PHYSICAL REVIEW D **83,** 073007 (2011)

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

Mariana Frank and Beste Korutlu
Department of Physics, Concordia University, 7141 Sherbrooke Street West, Montreal, H4B 1R6 Quebec, Canada
(Received 18 January 2011; published 29 April 2011)

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 nonrenormalizable 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 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 doubly charged Higgs boson is light and show that the parameter space for such Higgs masses and mixings is very restrictive, thus making the model more predictive.

DOI: 10.1103/PhysRevD.83.073007 PACS numbers: 12.15.Ji, 12.60.Cn, 12.60.Fr, 12.60.Jv

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 the 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 and 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 hierachy 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 nonrenormalizable 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 one can have neutral and charged Higgs bosons that conserve flavor below 1 TeV, while the flavor-violating bosons

 © 2011 American Physical Society

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 Sec. II, we summarize the particular LRSUSY model we use, with emphasis on the Higgs structure. In the following section, Sec. III, we present analytic formulas for the mass matrices in the neutral, singly charged, and doubly charged sectors, including first-order loop corrections. In Sec. 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 Sec. V, as well as presenting plots for masses consistent with the constraints. We summarize our findings and conclude in Sec. 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 the 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 as 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, $SU(2)_R$ triplets were introduced to provide the seesaw mechanism for neutrino masses [2], while both left- and right-handed triplet Higgs bosons were considered for parity conservation. The model was described extensively in several previous works [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 [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.* [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

$$\begin{aligned} W = {} & Y_u Q^T \tau_2 \Phi_1 \tau_2 Q^c + Y_d Q^T \tau_2 \Phi_2 \tau_2 Q^c \\ & + Y_\nu L^T \tau_2 \Phi_1 \tau_2 L^c + Y_\ell L^T \tau_2 \Phi_2 \tau_2 L^c + \text{H.c.} \\ & + i f L^{cT} \tau_2 \Delta^c L^c + S[\lambda \operatorname{Tr}(\Delta^c \bar{\Delta}^c) \\ & + \lambda_{ij} \operatorname{Tr}(\Phi_i^T \tau_2 \Phi_j \tau_2) - \mathcal{M}_R^2] + W', \end{aligned} \tag{2.1}$$

where

$$\begin{aligned} W' = {} & [M_\Delta \operatorname{Tr}(\Delta^c \bar{\Delta}^c)] + \mu_{ij} \operatorname{Tr}(\Phi_i^T \tau_2 \Phi_j \tau_2) \\ & + \mathcal{M}_S S^2 + \lambda_S S^3. \end{aligned} \tag{2.2}$$

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 μ term, as explained below.

The model is minimal in the following sense: Δ^c and $\bar{\Delta}^c$ fields are needed for breaking $SU(2)_R \otimes U(1)_{B-L}$ symmetry without R-parity violation, and the two bidoublets Φ_1 and Φ_2 are needed to generate the quark and lepton masses and Cabibbo-Kobayashi-Maskawa (CKM) mixings. The singlet field S is introduced to so that $SU(2)_R \otimes U(1)_{B-L}$ symmetry breaking occurs in the supersymmetric limit. For parity invariance, two left-handed triplet Higgses Δ and $\bar{\Delta}$ are sometimes introduced. We do not include them here, as we wish to construct a minimal model. If included, the VEVs of the left-handed triplet fields Δ and $\bar{\Delta}$, which

TABLE I. Higgs sector in the minimal supersymmetric left-right model.

Higgs Field	Matrix Representation	Vacuum Expectation Values
$\Delta^c(1, 1, 3, -2)$	$\begin{pmatrix} \frac{\delta^{c^-}}{\sqrt{2}} & \delta^{c^0} \\ \delta^{c^{--}} & -\frac{\delta^{c^-}}{\sqrt{2}} \end{pmatrix}$	$\begin{pmatrix} 0 & v_R \\ 0 & 0 \end{pmatrix}$
$\bar{\Delta}^c(1, 1, 3, 2)$	$\begin{pmatrix} \frac{\bar{\delta}^{c^+}}{\sqrt{2}} & \bar{\delta}^{c^{++}} \\ \bar{\delta}^{c^0} & -\frac{\bar{\delta}^{c^+}}{\sqrt{2}} \end{pmatrix}$	$\begin{pmatrix} 0 & 0 \\ \bar{v}_R & 0 \end{pmatrix}$
$\Phi_1(1, 2, 2, 0)$	$\begin{pmatrix} \phi_1^+ & \phi_2^0 \\ \phi_1^0 & \phi_2^- \end{pmatrix}$	$\begin{pmatrix} 0 & \kappa_1' \\ \kappa_1 & 0 \end{pmatrix}$
$\Phi_2(1, 2, 2, 0)$	$\begin{pmatrix} \chi_1^+ & \chi_2^0 \\ \chi_1^0 & \chi_2^- \end{pmatrix}$	$\begin{pmatrix} 0 & \kappa_2 \\ \kappa_2' & 0 \end{pmatrix}$
$S(1, 1, 1, 0)$	S_0	$\langle S \rangle$

determine the tree-level left-handed neutrino masses, must be extremely small and are assumed to be zero. In this case, the left-handed triplet fields decouple, and their addition amounts only to the proliferation of Higgs masses and representations in this model.

The charge is defined as

$$Q = I_{3L} + I_{3R} + \frac{B-L}{2}. \tag{2.3}$$

The VEVs of the Higgs fields in this model needed to break the symmetries, as described above, are given in Table I. If we assume that the VEVs of the bidoublet Higgs are real, the fermion mass matrices become Hermitian. In the supersymmetric limit, the VEV of the singlet S Higgs boson is zero, but, after SUSY breaking, $\langle S\rangle \sim m_{\rm SUSY}$. Thus, the μ term for the bidoublet Φ will arise from the coupling λ_{ij}, with a magnitude of order $m_{\rm SUSY}$ [11]. In the SUSY limit,

$$|v_R| = |\bar{v}_R|, \qquad \lambda v_R \bar{v}_R = \mathcal{M}_R^2, \qquad \langle S\rangle = 0. \tag{2.4}$$

The VEV of the S field, generated after SUSY breaking, arises from linear terms in SUSY breaking

$$V_{\rm soft} = A_\lambda \lambda S\,{\rm Tr}(\Delta^c\bar{\Delta}^c) - C_\lambda \mathcal{M}_R^2 S + {\rm H.c.} \tag{2.5}$$

Minimization of the resulting potential yields $\langle S^*\rangle = \frac{1}{2\lambda} \times (C_\lambda - A_\lambda)$, which is of order $m_{\rm SUSY}$. If the coupling λ 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 remain below 1 TeV. Consistency of the model (a nonvanishing CKM mixing angle) requires the asymmetry $\mu_{12} = \mu_{21}$.

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

$$\begin{aligned} V_F &= |\lambda\,{\rm Tr}(\Delta^c\bar{\Delta}^c) + \lambda_{ij}\,{\rm Tr}(\Phi_i^T\tau_2\Phi_j\tau_2) - \mathcal{M}_R^2|^2 \\ &\quad + \lambda^2|S|^2|{\rm Tr}(\Delta^c\Delta^{c\dagger}) + {\rm Tr}(\bar{\Delta}^c\bar{\Delta}^{c\dagger})|, \\ V_{\rm soft} &= M_1^2\,{\rm Tr}(\Delta^{c\dagger}\Delta^c) + M_2^2\,{\rm Tr}(\bar{\Delta}^{c\dagger}\bar{\Delta}^c) + M_3^2\,{\rm Tr}(\Phi_1^\dagger\Phi_1) \\ &\quad + M_4^2\,{\rm Tr}(\Phi_2^\dagger\Phi_2) + M_S^2|S|^2 \\ &\quad + \{A_\lambda\lambda S\,{\rm Tr}(\Delta^c\bar{\Delta}^c) - C_\lambda\mathcal{M}_R^2 S + {\rm H.c.}\}, \\ V_D &= \frac{g_L^2}{8}\sum_i |{\rm Tr}(\Phi_a\tau_i^T\Phi_b^\dagger)|^2 + \frac{g_R^2}{8}\sum_i |{\rm Tr}(2\Delta^{c\dagger}\tau_i\Delta^c \\ &\quad + 2\bar{\Delta}^{c\dagger}\tau_i\bar{\Delta}^c + \Phi_a\tau_i^T\Phi_b^\dagger)|^2 \\ &\quad + \frac{g'^2}{2}|{\rm Tr}(-\Delta^{c\dagger}\Delta^c + \bar{\Delta}^{c\dagger}\bar{\Delta}^c)|^2. \end{aligned} \tag{2.6}$$

We use this potential and proceed in 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

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

to obtain masses and compositions of the Higgs bosons. However, this procedure does not lead to the correct minimum of the potential. The reason is that all the terms in the scalar potential are identical for the configurations in which VEVs are given to the neutral right-handed triplet Higgs, except for the D term, which is lower for the charge-breaking configuration

$$\langle\Delta^c\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & v_R \\ v_R & 0\end{pmatrix}, \qquad \langle\bar{\Delta}^c\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & \bar{v}_R \\ \bar{v}_R & 0\end{pmatrix}. \tag{2.7}$$

Previous solutions suggested are breaking R parity, which would have the attractive feature that $v_R \sim 1$ TeV, but which abandons the lightest supersymmetric particle as the candidate for dark matter [7]; or, they are 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 the effects of introducing a one-loop Coleman-Weinberg effective potential generated by a one-family right-chiral neutrino to the Δ^c field:

$$V_{\rm eff}^{\rm 1\text{-}loop} = \frac{1}{16\pi^2}\sum_i (-1)^{2s}(2s+1)M_i^4\left[\ln\left(\frac{M_i^2}{\mu^2}\right) - \frac{3}{2}\right]. \tag{2.8}$$

Expanding this potential in the limit in which the SUSY-breaking parameters are small with respect to the triplet VEVs $(v_R, \bar{v}_R)$, one obtains an effective form, in terms of the small parameter

$$x = \frac{{\rm Tr}(\Delta^c\Delta^c){\rm Tr}(\Delta^{c\dagger}\Delta^{c\dagger})}{[{\rm Tr}(\Delta^{c\dagger}\Delta^c)]^2}.$$

The one-loop potential becomes

$$\begin{aligned} V_{\rm eff}^{\rm 1\text{-}loop} \simeq &-\frac{|f|^2 m_{L^c}^2\,{\rm Tr}(\Delta^c\Delta^c){\rm Tr}(\Delta^{c\dagger}\Delta^{c\dagger})}{128\pi^2|v_R|^2} \\ &\times\left\{(a_1 - a_2)g_R^2\left(2\ln\frac{|fv_R|^2}{\mu^2} + \ln x - 2\ln 2 - 2\right)\right. \\ &\left. - [2 + (a_1 + a_2)g_{B-L}^2](\ln x - 2\ln 2)\right\}. \end{aligned} \tag{2.9}$$

Here, a_1 and a_2 are coefficients which vanish in the SUSY limit (when D terms vanish), and $m_{L^c}^2$ are soft right-handed scalar lepton masses. The effect of this potential is to 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. Whereas before, the global minimum contained at least one doubly charged Higgs boson with zero or negative mass, after one-loop corrections, all the masses are 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 arbitrary sneutrino VEVs. In the next section, we study explicitly the implications for the Higgs masses in this model. Before continuing, we wish to point out that, should the model have included left-handed triplet Higgs bosons, their mass would remain negative and could not be fixed by the first-order loop corrections. A left-handed counterpart of the one-loop correction would not work, as the VEVs of these fields v_L is zero, or very small. Thus, one would have to consider higher-order corrections or additional Higgs representations.

