Electron/Muon Specific Two Higgs Doublet Model

Yuji Kajiyama,1 Hiroshi Okada,2 and Kei Yagyu3

1 Akita Highschool, Tegata-Nakadai 1, Akita, 010-0851, Japan
2 School of Physics, KIAS, Seoul 130-722, Korea
3 Department of Physics, National Central University, Chungli, Taiwan 32001, ROC

We discuss two Higgs doublet models with a softly-broken discrete $S_3$ symmetry, where the mass matrix for charged-leptons is predicted as the diagonal form in the weak eigenbasis of lepton fields. Similar to an introduction of $Z_2$ symmetry, the tree level flavor changing neutral current can be forbidden by imposing the $S_3$ symmetry to the model. Under the $S_3$ symmetry, there are four types of Yukawa interactions depending on the $S_3$ charge assignment to right-handed fermions. We find that extra Higgs bosons can be muon and electron specific in one of four types of the Yukawa interaction. This property does not appear in any other two Higgs doublet models with a softly-broken $Z_2$ symmetry. We discuss the phenomenology of the muon and electron specific Higgs bosons at the Large Hadron Collider; namely we evaluate allowed parameter regions from the current Higgs boson search data and discovery potential of such a Higgs boson at the 14 TeV run.

PACS numbers:
1. INTRODUCTION

A Higgs boson has been discovered at the CERN Large Hadron Collider (LHC) \cite{1, 2}, whose properties; e.g., mass, spin, CP and observed number of events are consistent with those of the Higgs boson predicted in the Standard Model (SM). The SM-like Higgs boson also appears in Higgs sectors extended from the SM one, so that there are still various possibilities non-minimal Higgs sectors. They are often introduced in models beyond the SM which have been considered to explain problems unsolved within the SM such as the neutrino oscillation, dark matter (DM) and baryon asymmetry of the Universe.

In addition to the above problems, one of the deepest mystery in the SM is the flavor structure. In the SM, all the masses of charged fermions are accommodated by the vacuum expectation value (VEV) of the Higgs doublet field through Yukawa interactions. However, there are redundant number of parameters to obtain physical observables; i.e., the Yukawa couplings are given by general $3 \times 3$ complex matrices (totally 18 degrees of freedom) for each up-type and down-type quarks and charged-leptons. In fact, only three independent parameters are enough in the charged-leptons sector to describe the masses of $e$, $\mu$ and $\tau$. In order to constrain the structure of Yukawa interactions, Non-Abelian discrete symmetries have been introduced such as based on the $S_3$ \cite{3, 4} and $A_4$ \cite{5} groups. Usually, in a model with such a discrete symmetry, the Higgs sector is extended to be the multi-doublet structure. Therefore, phenomenological studies for the extended Higgs sector with multi-doublet structure are important to probe such a model.

In this paper, we discuss two Higgs doublet models (THDMs) with the $S_3$ symmetry as the simplest realization of the diagonalized mass matrix for the charged-leptons without introducing any unitary matrices. This can be achieved by assigning the first and second generation lepton fields to be the $S_3$ doublet$^1$.

In general, there appears the flavor changing neutral current (FCNC) via a neutral Higgs boson mediation at the tree level in two Higgs doublet models (THDMs), which is strictly constrained by flavor experiments. Usually, such a tree level FCNC is forbidden by introducing a discrete $Z_2$ symmetry \cite{6} to realize the situation where one of two Higgs doublet fields couples to each fermion. In our model, this situation is realized in terms of the $S_3$ flavor symmetry. The Yukawa interaction among the Higgs doublet fields and fermions can be classified into four types depending on the

