( 2.49^{+0.81}_{-0.79} \) MeV | \( 5.05^{+0.75}_{-0.95} \) MeV | \( 101^{+29}_{-21} \) MeV |
| \( m_c(m_c) \) | \( m_b(m_b) \) | \( m_t(m_t) \) |
| \( 1270^{+70}_{-90} \) MeV | \( 4190^{+180}_{-60} \) MeV | \( (172 \pm 0.9 \pm 1.3) \times 10^3 \) MeV |

**TABLE III.** Quark masses
we get
\[ \Delta M_K = \frac{6.9269 \times 10^{-7} A_{16}^{2*} + 2.0088 \times 10^{-7} B_{16}^{2*}}{M_i^2} (1 + \tan \beta^2), \] (4.20)

and
\[ \epsilon_K = \frac{9.9975 \times 10^6 A_{16}^{2*} - 9.8616 \times 10^9 A_{16}^{*} B_{16}^{*} + 2.8993 \times 10^7 B_{16}^{2*}}{M_i^2} (1 + \tan \beta^2). \] (4.21)

By comparing the calculated expressions with their experimental values, we obtain on the sources of flavor and CP violation in the LRSUSY.

The experimental value for the mass difference of \( K_L \) and \( K_S \) is given by [21]
\[ |\Delta M_K| = M_{K_L} - M_{K_S} = (3.483 \pm 0.006) \times 10^{-12} \text{ MeV} \] (4.22)
and indirect CP violation in \( K \to \pi \pi \) [22] and in \( K \to \pi \ell \nu \) decays is given by [21]
\[ |\epsilon_K| = (2.228 \pm 0.011) \times 10^{-3} \] (4.23)

We give below the analytical expressions for the constraints on the parameters in the neutral scalar and pseudoscalar mixing from \( K \) meson mixing. Taking the lightest neutral Higgs mass to be \( M_{H^0} = M_{H^0} = M_i \), the value of \( \Delta M_K = 3.483 \times 10^{-15} \text{ GeV} \) yields the constraint
\[ M_i^2 \geq (1.9888 \times 10^8 A_{16}^{2*} + 5.7675 \times 10^8 B_{16}^{2*}) (1 + \tan \beta^2) \text{ GeV}^2 \] (4.24)
while the value of \( \epsilon_K = 2.228 \times 10^{-3} \) [21] yields the constraint
\[ M_i^2 \geq (4.4872 \times 10^9 A_{16}^{2*} - 4.4262 \times 10^{-6} A_{16}^{*} B_{16}^{*} + 1.3013 \times 10^{10} B_{16}^{2*}) (1 + \tan \beta^2) \text{ GeV}^2 \] (4.25)

In the above expressions we assumed that the lightest Higgs mass provides the dominant contribution, and neglected the rest, while in our numerical evaluations we have summed over all mass contributions, as in (4.20) and (4.21). These become, for example, when \( \tan \beta = 10 \)
\[ M_i^2 \geq (2.0087 \times 10^{10} A_{16}^{2*} + 5.8251 \times 10^{10} B_{16}^{2*}) \text{ GeV}^2 \] (4.26)
and
\[ M_i^2 \geq (4.5320 \times 10^{11} A_{16}^{2*} - 4.4704 \times 10^{-4} A_{16}^{*} B_{16}^{*} + 1.3143 \times 10^{12} B_{16}^{2*}) \text{ GeV}^2 \] (4.27)

We tried varying the lightest relative masses in the scalar and pseudoscalar sector and found that the results do not change.
C. \( B^0_d - \bar{B}^0_d \) Mixing

We proceed the same way as for \( K^0 - \bar{K}^0 \) mixing to evaluate the constraints from the \( B^0_d, B^0_s \) meson mixing. We use again four quark operators \( Q_1, Q_2, \bar{Q}_1, \) and \( \bar{Q}_2 \) defined previously. Setting as before the Higgs mass to be equal to the lightest scalar mass \( M_{H^0} = M_i \) the expression for \( \Delta M_{B_d} \) becomes

\[
\Delta M_{B_d} = \frac{(9.4139 \times 10^{-6} A_{i6}^2 + 3.6405 \times 10^{-5} B_{i6}^2)(1 + \tan^2 \beta)}{M_i^2} \text{ GeV}^3 \quad (4.28)
\]