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 flavor-changing neutral current (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].

After minimizing the Higgs potential, as in the previous section, we evaluate the masses by taking into account corrections induced by the heavy Majorana neutrino Yukawa couplings. We shall give the expressions for these contributions explicitly. We proceed to give the masses obtained by minimizing the potential. For simplicity, we use the abbreviations

$$\kappa_{\rm dif}^2 = \kappa_1^2 - \kappa_2^2, \tag{3.1}$$

$$\rho_{\rm dif}^2 = v_R^2 - \bar{v}_R^2 + \tfrac{1}{2}(\kappa_1^2 - \kappa_2^2), \tag{3.2}$$

$$Y = A_\lambda \lambda S + \lambda(-M_R^2 - 2\lambda_{21}\kappa_1\kappa_2 + \lambda v_R \bar{v}_R), \tag{3.3}$$

$$M = 2\lambda_{21}(-M_R^2 - 2\lambda_{21}\kappa_1\kappa_2 + \lambda v_R \bar{v}_R), \tag{3.4}$$

$$f(\epsilon) = \epsilon\left(\frac{M}{2\lambda_{21}} - 2\lambda_{21}\kappa_1\kappa_2 - \epsilon\kappa_1\kappa_2\right), \tag{3.5}$$

$$g(\epsilon) = \epsilon\lambda\kappa_1\kappa_2, \tag{3.6}$$

$$h(\epsilon) = \epsilon\kappa_1\kappa_2(4\lambda_{21} + \epsilon), \tag{3.7}$$

with $\epsilon = \mu_{21} - \mu_{12}$ small but nonzero 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 $(\delta^{c--*}, \bar{\delta}^{c++})$ fields

$$M^2_{\delta^{c--*}\,\bar{\delta}^{c++}} = \begin{pmatrix} -2g_R^2\rho_{\rm dif}^2 - \frac{\bar{v}_R}{v_R}Y' & Y' \\ Y' & 2g_R^2\rho_{\rm dif}^2 - \frac{v_R}{\bar{v}_R}Y' \end{pmatrix}, \tag{3.8}$$

where $Y' = Y - g(\epsilon)$. From these expressions, we can find the exact analytic forms for the doubly charged Higgs masses. Setting $\tan\delta = \frac{\bar{v}_R}{v_R}$, these are

$$M^2_{H^{++}_{1,2}} = -\frac{Y'}{\sin 2\delta} \pm \sqrt{4g_R^4\rho_{\rm dif}^4 + \frac{Y'^2}{\sin^2 2\delta} - 4Y' g_R^2\rho_{\rm dif}^2 \cot 2\delta}. \tag{3.9}$$

It is clear that one must require $Y' < 0$, but, even so, one of the mass eigenvalues will be negative. It is thus essential that we include the first-order correction to the doubly charged Higgs masses, which arise from derivatives of the quartic potential (2.9) with respect to the doubly charged Higgs boson fields. The corrected (Mass^2) matrix elements are

$$\begin{aligned} M^2_{\delta^{c--*}\,\delta^{c--}} &= -\frac{f^2 m^2_{L^c}}{16\pi^2}\left[a_1 g_R^2\left(2\ln\left(\frac{|f v_R|}{\mu}\right) - 1\right) + \ln 2(2 - a_1(g_R^2 - g_{B-L}^2))\right] - 2g_R^2\rho_{\rm dif}^2 - \frac{\bar{v}_R}{v_R}Y', \\ M^2_{\bar{\delta}^{c++}\,\delta^{c--}} &= Y', \\ M^2_{\delta^{c--*}\,\bar{\delta}^{c++*}} &= Y', \\ M^2_{\bar{\delta}^{c++}\,\bar{\delta}^{c++*}} &= 2g_R^2\rho_{\rm dif}^2 - \frac{\bar{v}_R}{v_R}Y', \end{aligned} \tag{3.10}$$

yielding a positive correction to the masses for $m^2_{L^c} < 0$. The first-order correction is not finite at $x = 0$; however, the divergence is very mild (logarithmic), and higher-order effects cure it without altering the masses significantly [13].

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 (ϕ_1^+, χ_2^{-*}) fields and one four-by-four matrix for $(\delta^{c^+}, \bar{\delta}^{c^{-*}}, \phi_2^{-*}, \chi_1^+)$ fields, respectively,

$$M^2_{\phi_1^+,\chi_2^{-*}} = \begin{pmatrix} \frac{\kappa_2}{\kappa_1}M' & M' \\ M' & \frac{\kappa_1}{\kappa_2}M' \end{pmatrix}, \tag{3.11}$$

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

$$M^2_{\delta^{c-*}\,\delta^{c-}} = g_R^2 v_R^2 - g_R^2\rho_{\rm dif}^2 - \frac{\bar{v}_R}{v_R}Y', \tag{3.12}$$

$$M^2_{\delta^{c-*}\,\bar{\delta}^{c+*}} = -g_R^2 v_R \bar{v}_R + Y', \tag{3.13}$$

$$M^2_{\delta^{c-*}\phi_2^-} = -\frac{1}{\sqrt{2}} g_R^2 \kappa_1 v_R, \quad (3.14)$$

$$M^2_{\delta^{c-*}\chi_1^{+*}} = -\frac{1}{\sqrt{2}} g_R^2 \kappa_2 v_R, \quad (3.15)$$

$$M^2_{\bar{\delta}^{c+}\bar{\delta}^{c+*}} = g_R^2 \bar{v}_R^2 + g_R^2 \rho_{\rm dif}^2 - \frac{v_R}{\bar{v}_R} Y', \quad (3.16)$$

$$M^2_{\bar{\delta}^{c+}\phi_2^-} = \frac{1}{\sqrt{2}} g_R^2 \kappa_1 \bar{v}_R, \quad (3.17)$$

$$M^2_{\bar{\delta}^{c+}\chi_1^{+*}} = \frac{1}{\sqrt{2}} g_R^2 \kappa_2 \bar{v}_R, \quad (3.18)$$

$$M^2_{\phi_2^{-*}\phi_2^-} = \frac{1}{2}\kappa_1^2(g_L^2 + g_R^2) - \frac{1}{2} g_L^2 \kappa_{\rm dif}^2 - g_R^2 \rho_{\rm dif}^2 + \frac{\kappa_2}{\kappa_1} M', \quad (3.19)$$

$$M^2_{\phi_2^{-*}\chi_1^{+*}} = \tfrac{1}{2}\kappa_1\kappa_2(g_L^2 + g_R^2) + M', \quad (3.20)$$

$$M^2_{\chi_1^+\chi_1^-} = \frac{1}{2}\kappa_2^2(g_L^2 + g_R^2) + \frac{1}{2} g_L^2 \kappa_{\rm dif}^2 + g_R^2 \rho_{\rm dif}^2 + \frac{\kappa_1}{\kappa_2} M'. \quad (3.21)$$

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 $(\phi_2^{0r}, \chi_1^{0r})$ fields and one five-by-five matrix for $(\delta^{c^{0r}}, \bar{\delta}^{c^{0r}}, \phi_1^{0r}, \chi_2^{0r}, S^{0r})$ fields, respectively,

$$M^2_{\phi_2^{0r},\chi_1^{0r}} = \begin{pmatrix} -\frac{1}{2} g_L^2 \kappa_{\rm dif}^2 - g_R^2 \rho_{\rm dif}^2 + \frac{\kappa_2}{\kappa_1} M' & -M' \\ -M' & \frac{1}{2} g_L^2 \kappa_{\rm dif}^2 + g_R^2 \rho_{\rm dif}^2 + \frac{\kappa_1}{\kappa_2} M' \end{pmatrix}. \quad (3.22)$$