\footnote{1 Our $S_3$ charge assignments for the quarks and Higgs doublet fields are different from those in the previous studies for $S_3$ models \cite{3}. Usually, all the quarks, leptons and Higgs doublet fields are embedded in the $S_3$ doublet plus singlet. However, we treat that the quark sector is the same as in the SM assuming the quark fields to be the singlet, because it is suitable and economical to explain the observed SM-like Higgs boson at the LHC.}
\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|c|}
\hline
Particle & $Q_i$ & $L_u$ & $L_\tau$ & $u_{iR}$ & $d_{iR}$ & $e_{uR}$ & $\tau_R$ & $\Phi_1$ & $\Phi_2$ \\
\hline
$SU(2)_L \times U(1)_Y$ & $\frac{1}{2}, \frac{1}{6}$ & $2, -\frac{1}{2}$ & $2, -\frac{1}{2}$ & $1, \frac{2}{3}$ & $1, -\frac{1}{3}$ & $1, -1$ & $1, -1$ & $2, \frac{1}{2}$ & $2, \frac{1}{2}$ \\
$S_3$ & 1 & 2 & 1 & 1' & 1 or 1' & 2 & 1 or 1' & 1 & 1' \\
\hline
\end{tabular}
\caption{The particle contents and their charge assignment of the $SU(2)_L \times U(1)_Y \times S_3$ symmetry.}
\end{table}

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|c|c|c|}
\hline
Particle & $u_{iR}$ & $d_{iR}$ & $\tau_R$ & $\xi_u$ & $\xi_d$ & $\xi_\tau$ \\
\hline
Type-I & 1' & 1' & 1' & cot $\beta$ & cot $\beta$ & cot $\beta$ \\
Type-II & 1' & 1 & 1 & cot $\beta$ & $-\tan \beta$ & $-\tan \beta$ \\
Type-X & 1' & 1' & 1 & cot $\beta$ & cot $\beta$ & $-\tan \beta$ \\
Type-Y & 1' & 1 & 1' & cot $\beta$ & $-\tan \beta$ & cot $\beta$ \\
\hline
\end{tabular}
\caption{Four patterns of the assignment of $S_3$ charges to the right-handed fermions, and $\xi_f$ factors appearing in Eq. (II.6).}
\end{table}

$S_3$ charge assignments to the right-handed fermions. Similar classification has been defined in THDMs with a softly-broken $Z_2$ symmetry \cite{7, 8}. A comprehensive review for the THDMs with the softly-broken $Z_2$ symmetry has been given in Ref. \cite{9}.

We find that extra neutral and charged Higgs bosons can be muon and electron specific; namely, they can mainly decay into $\mu^+ \mu^-$ or $e^+ e^-$ and $\mu^\pm \nu$ or $e^\pm \nu$, respectively, in one of four types of the Yukawa interaction. This phenomena cannot be seen in any other THDMs without the tree level FCNC such as the softly-broken $Z_2$ symmetric version. We show excluded parameter regions from the current LHC data in this scenario. We then evaluate discovery potential of signal events from these extra Higgs bosons at the LHC with the collision energy to be 14 TeV.

This paper is organized as follows. In Sec. II, we define the particle content and give the Lagrangian in our model. The mass matrices for the charged-leptons and neutrinos are then calculated. The Higgs boson interactions are also discussed in this section. In Sec. III, we discuss the collider phenomenology, especially focusing on the muon and electron specific Higgs bosons in the Type-I $S_3$ model. We give a summary and conclusion of this paper in Sec. IV. The explicit form of the Higgs potential and mass matrices for the scalar fields are given in Appendix A.
II. THE MODEL

A. Charge assignment

We discuss the THDM with the softly-broken discrete $S_3$ symmetry. In the $S_3$ group, there are the following irreducible representations; two singlets $1$ (true-singlet) and $1'$ (pseudo-singlet) and doublet $2$ (see Ref. [10]). Particle contents are shown in Table I. The $i$-th generation of left-handed quarks $Q_i$ are assigned to be $S_3$ true-singlet, while the right-handed up type quarks $u_{iR}$ and down type quarks $d_{iR}$ are assigned to be $S_3$ true- or pseudo-singlet. The left- (right-) handed electron and muon $L_a$ ($e_{aR}$) are embedded as the doublet representation of the $S_3$ symmetry. Both left-handed and right-handed tau leptons $L_\tau$ and $\tau_R$, respectively, are singlets under $S_3$. The isospin doublet Higgs fields $\Phi_1$ and $\Phi_2$ are transformed as $S_3$ true- or pseudo-singlet.

