Exploring the charged Higgs bosons in the left-right symmetric model

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‡(Dated: February 26, 2018)

We explore constraints on the charged Higgs sector in the left-right symmetric model from the present experimental data. Due to the different Yukawa structure, the allowed parameter space of the charged Higgs boson in the LR model is different from that of the two Higgs doublet model. We find that the constraint from $t \to bH^\pm$ decay at Tevatron is most significant, while $B \to \tau\nu$ decay could provide no constraint on the LR model. Bounds from $e^-e^+ \to H^-H^+$ process at LEP are similar to those in the two Higgs doublet model.

I. INTRODUCTION

Understanding the mechanism of the electroweak symmetry breaking (EWSB) is one of the most important motivations of new physics beyond the standard model (SM). The nature of the EWSB will be experimentally studied at the CERN Large Hadron Collider (LHC) and the $e^-e^+$ linear collider (ILC) in the future. In the SM, one neutral Higgs boson exists as a result of the EWSB, of which mass is not predicted in the theoretical framework. In many models of new physics beyond the SM, more symmetries are involved and the Higgs sector should be extended to break larger symmetry. Generically, the charged Higgs bosons are predicted by models with extended Higgs sector although they do not exist within the SM. Thus the observation of the charged Higgs boson is clearly a direct evidence of the new physics. The charged Higgs boson in the two Higgs doublet (2HD) model or the minimal supersymmetric standard model (MSSM) has been examined through the pair production at the CERN Large Electron Positron Collider (LEP) [1] and top quark decay process $t \to bH^\pm$ at Tevatron [2]. The absence of the observed charged Higgs boson so far derives constraints on $(\tan \beta, m_{H^\pm})$ parameter space for these 2HD type models. Recent measurement of $\text{Br}(B^{\pm} \to \tau\nu)$ by Belle provides indirect constraints [3] since $B^{\pm} \to \tau\nu$ channel is sensitive to the annihilation diagram mediated by the charged Higgs boson in the 2HD model [3]. The phenomenology of the charged Higgs boson has been widely studied at the LHC [5, 6, 7, 8, 9, 10, 11].

Inspired by the custodial symmetry in the Higgs potential, additional SU(2) gauge symmetry attracts much interest as an underlying structure of the new physics [12, 13]. The left-right (LR) symmetric model based on the gauge symmetry, $\text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L}$, is one of the attractive extensions of the SM with the additional $\text{SU}(2)_R$ symmetry [14]. In the minimal version of the LR model with the manifest left-right symmetry, the parity is an exact symmetry of the Lagrangian and spontaneously broken along with the gauge symmetry breaking. The triplet Higgs boson $\Delta_R$ is introduced to break the $\text{SU}(2)_R$ symmetry and another triplet $\Delta_L$ introduced as a result of left-right symmetry. The scale of the $\text{SU}(2)_R$ symmetry breaking is required to be much higher than the electroweak scale since the masses of the heavy gauge bosons are constrained by experiments [15, 16, 17, 18]. The right-handed fermions transform as doublets under $\text{SU}(2)_R$ and singlets under $\text{SU}(2)_L$ and the left-handed fermions behave reversely in this model. Thus a bidoublet Higgs field is introduced for the Yukawa couplings and the EWSB. The triplet Higgs fields violate the lepton number and baryon number and do not allow the ordinary Yukawa coupling terms. Consequently, the weak scale phenomenology of the Higgs sector is principally determined by the bidoublet Higgs fields, and the dominant field contents are similar to those of the 2HD model: a pair of the charged Higgs boson and three neutral Higgs bosons. However, the structure of the Yukawa couplings and potential of the bidoublet Higgs fields are much different from those of the 2HD model. It leads to the different phenomenology involving charged Higgs bosons and different constraints on the parameter space by experiments from those in the 2HD type models.

In this paper, we examine the constraints on the charged Higgs sector of the LR model using the present experimental results at LEP, Tevatron and B-factory. This paper is organized as follows: In section 2, the LR model is briefly

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reviewed, focusing on the Higgs sector. The analysis of $t \to bH^\pm$ process at Tevatron, $H^\pm$ pair production at LEP, and $B \to \tau\bar{\nu}$ decay at Belle in the LR model are presented in section 3. Finally we conclude in section 4.