The elements of the five-by-five matrix are

$$M^2_{\delta^{c0r}\delta^{c0r}} = 2v_R^2(g_{B-L}^2 + g_R^2) + \lambda^2 \bar{v}_R^2 - \frac{\bar{v}_R}{v_R} Y', \quad (3.23)$$

$$M^2_{\delta^{c0r}\bar{\delta}^{c0r}} = -2v_R\bar{v}_R(g_{B-L}^2 + g_R^2) + \lambda^2 v_R \bar{v}_R + Y', \quad (3.24)$$

$$M^2_{\delta^{c0r}\phi_1^{0r}} = g_R^2 \kappa_1 v_R - 2\lambda\lambda_{21}\kappa_2 \bar{v}_R - 2\frac{\bar{v}_R}{\kappa_1} g(\epsilon), \quad (3.25)$$

$$M^2_{\delta^{c0r}\chi_2^{0r}} = -g_R^2 \kappa_2 v_R - 2\lambda\lambda_{21}\kappa_1 \bar{v}_R - \frac{\bar{v}_R}{\kappa_2} g(\epsilon), \quad (3.26)$$

$$M^2_{\delta^{c0r}S^{0r}} = 2\lambda^2 S v_R + A_\lambda \lambda \bar{v}_R, \quad (3.27)$$

$$M^2_{\bar{\delta}^{c0r}\bar{\delta}^{c0r}} = 2(g_{B-L}^2 + g_R^2)\bar{v}_R^2 + \lambda^2 v_R^2 - \frac{v_R}{\bar{v}_R} Y', \quad (3.28)$$

$$M^2_{\bar{\delta}^{c0r}\phi_1^{0r}} = -g_R^2 \kappa_1 \bar{v}_R - 2\lambda\lambda_{21}\kappa_2 v_R - \frac{v_R}{\kappa_1} g(\epsilon), \quad (3.29)$$

$$M^2_{\bar{\delta}^{c0r}\chi_2^{0r}} = g_R^2 \kappa_2 \bar{v}_R - 2\lambda\lambda_{21}\kappa_1 v_R - \frac{v_R}{\kappa_2} g(\epsilon), \quad (3.30)$$

$$M^2_{\bar{\delta}^{c0r}S^{0r}} = 2\lambda^2 S \bar{v}_R + A_\lambda \lambda v_R, \quad (3.31)$$

$$M^2_{\phi_1^{0r}\phi_1^{0r}} = \frac{1}{2}\kappa_1^2(g_L^2 + g_R^2) + 4\lambda_{21}^2\kappa_2^2 + \frac{\kappa_2}{\kappa_1}[M' + h(\epsilon)], \quad (3.32)$$

$$M^2_{\phi_1^{0r}\chi_2^{0r}} = -\tfrac{1}{2}\kappa_1\kappa_2(g_L^2 + g_R^2) + 4\lambda_{21}^2\kappa_1\kappa_2 - [M' - h(\epsilon)], \quad (3.33)$$

$$M^2_{\phi_1^{0r}S^{0r}} = 0, \quad (3.34)$$

$$M^2_{\chi_2^{0r}\chi_2^{0r}} = \frac{1}{2}\kappa_2^2(g_L^2 + g_R^2) + 4\lambda_{21}^2\kappa_1^2 + \frac{\kappa_1}{\kappa_2}[M' + h(\epsilon)], \quad (3.35)$$

$$M^2_{\chi_2^{0r}S^{0r}} = 0, \quad (3.36)$$

$$M^2_{S^{0r}S^{0r}} = M_S^2 + \lambda^2(v_R^2 + \bar{v}_R^2). \quad (3.37)$$

Mass matrices for the neutral pseudoscalar Higgs fields are similar, of block diagonal form of one two-by-two matrix for $(\phi_2^{0i}, \chi_1^{0i})$ fields and one five-by-five matrix for $(\delta^{c^{0i}}, \bar{\delta}^{c^{0i}}, \phi_1^{0i}, \chi_2^{0i}, S^{0i})$ fields, respectively,

$$M^2_{\phi_2^{0i},\chi_1^{0i}} = \begin{pmatrix} -\frac{1}{2} g_L^2 \kappa_{\rm dif}^2 - g_R^2 \rho_{\rm dif}^2 + \frac{\kappa_2}{\kappa_1} M' & M' \\ M' & \frac{1}{2} g_L^2 \kappa_{\rm dif}^2 + g_R^2 \rho_{\rm dif}^2 + \frac{\kappa_1}{\kappa_2} M' \end{pmatrix}. \quad (3.38)$$

The elements of the five-by-five matrix are

$$M^2_{\delta^{c0i}\delta^{c0i}} = \lambda^2 \bar{v}_R^2 - \frac{\bar{v}_R}{v_R} Y', \quad (3.39)$$

$$M^2_{\delta^{c0i}\bar{\delta}^{c0i}} = \lambda^2 v_R \bar{v}_R - Y', \quad (3.40)$$

$$M^2_{\delta^{c0i}\phi_1^{0i}} = -2\lambda\lambda_{21}\kappa_2 \bar{v}_R - \frac{\bar{v}_R}{\kappa_1} g(\epsilon), \quad (3.41)$$

$$M^2_{\delta^{c0i}\chi_2^{0i}} = -2\lambda\lambda_{21}\kappa_1 \bar{v}_R - \frac{\bar{v}_R}{\kappa_2} g(\epsilon), \quad (3.42)$$

$$M^2_{\delta^{c0i}S^{0i}} = -A_\lambda \lambda \bar{v}_R, \quad (3.43)$$

$$M^2_{\bar{\delta}^{c0i}\bar{\delta}^{c0i}} = \lambda^2 v_R^2 - \frac{v_R}{\bar{v}_R} Y', \tag{3.44}$$

$$M^2_{\bar{\delta}^{c0i}\phi_1^{0i}} = -2\lambda\lambda_{21}\kappa_2 v_R - \frac{v_R}{\kappa_1} g(\epsilon), \tag{3.45}$$

$$M^2_{\bar{\delta}^{c0i}\chi_2^{0i}} = -2\lambda\lambda_{21}\kappa_1 v_R - \frac{v_R}{\kappa_2} g(\epsilon), \tag{3.46}$$

$$M^2_{\bar{\delta}^{c0i}S^{0i}} = -A_\lambda \lambda v_R, \tag{3.47}$$

$$M^2_{\phi_1^{0i}\phi_1^{0i}} = 4\lambda_{21}^2\kappa_2^2 + \frac{\kappa_2}{\kappa_1}[M' + h(\epsilon)], \tag{3.48}$$

$$M^2_{\phi_1^{0i}\chi_2^{0i}} = 4\lambda_{21}^2\kappa_2\kappa_2 + [M' + h(\epsilon)], \tag{3.49}$$

$$M^2_{\phi_1^{0i}S^i} = 0, \tag{3.50}$$

$$M^2_{\chi_2^{0i}\chi_2^{0r}} = 4\lambda_{21}^2\kappa_1^2 + \frac{\kappa_1}{\kappa_2}[M' + h(\epsilon)], \tag{3.51}$$

$$M^2_{\chi_2^{0i}S^{0i}} = 0, \tag{3.52}$$

$$M^2_{S^{0i}S^{0i}} = M_S^2 + \lambda^2(v_R^2 + \bar{v}_R^2). \tag{3.53}$$

The one-loop correction to the Higgs boson masses is significant for doubly charged bosons only.

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 [14]. 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

$$\mathcal{L}_Y = \bar{d}_L Y_u \phi_2^0 d_R + \bar{d}_L Y_d \chi_2^0 d_R + \bar{u}_L Y_u \phi_1^0 u_R + \bar{u}_L Y_d \chi_1^0 u_R + \text{H.c.}, \tag{4.1}$$

where Y_u and Y_d are 3×3 Hermitian matrices in flavor space. When the bidoublets acquire the VEV, as in Table I, with κ_1, κ_2, κ_1', and κ_2' real, the up- and the down-type quark mass matrices are given by

$$M_u = Y_u\kappa_1 + Y_d\kappa_2', \qquad M_d = Y_u\kappa_1' + Y_d\kappa_2. \tag{4.2}$$

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

$$\mathcal{L}_Y^N(d) = \frac{[d_L^{i*} M_u^{ij} d_R^j(\kappa_2\phi_2^0 - \kappa_2'\chi_2^0) + d_L^{i*} M_d^{ij} d_R^j(\kappa_1\chi_2^0 - \kappa_1'\phi_2^0)]}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} + \frac{[d_R^{j*} M_u^{ij*} d_L^i(\kappa_2\phi_2^{0*} - \kappa_2'\chi_2^{0*}) + d_R^{j*} M_d^{ij*} d_L^i(\kappa_1\chi_2^{0*} - \kappa_1'\phi_2^{0*})]}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'}. \tag{4.3}$$

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

$$M_u^{ij} = U_u^{ik}\hat{M}_u^{km}W_u^{jm*}\delta^{km}, \qquad M_d^{ij} = U_d^{ik}\hat{M}_d^{km}W_d^{jm*}\delta^{km}, \tag{4.4}$$

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

$$d_L^i \to U_d^{ij} d_L^j, \qquad d_R^i \to W_d^{ij} d_R^j. \tag{4.5}$$

We define $U_d^{ji*}U_u^{ik} = V_L^{jk}$ and $W_u^{lj*}W_d^{jm} = V_R^{lm}$, 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

$$\begin{aligned}\mathcal{L}_Y^N(d) = {}& \frac{d_L^{n*}V_L^{kn*}\hat{M}_u^{km}V_R^{ml}d_R^l\delta^{km}(\kappa_2\phi_2^0 - \kappa_2'\chi_2^0)}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} \\ &+ \frac{d_L^{n*}\delta^{nk}\hat{M}_d^{km}\delta^{ml}d_R^l\delta^{km}(\kappa_1\chi_2^0 - \kappa_1'\phi_2^0)}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} \\ &+ \frac{d_R^{n*}V_R^{mn*}\hat{M}_u^{km*}V_L^{kl}d_L^l\delta^{km}(\kappa_2\phi_2^{0*} - \kappa_2'\chi_2^{0*})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} \\ &+ \frac{d_R^{n*}\delta^{nm}\hat{M}_d^{km*}\delta^{kl}d_L^l\delta^{km}(\kappa_1\chi_2^{0*} - \kappa_1'\phi_2^{0*})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'},\end{aligned} \tag{4.6}$$

where the up and down mass matrices are Hermitian, since the VEVs of bidoublets are taken to be real. For simplicity, we assume $V_L = V_R = V$. The fields ϕ_2^0 and χ_2^0 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

$$\mathcal{L}_{\rm FCNC}(d) = \frac{d_L^{n*}V^{kn*}\hat{M}_u^{kk}V^{kl}d_R^l(\kappa_2\phi_2^{0r} - \kappa_2'\chi_2^{0r})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} + \frac{id_L^{n*}V^{kn*}\hat{M}_u^{kk}V^{kl}d_R^l(\kappa_2\phi_2^{0i} - \kappa_2'\chi_2^{0i})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} + \frac{d_R^{n*}V^{kn*}\hat{M}_u^{kk*}V^{kl}d_R^l(\kappa_2\phi_2^{0r} - \kappa_2'\chi_2^{0r})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} - \frac{id_R^{n*}V^{kn*}\hat{M}_u^{kk*}V^{kl}d_R^l(\kappa_2\phi_2^{0i} - \kappa_2'\chi_2^{0i})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'}, \tag{4.7}$$

where ϕ_2^{0r} and χ_2^{0r} are the two of the nine bare scalar fields, and ϕ_2^{0i} and χ_2^{0i} are the two of the nine bare pseudoscalar fields appearing in the 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 $\lambda = |V_{us}| = 0.2246 \pm 0.0012$ [15]:

$$V = \begin{pmatrix} 1 - \frac{\lambda^2}{2} & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda & 1 - \frac{\lambda^2}{2} & A\lambda^2 \\ A\lambda^3(1 - \rho - i\eta) & -A\lambda^2 & 1 \end{pmatrix} + \mathcal{O}(\lambda^4). \tag{4.8}$$