We can define four independent patterns of the charge assignment for $u_{iR}$, $d_{iR}$ and $\tau_R$ in the $S_3$ symmetric THDMs. We call them as Type-I, Type-II, Type-X and Type-Y $S_3$ models, and the $S_3$ charge assignment in each model is listed in Table II. This charge assignment is the analogy of that of a softly-broken $Z_2$ symmetry in the THDMs [13].

B. Higgs Potential

The softly-broken $S_3$ symmetric Higgs potential is given as

$$V = m_1^2 \Phi_1^\dagger \Phi_1 + m_2^2 \Phi_2^\dagger \Phi_2 + \left[ m_3^2 \Phi_1^\dagger \Phi_2 + \text{h.c.} \right]$$

$$+ \frac{1}{2} \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \frac{1}{2} \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1)(\Phi_2^\dagger \Phi_2) + \lambda_4 |\Phi_1^\dagger \Phi_2|^2 + \frac{1}{2} \left[ \lambda_5 (\Phi_1^\dagger \Phi_2)^2 + \text{h.c.} \right],$$

where the doublet Higgs fields can be parameterized as

$$\Phi_\alpha = \begin{pmatrix} w_\alpha^+ \\ \frac{1}{\sqrt{2}}(h_\alpha + v_\alpha + iz_\alpha) \end{pmatrix}, \quad \alpha = 1, 2,$$

where $v_\alpha$ are the VEVs of the doublet Higgs fields, and they satisfy $v^2 \equiv v_1^2 + v_2^2 = 1/\sqrt{2}G_F = (246 \text{ GeV})^2$. The ratio of the two VEVs can be parameterized by $\tan \beta \equiv v_2/v_1$ as usual in THDMs. Although among the parameters in the potential, $m_3^2$ and $\lambda_5$ are complex in general, we assume the CP-conservation in the Higgs potential for simplicity. We note that we can retain the $Z_2$ symmetry as the subgroup of $S_3$ by taking $m_3^2 = 0$. However, the potential without the $m_3^2$ term

---

2 The Type-X and Type-Y THDMs are respectively referred as the lepton-specific [11] and flipped [12] THDMs.
results non-decoupling theory; namely, all the masses of Higgs bosons are determined by the Higgs VEV times λ couplings. In the following, we consider the case with \( m_3^2 \neq 0 \).

The mass eigenstates for the CP-odd, singly-charged and CP-even Higgs bosons from the doublet fields are given by the following orthogonal matrices as

\[
\begin{pmatrix}
z_1 \\
z_2
\end{pmatrix}
= \begin{pmatrix}
c_\beta & -s_\beta \\
s_\beta & c_\beta
\end{pmatrix}
\begin{pmatrix}
G^0 \\
A
\end{pmatrix},
\begin{pmatrix}
w_1^+ \\
w_2^+
\end{pmatrix}
= \begin{pmatrix}
c_\beta & -s_\beta \\
s_\beta & c_\beta
\end{pmatrix}
\begin{pmatrix}
G^+ \\
H^+
\end{pmatrix},
\begin{pmatrix}
h_1 \\
h_2
\end{pmatrix}
= \begin{pmatrix}
c_\alpha & -s_\alpha \\
s_\alpha & c_\alpha
\end{pmatrix}
\begin{pmatrix}
H \\
h
\end{pmatrix},
\]

where \( G^\pm \) and \( G^0 \) are the Nambu-Goldstone bosons which are absorbed by the longitudinal component of \( W^\pm \) and \( Z \). Because the potential given in Eq. (II.1) is the completely same form as in the softly-broken \( Z_2 \) symmetric THDMs, the mass formulae are also the same form. The detailed formulae for the masses of the physical Higgs bosons can be seen in Ref. [14], for example.

C. Yukawa Lagrangian

The renormalizable Yukawa Lagrangian under the \( S_3 \) invariance is given by

\[
-\mathcal{L}_Y = y_1^e (\bar{L}_1 e_2 R + \bar{L}_2 e_1 R) \Phi_1 + y_2^e (\bar{L}_1 e_2 R - \bar{L}_2 e_1 R) \Phi_2 + \text{h.c.}
+ y_1^d Q_i (i \tau_2 \Phi_u^* u_j R + y_2^d Q_i \Phi_d d_j R + y_2^d \bar{L}_\tau \Phi_\tau R + \text{h.c.}),
\]

where \( \Phi_{u,d,\tau} \) are \( \Phi_1 \) or \( \Phi_2 \) depending on the \( S_3 \) charge assignment of \( u_{iR}, d_{iR} \) and \( \tau_R \) as listed in Table II.