II. THE LEFT-RIGHT SYMMETRIC MODEL

The Higgs sector of the minimal LR model consists of a bidoublet Higgs field $\phi(2, 2, 0)$ and two triplet Higgs fields $\Delta_L(3, 1, 2)$ and $\Delta_R(1, 3, 2)$ represented by

$$\phi = \left( \begin{array}{c} \phi^0_L \\ \phi^+_R \\ \phi^0_R \end{array} \right), \quad \Delta_{L,R} = \frac{1}{\sqrt{2}} \left( \begin{array}{c} \phi_0^L \\ \sqrt{2}\phi_0^R \\ -\phi_0^{+R} \end{array} \right),$$

under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group. The general Higgs potential has been analyzed in [19, 20, 21]. Minimizing condition of the Higgs potential is presented in Eq. (10) - (15) of Ref. [21]. For simplicity in this work, we assume no CP violation in the Higgs sector. The gauge symmetries are spontaneously broken by the vacuum expectation values (VEV)

$$\langle \phi \rangle = \frac{1}{\sqrt{2}} \left( \begin{array}{c} k_1 \\ 0 \\ k_2 \end{array} \right), \quad \langle \Delta_{L,R} \rangle = \frac{1}{\sqrt{2}} \left( \begin{array}{c} 0 \\ v_{L,R} \\ 0 \end{array} \right),$$

which lead to the charged gauge boson masses defined by

$$\mathcal{L}_M = (W^+_{L,R}, W^+_{L,R}^\dagger) M_{W^\pm}^2 \left( \begin{array}{c} W^-_L \\ W^-_R \end{array} \right),$$

where

$$M_{W^\pm}^2 = \frac{g^2}{4} \left( \begin{array}{cc} k_+^2 + 2v_L^2 & -2k_1k_2 \\ -2k_1k_2 & k_+^2 + 2v_R^2 \end{array} \right),$$

with $k_+^2 = |k_1|^2 + |k_2|^2$. Since $SU(2)_R$ symmetry should be broken by $v_R$ at the higher scale than the electroweak scale $\sim k_+$, we have $k_{1,2} \ll v_R$. Actually $v_L$ is irrelevant for the symmetry breaking and just introduced in order to manifest the left-right symmetry. If neutrino masses are derived by the see-saw mechanism, $m_\nu \sim v_L + k_+^2/v_R$, $v_R$ should be very large $\sim 10^{11}$ GeV. Then the heavy gauge bosons are too heavy to be produced at the accelerator experiments and the SU(2)$_R$ structure is hardly probed in the laboratory. Thus we assume $v_R$ to be only moderately large for the heavy gauge bosons to be studied at LHC. Since $v_L$ is less than a generic neutrino mass from the see-saw relation, it should be very small and close to 0. This is achieved when the quartic couplings of $(\phi\phi\Delta_{L,R})$-type terms in the Higgs potential are set to be zero [21, 22]. We adopt this limit here.

We introduce the parameters $\xi = k_2/k_1$ and $\epsilon = k_1/v_R$. Since the parameter $\xi$ is the ratio of two VEVs for the EWSB, it is corresponding to $\tan \beta$ in the 2HD model. It is clear that $\epsilon \ll 1$. We diagonalize the mass matrix of the charged gauge bosons by a unitary transform

$$\left( \begin{array}{c} W^\pm_L \\ W^\pm_R \end{array} \right) = \left( \begin{array}{cc} \cos \zeta & \sin \zeta \\ -\sin \zeta & \cos \zeta \end{array} \right) \left( \begin{array}{c} W_1^\pm \\ W_2^\pm \end{array} \right),$$

with the mixing angle

$$\tan 2\zeta = \frac{2k_1k_2}{v_L^2 - v_R^2} \approx -2\epsilon^2 \xi,$$

which yield the masses

$$M_{W_1}^2 \approx \frac{g^2}{4} |k_+|^2 \left( 1 - \epsilon^2 \frac{2\xi^2}{1 + \xi^2} \right),$$

$$M_{W_2}^2 \approx \frac{g^2}{4} 2v_R^2 \left( 1 + \epsilon^2 \frac{1 + \xi^2}{2} \right),$$

in the leading order of $\epsilon$. We identify $W_1 \equiv W_{SM}$ and let $W_2 \equiv W'$ hereafter.
The Yukawa couplings for quark sector is written by

