New mechanism for the top-bottom mass hierarchy

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We propose a mechanism to generate hierarchy between masses of the top and bottom quarks without fine tuning of the Yukawa coupling constants in the context of the two Higgs doublet model (THDM). In the THDM with a discrete symmetry, there exists the vacuum only where the top quark receives the mass of the order of the electroweak symmetry breaking scale v (≈ 246 GeV), while the bottom quark remains massless. By introducing a small soft-breaking parameter \( m_2^2 \) of the discrete symmetry, the bottom quark perturbatively acquires a nonzero mass. We show a model in which the small \( m_2^2 [\sim v^2/(4\pi)^2] \) is generated by the dynamics above the cutoff scale of the THDM. The ratio \( \tan \beta \) of the two vacuum expectation values is necessarily very large; i.e., \( \tan \beta \sim m_t/m_b \). We also find a salient relation, \( 1/\tan \beta \simeq m_2^2/m_H^2 \), where \( m_H \) is the mass of the extra CP-even Higgs boson. Our scenario yields some specific features that can be tested in future collider experiments.

PACS numbers: 12.60.-i, 12.60.Fr, 14.80.Cp

[ November 5, 2018 ]

I. Introduction

The measured quark mass spectrum shows a specific feature. Only the top quark has the mass of the order of the electroweak symmetry breaking (EWSB) scale \( v = (\sqrt{2} G_F)^{-1/2} \approx 246 \text{ GeV} \), while masses of the other quarks are much smaller. The top quark mass is 174 GeV (≈ \( v/\sqrt{2} \)), while the bottom quark, the second heaviest, has the mass of 4.2 GeV (≈ \( v \)). In the Standard Model (SM), however, the unique Higgs doublet field \( \Phi_{SM} \) is responsible for the EWSB and gives masses of all quarks via the Yukawa interactions; i.e., \( m_f \simeq y_f \langle \Phi_{SM} \rangle \) with \( \langle \Phi_{SM} \rangle = (0, v/\sqrt{2})^T \). Therefore, the observed mass spectrum is obtained only by assuming unnatural hierarchy among the Yukawa coupling constants \( y_f \). For instance, the hierarchy \( y_b/y_t \simeq 1/40 \) must be required for the top and bottom quarks. Nevertheless, no explanation for such fine tuning is given in the SM.

In this paper, we propose an alternative scenario in which the quark mass spectrum is reproduced without fine tuning in magnitude of the Yukawa coupling constants. We study the hierarchy between \( m_t \) and \( m_b \) under the assumption of \( y_t \approx y_b \approx O(1) \). In order to realize \( m_b/m_t \approx 1/40 \) in a natural way, we consider the two Higgs doublet model (THDM) with \( \Phi_1 \) and \( \Phi_2 \), imposing the discrete \( Z_2 \) symmetry\(^{2} \) under the transformation

\[
\Phi_1 \to -\Phi_1, \quad \Phi_2 \to +\Phi_2
\]

as well as

\[
\begin{pmatrix}
t \\ b
\end{pmatrix}_L \to \begin{pmatrix}
t \\ b
\end{pmatrix}_L, \quad t_R \to +t_R, \quad b_R \to -b_R.
\]

Due to the \( Z_2 \) symmetry, only \( \Phi_1 \) couples to the bottom quark while \( \Phi_2 \) does to the top quark. The hierarchy \( m_t \gg m_b \) is then equivalent to \( v_2 \gg v_1 \), where \( \langle \Phi_{1,2} \rangle = (0, v_{1,2}/\sqrt{2})^T \). We note that there exists the vacuum with \( v_1 = 0 \) and \( v_2 = v \) when the \( Z_2 \) symmetry is exact. A nonzero value of \( v_1 \) (≈ \( v_2 \)) is induced as a perturbation of a small soft-breaking parameter \( m_2^2 \) for the \( Z_2 \) symmetry. The small \( m_2^2 [\sim v^2/(4\pi)^2] \) is generated by the dynamics above the cutoff scale of the THDM. We find a salient relation, \( 1/\tan \beta \simeq m_2^2/m_H^2 \), where \( m_H \) is the mass of the extra CP-even Higgs boson. Consequently, we obtain \( m_b/m_t \ll 1 \). This scenario is extended to include the first two generation quarks.