The charged-lepton mass matrix defined by \((\bar{e}_L, \bar{\mu}_L, \bar{\tau}_L) M_\ell (e_R, \mu_R, \tau_R)^T\), under the identifications of the lepton fields as \( L_1 = L_e, L_2 = L_\mu, e_1 R = \mu_R, e_2 R = e_R \), can be obtained in the diagonal form by

\[
M_\ell = \frac{1}{\sqrt{2}} \text{diag}(y_1^e v_1 + y_2^e v_2, y_1^e v_1 - y_2^e v_2, y^\tau v_\tau),
\]

where \( v_\tau \) is either \( v_1 \) or \( v_2 \).

The quarks masses and mixings are obtained as the same way in the SM. As already mentioned in the Introduction, this treatment is different from that in the previous \( S_3 \) models in which the part of Yukawa Lagrangian is given by the \( S_3 \) singlet from \( 2 \times 2 \times 2 \), where each \( 2 \) denotes the left-handed quark, right-handed quark and Higgs doublet fields. In such a model, there are predictions in the quark sector such as the Cabibbo mixing angle. In our model, we choose singlet
representations for all the quark fields and Higgs doublet fields, so that there is no such a prediction. However, by this assignment, the minimal content for the Higgs sector; i.e., two Higgs doublet fields can be realized within the framework of $S_3$ with the diagonalized charged-lepton mass matrix and the SM-like Higgs boson which is necessary to explain the observed Higgs boson at the LHC as will be discussed in the next subsection.

The Yukawa interactions are given in the mass eigenbasis for the physical Higgs bosons as

$$-L_Y^{\text{int}} = \frac{m_\mu}{v} \left\{ -\frac{1}{2}(\tan \beta + \cot \beta)c_{\beta-\alpha}\bar{e}eh + \left[ s_{\beta-\alpha} - \frac{1}{2}(\tan \beta - \cot \beta)c_{\beta-\alpha} \right] \bar{\mu}h ight. \\
+ \frac{1}{2}(\tan \beta + \cot \beta)s_{\beta-\alpha}\bar{e}eH + \left[ c_{\beta-\alpha} + \frac{1}{2}(\tan \beta - \cot \beta)s_{\beta-\alpha} \right] \bar{\mu}H \\
- \frac{i}{2}(\tan \beta + \cot \beta)\bar{e}\gamma_5eA - \frac{i}{2}(\tan \beta - \cot \beta)\bar{\mu}\gamma_5A \\
- \frac{1}{\sqrt{2}} \left[ (\tan \beta + \cot \beta)\bar{\nu}_P \bar{P}R eH^+ + (\tan \beta - \cot \beta)\bar{\nu}_\mu \bar{P}R H + h.c. \right] \right\} \\
\sum_{f=u,d,\tau} \frac{m_f}{v} \left[ (s_{\beta-\alpha} + \xi_f c_{\beta-\alpha})\bar{f}h + (c_{\beta-\alpha} - \xi_f s_{\beta-\alpha})\bar{f}H - 2iI_f \xi_f \bar{f}\gamma_5fA \right] \\
+ \left[ \frac{\sqrt{2}V_{ud}}{v} \bar{m}_d \xi_d P_R - m_u \xi_u P_L \right] d H^+ + \frac{\sqrt{2}m_\tau \xi_\tau}{v} \bar{P}_R \tau H + h.c. \right].$$

where the electron mass is neglected in the above expression, and $I_f = +1/2$ ($-1/2$) for $f = u$ ($d, \tau$). The $\xi_f$ factors are listed in Table III.

The $hVV$ and $HVV$ ($V = W^\pm, Z$) coupling constants are given by $\sin(\beta - \alpha) \times g_{hVV}^{\text{SM}}$ and $\cos(\beta - \alpha) \times g_{hVV}^{\text{SM}}$ with $g_{hVV}^{\text{SM}}$ being the coupling constant of the SM Higgs boson and gauge bosons. Thus, when we take the limit of $\sin(\beta - \alpha) = 1$, $h$ has the same coupling constants with the gauge bosons and fermions (see Eq. (II.6)) as those in the SM Higgs boson.