\[ - \mathcal{L} = \bar{\Psi}_L \begin{pmatrix} \mathcal{F} + \mathcal{G} \end{pmatrix} \Psi_R + H.c., \tag{8} \]

where \( \Psi^i = (\hat{U}, \hat{D})^\dagger \) is the flavour eigenstates, \( \hat{\phi} = \tau_2 \phi^+ \tau_2 \), and \( \mathcal{F}, \mathcal{G} \) are 3 \times 3 Yukawa coupling matrices. We rotate \( \hat{U} \) and \( \hat{D} \) into the mass eigenstates by unitary transforms

\[
\begin{align*}
\hat{U}_{L,R} &= V^U_{L,R} U_{L,R}, \\
\hat{D}_{L,R} &= V^D_{L,R} D_{L,R},
\end{align*}
\tag{9}
\]

and define Cabibbo-Kobayashi-Maskawa (CKM) matrix \( V^{CKM}_L = V^U_L \dagger V^D_R \). Here we assume the manifest left-right symmetry, \( V^{CKM}_L = V^{CKM}_R \). Solving the equations for the diagonal mass matrices,

\[
\begin{align*}
\frac{1}{\sqrt{2}} V^U_{\alpha} (\mathcal{F} k_1 + \mathcal{G} k_2) V^U_{\beta} &= \mathcal{M}^U, \\
\frac{1}{\sqrt{2}} V^D_{\alpha} (\mathcal{F} k_2 + \mathcal{G} k_1) V^D_{\beta} &= \mathcal{M}^D,
\end{align*}
\tag{10}
\]

the Yukawa coupling matrices \( \mathcal{F} \) and \( \mathcal{G} \) are given by

\[
\begin{align*}
\mathcal{F} &= \frac{\sqrt{2}}{k^2} \left( k_1 V^U_1 \mathcal{M}^U_{11} V^U_{11} - k_2 V^D_1 \mathcal{M}^D_{11} V^D_{11} \right), \\
\mathcal{G} &= \frac{\sqrt{2}}{k^2} \left( -k_2 V^U_1 \mathcal{M}^U_{11} V^U_{11} + k_1 V^D_1 \mathcal{M}^D_{11} V^D_{11} \right),
\end{align*}
\tag{11}
\]

where \( k^2 = |k_1|^2 - |k_2|^2 \). If \( \xi = 1 \), these solutions for the Yukawa coupling matrices given in Eq. (11) no more hold. Actually Eq. (10) is overdetermined and we have to treat it in a separate way. On the other hand, \( \xi = 1 \) implies the maximal LR mixing, which is phenomenologically disfavored. When \( M_{W'} = 1 \, \text{TeV}, |\tan 2\xi| \sim 0.01 \) while the indirect constraints on the mixing angle \( \xi \) derive the bound \( |\xi| < 10^{-3} \) \cite{22, 23}. Although small \( \xi \) is preferred in order to generate the ratio \( m_\ell/m_t, \xi > 1 \) region cannot be excluded in general. We do not consider the \( \xi = 1 \) case in this work.

Taking the limit that the quartic couplings for \( (\phi \phi \Delta_L \Delta_R) \) terms and \( v_L \) go to 0, the charged Higgs boson mass matrix is given in the basis of \( (\phi_1^+, \phi_2^+, \delta_R^+, \delta_L^+) \) by

\[
\begin{pmatrix}
\frac{m_+^2}{\sqrt{2} m_r^2} & m_\pm^2 \xi & \frac{1}{\sqrt{2}} m_r^2 (1-\xi^2) & 0 \\
\frac{m_\pm^2}{\sqrt{2} m_r^2} & \frac{m_\pm^2}{\sqrt{2} m_r^2} & \frac{1}{\sqrt{2}} m_r^2 (1-\xi^2) & 0 \\
\frac{1}{\sqrt{2}} m_r^2 (1-\xi^2) & \frac{1}{\sqrt{2}} m_r^2 (1-\xi^2) & \frac{m_r^2}{2} (1-\xi^2)^2 & 0 \\
0 & 0 & 0 & m_r^{(+)2}
\end{pmatrix}
\tag{12}
\]