For $\lambda = 0.2246$, $A = 0.832$, $\rho = 0.130$, and $\eta = 0.350$ [16,17],

$$V^{kd*}\hat{M}_u^{kk}V^{ks} = (m_u - m_c)\left(\lambda - \frac{\lambda^3}{2}\right) - m_t A^2\lambda^5(1 - \rho + i\eta). \tag{4.9}$$

We express the bare scalar $\psi^{0r^T} = (\delta^{c0r}\bar{\delta}^{c0r}\phi_1^{0r}\phi_2^{0r}\chi_1^{0r}\chi_2^{0r}S^{0r})$ and pseudoscalar Higgs fields $\psi^{0i^T} = (\delta^{c0i}\bar{\delta}^{c0i}\phi_1^{0i}\phi_2^{0i}\chi_1^{0i}\chi_2^{0i}S^{0i})$ as physical CP-even Higgs fields $H^{0r^T} = (H_1^{0r}H_2^{0r}H_3^{0r}H_4^{0r}H_5^{0r}H_6^{0r}H_7^{0r})$ and physical CP-odd Higgs fields $H^{0i^T} = (H_1^{0i}H_2^{0i}H_3^{0i}H_4^{0i}H_5^{0i}H_6^{0i}H_7^{0i})$. Call A_{ij} the transformation matrix which transforms the bare scalar fields into the physical CP-even ones and B_{ij} the matrix which transforms the bare pseudoscalar 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

$$\begin{aligned} \mathcal{L}_{\rm FCNC}^{\Delta S=2}(d) = {}& \frac{m_t\lambda}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'}\left(\left[\left(\frac{m_u}{m_t} - \frac{m_c}{m_t}\right)\left(1 - \frac{\lambda^2}{2}\right) - A^2\lambda^4(1 - \rho)\right](\kappa_2 A_{i4}^* - \kappa_2' A_{i6}^*)H_i^{0r}(\bar{d}P_R s + \bar{d}P_L s)\right. \\ & \left. + A^2\lambda^4\eta(\kappa_2 B_{i4}^* - \kappa_2' B_{i6}^*)H_i^{0i}(\bar{d}P_R s - \bar{d}P_L s)\right) + \frac{im_t\lambda}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'}\left(\left[\frac{m_u}{m_t} - \frac{m_c}{m_t}\right)\left(1 - \frac{\lambda^2}{2}\right) - A^2\lambda^4(1 - \rho)\right] \\ & \left. \times (\kappa_2 B_{i4}^* - \kappa_2' B_{i6}^*)H_i^{0i}(\bar{d}P_R s - \bar{d}P_L s) - A^2\lambda^4\eta(\kappa_2 A_{i4}^* - \kappa_2' A_{i6}^*)H_i^{0r}(\bar{d}P_R s + \bar{d}P_L s)\right). \end{aligned} \tag{4.10}$$

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

$$\mathcal{L}_Y^N(u) = u_L^{i*}Y_u^{ij}\phi_1^0 u_R^j + u_L^{i*}Y_d^{ij}\chi_1^0 u_R^j + u_R^{j*}\phi_1^{0*}Y_u^{ji*}u_L^i + u_R^{j*}\chi_1^{0*}Y_d^{ji*}u_L^i. \tag{4.11}$$

We use the same substitutions as before and express the Lagrangian in terms of the complex fields ϕ_2^0 and χ_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

$$\begin{aligned} \mathcal{L}_{\rm FCNC}(u) = {}& \frac{u_L^{n*}V^{nk}\hat{M}_d^{kk}V^{lk*}u_R^l(\kappa_2\phi_1^{0r} - \kappa_2'\chi_1^{0r})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} + \frac{iu_L^{n*}V^{nk}\hat{M}_d^{kk}V^{lk*}u_R^l(\kappa_2\phi_1^{0i} - \kappa_2'\chi_1^{0i})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} \\ & + \frac{u_R^{n*}V^{nk}\hat{M}_d^{kk*}V^{lk*}u_R^l(\kappa_2\phi_1^{0r} - \kappa_2'\chi_1^{0r})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'} - \frac{iu_R^{n*}V^{nk}\hat{M}_d^{kk*}V^{lk*}u_R^l(\kappa_2\phi_1^{0i} - \kappa_2'\chi_1^{0i})}{\kappa_1\kappa_2 - \kappa_1'\kappa_2'}, \end{aligned} \tag{4.12}$$

where ϕ_1^{0r} and χ_1^{0r} are the two of the nine bare scalar fields, and ϕ_1^{0i} and χ_1^{0i} are the two of the nine bare pseudoscalar fields appearing in the LRSUSY Lagrangian. The $u - c$ coupling in Eq. (4.12) allows a $\Delta C = 2$ transition at tree level. Inserting $V^{uk}\hat{M}_u^{kk}V^{ck*}$ in terms of Wolfenstein parameters,

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

and using physical states instead of ϕ_1^{0r} and χ_1^{0r}, we obtain the explicit form of the Lagrangian responsible for FCNC in the up sector

$$\mathcal{L}^{\Delta C=2}_{\rm FCNC}(u) = \frac{m_b\lambda}{\kappa_1\kappa_2-\kappa_1'\kappa_2'}\bigg(\bigg[\bigg(\frac{m_s}{m_b}-\frac{m_d}{m_b}\bigg)\bigg(1-\frac{\lambda^2}{2}\bigg)+A^2\lambda^4\rho\bigg](\kappa_2 A^*_{i3}-\kappa_2' A^*_{i5})H_i^{0r}(\bar{u}P_R c+\bar{u}P_L c)$$
$$+A^2\lambda^4\eta(\kappa_2 B^*_{i3}-\kappa_2' B^*_{i5})H_i^{0i}(\bar{u}P_R c-\bar{u}P_L c)\bigg)+\frac{im_b\lambda}{\kappa_1\kappa_2-\kappa_1'\kappa_2'}\bigg(\bigg[\bigg(\frac{m_s}{m_b}-\frac{m_d}{m_b}\bigg)\bigg(1-\frac{\lambda^2}{2}\bigg)+A^2\lambda^4\rho\bigg]$$
$$\times(\kappa_2 B^*_{i3}-\kappa_2' B^*_{i5})H_i^{0i}(\bar{u}P_R c-\bar{u}P_L c)-A^2\lambda^4\eta(\kappa_2 A^*_{i3}-\kappa_2' A^*_{i5})H_i^{0r}(\bar{u}P_R c+\bar{u}P_L c)\bigg). \quad (4.14)$$

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. ϵ_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.

$$\mathrm{Re}\langle\bar{K}^0|H_{\rm eff}|K^0\rangle=\frac{m_t^2\lambda^2}{4M_i^2(\kappa_1\kappa_2-\kappa_1'\kappa_2')^2}\bigg\{\bigg[\bigg(\frac{m_u}{m_t}-\frac{m_c}{m_t}\bigg)(2-\lambda^2)-2A^2\lambda^4(1-\rho)\bigg]$$
$$\times([(\kappa_2A^*_{i4}-\kappa_2'A^*_{i6})^2-(\kappa_2B^*_{i4}-\kappa_2'B^*_{i6})^2](\langle\tilde{Q}_1(\mu)\rangle+\langle Q_1(\mu)\rangle)$$
$$+[(\kappa_2A^*_{i4}-\kappa_2'A^*_{i6})^2+(\kappa_2B^*_{i4}-\kappa_2'B^*_{i6})^2](\langle\tilde{Q}_2(\mu)\rangle+\langle Q_2(\mu)\rangle))+4A^4\lambda^8\eta^2([(\kappa_2A^*_{i4}-\kappa_2'A^*_{i6})^2$$
$$+(\kappa_2B^*_{i4}-\kappa_2'B^*_{i6})^2](\langle\tilde{Q}_1(\mu)\rangle+\langle Q_1(\mu)\rangle)+[(\kappa_2A^*_{i4}-\kappa_2'A^*_{i6})^2-(\kappa_2B^*_{i4}-\kappa_2'B^*_{i6})^2](\langle\tilde{Q}_2(\mu)\rangle+\langle Q_2(\mu)\rangle))\bigg\}, \quad (4.15)$$

and

$$\mathrm{Im}\langle K^0|H_{\rm eff}|\bar{K}^0\rangle$$
$$=\frac{im_t^2\lambda^2}{4M_i^2(\kappa_1\kappa_2-\kappa_1'\kappa_2')^2}\bigg[\bigg(\frac{m_u}{m_t}-\frac{m_c}{m_t}\bigg)(2-\lambda^2)A^2\lambda^4\eta$$
$$-2A^4\lambda^8(1-\rho)\eta]([(\kappa_2B^*_{i4}-\kappa_2'B^*_{i6})^2$$
$$-(\kappa_2A^*_{i4}-\kappa_2'A^*_{i6})^2\bigg](\langle\tilde{Q}_1(\mu)\rangle+\langle Q_1(\mu)\rangle)$$
$$-[(\kappa_2B^*_{i4}-\kappa_2'B^*_{i6})^2+(\kappa_2A^*_{i4}-\kappa_2'A^*_{i6})^2]$$
$$\times(\langle\tilde{Q}_2(\mu)\rangle+\langle Q_2(\mu)\rangle)). \quad (4.16)$$

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

$$Q_1=(\bar{q}_1^\alpha P_L q_2^\alpha)\otimes(\bar{q}_1^\beta P_L q_2^\beta),$$
$$\tilde{Q}_1=(\bar{q}_1^\alpha P_R q_2^\alpha)\otimes(\bar{q}_1^\beta P_R q_2^\beta),$$
$$Q_2=(\bar{q}_1^\alpha P_L q_2^\alpha)\otimes(\bar{q}_1^\beta P_R q_2^\beta),$$
$$\tilde{Q}_2=(\bar{q}_1^\alpha P_R q_2^\alpha)\otimes(\bar{q}_1^\beta P_L q_2^\beta), \quad (4.17)$$

where α and β are the color indices. The matrix elements are [18]

$$\langle Q_1(\mu)\rangle=-\frac{5}{24}\bigg(\frac{m_a}{m_{q_1}(\mu)+m_{q_2}(\mu)}\bigg)^2 m_aF_a^2B_1(\mu),$$
$$\langle Q_2(\mu)\rangle=\frac{1}{4}\bigg(\frac{m_a}{m_{q_1}(\mu)+m_{q_2}(\mu)}\bigg)^2 m_aF_a^2B_2(\mu), \quad (4.18)$$

where $a=K$, B_d, B_s, and 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 [19] calculated in the naïve dimensional regularization scheme for an energy scale μ. 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 those for the quark mass values are summarized in Table III. The same expressions for the operators Q_1 and Q_2 are valid for the operators $\tilde{Q}_1$ and $\tilde{Q}_2$.