We find that the extra Higgs bosons almost decouple with the weak gauge bosons in our model. Moreover, the extra Higgs bosons as well as the SM-like one turn out to have masses of the order of \( v \). The Higgs bosons with such masses are expected to be discovered at the CERN LHC because of the large value of \( \tan \beta \). The characteristics of our scenario can further be tested by precision measurement at future linear colliders (LC’s).\(^{4} \)

II. Minimal Model

The Lagrangian of the THDM with the softly-broken \( Z_2 \) symmetry is described as

\[
\mathcal{L} = \mathcal{L}_{\text{kin}} + \mathcal{L}_Y - V,
\]

where \( \mathcal{L}_{\text{kin}} \) and \( \mathcal{L}_Y \) are the kinetic and Yukawa interaction terms, respectively. The Higgs potential \( V \) is given by

\[
V = m_t^2 |\Phi_1|^2 + m_2^2 |\Phi_2|^2 - m_2^4 \Phi_1^3 + (\text{h.c.})
+ \lambda_1 |\Phi_1|^4 + \lambda_2 |\Phi_2|^4 + 2\lambda_3 |\Phi_1|^2 |\Phi_2|^2
+ 2\lambda_4 |\Phi_1|^2 |\Phi_2|^2 + \left[ \lambda_5 \left( \Phi_2 \Phi_1^* \right)^2 + (\text{h.c.}) \right],
\]

where \( m_t^2, m_2^2 \) and \( \lambda_1 \) to \( \lambda_5 \) are real, while \( m_2^2 \) and \( \lambda_5 \) are complex. The Higgs doublet fields \( \Phi_i \) (\( i = 1, 2 \)) with hypercharge \( Y = 1/2 \) are parameterized by

\[
\Phi_i = \left[ \frac{1}{\sqrt{2}} \left( v_i + h_i + ia_i \right) \right],
\]
where the vacuum expectation values (VEV’s) \( v_i \) \((i = 1, 2)\) satisfy \( v_1^2 + v_2^2 = v^2 \). The mass matrices for the Higgs bosons are diagonalized by mixing angles \( \alpha \) and \( \beta \). We then obtain five physical scalar states, \( h \) and \( H \) (CP-even), \( A \) (CP-odd), and \( H^\pm \) (charged), as well as three Nambu-Goldstone (NG) bosons, \( \phi^0 \) and \( \phi^\pm \).

We consider only the top and bottom quarks among fermions at first. We discuss the extension for the other quarks later on. In order to describe the assumption of \( y_t \approx y_b \), we introduce the global \( SU(2)_R \) symmetry, in addition to the \( SU(2)_L \) gauge symmetry:

\[
q_{L,R} \rightarrow q'_{L,R} = U_{L,R} q_{L,R}, \quad M_{21} \rightarrow M'_{21} = U_L M_{21} U_R^T ,
\]

where \( q_{L,R} \equiv (t_{L,R}, b_{L,R}) \) and \( U_{L,R} \in SU(2)_{L,R} \), respectively. The \( 2 \times 2 \) matrix \( M_{21} \) is defined by

\[
M_{21} \equiv (\Phi_2, \Phi_1) , \quad \text{with} \quad \Phi_2 = i\tau_2 \Phi_2. \tag{8}
\]

The \( Z_2 \) symmetry can be expressed in terms of \( q_{L,R} \) and \( M_{21} \) by

\[
q_L \rightarrow q'_L = q_L, \quad q_R \rightarrow q'_R = \tau_3 q_R, \quad M_{21} \rightarrow M'_{21} = M_{21} \tau_3.
\]

The Yukawa interaction then is written as

\[
\mathcal{L}_Y = -y q_L M_{21} q_R + (\text{h.c.}), \tag{11}
\]

with \( y \equiv y_t = y_b \). We also set

\[
\lambda_1 = \lambda_2 = \lambda_3 (\equiv \lambda) \tag{12}
\]

in Eq. \[12\] to realize the \( SU(2)_R \) symmetry in quartic interactions. The Higgs potential then is expressed by

\[
V(M_{21}) = \frac{1}{2} m^2 \text{tr}(M_{21}^2 M_{21}) - \frac{1}{2} \Delta_{12} \text{tr}(M_{21}^2 M_{21} \tau_3) - [m_2^2 \det M_{21} + (\text{h.c.})] + \lambda \left[ \text{tr}(M_{21}^2 M_{21}) \right]^2 + 2\lambda_4 \det(M_{21}^2 M_{21}) + [\lambda_5 (\det M_{21})^2 + (\text{h.c.})], \tag{13}
\]

where \( m^2 = m_1^2 + m_3^2 \) and \( \Delta_{12} = m_1^2 - m_2^2 \). The \( Z_2 \) symmetry is softly broken by the mass term of \( m_3^2 \). A non-zero value of \( \Delta_{12} \) measures the soft breaking of the global \( SU(2)_R \) symmetry. In order to evade explicit CP violation, we choose the phases in \( m_3^2 \) and \( \lambda_5 \) to be zero.