where \( m_\pm^2 = \alpha_3 v_r^2/2(1-\xi^2) \) with the quartic coupling \( \alpha_3 \) for \( \text{Tr}(\phi^4 \phi \Delta_L \Delta^R_L) + \text{Tr}(\phi^4 \phi \Delta_R \Delta^R_L) \) term \cite{21}. This limit is warranted by the approximate horizontal U(1) symmetry \cite{24} and the see-saw picture for light neutrino masses. Higgs boson masses are not affected by taking this limit \cite{21}. If \( \xi > 1 \), \( \alpha_3 \) should be negative to avoid the dangerous negative mass square of scalar fields. Note that \( \delta_L^+ \) field decouples from other three charged Higgs fields with mass \( m_r^{(+)2} \) and is irrelevant for our phenomenological discussion here since it comes from the Higgs triplet \( \Delta_L \). By an appropriate unitary transform \( V \), we diagonalize the mass matrix \( V^\dagger M_r^2 V = M_r^{diag} \) to lead to the transform \( (\phi_1^+, \phi_2^+, \delta_R^+, \delta_L^+) \rightarrow (G_{1L}^+, G_{2L}^+, H^+, \delta_L^+) \), where \( G_{1,2} \) are Goldstone bosons and \( H^+ \) is the charged Higgs boson of which mass is given by

\[
m_m^2 = m_+^2 (1+\xi^2) \left( 1 + \frac{1}{2} \xi^2 \right) + \mathcal{O}(\xi^4), \tag{13}\]

which couples to the quarks and leptons.

Finally we obtain the \( H^\pm \)-quark–quark couplings

\[
-\mathcal{L} = \bar{D}(g_L P_L + g_R P_R) U H^- + H.c., \tag{14}\]
where the couplings are defined by

\[
\begin{align*}
g_L &= \sqrt{2\sqrt{2}G_F V_{UD}^*} \left( m_U \frac{1 + \xi^2}{|1 - \xi^2|} - m_D \frac{2\xi}{|1 - \xi^2|} \right) \left( 1 - \frac{1}{4} \xi^2 (1 + \xi^2) \right) + \mathcal{O}(\xi^4), \\
g_R &= \sqrt{2\sqrt{2}G_F V_{UD}^*} \left( m_U \frac{2\xi}{1 - \xi^2} - m_D \frac{1 + \xi^2}{|1 - \xi^2|} \right) \left( 1 - \frac{1}{4} \xi^2 (1 + \xi^2) \right) + \mathcal{O}(\xi^4).
\end{align*}
\]

We also have the lepton Yukawa couplings involving the lepton number violating terms

\[
\begin{align*}
- \mathcal{L} &= f_{ij} \bar{\Psi}_L^i \phi \Psi_R^j + g_{ij} \bar{\Psi}_L^i \tilde{\phi} \Psi_R^j + \text{H.c.} + i(h_M)_{ij} \left( \Psi_L^T C \Delta \xi \Psi_L^j + \Psi_R^T C \Delta \xi \Psi_R^j \right) + \text{H.c.},
\end{align*}
\]

where \( f \) and \( g \) are \( 3 \times 3 \) Yukawa coupling matrices for the Dirac masses while \( h_M \) is a \( 3 \times 3 \) Yukawa coupling matrix to yield the lepton number violating Majorana masses. Focusing on the charged Higgs boson coupling here, we ignore the masses of neutrinos. We write the Yukawa couplings for leptons by

\[
- \mathcal{L} = g_l (\bar{L}_L^i, \nu_l) H^- + \text{H.c.}
\]

where \( g_l = \sqrt{2\sqrt{2}G_F} \cdot m_i \frac{2\xi}{|1 - \xi^2|} \left( 1 - \frac{1}{4} \xi^2 (1 + \xi^2) \right) + \mathcal{O}(\xi^4). \)

### III. CONSTRAINTS FROM EXPERIMENTS

#### A. Top decays at Tevatron

The charged Higgs boson of which mass is \( m_{H^\pm} < m_t - m_b \) can be produced through the top quark decay \( t \to bH^\pm \) competing with the SM decay \( t \to bW^\pm \). The production of such a light charged Higgs boson has been examined using an integrated luminosity of 193 pb\(^{-1}\) data of CDF collaboration at Tevatron in the MSSM framework \(^2\). Considering \( t\bar{t} \) production, the expected number of events for observed channel should be modified in the presence of the charged Higgs boson, which depends on the top and Higgs branching ratios. Final states of \( t\bar{t} \) events consist of four channels; all-hadronic channel, lepton+jet channel, dilepton channel, and lepton+\( \tau \) channel. The expected number of events in channel \( k \) is given by