Substituting $\mu=2$ GeV in the expressions for ΔM_K and the CP-violating parameter ϵ_K given below,

TABLE II. QCD parameters used for meson mixing.

	$K^0-\bar{K}^0$	$B^0_d-\bar{B}^0_d$	$B^0_s-\bar{B}^0_s$	$D^0-\bar{D}^0$
μ	2 GeV	m_b	m_b	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 III. Quark masses.

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

$$\Delta M_K = 2\,\mathrm{Re}\langle \bar{K}^0 | H_{\rm eff} | K^0 \rangle,$$
$$\Delta \epsilon_K = \frac{1}{\sqrt{2}\Delta M_K}\, \mathrm{Im}\langle \bar{K}^0 | H_{\rm eff} | K^0 \rangle, \quad (4.19)$$

we get

$$\Delta M_K = \frac{6.9269\times 10^{-7} A_{i4}^{2*} + 2.0088\times 10^{-7} B_{i4}^{2*}}{M_i^2}(1+\tan\beta^2) \quad (4.20)$$

and

$$\epsilon_K = \frac{9.9975\times 10^{6} A_{i4}^{2*} - 9.8616\times 10^{-9} A_{i4}^{*} B_{i4}^{*} + 2.8993\times 10^{7} B_{i4}^{2*}}{M_i^2}(1+\tan\beta^2). \quad (4.21)$$

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 [20]

$$|\Delta M_K| = M_{K_L} - M_{K_S} = (3.483 \pm 0.006) \times 10^{-12}\ \text{MeV}, \quad (4.22)$$

and indirect CP violation in $K \to \pi\pi$ [21] and in $K \to \pi l \nu$ decays is given by [20]

$$|\epsilon_K| = (2.228 \pm 0.011) \times 10^{-3}. \quad (4.23)$$

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 $M_{H_i^{0r}} = M_{H_i^{0i}} = M_i$, the value of $\Delta M_K = 3.483 \times 10^{-15}$ GeV yields the constraint

$$M_i^2 \geq (1.9888\times 10^{8} A_{i4}^{2*} + 5.7675\times 10^{8} B_{i4}^{2*})(1 + \tan\beta^2)\ \text{GeV}^2, \quad (4.24)$$

while the value of $\epsilon_K = 2.228 \times 10^{-3}$ [20] yields the constraint

$$M_i^2 \geq (4.4872\times 10^{9} A_{i4}^{2*} - 4.4262\times 10^{-6} A_{i4}^{*} B_{i4}^{*} + 1.3013\times 10^{10} B_{i4}^{2*})(1 + \tan\beta^2)\ \text{GeV}^2. \quad (4.25)$$

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 constraints become, for example, when $\tan\beta = 10$:

$$M_i^2 \geq (2.0087\times 10^{10} A_{i4}^{2*} + 5.8251\times 10^{10} B_{i4}^{2*})\ \text{GeV}^2 \quad (4.26)$$

and

$$M_i^2 \geq (4.5320\times 10^{11} A_{i4}^{2*} - 4.4704\times 10^{-4} A_{i4}^{*} B_{i4}^{*} + 1.3143\times 10^{12} B_{i4}^{2*})\ \text{GeV}^2. \quad (4.27)$$

We tried varying the lightest relative masses in the scalar and pseudoscalar sector and found that the results do not change.

C. $B_d^0 - \bar{B}_d^0$ mixing

We proceed the same way as for $K^0 - \bar{K}^0$ mixing to evaluate the constraints from the B_d^0, B_s^0 meson mixing. We use again four quark operators Q_1, Q_2, $\tilde{Q}_1$, and $\tilde{Q}_2$, defined previously. Setting as before the Higgs mass to be equal to the lightest scalar mass $M_{H_i^{0r}} = M_{H_i^{0i}} = M_i$, the expression for ΔM_{B_d} becomes

$$\Delta M_{B_d} = \frac{(9.4139\times 10^{-6} A_{i4}^{2*} + 3.6405\times 10^{-5} B_{i4}^{2*})(1+\tan^2\beta)}{M_i^2}\ \text{GeV}^3. \quad (4.28)$$

Using the experimental value of $\Delta M_{B_d} = 3.337 \times 10^{-13}$ GeV [20], we obtain, assuming, as before, dominance by the lightest mass

$$M_i^2 \geq (2.8211\times 10^{7} A_{i4}^{2*} + 1.6909\times 10^{8} B_{i4}^{2*}) \times (1+\tan^2\beta)\ \text{GeV}^2, \quad (4.29)$$

which becomes, for $\tan\beta = 10$,

$$M_i^2 \geq (2.8493\times 10^{9} A_{i4}^{2*} + 1.1019\times 10^{10} B_{i4}^{2*})\ \text{GeV}^2. \quad (4.30)$$

D. $B_s^0 - \bar{B}_s^0$ mixing

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

$$\Delta M_{B_s} = \frac{(4.2314 \times 10^{-4} A_{i4}^{*2} + 1.6469 \times 10^{-3} B_{i4}^{*2})(1 + \tan^2\beta)}{M_i^2}\ \text{GeV}^3. \tag{4.31}$$

Using the experimental value of $\Delta M_{B_d} = 117 \times 10^{-13}$ GeV [20,22],

$$M_i^2 \geq (3.6166 \times 10^7 A_{i4}^{2*} + 1.4076 \times 10^8 B_{i4}^{2*})(1 + \tan^2\beta)\ \text{GeV}^2 \tag{4.32}$$

or, for $\tan\beta = 10$,

$$M_i^2 \geq (3.6528 \times 10^9 A_{i4}^{2*} + 1.4217 \times 10^{10} B_{i4}^{2*})\ \text{GeV}^2. \tag{4.33}$$

E. $D^0 - \bar{D}^0$ mixing

In Subsec. IV 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 pseudoscalar Higgs fields.

$$\begin{aligned}\text{Re}\langle\bar{D}^0|H_{\text{eff}}|D^0\rangle = \frac{m_b^2\lambda^2}{4M_i^2(\kappa_1\kappa_2 - \kappa_1'\kappa_2')^2}\Bigg\{\Bigg[\bigg(\frac{m_s}{m_b} - \frac{m_d}{m_b}\bigg)(2-\lambda^2) + 2A^2\lambda^4\rho\Bigg]^2([(\kappa_2 A_{i3}^* - \kappa_2' A_{i5}^*)^2 - (\kappa_2 B_{i3}^* - \kappa_2' B_{i5}^*)^2] \\ \times (\langle\tilde{Q}_1(\mu)\rangle + \langle Q_1(\mu)\rangle) + [(\kappa_2 A_{i3}^* - \kappa_2' A_{i5}^*)^2 + (\kappa_2 B_{i3}^* - \kappa_2' B_{i5}^*)^2](\langle\tilde{Q}_2(\mu)\rangle + \langle Q_2(\mu)\rangle)) \\ + 4A^4\lambda^8\eta^2([(\kappa_2 A_{i3}^* - \kappa_2' A_{i5}^*)^2 + (\kappa_2 B_{i3}^* - \kappa_2' B_{i5}^*)^2](\langle\tilde{Q}_1(\mu)\rangle + \langle Q_1(\mu)\rangle) \\ + [(\kappa_2 A_{i3}^* - \kappa_2' A_{i5}^*)^2 - (\kappa_2 B_{i3}^* - \kappa_2' B_{i5}^*)^2](\langle\tilde{Q}_2(\mu)\rangle + \langle Q_2(\mu)\rangle)\Bigg\},\end{aligned} \tag{4.34}$$

where Q_1, Q_2, $\tilde{Q}_1$, and $\tilde{Q}_2$ are the four quark operators defined as before; the mass difference $\Delta M_D = 2\,\text{Re}\langle\bar{D}^0|H_{\text{eff}}|D^0\rangle$ is obtained as

$$\Delta M_D = \frac{5.2816 \times 10^{-10} A_{i5}^{2*} + 5.8097 \times 10^{-9} B_{i5}^{2*}}{M_i^2} \times \frac{(1 + \tan\beta^2)}{\tan\beta^2}\ \text{GeV}^3. \tag{4.35}$$

Comparing the calculated expression with the experimental value [20]

$$|\Delta M_D| = M_{D_1^0} - M_{D_2^0} = (1.573\,13) \times 10^{-17}\ \text{MeV}, \tag{4.36}$$

we obtain

$$M_i^2 \geq \frac{(3.3574 \times 10^{10} A_{i5}^{2*} + 3.6931 \times 10^{11} B_{i5}^{2*})(1 + \tan\beta^2)}{\tan\beta^2}\ \text{GeV}^2, \tag{4.37}$$

which becomes, for $\tan\beta = 10$,

$$M_i^2 \geq (3.3909 \times 10^{10} A_{i5}^{2*} + 3.7300 \times 10^{11} B_{i5}^{2*})\ \text{GeV}^2. \tag{4.38}$$

V. NUMERICAL RESULTS AND DISCUSSION

The FCNC tree-level diagrams are mediated by the physical scalar fields H_2^0 and H_6^0 and the pseudoscalars A_1^0 and A_4^0. These fields are linear superpositions of the χ_1^{0r} or ϕ_2^{0r} (χ_1^{0i} or ϕ_2^{0i}, respectively, for the pseudoscalars) components from the bidoublet Higgs.