Using the experimental value of \( \Delta M_{B_d} = 3.337 \times 10^{-13} \text{ GeV} \) [21], we obtain, assuming as before dominance by the lightest mass

\[
M_i^2 \geq (2.8211 \times 10^7 A_{i6}^2 + 1.6909 \times 10^8 B_{i6}^2)(1 + \tan^2 \beta) \text{ GeV}^2 \quad (4.29)
\]

which becomes, for \( \tan \beta = 10 \)

\[
M_i^2 \geq (2.8493 \times 10^9 A_{i6}^2 + 1.1019 \times 10^{10} B_{i6}^2) \text{ GeV}^2 \quad (4.30)
\]

D. \( B^0_s - \bar{B}^0_s \) Mixing

We proceed exactly as in the previous subsection, substituting \( s \) instead of \( d \) quark. The parameters for \( B^0_s - \bar{B}^0_s \) mixing are given in Table II.

\[
\Delta M_{B_s} = \frac{(4.2314 \times 10^{-4} A_{i6}^2 + 1.6469 \times 10^{-3} B_{i6}^2)(1 + \tan^2 \beta)}{M_i^2} \text{ GeV}^3 \quad (4.31)
\]

Using the experimental value of \( \Delta M_{B_d} = 117 \times 10^{-13} \text{ GeV} \) [21, 23]

\[
M_i^2 \geq (3.6166 \times 10^7 A_{i6}^2 + 1.4076 \times 10^8 B_{i6}^2)(1 + \tan^2 \beta) \text{ GeV}^2 \quad (4.32)
\]

or, for \( \tan \beta = 10 \)

\[
M_i^2 \geq (3.6528 \times 10^9 A_{i6}^2 + 1.4217 \times 10^{10} B_{i6}^2) \text{ GeV}^2 \quad (4.33)
\]

E. \( D^0 - \bar{D}^0 \) Mixing

In subsection A, we evaluated the real and imaginary parts of the \( D^0 - \bar{D}^0 \) transition. We assume as before a common mass for scalar and pseudo-scalar Higgs fields.

\[
\text{Re}\langle \bar{D}^0 | H_{eff} | D^0 \rangle = \frac{m_b^2 \lambda^2}{4M_i^2(\kappa_1 \kappa_2 - \kappa_1' \kappa_2')^2} \left\{ \left[ \frac{m_b}{m_b} - \frac{m_d}{m_b} \right](2 - \lambda^2) + 2A^2 \lambda^4 \rho \right\}^2
\]
Comparing the calculated expression with the experimental value \[21\] we obtain

\[
\times \left[\left((\kappa_2 A^*_{i5} - \kappa'_2 A^*_{i7})^2 - (\kappa_2 B^*_{i5} - \kappa'_2 B^*_{i7})^2\right)\left(\langle \tilde{Q}_1(\mu) \rangle + \langle Q_1(\mu) \rangle\right)
\right.
\]
\[
+ \left[\left((\kappa_2 A^*_{i5} - \kappa'_2 A^*_{i7})^2 + (\kappa_2 B^*_{i5} - \kappa'_2 B^*_{i7})^2\right)\left(\langle \tilde{Q}_2(\mu) \rangle + \langle Q_2(\mu) \rangle\right)\right]
\]
\[
+ 4A^4\lambda_1^2\eta^2 \left[\left((\kappa_2 A^*_{i5} - \kappa'_2 A^*_{i7})^2 + (\kappa_2 B^*_{i5} - \kappa'_2 B^*_{i7})^2\right)\left(\langle \tilde{Q}_1(\mu) \rangle + \langle Q_1(\mu) \rangle\right)\right]
\]
\[
\left.\left.\left.+ \left[\left((\kappa_2 A^*_{i5} - \kappa'_2 A^*_{i7})^2 - (\kappa_2 B^*_{i5} - \kappa'_2 B^*_{i7})^2\right)\left(\langle \tilde{Q}_2(\mu) \rangle + \langle Q_2(\mu) \rangle\right)\right]\right\}\right]
\]
\[\tag{4.34}\]

where \(Q_1, Q_2, \tilde{Q}_1,\) and \(\tilde{Q}_2\) are the four quark operators defined as before, the mass difference \(\Delta M_D = 2\text{Re}(\tilde{D}^0|H_{eI}|D^0)\) is obtained as

\[
\Delta M_D = \frac{5.2816 \times 10^{-10} A^2_{27} + 5.8097 \times 10^{-9} B^2_{27} (1 + \tan^2 \beta)}{\tan \beta^2} \text{GeV}^3. \tag{4.35}\]