\[
\mu_k = \sigma_{t\bar{t}}^{\text{prod}} A_k + n_k^{\text{back}},
\]

where \( \sigma_{t\bar{t}}^{\text{prod}} \) is the \( t\bar{t} \) production cross section and \( n_k^{\text{back}} \) the number of SM-expected background events. The detector acceptance \( A_k \) is defined by

\[
A_k = \sum_{i,j} B_i \cdot \bar{B}_j \cdot \epsilon_{ij,k},
\]

where \( B_i \) (\( \bar{B}_j \)) denotes the branching ratios of the \( t \) (\( \bar{t} \)) quark decaying into \( i(j) \)-th modes, which are \( 1 \) \( t \to W^+b, 2 \) \( t \to H^+b, H^+ \to c\bar{s}, 3 \) \( t \to H^+b, H^+ \to \tau\nu, 4 \) \( t \to H^+b, H^+ \to t\bar{b}, 5 \) \( t \to H^+b, H^+ \to W^+h^0, h^0 \to bb \). The efficiencies times integrated luminosity in channel \( k \), \( \epsilon_{ij,k} \) are obtained from the Monte Carlo simulation of \( t\bar{t} \) events depending on the parameters of the model and presented at Ref. \(^{24}\). The exclusion region on the model parameter space is obtained by the absence of the observed charged Higgs boson using the number of events in the CDF data, which implies that \( \Gamma(t \to H^+b) < \Gamma_0 \), where the CDF exclusion limit \( \Gamma_0 \) is a function of model parameters.

In the LR model, the relevant model parameters are the charged Higgs mass \( m_{H^\pm} \) and the ratio of VEVs for EWSB \( \xi \equiv k_2/k_1 \) which is corresponding to \( \tan \beta \) in the 2HD type model. Since the dependence of the Yukawa coupling on \( \xi \) is different from that on \( \tan \beta \), the phenomenology of the charged Higgs boson in the LR model is also different from that in the 2HD model. As shown in the Eq. (15), Yukawa couplings depend upon \( \xi \) in the form of \( (1 + \xi^2)/(1 - \xi^2) \) or \( \xi/|1 - \xi^2| \), while those in the 2HD model are proportional to \( \tan \beta \) or \( 1/\tan \beta \). Instead of using the number of events observed in the data, we just read out the bound \( Br_{[0]} = \Gamma_{0}/[\Gamma_{t}^{\text{SM}} + \Gamma(t \to H^\pm b)] \) by comparing the contour plot of \( Br(t \to bH^\pm) \) on the \( (\tan \beta, m_{H^\pm}) \) plane presented in Ref. \(^{24}\) to the exclusion limits given as the results of Ref. \(^{2}\) for \( m_{H^\pm} = 80 - 160 \) GeV. Using this bound, we vary \( \xi \) and set the condition \( Br_{[0]}(t \to H^\pm b) < Br_{[0]} \) for each
$m_{H^\pm}$ to obtain the exclusion limits on $(\xi, m_{H^\pm})$. The exclusion region on $(\xi, m_{H^\pm})$ plane is depicted in Fig. 1 where the SM total decay width $\Gamma^{SM}_{\text{SM}} = 1.42$ GeV [25]. We can find that the exclusion region is much different from that of $(\tan \beta, m_{H^\pm})$ plane in the 2HD model presented in Ref [2]. We find that there is a lower bound on $m_{H^\pm} > 145$ GeV in the LR model. Note that $m_{H^\pm}$ is proportional to $v_R$, while $m_t$ is determined by $k_+$. Thus the parameter region $m_{H^\pm} < m_t - m_b$ examined here denotes a very small $\alpha_3$ region.

![Excluded region on $(\xi, m_{H^\pm})$ parameter space by Tevatron and LEP data at 95% C.L.](image)

**FIG. 1:** Excluded region on $(\xi, m_{H^\pm})$ parameter space by Tevatron and LEP data at 95% C.L..

### B. Pair production at LEP

At the $e^- e^+$ collider, the most promising channel of $H^\pm$ production is the tree level pair production mediated by neutral gauge bosons. There exist three neutral gauge fields in the LR model, $B_\mu$, $W^3_{L\mu}$ and $W^3_{R\mu}$. We diagonalize the mass matrix by an orthogonal transform to produce one massless photon, one massive gauge boson identical to $Z$ boson, and one new heavy gauge boson $Z'$. In the orthogonal transform, we need three mixing angles, $\theta_W$, $\theta_R$, and $\theta_\xi$, among which $\theta_W$ is identical to the Weinberg angle in the SM. We use the notation of angles introduced in Ref. [17] where the constraints on the mixing angles have been studied in detail. Note that $\xi$ is replaced by $\theta_\xi$ in this paper.