We have introduced the global \( SU(2)_R \) symmetry only for the description of \( y_t = y_b \) in terms of a symmetry. The Higgs potential also becomes simple since this symmetry requires the relation \[12\]. Our main results, however, turn out to be unchanged even when this relation is relaxed to some extent. Cases without \( SU(2)_R \) as well as those with CP violation will be discussed in details elsewhere [8].

Let us consider the effective potential \( V(\langle M_{21} \rangle) \) to study the vacuum structure. By using \( SU(2)_L \) and \( U(1)_Y \), the VEV’s in the THDM can be generally parameterized as

\[
\langle M_{21} \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_2 & v_E \\ v_1 & iv_A \end{pmatrix} . \tag{14}
\]

Spontaneous breakdown of \( U(1)_\text{EM} \) and the CP symmetry occurs if \( v_E \neq 0 \) and \( v_1 v_A \neq 0 \), respectively. We can easily show that the spontaneous \( U(1)_\text{EM} \) breaking cannot occur at the tree level in our model. The conditions for CP conservation are studied in Ref. [9]. The effective potential is bounded from below by the requirement of the vacuum stability [10], which leads to

\[
\lambda > 0, \quad 2\lambda + \lambda_4 - |\lambda_5| > 0. \tag{15}
\]

We investigate details of the vacuum structure of our model in the tree level approximation. We first study the case with \( m_3^2 = 0 \) where the discrete \( Z_2 \) symmetry is exact. We next include effects of \( m_3^2 \neq 0 \).

For \( m_3^2 = 0 \), the effective potential \( V(\langle M_{21} \rangle) \) is given by

\[
V(\langle M_{21} \rangle) = \frac{m_1^2}{2} (v_1^2 + v_A^2) + m_2^2 v_2^2 + \frac{\lambda}{4} (v_1^2 + v_A^2 + v_2^2) + \frac{\lambda_4}{2} (v_1^2 + v_A^2) v_2^2 + \frac{\lambda_5}{2} (v_1^2 - v_A^2) v_2^2 , \tag{16}
\]

where we used Eq. \[14\] with \( v_E = 0 \). The VEV’s, \( v_1 \), \( v_2 \), and \( v_A \), are determined by the stationary conditions \( \partial V(\langle M_{21} \rangle) / \partial v_i = 0 \), \((i = 1, 2, A)\). Since spontaneous CP violation does not occur for \( m_3^2 = 0 \), three types of the nontrivial vacuum are possible [10]:

- \( v_1 = v_A = 0, v_2 \neq 0 \),
- \( v_1 v_2 \neq 0, v_A = 0 \),
- \( v_A v_2 \neq 0, v_1 = 0 \).

In Fig. [1], the area (I) corresponds to the vacuum (a), while the areas (II) and (III) do to the vacua (b) and (c), respectively. Due to the vacuum stability conditions [10], there does not exist the stable vacuum out of the three areas. Performing the transformation \( \Phi_1 \to e^{i\pi/2} \Phi_1 \) to the nontrivial vacuum (b), we obtain the vacuum (c). The transformation corresponds to \( \lambda_5 \rightarrow -\lambda_5 \) in the Higgs potential with \( m_3^2 = 0 \). The area (III) is thus the mirror image of the area (II).