Comparing the calculated expression with the experimental value \[21\]

\[|\Delta M_D| = M_{D_9} - M_{D_0} = (1.57313) \times 10^{-17} \text{MeV}, \tag{4.36}\]

we obtain

\[
M_i^2 \geq \frac{(3.3574 \times 10^{10} A^2_{27} + 3.6931 \times 10^{11} B^2_{27})(1 + \tan^2 \beta)}{\tan \beta^2} \text{GeV}^2, \tag{4.37}\]

which becomes for \(\tan \beta = 10,\)

\[
M_i^2 \geq (3.3909 \times 10^{10} A^2_{27} + 3.7300 \times 10^{11} B^2_{27}) \text{GeV}^2. \tag{4.38}\]

V. NUMERICAL RESULTS AND DISCUSSION

The FCNCs tree-level diagrams are mediated by the physical scalar fields \(H^0_3\) and \(H^0_9\), and the pseudoscalars \(A^0_2\) and \(A^0_7\). These fields are linear superpositions of the \(\chi^0r\) or \(\phi^0r\) (\(\chi^0i\) or \(\phi^0i\), respectively, for the pseudoscalars) components from the bidoublet Higgs.

As the fields \(H^0_3\) and \(H^0_9\) must be heavy, the light neutral scalars would likely be linear combinations of the complimentary \(\chi^0r\) or \(\phi^0r\) components from the bidoublets, or \(\delta^0r, \delta^0r, \delta^0r,\) and \(\delta^0r\) from the triplet Higgs. We set \(v_R\) in the interval obtained from the requirement that the doubly charged Higgs are light \((3 - 10 \text{ TeV})\). Varying \(v_R\) outside this range adversely affects the masses of the lightest doubly charged Higgs, and some of the light neutral and singly charged scalars.

The mass of the lightest scalar field \(H^0_1\) (the SM-like) changes at most a few GeV, if we vary any of the parameters, whereas the second lightest scalar field \(H^0_2\) is highly dependent
on the changes in the parameter $v_R$. Both of these fields can be light, as our numerical explorations indicate. Similarly, the lightest pseudoscalar field $A^0_1$ behaves like the second lightest neutral scalar field and is also effected by the changes in $v_R$. $H^0_1$ is SM-like, and the parameter that seems to affect $H^0_1$ mass the most is the $\lambda_{21}$ coupling. (This parameter is the coupling that generates the $\mu_{21} = \lambda_{21} \langle S \rangle$ Higgsino coupling). The dependence is not smooth, but varying $\lambda_{21}$ in the interval $0.01 - 1$ produces a 30% change in $M_{H^0_1}$.

The tree-level flavor-changing neutral currents in the down-quark sector are governed by $H^0_9$ and $A^0_7$. The mass values of the fields $H^0_9$ and $A^0_7$ are the same, and they are dependent on the parameters $\lambda_{21}$, $v_R$, $\tan \beta$ and $M_R$. Numerical investigation reveals that only $\tan \beta$, $M_R$ and $\lambda_{21}$ can affect the $H^0_9$ and $A^0_7$ masses significantly. For instance, if $\lambda_{21}$ increases from 0.01 to 1, the $H^0_9$ mass increases almost 10 times. These masses are also slightly dependent on the parameters $M_R$ and $\tan \beta$ such that when they increase, mass values of these physical fields also increase. The dependence of the $H^0_9$ mass on the parameter $M_R$ is more dominant than on $\tan \beta$. Requiring $M_R \sim 100$ TeV insures that Higgs-mediated FCNCs in K and B neutral mesons are suppressed to levels consistent with experimental data. The variations of $H^0_9$ mass with these parameters are shown in Fig. 1.

The fields $H^0_3$ and $A^0_2$ are responsible for flavor-changing neutral currents in the up-quark sector. Their masses are the same (as one can infer from the mass matrices in Section (III)), and although they depend in principle on $v_R$, $\tan \beta$ and $\lambda_{21}$, the only significant dependence is on $v_R$, such that if $v_R$ increases from 3 to 10 TeV, their mass values increase approximately 5 times. The mass also varies with the ratio $\tan \delta = \bar{v}_R/v_R$, while almost independent of the changes in the other parameters. The parameter dependence is shown in Fig. 2 where we plot the explicit $v_R$ dependence for three values of $\tan \delta$, as well as a more extensive illustration of the $v_R - \bar{v}_R$ dependence in a contour plot. $D^0 - \bar{D}^0$ mixing constraints require $v_R \geq 3$ TeV. While the dependence on both $\tan \beta$ and $\lambda_{21}$ is very weak, the dependence on $v_R$ is almost linear.