We here comment on the new contributions to the muon anomalous magnetic moment ($g - 2$) from the additional scalar boson loops. In our model in the case of $\sin(\beta - \alpha) = 1$, the $H$, $A$ and $H^\pm$ loop contributions are calculated by using the formula given in Ref. 15 as

$$\Delta a_\mu = \frac{1}{32\pi^2} \frac{m_\mu^2}{v^2} \frac{g^2}{16\pi^2} \frac{\sqrt{2}m_\mu}{v} \left\{ F_1(m_H^2/m_\mu^2) + F_2(m_H^2/m_\mu^2) \right\}$$

where

$$F_1(x) = \frac{1 - 4x + 3x^2 - 2x^2 \ln x}{2(1 - x)^3}, \quad F_2(x) = \frac{(1 - x)(2x^2 + 5x - 1) + 6x^2 \ln x}{6(1 - x)^4}.$$
give destructive contributions to the $H$ loop contribution. On the other hand, the discrepancy of the measured muon $g - 2$ from the SM prediction is roughly given as $3 \times 10^{-9}$ [17, 18] which is two orders of magnitude larger than the above result with $m_H = 150$ GeV and $\tan \beta = 100$. Therefore, it is difficult to compensate the discrepancy by the additional scalar boson loop contributions in our model similar to the Type-II THDM.

III. PHENOMENOLOGY AT THE LHC

In this section, we discuss the phenomenology of the Higgs bosons at the LHC. We consider the case with $\sin(\beta - \alpha) = 1$ in which $h$ can be regarded as the SM-like Higgs boson with the mass of 126 GeV, because the current Higgs boson search data at the LHC suggest that the observed Higgs boson is consistent with the SM Higgs boson. We then focus on collider signatures from the extra Higgs bosons; i.e., $H$, $A$ and $H^\pm$ at the LHC.

A. The $\mu$ and $e$ specific Higgs bosons

In all the $S_3$ models defined in Table II, the coupling constants of the extra Higgs bosons with $\mu$ and $e$ are respectively proportional to $(\tan \beta - \cot \beta)$ and $(\tan \beta + \cot \beta)$ as seen in Eq. (II.6). Thus, the extra Higgs bosons are expected to be $\mu$ and $e$ specific in large or small $\tan \beta$ regions\(^3\). However, this feature is hidden in the Type-II, Type-X and Type-Y $S_3$ models, because at least one of the bottom or tau Yukawa couplings is also enhanced as getting larger values of $\tan \beta$. Therefore, phenomenology in the Type-II, Type-X and Type-Y $S_3$ models are almost the same as those in the Type-II, Type-X and Type-Y THDMs with the softly-broken $Z_2$ symmetry, respectively. Studies for collider signatures using data of 126 GeV Higgs boson at the LHC have been analyzed in Refs. [14, 22] in the softly-broken $Z_2$ symmetric THDMs. Only in the Type-I $S_3$ model, all the Yukawa couplings of the extra Higgs bosons are suppressed by $\cot \beta$, so that the $\mu$ and $e$ specific nature is maintained.

We would like to emphasize that appearance of the $\mu$ and $e$ specific extra Higgs bosons does not appear in the other THDMs without the tree level FCNC; e.g., the $Z_2$ symmetric version and the THDMs with Yukawa alignments discussed in Ref. [23]. In such a THDM, the interaction matrices among a Higgs boson and fermions are proportional to the fermion mass matrices. Therefore,

\(^3\) Cases with small $\tan \beta$; i.e., $\tan \beta \lesssim 1$ is typically disfavored by the $B$ physics data such as the $b \to s\gamma$ process [13, 21].
FIG. 1: Decay branching ratio of $H$ as a function of $\tan \beta$ in the case with $\sin(\beta - \alpha) = 1$. In the left and right panel, the mass of $H$ is taken to be 150 GeV and 350 GeV, respectively.