In order to realize \( m_b/m_t \ll 1 \) without fine tuning of Yukawa couplings, we choose the vacuum (a) which leads to

\[
m_t = \frac{1}{\sqrt{2}} y v, \quad m_b = 0 , \tag{17}
\]

because of \( v_2 = v \). Although the bottom quark may receive a small mass even in the vacuum (b), the parameters of the Higgs potential must be very close to the boundary between the areas of (I) and (II). This is fine
tuning in a sense, so that we avoid such a case. The vacuum (a) for $m_2^2 = 0$ is realized when\(^1\)

\[
m_2^2 < -|m_1^2|, \quad -\frac{\lambda_1}{\lambda} - \frac{\Delta_{12}}{-m_2^2} < \frac{\lambda_5}{\lambda} + \frac{\Delta_{12}}{-m_2^2}, \quad (18)
\]

or

\[
m_1^2 \geq -m_2^2 > 0. \quad (19)
\]

Only the doublet $\Phi_2$ is responsible for the EWSB in the vacuum (a). The doublet fields $\Phi_1$ and $\Phi_3$ do not mix for $m_2^2 = 0$ because of the remaining $Z_2$ symmetry after the EWSB, $\Phi_1 \rightarrow -\Phi_1$. The mass formulae of the physical Higgs bosons are

\[
\begin{align*}
m_h^2 &= 2\lambda v^2, \\
m_{H^\pm}^2 &= \Delta_{12}, \\
m_H^2 &= \Delta_{12} + (\lambda_4 + \lambda_5)v^2, \\
m_A^2 &= \Delta_{12} + (\lambda_4 - \lambda_5)v^2.
\end{align*} \quad (20 - 23)
\]

When $\Delta_{12} = 0$, the charged Higgs bosons become the extra NG bosons associated with the breaking of the exact $SU(2)_R$ symmetry.

We now switch on a small soft-breaking parameter $m_3^2(\ll v^2)$ of the discrete $Z_2$ symmetry. We do not consider the possibility of spontaneous CP violation.\(^2\) A nonzero $v_1$ is necessarily induced for $m_3^2 \neq 0$ from the stationary condition. As a perturbation from the vacuum (a) with $m_2^2 = 0$, we consequently obtain

\[
\frac{v_1}{v_2} = \frac{1}{\tan \beta} = \frac{m_2^2}{m_H^2} \left\{ 1 + O \left( \frac{m_3^2}{v^4} \right) \right\}, \quad (24)
\]

where we used the tree-level mass formula in Eq. (22). Because of $v_1^2 + v_2^2 = v^2$, the expression for $v_2$ is slightly modified to $v_2 = v[1 - O(m_3^2/v^4)]$ from $v_2 = v$. The masses of the top and bottom quarks are given by

\[
m_t \simeq \frac{1}{\sqrt{2}} y v, \quad m_b = \frac{1}{\sqrt{2}} y v_1, \quad (25)
\]

so that the bottom quark finally obtains the small mass. The mass hierarchy of $m_t$ and $m_b$ then is deduced from Eqs. (24) and (25) without fine tuning of the Yukawa coupling constants; i.e., $m_t/m_b = \tan \beta$. With nonzero $m_3^2$ the Higgs doublets $\Phi_1$ and $\Phi_2$ do mix. The mixing angle $\beta - \alpha$ is expressed as

\[
\sin(\beta - \alpha) = 1 - \left( \frac{m_H^2 - m_{H^\pm}^2}{m_H^2 - m_A^2} \right)^2 \frac{2}{\tan^2 \beta} + O \left( \frac{m_3^2}{\lambda v^4} \right), \quad (26)
\]

where Eqs. (20 - 23) and $\tan \beta \gg 1$ are used. From Eqs. (20) and (26), the property of the CP-even Higgs $h$ is similar to the SM one. We note that Higgs boson masses in Eqs. (20 - 23) receive corrections of $O(m_3^2/v^4)$. These corrections, however, do not affect the expressions in Eqs. (24) and (26).

Let us estimate the typical size of the masses of the extra Higgs bosons. The value of $\tan \beta$ is fixed by $\tan \beta = m_t/m_b \sim 40$. On the other hand, the small value of $m_3^2(\ll v^2)$ can be interpreted as $m_3^2 \simeq v^2/(4\pi)^2$. In the next section, we shall present a concrete model in which such a small $m_3^2(\simeq v^2/(4\pi)^2)$ is radiatively induced by the dynamics above the cutoff scale of the THDM. From Eq. (24), the mass of $H$ is expressed as

\[
m_H^2 \simeq m_3^2 \tan \beta. \quad (27)
\]

Therefore, the size of $m_H$ is at most of the order of $v$. Furthermore, the masses of $A$ and $H^\pm$ are also the same order because of the relations

\[
m_{H^\pm}^2 = m_H^2 - (\lambda_4 + \lambda_5)v^2, \quad m_A^2 = m_H^2 - 2\lambda_5v^2, \quad (28)
\]

which are obtained from Eqs. (21 - 28).