From the approximate analytical expressions in Section (III), the mass of the lightest doubly charged physical field $H^\pm \pm_1$ depends on $v_R$, $\lambda$ and $M_R$. Analysis shows that only the dependence on $v_R$ is significant. However, the exact mass also depends on $\bar{v}_R$ through the ratio $\tan \delta = \bar{v}_R/v_R$. As before we show, in Fig. 3 the dependence of these two parameters as a contour plot in the $v_R - \bar{v}_R$ plane. The mass of $H^\pm \pm_1$ increases with the increasing values of $v_R$, as shown on the right hand side of Fig. 3 for three values of $\tan \delta$, while it is basically
FIG. 1. The variation of the FCNC neutral Higgs $H_0^0$ mass with the parameters of the LRSUSY model. $H_0^0$ induces tree-level FCNC in the down-quark sector. Shown are: contour plots in the $M_R - \tan \beta$ plane, the variation of $M_{H_0^0}$ with $M_R$, and with $\lambda_{21}$, for two values of $\tan \beta$. Masses are given in GeV.

FIG. 2. The variation of the FCNC neutral Higgs $H_3^0$ mass with the parameters of the LRSUSY model. $H_3^0$ induces tree-level FCNC in the up-quark sector. To the left, a contour plot in the $v_R - \bar{v}_R$ plane and, at the right, as a function of $v_R$ for three values of $\tan \delta = \bar{v}_R/v_R$. Masses are given in GeV.

independent on $M_R$. One can see that for $v_R \sim 3.5$ TeV the doubly charged Higgs boson mass is light for all values of $\tan \delta$, while for $v_R = 10$ TeV the mass is highly dependent on $\bar{v}_R/v_R$. For example, when we change $v_R$ from 3 to 10 TeV, the $H_1^{\pm \pm}$ mass values increase approximately 4 times. The effect of varying the other parameters is negligible for the
FIG. 3. The masses of the lightest doubly-charged Higgs boson as a contour plot in the $v_R - \bar{v}_R$ plane (left) and as a function of $v_R$ for three values of $\tan \delta = \bar{v}_R/v_R$ (right). Masses are given in GeV.

The lightest doubly charged Higgs, whereas the mass of the heavier doubly charged Higgs $H_2^{\pm\pm}$ depends almost exclusively on $M_R$.

The lightest singly charged physical field $H_1^\pm$ mass can also depend in principle on $v_R$, $\lambda$, $\tan \beta$ and $M_R$. However, upon inspection, the only significant dependence is on $v_R$, much like the doubly charged Higgs boson. The reason is that, requiring one doubly charged Higgs to be light makes the lightest singly charged Higgs to be an eigenstate of the $2 \times 2$ mass matrix $M_{\delta+\delta-}^2$. There is a slight difference between the lightest singly charged and lightest doubly charged Higgs boson, but the difference comes from SM-like observables in $g_L^2\kappa_{dij}^2$, and it is overwhelmed by parameters proportional to $v_R$ and $M_R$. Its mass also increases with the increasing values of $v_R$, and depends on the ratio $\bar{v}_R/v_R$ in much the same way as the mass of the doubly charged Higgs does. Thus we do not show the dependence separately.

Finally, we present two explicit numerical scenarios for the Higgs masses, which obey the constraints from meson mixings: one for $v_R = 3.5$ TeV and $\tan \beta = 10$, the other for $v_R = 5$ TeV and $\tan \beta = 50$. The other parameters in both scenarios are taken to be $\tan \delta \equiv \bar{v}_R/v_R = 1/1.05$, $M_R = 100$ TeV, $\lambda = 1$, $\lambda_{21} = 1$, $C_\lambda = 2.5$ TeV, $\langle S \rangle = 1$ TeV, $M_S = 1$ TeV. We give masses and compositions in terms of the bare states. One can see that, except for raising the lightest neutral Higgs mass, increasing $\tan \beta$ has little effect on
the spectrum. However raising $v_R$ increases the mass of the lighter non-SM-like Higgs bosons in the neutral scalar and pseudoscalar sector, as well as in the singly and doubly charged Higgs sectors. While we did not prove in general that the model conserves $R$-parity, the numerical results obtained from minimizing the masses confirm the results of [11]. Both of these scenarios allow for a pair of light flavor-conserving neutral scalar Higgs bosons (one SM-like, one mostly triplet $SU(2)_L$); as well as for one light singly charged Higgs and a pair of doubly charged Higgs bosons. The FCNC Higgs responsible for mixing in the up ($D^0 - \bar{D}^0$) or down ($K^0 - \bar{K}^0$ and $B^0_{d,s} - \bar{B}^0_{d,s}$) quark sectors are heavy and satisfy the experimental constraints in each sector.