As the fields H_2^0 and H_6^0 must be heavy, the light neutral scalars would likely be linear combinations of the complimentary χ_2^{0r} or ϕ_1^{0r} components from the bidoublets, or δ^{0r}, $\bar{\delta}^{0r}$, $\delta^{c^{0r}}$, and $\bar{\delta}^{c^{0r}}$ from the triplet Higgs. We set v_R in the interval obtained from the requirement that the doubly charged Higgses are light. 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_1^0 (SM-like) changes, at most, a few GeV, if we vary any of the parameters, whereas the second-lightest scalar field H_2^0 is highly dependent on the changes in the parameter v_R. Similarly, the lightest pseudoscalar field A_1^0 behaves like the second-lightest neutral scalar field and is also affected by the changes in v_R. H_1^0 is SM-like, and the parameter that seems to affect the H_1^0 mass the most is the λ_{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 $M_{H_1^0}$.

The tree-level flavor-changing neutral currents in the down-quark sector are governed by H_6^0 and A_4^0. The mass values of the fields H_6^0 and A_4^0 are the same, and they are dependent on the parameters λ_{21}, v_R, λ, $\tan\beta$, and M_R. Numerical investigation reveals that only $\tan\beta$, λ_{21}, and

M_R can affect the H_6^0 and A_4^0 masses significantly, while there is practically no variation with v_R. These masses are dependent on the parameters M_R, λ_{21}, and $\tan\beta$, such that, when they increase, mass values of these physical fields also increase. The dependence of the H_6^0 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_6^0 mass with these parameters are shown in Fig. 1.

The fields H_2^0 and A_1^0 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 Sec. III), and, although they depend, in principle, on v_R, $\tan\beta$, and λ_{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 3–12 times. The mass also varies with the ratio $\tan\delta = \bar{v}_R/v_R$, while it is 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 λ_{21} is very weak, the dependence on v_R is almost linear.

From the approximate analytical expressions in Sec. III, the mass of the lightest doubly charged physical field $H_1^{\pm\pm}$ depends on v_R, λ, and M_R, as well as on the soft slepton mass $m_{L^c}^2$. Analysis shows that only the dependence on v_R and $m_{L^c}^2$ 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 parameters as a plot, as a function of v_R for different values of $\tan\delta$, as well as contour plot in the $v_R - m_{L^c}^2$ plane. The mass of $H_1^{\pm\pm}$ 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 independent on M_R. One can see that 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. Of course, in all cases, different $m_{L^c}^2 < 0$ are needed to keep the masses positive. Within the parameter space considered, $m_{L^c} \in (4.5i\text{–}10i)$ TeV. The effect of varying the other parameters is negligible for the lightest doubly charged Higgs, whereas the mass of the heavier doubly charged Higgs $H_2^{\pm\pm}$ depends almost exclusively on M_R.

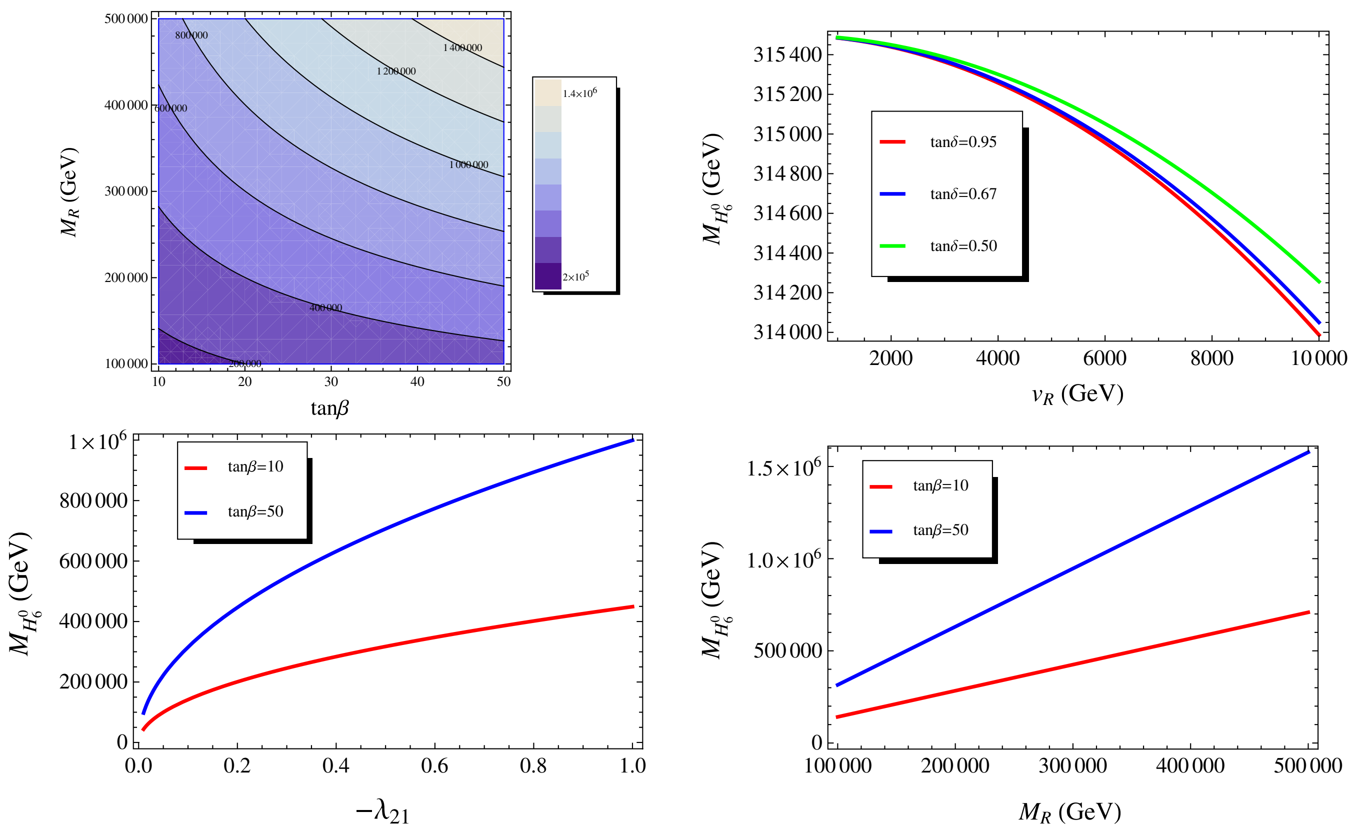

FIG. 1 (color online). The variation of the FCNC neutral Higgs H_6^0 mass with the parameters of the LRSUSY model. H_6^0 induces tree-level FCNC in the down-quark sector. Shown are contour plots in the $M_R - \tan\beta$ plane and the mass dependence on v_R for three values of $\tan\delta = \bar{v}_R/v_R$ (top row) and the mass dependence on λ_{21} and M_R for $\tan\beta = 10$ or $\tan\beta = 50$ (bottom row). Masses are given in GeV.

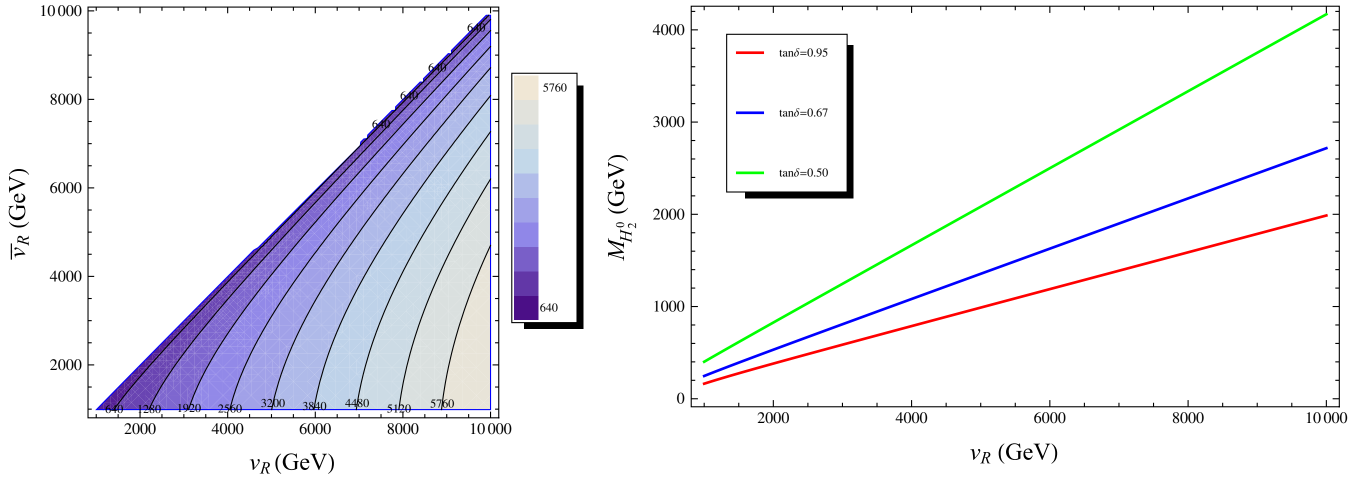

FIG. 2 (color online). The variation of the FCNC neutral Higgs H_2^0 mass with the parameters of the LRSUSY model. H_2^0 induces tree-level FCNC in the up-quark sector. To the left, a contour plot in the $v_R - \bar{v}_R$ plane; and, at the right, a plot as a function of v_R for three values of $\tan\delta = \bar{v}_R/v_R$. Masses are given in GeV.

The lightest singly charged physical field $H_1^{\pm}$ mass corresponds to the MSSM-like charged Higgs boson. The singly charged state that is tripletlike is H_2^+ and is heavy. The orthogonal combination of tripletlike charged Higgs bosons is the Goldstone boson $G_2^{\pm}$ responsible for giving mass to $W_R^{\pm}$ bosons. The other charged Higgses which come from bidoublet components are heavy, a consequence of requiring the mass parameters to satisfy FCNC bounds.