We have found Eqs. (21) and (28), assuming the softly broken $SU(2)_R$ symmetry; i.e., $\lambda_1 = \lambda_2 = \lambda_3(= \lambda)$. We now give comments on the case with $\lambda_1 \neq \lambda_2 \neq \lambda_3$, relaxing the $SU(2)_R$ symmetry. First, it can be shown that Eq. (24) does not change. Second, although Eq. (26) is slightly modified, the essential result of $\sin(\beta - \alpha) = 1 - O(\tan^{-2}\beta)$ still holds. Finally, the masses of the extra Higgs bosons remain $O(v)$ even for $\lambda_1 \neq \lambda_2 \neq \lambda_3$, because Eq. (25) turns out to be unchanged as well.

III. A MECHANISM FOR SMALL $m_3^2$

We discuss an example where the small $m_3^2$ is generated radiatively in the low energy scale. Let us consider a model with a complex scalar field $S$ which is a $SU(2)_L$ singlet without $U(1)_{Y'}$ charge. The Lagrangian is given by

\[
\mathcal{L} = \mathcal{L}_{kin} - V_{\Phi} - V_S - V_{\Phi S}, \quad (29)
\]
where $L_{\text{kin}}$ represents the kinetic term and $V_\Phi$ is the $Z_2$ symmetric part of the THDM potential [11] with $m_3^2 = 0$. The potential $V_S$ for the complex scalar $S$ and the interaction term $V_{Z_2}$ between $S$ and $\Phi_{1,2}$ are given by

$$V_S = M_3^2 S^\dagger S + \kappa (S^\dagger S)^2 + V_{Z_2}, \quad M_3^2 > 0, \quad (30)$$

with

$$V_{Z_2} = \frac{\eta}{\Lambda^{2\ell-2}} (S^{2n} + \text{h.c.}), \quad \eta \sim O(1), \quad (31)$$

and

$$V_{Z_2} = \frac{\xi}{\Lambda^{2\ell-2}} (S^{2n} + \text{h.c.}), \quad \xi \sim O(1), \quad (32)$$

respectively. In Eqs. (31) and (32), $\Lambda$ denotes the cutoff scale of the model. We now set $n = 1$ (case A) or $n = \ell$ (case B) with $\ell \geq 1$. We note that $V_S$ has the $Z_{2n}$ symmetry under $S \rightarrow e^{i\pi} S$, while $V_\Phi$ is $Z_2$ invariant under the transformation $\Phi_1 \rightarrow -\Phi_1$, $\Phi_2 \rightarrow +\Phi_2$. The interaction term $V_{Z_2}$ explicitly breaks both $Z_{2n}$ and $Z_2$. Some invariant terms under $Z_{2n}$ and $Z_2$ are not explicitly included here, as they are irrelevant to our conclusion.

Supposing that $M_S(\sim \Lambda)$ is much larger than the EWSB scale, we integrate out the field $S$ and thereby obtain the THDM with the softly-broken $Z_2$ symmetry ($m_3^2 \neq 0$) as the low-energy effective theory. From the Feynman diagrams depicted in Fig. 2 we estimate

$$m_3^2 \sim \xi \eta \frac{\Lambda}{(4\pi)^{2\ell}} M_S^2, \quad \text{for Case A}, \quad (33)$$

$$m_3^2 \sim \xi \eta \frac{\Lambda}{(4\pi)^{2(\ell-1)}} M_S^2, \quad \text{for Case B}. \quad (34)$$

For example, we can obtain $m_3^2 \sim v^2/(4\pi)^2$ for $\ell = 2$, if we take the cutoff $M_S = 4\pi v$ for Case A or $M_S = (4\pi)^2 v$ for Case B. For $\ell = 2$, we do not need higher dimensional operators except for the $Z_2$ breaking term $V_{Z_2}$.

We may consider other possibilities to obtain small $m_3^2$ values based on many ideas such as Topcolor instanton [11] and large extra dimensions [12]. Also useful is a model which provides effectively $(\Phi_1^2 \Phi_2^2)$ with a coefficient $\sim O(1)$ while prohibits the hard breaking terms of the $Z_2$ symmetry such as $(\Phi_1^2 \Phi_1^2)$.