VI. SUMMARY AND CONCLUSION

We analyzed the Higgs sector of a minimal left-right supersymmetric model with automatic $R$-parity violation. Symmetries of the model forbid explicit $R$-parity violation. Inclusion of the effects of the Yukawa coupling of the heavy Majorana neutrino insures a global minimum which is charge conserving, thus avoiding spontaneous $R$-parity breaking or the need to introduce higher dimensional terms.

The Higgs sector contains four doubly charged Higgs, six singly charged Higgs fields, nine neutral scalar fields, and seven pseudoscalar fields (in addition to two neutral Goldstone bosons, and two charged ones). One would expect that, with so many free parameters in the Lagrangian, and so many free masses, almost any scenario is possible for the Higgs masses in this model. We show that the requirement that 1) there is a light neutral scalar Higgs boson, flavor conserving, which is the counterpart to the SM Higgs boson; 2) there exist at least one light doubly charged Higgs boson (as it is interesting for phenomenology); and 3) the flavor-violating neutral Higgs bosons satisfy the constraints imposed by the experimental data from $K^0 - \bar{K}^0$, $D^0 - \bar{D}^0$, and $B^0_{d,s} - \bar{B}^0_{d,s}$ mixings, makes the Higgs sector fairly predictive and fixes some of the parameters in a narrow range. The masses of the light neutral and doubly charged Higgs bosons depend on very few parameters. For instance, we find that requirement 1) and 2) are related, and satisfied by $v_R \in (3, 10)$ TeV range. Assuming $v_R \sim \bar{v}_R$ and $g_L = g_R$, this predicts masses for the $W_R$ around $4 - 13$ TeV (assuming negligible mixing with $W_L$), and for $Z_R$ bosons in the $3 - 10$ TeV range. Thus, while the model can allow for light neutral, singly and doubly charged Higgs bosons, it predicts new
| Particle | Mass (GeV) | Composition |
|----------|------------|-------------|
| $H_1^0$ | 100.6 | $0.015\delta^{cr} + 0.014\delta^{dr} + 0.099\phi_1^{0r} + 0.995\chi_2^{0r} - 0.010\delta^{0r}$ |
| $H_2^0$ | 151.9 | $0.724\delta^{0r} + 0.690\delta^{0r}$ |
| $H_3^0$ | 680.9 | $1.000\phi_2^{0r} + 1.000\chi_1^{0r}$ |
| $H_4^0$ | 3433.4 | $-0.522\delta^{0r} - 0.497\delta^{0r} + 0.002\phi_1^{0r} + 0.032\chi_2^{0r} + 0.693\delta^{0r}$ |
| $H_5^0$ | 5997.1 | $0.502\delta^{0r} + 0.478\delta^{0r} - 0.001\phi_1^{0r} - 0.007\chi_2^{0r} + 0.721\delta^{0r}$ |
| $H_6^0$ | 141419.5 | $-0.690\delta^{0r} + 0.724\delta^{0r}$ |
| $H_7^0$ | 141537.9 | $-0.690\delta^{0r} + 0.724\delta^{0r}$ |
| $H_8^0$ | 449294.1 | $1.000\phi_1^{0r} + 1.000\chi_2^{0r}$ |
| $H_9^0$ | 449294.8 | $1.000\phi_2^{0r} - 0.100\chi_1^{0r}$ |
| $A_1^0$ | 151.9 | $0.724\delta^{0i} - 0.690\delta^{0i}$ |
| $A_2^0$ | 680.9 | $-0.100\phi_2^{0i} + 1.000\chi_1^{0i}$ |
| $A_3^0$ | 4935.7 | $1.000\delta^{0i}$ |
| $A_4^0$ | 141419.5 | $-0.690\delta^{0i} - 0.724\delta^{0i}$ |
| $A_5^0$ | 141502.0 | $0.690\delta^{0i} + 0.724\delta^{0i}$ |
| $A_6^0$ | 449294.1 | $1.000\phi_1^{0i} + 1.000\chi_2^{0i}$ |
| $A_7^0$ | 449294.8 | $1.000\phi_2^{0i} + 1.000\chi_1^{0i}$ |
| $H_1^+$ | 152.9 | $0.724\delta^{+} + 0.690\delta^{+}$ |
| $H_2^+$ | 690.2 | $-0.018\delta^{-} - 0.018\delta^{-} - 0.099\delta^{-} + 0.995\chi_1^{+}$ |
| $H_3^+$ | 141419.5 | $-0.690\delta^{+} + 0.724\delta^{+}$ |
| $H_4^+$ | 141454.7 | $0.690\delta^{-} - 0.724\delta^{+}$ |
| $H_5^+$ | 449294.3 | $0.995\phi_1^{+} + 1.000\chi_2^{-}$ |
| $H_6^+$ | 449294.8 | $0.995\phi_2^{-} - 0.100\chi_1^{+}$ |
| $H_1^{++}$ | 153.9 | $0.724\delta^{++} + 0.690\delta^{--}$ |
| $H_2^{++}$ | 216.3 | $0.724\delta^{--} + 0.690\delta^{++}$ |
| $H_3^{++}$ | 141419.5 | $-0.690\delta^{++} + 0.724\delta^{--}$ |
| $H_4^{++}$ | 141419.6 | $-0.690\delta^{--} + 0.724\delta^{++}$ |
| $G_1^0$ | 0 | $-0.721\delta^{0i} + 0.686\delta^{0i} + 0.010\phi_1^{0i} - 0.995\chi_2^{0i}$ |
| $G_2^0$ | 0 | $0.069\delta^{0i} - 0.066\delta^{0i} + 0.099\phi_1^{0i} - 0.990\chi_2^{0i}$ |
| $G_1^+$ | 0 | $0.100\phi_2^{+} - 0.995\chi_2^{-}$ |
| $G_2^+$ | 0 | $0.724\delta^{-} + 0.690\delta^{+} - 0.003\phi_1^{+} + 0.025\chi_1^{+}$ |