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 $\tan\beta = 10$ (Table IV), the other for $v_R = 5$ TeV and $\tan\beta = 50$ (Table V). 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, $\lambda = 1$, $\lambda_{21} = -0.1$, $C_\lambda = 2.5$ TeV, $\langle S \rangle = 1$ TeV, and $M_S = 1$ 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 light flavor-conserving neutral scalar Higgs boson; one light doubly charged Higgs boson; and three other Higgs bosons with masses below 1 TeV, one neutral and two singly charged. 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

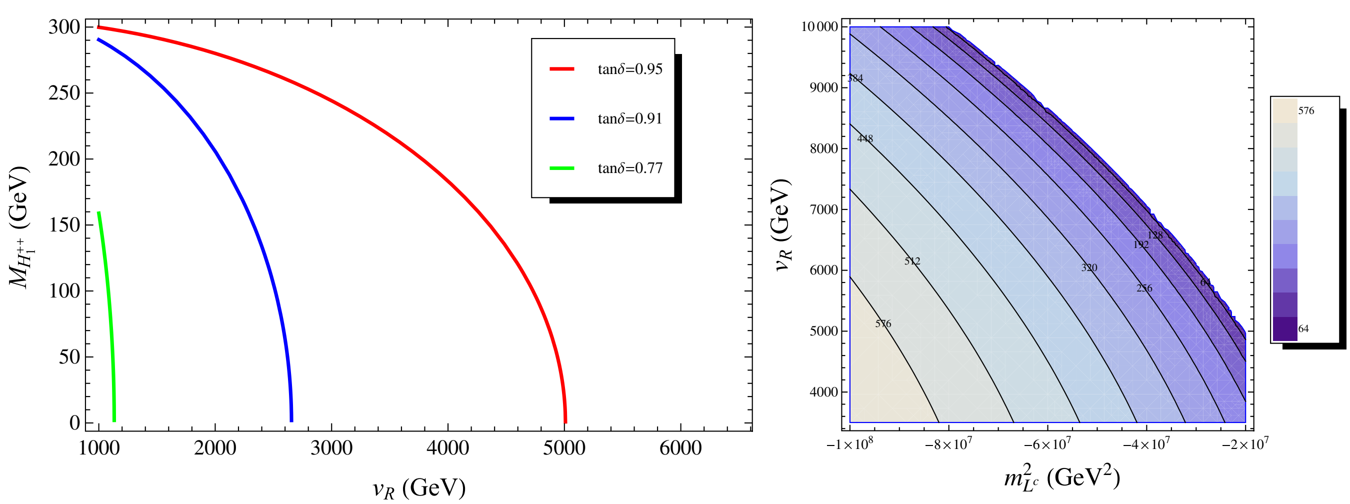

FIG. 3 (color online). The masses of the lightest doubly charged Higgs boson as a function of the v_R plane for three values of $\tan\delta = \bar{v}_R/v_R$ (left) and as a contour plot in the $v_R - m^2_{L^c}$ plane (right). Masses are given in GeV.

TABLE IV. Masses and compositions of physical Higgs fields and unphysical Goldstone bosons. Parameters are chosen as follows: $\tan\beta = 10$, $\tan\delta = 1/1.05$, $v_R = 3.5$ TeV, $M_R = 100$ TeV, $\lambda = 1$, $\lambda_{21} = -0.1$, $C_\lambda = 2.5$ TeV, $\langle S\rangle = 1$ TeV, $M_S = 1$ TeV, $m^2_{L^c} = -20$ TeV^2, and $f = 1$.

Particle	Mass (GeV)	Composition
H_1^0	111.6	$0.100\phi_1^{0r} + 0.995\chi_2^{0r}$
H_2^0	680.9	$-0.100\phi_2^{0r} - 1.000\chi_1^{0r}$
H_3^0	4557.4	$0.720\delta^{c^{0r}} + 0.685\bar{\delta}^{c^{0r}} + 0.001\chi_2^{0r} - 0.129S^{0r}$
H_4^0	11 140.6	$0.045\delta^{c^{0r}} + 0.043\bar{\delta}^{c^{0r}} + 0.998S^{0r}$
H_5^0	141 537.9	$0.690\delta^{c^{0r}} - 0.724\bar{\delta}^{c^{0r}}$
H_6^0	141 686.6	$-1.000\phi_2^{0r} + 0.100\chi_1^{0r}$
H_7^0	141 688.2	$0.001\delta^{c^{0r}} - 0.001\bar{\delta}^{c^{0r}} + 1.000\phi_1^{0r} - 0.100\chi_2^{0r}$
A_1^0	680.9	$0.100\phi_2^{0i} - 1.000\chi_1^{0i}$
A_2^0	11 106.8	$1.000S^{0i}$
A_3^0	141 502.0	$0.690\delta^{c^{0i}} + 0.724\bar{\delta}^{c^{0i}} - 0.003\phi_2^{0i}$
A_4^0	141 686.6	$-1.000\phi_2^{0i} - 0.100\chi_1^{0i}$
A_5^0	141 688.2	$-0.002\delta^{c^{0i}} - 0.002\bar{\delta}^{c^{0i}} - 1.000\phi_1^{0i} - 0.100\chi_2^{0i}$
H_1^+	690.2	$-0.018\delta^{c^{-*}} - 0.018\bar{\delta}^{c^+} - 0.099\phi_2^{-*} + 0.995\chi_1^+$
H_2^+	141 454.6	$0.690\delta^{c^{-*}} - 0.724\bar{\delta}^{c^+}$
H_3^+	141 686.6	$-0.995\phi_1^+ - 0.100\chi_2^{-*}$
H_4^+	449 688.2	$-0.995\phi_2^{-*} - 0.100\chi_1^+$
H_1^{++}	217.9	$-0.724\delta^{c^{--*}} - 0.690\bar{\delta}^{c^{++}}$
H_2^{++}	141 419.9	$-0.690\delta^{c^{--*}} + 0.724\bar{\delta}^{c^{++}}$
G_1^0	0	$0.568\delta^{c^{0i}} + -0.540\bar{\delta}^{c^{0i}} + 0.062\phi_1^{0i} - 0.617\chi_2^{0i}$
G_2^0	0	$-0.449\delta^{c^{0i}} + 0.428\bar{\delta}^{c^{0i}} + 0.078\phi_1^{0i} - 0.780\chi_2^{0i}$
G_1^+	0	$0.100\phi_1^+ - 0.995\chi_2^{-*}$
G_2^+	0	$-0.724\delta^{c^{-*}} - 0.690\bar{\delta}^{c^+} + 0.003\phi_2^{-*} - 0.025\chi_1^+$

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, $\lambda = 1$, $\lambda_{21} = -0.1$, $C_\lambda = 2.5$ TeV, $\langle S\rangle = 1$ TeV, $M_S = 1$ TeV, $m^2_{L^c} = -30$ TeV^2, and $f = 1$.

Particle	Mass (GeV)	Composition
H_1^0	113.6	$0.020\phi_1^{0r} + 1.000\chi_2^{0r}$
H_2^0	998.6	$-0.020\phi_2^{0r} - 1.000\chi_1^{0r}$
H_3^0	6797.1	$0.714\delta^{c^{0r}} + 0.680\bar{\delta}^{c^{0r}} - 0.168S^{0r}$
H_4^0	12 214.8	$0.068\delta^{c^{0r}} + 0.061\bar{\delta}^{c^{0r}} + 0.996S^{0r}$
H_5^0	141 575.2	$-0.690\delta^{c^{0r}} + 0.724\bar{\delta}^{c^{0r}}$
H_6^0	315 121.9	$-1.000\phi_2^{0r} + 0.020\chi_1^{0r}$
H_7^0	315 123.5	$-1.000\phi_1^{0r} + 0.020\chi_2^{0r}$
A_1^0	998.6	$0.020\phi_2^{0i} - 1.000\chi_1^{0i}$
A_2^0	12 152.2	$1.000S^{0i}$
A_3^0	141 502.0	$0.690\delta^{c^{0i}} + 0.724\bar{\delta}^{c^{0i}}$
A_4^0	315 121.9	$-1.000\phi_2^{0i} - 0.020\chi_1^{0i}$
A_5^0	315 123.5	$-1.000\phi_1^{0i} - 0.020\chi_2^{0i}$
H_1^+	995.3	$-0.013\delta^{c^{-*}} - 0.012\bar{\delta}^{c^+} - 0.020\phi_2^{-*} + 1.000\chi_1^+$
H_2^+	141 405.3	$0.690\delta^{c^{-*}} - 0.724\bar{\delta}^{c^+}$
H_3^+	315 121.9	$-1.000\phi_1^+ - 0.020\chi_2^{-*}$
H_4^+	315 123.5	$-1.000\phi_2^{-*} - 0.020\chi_1^+$
H_1^{++}	215.3	$-0.724\delta^{c^{--*}} - 0.690\bar{\delta}^{c^{++}}$
H_2^{++}	141 334.2	$-0.690\delta^{c^{--*}} + 0.724\bar{\delta}^{c^{++}}$
G_1^0	0	$-0.138\delta^{c^{0i}} + 0.131\bar{\delta}^{c^{0i}} + 0.019\phi_1^{0i} - 0.961\chi_2^{0i}$
G_2^0	0	$0.710\delta^{c^{0i}} - 0.677\bar{\delta}^{c^{0i}} + 0.004\phi_1^{0i} - 0.981\chi_2^{0i}$
G_1^+	0	$0.020\phi_1^+ - 1.000\chi_2^{-*}$
G_2^+	0	$0.724\delta^{c^{-*}} + 0.690\bar{\delta}^{c^+} + 0.018\chi_1^+$

in each sector. This scenario is completely consistent with the Tevatron [23] and LHC data [24] on Higgs boson searches.