the branching fractions of $\mathcal{H} \to \mu\mu$ and $\mathcal{H} \to ee$ are suppressed by the factors of $(m_\mu/m_\tau)^2$ and $(m_e/m_\tau)^2$, respectively, compared to that of $\mathcal{H} \to \tau\tau$, where $\mathcal{H}$ denotes an extra neutral Higgs boson. If we consider the most general THDM, sometimes it is called as the Type-III THDM \cite{24}, in which both Higgs doublet fields couple to each fermion, such a proportionality between the matrices can be broken in general. In that case, the $\mu$ and $e$ specific extra Higgs bosons can be obtained by choosing parameters in the interaction matrix. The important point in our model is that we can explain the $\mu$ and $e$ specific nature as a consequence of the $S_3$ symmetry.

Therefore, measuring signatures from the $\mu$ and $e$ specific extra Higgs bosons can be useful to distinguish the other THDMs without the tree level FCNC.

B. Decays of extra Higgs bosons

We first evaluate the decay branching ratios of $H$, $A$ and $H^\pm$ in the Type-I $S_3$ model. In the following calculation, the running quark masses are taken to be $\bar{m}_b = 3.0$ GeV, $\bar{m}_c = 0.677$ GeV and $\bar{m}_s = 0.0934$ GeV. The top quark mass is set to be 173.1 GeV. The strong coupling constant $\alpha_s$ is fixed by 0.118. In Fig. 1 the decay branching fraction of $H$ is shown as a function of $\tan \beta$ in the case of $m_H = 150$ GeV (left panel) and 350 GeV (right panel). In the small $\tan \beta$ region, the main decay modes are $bb$ ($t\bar{t}$), while they are replaced by $\mu^+\mu^-$ and $e^+e^-$ when $\tan \beta$ is larger than about 10 (20) in the case of 150 GeV (350 GeV).

In Fig. 2 the decay branching fraction of $A$ is shown as a function of $\tan \beta$ in the case of $m_A = 150$ GeV (left panel) and 350 GeV (right panel). The $\tan \beta$ dependence of the branching fraction is not so different from that of $H$ in the case of 150 GeV. On the other hand, in the case
FIG. 2: Decay branching ratio of $A$ as a function of $\tan \beta$ in the case with $\sin(\beta - \alpha) = 1$. In the left and right panel, the mass of $A$ is taken to be 150 GeV and 350 GeV, respectively.

FIG. 3: Decay branching ratio of $H^+$ as a function of $\tan \beta$ in the case with $\sin(\beta - \alpha) = 1$. In the left and right panel, the mass of $H$ is taken to be 150 GeV and 350 GeV, respectively.

of $m_A = 350$ GeV, the meeting point of two curves for $t\bar{t}$ and $e^+e^-$ or $\mu^+\mu^-$ is shifted into the larger $\tan \beta$ value about 50, because the suppression of the decay rate of $A \to t\bar{t}$ due to the phase space function is weaker than that of $H$.

The branching fraction of $H^+$ is shown in Fig. 3 as a function of $\tan \beta$ in the case of $m_{H^+} = 150$ GeV (left panel) and 350 GeV (right panel). When $\tan \beta \lesssim 7$ (tan $\beta > 7$), the $H^+ \to \tau^+\nu$ ($H^+ \to \mu^+\nu$ and $e^+\nu$) decay is dominant in the case of $m_{H^+} = 150$ GeV. When $m_{H^+} = 350$ GeV, the main decay mode is changed from $t\bar{t}$ to $\mu^+\nu$ and $e^+\nu$ at $\tan \beta \simeq 65$.