**TABLE IV.** Masses and compositions of physical Higgs fields and unphysical Goldstone bosons.

Parameters are chosen as follows: $\tan \beta = 10$, $\tan \delta = \bar{v}_R/v_R = 1/1.05$, $v_R = 3.5$ TeV, $M_R = 100$ TeV, $\lambda = 1$, $\lambda_{21} = 1$, $C_\chi = 2.5$ TeV, $\langle S \rangle = 1$ TeV, $M_S = 1$ TeV.
| Particle | Mass (GeV) | Composition |
|----------|------------|-------------|
| $H_1^0$  | 112.9      | $0.002\delta^{0r} + 0.002\bar{\delta}^{0r} + 0.02\phi_1^{0r} + 1.000\chi_2^{0r} - 0.001S^{0r}$ |
| $H_2^0$  | 218.5      | $0.72\delta^{0r} + 0.690\delta^{0r}$ |
| $H_3^0$  | 998.6      | $0.02\phi_2^{0r} + 1.000\chi_1^{0r}$ |
| $H_4^0$  | 5562.6     | $-0.52\delta^{0r} - 0.497\delta^{0r} + 0.005\chi_2^{0r} + 0.6935^{0r}$ |
| $H_5^0$  | 8901.0     | $0.51\delta^{0r} + 0.494\delta^{0r} - 0.003\chi_2^{0r} + 0.697S^{0r}$ |
| $H_6^0$  | 141333.6   | $-0.690\delta^{0r} + 0.72\delta^{0r}$ |
| $H_7^0$  | 141575.2   | $-0.690\delta^{0r} + 0.72\delta^{0r}$ |
| $H_8^0$  | 999258.8   | $-1.00\phi_1^{0r} + 0.02\chi_2^{0r}$ |
| $H_9^0$  | 999258.3   | $1.00\phi_2^{0r} - 0.02\chi_1^{0r}$ |
| $A_1^0$  | 218.5      | $0.72\delta^{0i} - 0.690\delta^{0i}$ |
| $A_2^0$  | 998.6      | $-0.02\phi_2^{0i} + 1.000\chi_1^{0i}$ |
| $A_3^0$  | 6976.8     | $1.000\phi^{0i}$ |
| $A_4^0$  | 141334.6   | $-0.690\delta^{0i} - 0.72\delta^{0i}$ |
| $A_5^0$  | 141502.0   | $-0.690\phi^{0i} - 0.72\delta^{0i}$ |
| $A_6^0$  | 999252.8   | $1.00\phi_1^{0i} + 0.02\chi_2^{0i}$ |
| $A_7^0$  | 999258.3   | $1.00\phi_2^{0i} + 0.02\chi_1^{0i}$ |
| $H_1^+$  | 219.2      | $0.72\delta^{+} + 0.690\bar{\delta}^{-}$ |
| $H_2^+$  | 995.3      | $0.01\delta^{c+} + 0.01\bar{\delta}^{c+} + 0.02\phi_2^{-} - 1.000\chi_1^{+}$ |
| $H_3^+$  | 141334.6   | $-0.690\delta^{+} + 0.72\bar{\delta}^{-}$ |
| $H_4^+$  | 141405.3   | $0.690\delta^{c+} - 0.72\delta^{c+}$ |
| $H_5^+$  | 999258.3   | $1.00\phi_1^{+} + 0.02\chi_2^{+}$ |
| $H_6^+$  | 999259.8   | $1.00\phi_2^{-} + 0.02\chi_1^{+}$ |
| $H_1^{++}$ | 219.9     | $0.72\delta^{++} + 0.690\bar{\delta}^{--}$ |
| $H_2^{++}$ | 310.2      | $0.72\delta^{c++} + 0.690\bar{\delta}^{c++}$ |
| $H_3^{++}$ | 141334.6   | $-0.690\delta^{c--} + 0.72\delta^{c+}$ |