We write the charged Higgs boson couplings to neutral gauge bosons in terms of their physical states. The photon coupling measures the electric charge of the charged Higgs boson and is same as that in the 2HD model. Since the mass of $Z'$ is bounded by 630 GeV from direct search and 860 GeV from electroweak fit [28], the contribution of $Z'$ is suppressed to be 10% or less compared with those of photon or $Z$ boson. Ignoring the $Z'$ contribution, the cross
gauge bosons and charged Higgs bosons such that
\[ \frac{\sigma(e^+e^- \to H^-H^+)}{s} = \frac{\pi\alpha^2}{3s} \left( 1 + \frac{2g_Y^2g_W^2}{1 - m_Z^2/s} + \frac{(g_Y^2 + g_W^2)g_V^2}{(1 - m_Z^2/s)^2} \right) \beta_H^2, \]  
(21)
where \( \beta_H = \sqrt{1 - 4m_H^2/s} \) is the velocity of the Higgs boson. The \( ZH^+H^- \) coupling is expressed by the mixing of gauge bosons and charged Higgs bosons such that
\[ g_Y^2 = -g_L^2 \left( \cos \theta_L \cos \theta_W + \sin \theta_L \cos \theta_R - \cos \theta_L \sin \theta_W \sin \theta_R \right) - \frac{2(1 + \xi^2)}{2(1 + \xi^2) + \epsilon^2(1 - \xi^2)^2} + g_R \left( \cos \theta_L \sin \theta_W \cos \theta_R + \sin \theta_L \sin \theta_W \sin \theta_R \right) \frac{\epsilon^2(1 - \xi^2)^2}{2(1 + \xi^2) + \epsilon^2(1 - \xi^2)^2}. \]  
(22)

Since we consider the minimal model with \( g_L = g_R \), the mixing angles satisfy the relation \( \sin \theta_R = \tan \theta_W \). Moreover the mixing angle \( \theta_L \) is strongly constrained by the experiment \[17\]. Thus the contribution to \( g_Y^2 \) in the leading order of \( \epsilon \) and \( \theta_L \) is reduced to
\[ g_Y^2 = -\frac{\epsilon(1 - 2\sin \theta_W^2)}{2\cos \theta_W \sin \theta_W} + \mathcal{O}(\epsilon^2) + \mathcal{O}(\theta_L), \]  
(23)
which is equal to that of the 2HD model. Consequently the constraints on the charged Higgs sector of the LR model from LEP data via the pair production of \( H^\pm \) is same in the leading order of \( \epsilon \) as those in the 2HD model. The conservative bound from LEP data is depicted in Fig. 1 by quoting in Ref. \[29\].

C. \( B \to \tau \nu \) Constraints

The purely leptonic \( B \to \tau \nu \) decay proceeds via annihilation of \( B \) meson into \( W \) boson in the SM and also into \( H^\pm \) in the LR model and 2HD model. The first measurement of the branching ratio of \( B \to \tau \nu \) decay has been performed by Belle \[4\]
\[ \text{Br}(B^- \to \tau \bar{\nu}_\tau) = (1.79^{+0.56}_{-0.49} +0.39_{-0.46}) \times 10^{-4}, \]  
(24)
which leads to the stringent constraints on the parameter set of the 2HD model via
\[ \frac{\text{Br}^{2HD}(B \to \tau \nu)}{\text{Br}^{SM}(B \to \tau \nu)} = \left( 1 - \frac{m_B^2}{m_H^2} \tan^2 \beta \right)^2 = 1.13 \pm 0.51, \]  
(25)
assuming \( f_B \) \[30\] and \( |V_{ub}| \) \[31\] are known.