IV. QUARK MASS MATRICES

We discuss the extension of our model incorporating first two generation quarks. Can we reproduce the observed quark mass spectrum and the Kobayashi-Maskawa (KM) matrix?

Under the discrete symmetry [2], two types of Yukawa interactions are possible in the THDM, so called Model I and Model II [3]. The flavor changing neutral current (FCNC) then does not appear at the tree level [2]. Obviously Model I is inconsistent with our scenario, so that we here study Model II,

$$-\mathcal{L} = \sum_{i,j=1}^{3} \left( Y_{D}^{ij} q_{L}^{i} \Phi_{1} D_{R}^{j} + Y_{U}^{ij} q_{L}^{i} \Phi_{2} U_{R}^{j} \right) + \text{h.c.}, \quad (35)$$

where $q_{L}^{i}$ is the left-handed quark doublet of the $i$th generation, and $D_{R}^{j} = (d_{R}, s_{R}, b_{R})^T$ and $U_{R}^{j} = (u_{R}, c_{R}, t_{R})^T$. We then assume that matrices of the Yukawa coupling take the following forms,

$$Y_{U}^{ij} \sim Y_{D}^{ij} \sim y \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}, \quad y \sim O(1), \quad (36)$$

which lead to $m_t \gg m_c, m_u$, and $m_b \gg m_s, m_d$, and the KM matrix becomes approximately diagonal. We can numerically reproduce the data for the mass spectrum and the KM matrix [4], allowing fluctuations of the Yukawa coupling constants,

$$Y_{U}^{ij} = y \epsilon_{ij}^U, \quad Y_{D}^{ij} = y \epsilon_{ij}^D, \quad \text{with} \quad 0.5 < |\epsilon_{ij}^{U,D}| < 1.5. \quad (37)$$

Three comments are in order: (a) Although we can avoid hierarchy among Yukawa couplings, subtle cancelation among the $O(1)$ mass-matrix elements is required to obtain masses of light quarks. (b) We may adopt Model III [13] to our scenario, if the FCNC is suppressed by some mechanism. (c) It is possible to apply our scenario to the lepton sector. The $\tau$ lepton then receives the small mass due to the similar mechanism to the bottom quark. At the same time, however, the Dirac mass of the tau neutrino could be produced around $m_t$. To explain the tiny (Majorana) mass of the tau neutrino, additional mechanism such as the Seesaw [14] might be helpful.

V. SUMMARY AND DISCUSSIONS

We have proposed the mechanism to explain the mass hierarchy between the top and bottom quarks without
fine tuning, starting from the vacuum with \((v_1, v_2) = (0, v)\). Such a vacuum can exist when the \(Z_2\) symmetry is exact. The observed mass spectrum \(m_t \gg m_b \neq 0\) is realized via the small soft-breaking parameter \(m_3^2\) for the \(Z_2\) symmetry. We have presented the model in which a small \(m_3^2\) is induced from the underlying physics above the cutoff scale of the THDM.

The phenomenological implication is as follows. The size of \(\tan \beta\) corresponds to the ratio \(m_t/m_b \sim 40\). We have found the relation \(m_{H^\pm} \approx m_3^2 \tan \beta \sim O(v^2)\). Therefore, the masses of the extra Higgs bosons \(H, A\) and \(H^\pm\) are expected to be \(O(v)\). The THDM with such parameters is constrained by the theoretical considerations \([15, 16]\) as well as the available data. When \(m_{H^\pm} \approx m_H\), or \(m_{H^\pm} \approx m_A\), our model can satisfy the constraint from the LEP precision data \([7]\). The mass of the charged Higgs boson in our scenario may not conflict with the \(b \to s\gamma\) result \([17]\). The doublet \(\Phi_2\) is mainly responsible for the EWSB, so that we obtain \(\sin(\beta - \alpha) \approx 1\) in a good approximation.

In addition to the SM-like Higgs boson \(h\), all the extra Higgs bosons in our model are expected to be discovered at the LHC. Our prediction of \(\sin(\beta - \alpha) \approx 1\) can also be confirmed at the LHC and LC’s. Our scenario may further be tested by measuring the \(hhh\) coupling at future LC’s \([18]\). More detailed phenomenological analysis will be done elsewhere \([8]\).

Acknowledgements

The authors thank Yasuhiro Okada for useful comments.

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