| $H_4^{++}$ | 141334.7   | $-0.690\delta^{++} + 0.72\bar{\delta}^{c--}$ |
| $G_1^0$  | 0          | $-0.20\delta^{0i} + 0.190\bar{\delta}^{0i} - 0.01\phi_1^{0i} + 0.961\chi_2^{0i}$ |
| $G_2^0$  | 0          | $0.696\delta^{0i} - 0.663\bar{\delta}^{0i} - 0.006\phi_1^{0i} + 0.27\chi_2^{0i}$ |
| $G_1^+$  | 0          | $0.02\phi_1^{+} - 1.000\bar{\chi}_2^{+}$ |
| $G_2^+$  | 0          | $0.72\delta^{c+} + 0.690\bar{\delta}^{c+} - 0.001\phi_2^{-} + 0.018\chi_1^{+}$ |

**TABLE V.** Masses and compositions of physical Higgs fields and unphysical Goldstone bosons.

Parameters are chosen as follows: $\tan \beta = 50$, $\tan \delta = 1/1.05$, $v_R = 5$ TeV, $M_R = 100$ TeV, $\lambda = 1$, $\lambda_2 = 1$, $C_\lambda = 2.5$ TeV, $\langle S \rangle = 1$ TeV, $M_S = 1000$ GeV.
gauge bosons just outside the range $M_{W_R} < 2(4) \text{ TeV}$ which can be observed at the LHC with a luminosity of $0.1(30) \text{ fb}^{-1}$ \cite{24}.

The parameter $M_R$, associated with the singlet Higgs field in the superpotential, must be of $\mathcal{O}(100) \text{ TeV}$, which insures high masses for the FCNC Higgs.

Our analysis is important for two reasons: first, we have shown that a reasonable Higgs mass spectrum is possible in LRSUSY, without all Higgs masses being required to be heavy. We can require that the Higgs generating tree-level FCNC in the $K$, $D$ and $B$ mesons are heavy, but still obtain two light neutral Higgs bosons, one light pseudoscalar, one light singly charged Higgs boson, and a pair of light doubly charged Higgs bosons. Second, as most Higgs masses are sensitive to few parameters, the model is very predictive and free of additional parameters, such as the sneutrino VEVs or extra higher-dimensional terms. Third, the non-SM light Higgs are mostly triplet $SU(2)_L$ bosons and expected to decay copiously to leptons, but not to quarks, giving clear distinguishing signals for the model. This analysis can now form the basis of a consistent phenomenological study of signals from such a Higgs sector, including production and decay rates, and has implications for the masses of the additional gauge bosons, as well as for the right-handed neutrinos.

VII. ACKNOWLEDGMENTS

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