Finally, we comment on the scalar leptons and gaugino masses. In [11], the authors attempt a complete model building, incorporating general (approximate) constraints on doubly charged Higgs boson fields and scalar lepton masses, as functions of gaugino masses. Using two-loop MSSM renormalization group equations [25], the relations between these parameters are

$$M_{++}^2(m_Z) < \frac{24}{5b_1} M_{\tilde{1}}^2(m_Z)\left[\frac{\alpha_{B-L}^2(v_R)}{\alpha_{B-L}^2(m_Z)} - 1\right],$$
$$M_{\tilde{\tau}_R}^2(m_Z) < \frac{5}{6b_1} M_{\tilde{1}}^2(m_Z)\left[\frac{\alpha_{B-L}^2(v_R)}{\alpha_{B-L}^2(m_Z)} - 1\right],$$
$$M_{\tilde{\tau}_L}^2(m_Z) < \frac{3}{10b_1} M_{\tilde{1}}^2(m_Z)\left[\frac{\alpha_{B-L}^2(v_R)}{\alpha_{B-L}^2(m_Z)} - 1\right] + \frac{3}{2b_2} M_{\tilde{L}}^2(m_Z)\left[\frac{\alpha_L^2(v_R)}{\alpha_L^2(m_Z)} - 1\right], \tag{5.1}$$

where $M_{\tilde{1}}$ and $M_{\tilde{L}}$ are gaugino masses, b_1 and b_2 are renormalization group equation coefficients, M_{++} is the soft doubly charged Higgs mass, and the last two equations give bounds on the right and left tau slepton masses. We use the renormalization group equations for the left-right supersymmetric model with triplets and an arbitrary number of singlets [26] to evaluate the mass bounds.[1] In our case, taking v_R in the 3.5–10 TeV region, the limits become

$$M_{++}^2(m_Z) < \frac{1}{8} M_{\tilde{1}}^2(m_Z)\left[\frac{\alpha_{B-L}^2(v_R)}{\alpha_{B-L}^2(m_Z)} - 1\right] + M_{\tilde{R}}^2(m_Z)\left[\frac{\alpha_R^2(v_R)}{\alpha_R^2(m_Z)} - 1\right],$$
$$M_{\tilde{\tau}_R}^2(m_Z) < \frac{1}{32} M_{\tilde{1}}^2(m_Z)\left[\frac{\alpha_{B-L}^2(v_R)}{\alpha_{B-L}^2(m_Z)} - 1\right] + \frac{3}{16} M_{\tilde{R}}^2(m_Z)\left[\frac{\alpha_R^2(v_R)}{\alpha_R^2(m_Z)} - 1\right],$$
$$M_{\tilde{\tau}_L}^2(m_Z) < \frac{1}{32} M_{\tilde{1}}^2(m_Z)\left[\frac{\alpha_{B-L}^2(v_R)}{\alpha_{B-L}^2(m_Z)} - 1\right] + \frac{3}{16} M_{\tilde{L}}^2(m_Z)\left[\frac{\alpha_L^2(v_R)}{\alpha_L^2(m_Z)} - 1\right]. \tag{5.2}$$

The approximate bounds on the soft masses depend critically on the relationship between the $U(1)_{B-L}$, $SU(2)_L$, and $SU(2)_R$ gaugino masses. For instance, for $M_{\tilde{L}} = M_{\tilde{R}} = M_{\tilde{1}}$: $M_{++}(m_Z) \le 0.24 M_{\tilde{1}}(m_Z)$, $M_{\tilde{\tau}_R}(m_Z) \le 0.11 M_{\tilde{1}}(m_Z)$, and $M_{\tilde{\tau}_L}(m_Z) \le 0.11 M_{\tilde{1}}(m_Z)$; for $M_{\tilde{L}} = M_{\tilde{R}} = 2M_{\tilde{1}}$, the bounds become $M_{++}(m_Z) \le 0.4 M_{\tilde{1}}(m_Z)$, $M_{\tilde{\tau}_R}(m_Z) \le 0.2 M_{\tilde{1}}(m_Z)$, and $M_{\tilde{\tau}_L}(m_Z) \le 0.2 M_{\tilde{1}}(m_Z)$; while, for $2M_{\tilde{L}} = M_{\tilde{R}} = 4M_{\tilde{1}}$, the limits are $M_{++}(m_Z) \le 0.7 M_{\tilde{1}}(m_Z)$, $M_{\tilde{\tau}_R}(m_Z) \le 0.32 M_{\tilde{1}}(m_Z)$, and $M_{\tilde{\tau}_L}(m_Z) \le 0.2 M_{\tilde{1}}(m_Z)$. While no precise conclusions can be reached, the bounds push the gaugino mass parameter $M_{\tilde{1}}$ to be very large, which is consistent with soft slepton masses in the TeV range. Note that these mass bounds are only a rough estimate, as we include gaugino masses but neglect other terms. The purpose of our calculations was to show that a self-consistent Higgs sector can be obtained within the framework of the minimal model, leaving the door open for a more thorough exploration of the model.

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 two doubly charged Higgses, four singly charged Higgs fields, seven neutral scalar fields, and five 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 requirements that 1) there is a light neutral scalar Higgs boson, flavor-conserving, which is the counterpart to the SM Higgs boson; 2) there exists 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 $B_{d,s}^0 - \bar{B}_{d,s}^0$ mixings, make the Higgs sector fairly predictive and fix 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 requirements 1) and 2) are related and satisfied by the $v_R \in (3, 10)$ TeV range. One could obtain much heavier Higgs bosons by increasing v_R. As long as $m_{L^c}^2$ is also increased, one of the doubly charged Higgs bosons remains light, but the rest of the mass spectrum shifts to higher scales. 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 and singly and doubly charged Higgs bosons, it predicts new gauge bosons just outside the range $M_{W_R} < 2(4)$ TeV, which can be observed at the LHC with a luminosity of 1(30) fb^{-1} [27].

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

[1] These bounds are completely consistent with what we obtain using relations given in [11] for $M_{\tilde{1}} = M_{\tilde{L}} = M_{\tilde{R}}$.

estimates show that requiring some of the Higgs bosons to be light will likely push the scale of supersymmetry above 1 TeV.

Our analysis is important for two reasons: first, we have shown that a reasonable Higgs mass spectrum is possible in LRSUSY. This analysis shows that, except for these two bosons, the rest are heavy (with only three—one neutral scalar, one pseudoscalar and one singly charged) just below the TeV scale. 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. The best signal for this model from the Higgs sector remains the observation of a doubly charged Higgs boson, decaying copiously to charged leptons. Observation of a light non-SM Higgs (neutral or singly charged) will invalidate the model, at least within the present minimal prescription for the Higgs sector. 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 and of the right-handed neutrinos, as well as for the supersymmetric partners.

ACKNOWLEDGMENTS

We thank NSERC of Canada for partial financial support under Grant No. SAP01105354. We are grateful to K. S. Babu and Ayon Patra for illuminating discussions on the model presented here.

[1] R. N. Mohapatra and J. C. Pati, Phys. Rev. D **11**, 2558 (1975); G. Senjanovic and R. N. Mohapatra, *ibid.* **12**, 1502 (1975).

[2] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).

[3] J. F. Gunion, J. Grifols, A. Mendez, B. Kayser, and F. I. Olness, Phys. Rev. D **40**, 1546 (1989); J. A. Grifols, *ibid.* **18**, 2704 (1978); J. A. Grifols, A. Mendez, and G. A. Schuler, Mod. Phys. Lett. A **4**, 1485 (1989); N. G. Deshpande, J. F. Gunion, B. Kayser, and F. I. Olness, Phys. Rev. D **44**, 837 (1991); M. L. Swartz, *ibid.* **40**, 1521 (1989); R. Vega and D. A. Dicus, Nucl. Phys. **B329**, 533 (1990); K. Huitu, J. Maalampi, A. Pietila, and M. Raidal, *ibid.* **B487**, 27 (1997); A. Datta and A. Raychaudhuri, Phys. Rev. D **62**, 055002 (2000); G. Barenboim, M. Gorbahn, U. Nierste, and M. Raidal, *ibid.* **65**, 095003 (2002); K. Kiers, M. Assis, and A. A. Petrov, *ibid.* **71**, 115015 (2005); A. G. Akeroyd and M. Aoki, *ibid.* **72**, 035011 (2005).

[4] R. M. Francis, M. Frank, and C. S. Kalman, Phys. Rev. D **43**, 2369 (1991); K. Huitu, J. Maalampi, and M. Raidal, Phys. Lett. B **328**, 60 (1994); Nucl. Phys. **B420**, 449 (1994).

[5] R. N. Mohapatra and A. Rasin, Phys. Rev. Lett. **76**, 3490 (1996); Phys. Rev. D **54**, 5835 (1996); R. Kuchimanchi, Phys. Rev. Lett. **76**, 3486 (1996).

[6] K. S. Babu and S. M. Barr, Phys. Rev. D **48**, 5354 (1993); **50**, 3529 (1994); M. Frank, H. Hamidian, and K. Puolamäki, Phys. Lett. B **456**, 179 (1999); Phys. Rev. D **60**, 095011 (1999); For a review and further references, see, e.g., R. N. Mohapatra, arXiv:hep-ph/9801235.

[7] R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. D **48**, 4352 (1993).

[8] C. S. Aulakh, A. Melfo, and G. Senjanovic, Phys. Rev. D **57**, 4174 (1998).

[9] Z. Chacko and R. N. Mohapatra, Phys. Rev. D **58**, 015003 (1998).

[10] K. Huitu and J. Maalampi, Phys. Lett. B **344**, 217 (1995).

[11] K. S. Babu and R. N. Mohapatra, Phys. Lett. B **668**, 404 (2008).

[12] C. S. Aulakh, A. Melfo, A. Rasin, and G. Senjanovic, Phys. Rev. D **58**, 115007 (1998); C. S. Aulakh, K. Benakli, and G. Senjanovic, Phys. Rev. Lett. **79**, 2188 (1997).

[13] K. S. Babu (private communication).

[14] M. E. Pospelov, Phys. Rev. D **56**, 259 (1997).

[15] L. Wolfenstein, Phys. Rev. Lett. **51**, 1945 (1983).

[16] M. Bona *et al.* (UTfit Collaboration), J. High Energy Phys. 07 (2005) 028, and updates at http://www.utfit.org/.

[17] M. Bona *et al.* (UTfit Collaboration), J. High Energy Phys. 03 (2008) 049.

[18] A. J. Buras, S. Jager, and J. Urban, Nucl. Phys. **B605**, 600 (2001).

[19] D. Becirevic, V. Gimenez, G. Martinelli, M. Papinutto, and J. Reyes, J. High Energy Phys. 04 (2002) 025.

[20] K. Nakamura *et al.* (Particle Data Group), J. Phys. G **37**, 075021 (2010).

[21] J. H. Christenson, J. W. Cronin, V. L. Fitch, and R. Turlay, Phys. Rev. Lett. **13**, 138 (1964).

[22] A. Abulencial *et al.* (CDF Run II Collaboration), Phys. Rev. Lett. **97**, 062003 (2006); **97**, 242003 (2006); V. M. Abazov *et al.* (D0 Collaboration), *ibid.* **97**, 021802 (2006).

[23] T. Aaltonen *et al.* (CDF Collaboration and D0 Collaboration), arXiv:1103.3233.

[24] S. Chatrchyan *et al.* (CMS Collaboration), arXiv:1102.5429.

[25] S. P. Martin and M. T. Vaughn, Phys. Rev. D **50**, 2282 (1994).

[26] N. Setzer and S. Spinner, Phys. Rev. D **71**, 115010 (2005).

[27] A. Ferrari *et al.*, Phys. Rev. D **62**, 013001 (2000).