Contributions of the LR model to \( B \to \tau \nu \) decay consist of the heavy \( W \) boson and the charged Higgs boson mediation diagrams as well as the SM contribution. The transition amplitude is given by
\[ \mathcal{M} = \mathcal{M}_W + \mathcal{M}_{W'} + \mathcal{M}_H, \]  
(26)
where
\[ \mathcal{M}_W = -\sqrt{2}G_Ff_Bm_Vab\cos \zeta (\cos \zeta + \sin \zeta) \left( \bar{u}(p_r)P_Lv(p_\nu) \right), \]
\[ \mathcal{M}_{W'} = -\sqrt{2}G_Ff_Bm_Vab\sin \zeta (\sin \zeta - \cos \zeta) \cdot \frac{1}{2} \epsilon^2(1 + \xi^2) \left( \bar{u}(p_r)P_Lv(p_\nu) \right), \]
\[ \mathcal{M}_H = -\sqrt{2}G_Ff_Bm_Vab\frac{m_B^2}{m_H^2} \frac{2\zeta}{1 + \xi^2} \left( 1 - \frac{1}{2} (1 + \xi^2)^2 \right) \left( \bar{u}(p_r)P_Lv(p_\nu) \right). \]  
(27)

Since \( \sin \zeta \sim \epsilon^2, W' \) contribution is suppressed by \( \epsilon^4 \) and ignored here. Although Yukawa couplings are proportional to \( 1/(1 - \xi^2) \), \( \mathcal{M}_H \) involves the combination of left-handed and right-handed Yukawa couplings given in Eq. (15), \( g_L - g_R \), which is proportional to \( (1 - \xi^2)^2 \) and cancel the divergent factor \( 1/(1 - \xi^2) \). The ratio of branching ratios is given by
\[ \frac{\text{Br}^{LR}(B \to \tau \nu)}{\text{Br}^{SM}(B \to \tau \nu)} = \left( 1 + \frac{m_B^2}{m_H^2} \frac{2\zeta}{1 + \xi^2} - \frac{m_W^2}{m_W^2} \frac{2\zeta}{1 + \xi^2} \right)^2. \]  
(28)
The new physics effects involve the factor \( \xi/(1 + \xi^2) \) and \( \xi/(1 + \xi^2) \) instead of \( \tan \beta \) and \( \cot \beta \). Since these factors are at most 1 with varying \( \xi \), the new physics effects are suppressed by the mass ratio \( m_B^2/m_H^2 \) and \( m_W^2/m_W^2 \), without enhancement factors. Consequently the LR model effects are at most a few percent and the recent Belle measurement of Eq. (24) could not provide any bounds on the LR model parameters.
IV. CONCLUDING REMARKS

If we observe a charged Higgs boson, it is a clear evidence for existence of new physics beyond the SM. The next step is to find the underlying physics for the Higgs sector. In this work, we examine the charged Higgs sector of the LR model with the present experiments; $H^\pm$ pair production at LEP, top quark decay into $bH^\pm$ at Tevatron, and $B$ annihilation decay mediated by $H^\pm$ at $B$-factory. Observables for each experiments depend upon different couplings of the charged Higgs from each others. In the 2HD model, the $t\to bH^\pm$ decay rate is governed by the $tbH^\pm$ Yukawa coupling involving $m_t$, which is proportional to $\tan \beta$ inversely, while the $B\to \tau\nu$ decay depends on the $ubH^\pm$ Yukawa coupling involving $m_b$, proportional to $\tan \beta$. On the other hand, the pair production of the charged Higgs bosons at $e^-e^+$ collision is related to the gauge couplings of $H^\pm$.

In the LR model, the Yukawa couplings involve the factors of $\xi/|1-\xi^2|$ or $(1+\xi^2)/|1-\xi^2|$ instead of $\tan \beta$ or $1/\tan \beta$, while the gauge couplings for the charged Higgs boson are identical to those of the 2HD model in the leading order of $\epsilon=k_1/v_R$. Thus the exclusion region of the parameter space on $(\xi, m_{H^\pm})$ in the LR model is much different from that on $(\tan \beta, m_{H^\pm})$ in the 2HD model regarding processes involving Yukawa couplings. Constraints from the $H^\pm$ pair production through the gauge couplings are similar in both the LR model and the 2HD model and weaker than the Tevatron bounds obtained in this work. The annihilation decay $B\to \tau\nu$ turns to provide no additional constraint on the LR model parameters because it involves the combination of left-handed and right-handed Yukawa couplings to quarks. In conclusion, we present the constraints on $(\xi, m_{H^\pm})$ space for the light charged Higgs boson in the LR model from the present experimental data.

Acknowledgments

This work was supported by the BK21 program of Ministry of Education (K.Y.L.).

[1] The LEP Collaborations ALEPH, DELPHI, L3 and OPAL, LEP working group for Higgs boson searches, [hep-ex/0107031].
[2] A. Abulencia et al., CDF Collaboration, Phys. Rev. Lett. 96, 042003 (2006).
[3] W.-S. Hou, Phys. Rev. D 48, 2342 (1993).
[4] T. Browder, Rare $B$ Decays with “Missing Energy”, talk delivered at the 33rd International Conference on High Energy Physics (ICHEP 06), Moscow, Russia, 26 Jul - 2 Aug 2006
[5] ATLAS Collaboration, Report No. CERN/LHCC/99-15 (1999) Vol. 2.
[6] CMS Collaboration, Report No. CERN/LHCC/94-38 (1994).
[7] M. Hashemi, [hep-ph/0612104].
[8] A. Belyaev, D. Garcia, J. Guasch, and J. Sola, JHEP 06, 050 (2002).
[9] A. C. Bawa, C. S. Kim, and A. D. Martin, Z. Phys. C 47, 75 (1990).
[10] V. D. Barger, R. J. N. Phillips, and D. P. Roy, Phys. Lett. B 324, 236 (1994); S. Moretti and K. Odagiri, Phys. Rev. D 55, 5627 (1997).
[11] D.-W. Jung, K. Y. Lee, and H. S. Song, Phys. Rev. D 70, 117701 (2004).
[12] C. Csaki, C. Grojean, L. Pilo, and J. Terning, Phys. Rev. Lett. 92, 101802 (2004); Y. Nomura, JHEP 0111, 050 (2003).
[13] Z. Chacko, H.-S. Goh, and R. Harnik, JHEP 0601, 108 (2006); H.-S. Goh and S. Su, [hep-ph/0608330].
[14] R. N. Mohapatra and J. C. Pati, Phys. Rev. D 11, 2558 (1975); J. C. Pati and A. Salam, Phys. Rev. D 10, 275 (1974); Erratum ibid. D 11, 703 (1975); For reviews, see R. N. Mohapatra, Unification and Supersymmetry (Springer, New York, 1986).
[15] M. Czagon, J. Gluza, and M. Zralek, Phys. Lett. B 458, 355 (1999).
[16] K. Cheung, Phys. Lett. B 517, 167 (2001).
[17] J. Chay, K. Y. Lee, and S.-h. Nam, Phys. Rev. D 61, 035002 (2000).
[18] J. Erler, and P. Langacker, Phys. Lett. B 456, 68 (1999).
[19] J. F. Gunion, J. Grifols, A. Mendez, B. Kayser, and F. Olness, Phys. Rev. D 40, 1546 (1989).
[20] N. G. Deshpande, J. F. Gunion, B. Kayser, and F. Olness, Phys. Rev. D 44, 837 (1991).
[21] K. Kiers, M. Assis, and A. A. Petrov, Phys. Rev. D 71, 115015 (2005).
[22] J. Donoghue and B. Holstein, Phys. Lett. B 113, 383 (1982).
[23] L. Wolfenstein, Phys. Rev. D 29, 2130 (1984).
[24] O. Khasanov and G. Perez, Phys. Rev. D 65, 053007 (2002).
[25] A. Ghinculov, and Y.-P. Yao, Mod. Phys. Lett. A 15, 925 (2000).
[26] R. Eusebi, Ph. D. thesis, University of Rochester, 2005.
[27] A. Djouadi, J. Kalinowski, and P. M. Zerwas, Z. Phys. C 57, 569 (1993).
[28] W.-M. Yao et al., J. of Phys. G: Nucl. Part. Phys. 33, 1 (2006).
[29] R. Barate et al., ALEPH Collaboration, Phys. Lett. B 543, 1 (2002); J. Abdallah et al., DELPHI Collaboration, Euro. Phys. J. C 34, 399 (2004); P. Achard et al., L3 Collaboration, Phys. Lett. B 575, 208 (2003).

[30] A. Grey et al., HPQCD Collaboration, Phys. Rev. Lett. 95, 212001 (2005).

[31] Heavy Flavor Averaging Group, http://www.slac.stanford.edu/xorg/hfag/