We would like to mention that measuring almost the same branching fractions of $H/A \to e^+e^-$ and $H/A \to \mu^+\mu^-$ as well as those of $H^+ \to e^+\nu$ and $H^+ \to \mu^+\nu$ can be an evidence of the $S_3$ symmetric nature of the model; namely, the electron and muon are included in the same $S_3$
\[
\begin{array}{|c|c|c|c|c|}
\hline
m_{h_{\text{SM}}}[\text{GeV}] & \kappa & \tan \beta (H) & \tan \beta (A) & \tan \beta (H \text{ and } A) \\
\hline
110 & 5.1 & 5.0-16.8 & 3.3-28.1 & 3.0-33.3 \\
\hline
115 & 5.7 & 5.3-16.2 & 3.3-27.4 & 3.0-32.3 \\
\hline
120 & 9.2 & 6.6-12.2 & 4.0-22.1 & 3.6-26.1 \\
\hline
125 & 9.8 & 6.2-12.9 & 4.0-22.8 & 3.3-27.1 \\
\hline
130 & 10.8 & 6.3-13.5 & 4.0-23.4 & 3.3-27.7 \\
\hline
135 & 11.0 & 5.6-15.2 & 3.6-25.8 & 3.3-30.4 \\
\hline
140 & 16.8 & 6.0-13.5 & 4.0-23.4 & 3.3-28.1 \\
\hline
145 & 16.9 & 5.0-16.5 & 3.6-27.7 & 3.0-32.7 \\
\hline
150 & 22.1 & 4.6-17.8 & 3.3-29.7 & 3.0-35.0 \\
\hline
\end{array}
\]

TABLE III: $\kappa$ values and the excluded range of $\tan \beta$ with the 95% C.L. for each mass of the SM Higgs boson.

doublet.

C. Collider signatures

Next, we discuss signatures of the extra Higgs bosons at the LHC. The main production mode of $H$ and $A$ is the gluon fusion process, especially in the small $\tan \beta$ region. The cross section of this mode is suppressed by the factor of $\cot^2 \beta$, so that it does not use in the large $\tan \beta$ region. On the other hand, the cross section for the pair production processes $pp \to HA, H^\pm H$ and $H^\pm A$ do not depend on $\tan \beta$, so that they can be useful even in the large $\tan \beta$ region. We note that the vector boson fusion processes for $H$ and $A$ are vanished at the tree level in the scenario based on $\sin(\beta - \alpha) = 1$. Thus, we consider the signal events from the gluon fusion and the pair production processes.

From the gluon fusion process, the opposite-sign dimuon or dielectron signal can be considered as

\[
gg \to H/A \to \ell^+\ell^-,
\]

where $\ell^\pm$ are $e^\pm$ or $\mu^\pm$. The cross section for this process for $\ell^\pm = \mu^\pm$ is constrained by using the analysis of the search for the SM Higgs boson in the dimuon decay which has been performed from the ATLAS data with the collision energy to be 8 TeV and the integrated luminosity to be 20.7 fb$^{-1}$. The current 95% C.L. upper limit for the cross section $\sigma(pp \to h \to \mu^+\mu^-)$95% C.L. is given by $\sigma(pp \to h \to \mu^+\mu^-)_{\text{SM}} \times \kappa$, where $\sigma(pp \to h \to \mu^+\mu^-)_{\text{SM}}$ is the SM prediction of the cross
TABLE IV: Cross sections for the $H A$, $H^+ H$ and $H^- H$ productions for each fixed value of $m_A$ with the collision energy to be 7 TeV (14 TeV). The masses of $H$ and $H^\pm$ are taken to be the same as $m_A$. The $H^\pm A$ production cross sections are the same as those of $H^\pm H$.

section of the $pp \to h \to \mu^+ \mu^-$ process. The $\kappa$ values are listed for each mass of the SM Higgs boson $m_{h_{SM}}$ in Table III. In the $S_3$ model, this cross section with the $H$ and $A$ mediations can be calculated by

$$\sigma(gg \to H/A \to \mu^+ \mu^-) = \sigma(gg \to h)_{SM} \frac{\Gamma(gg \to H/A)}{\Gamma(gg \to h)_{SM}} \times BR(H/A \to \mu^+ \mu^-),$$

(III.2)

where $\sigma(gg \to h)_{SM}$ is the gluon fusion cross section for the SM Higgs boson, $\Gamma(gg \to h)_{SM}$ is the decay rate of the SM Higgs boson $[H/A]$ into two gluons, and $\text{BR}(H/A \to \mu^+ \mu^-)$ is the branching fraction of the dimuon decay of $H/A$. In order to obtain the cross section from Eq. (III.2), the masses of $H$ and $A$ are taken to be the same as that of the SM Higgs boson. We use the value of $\sigma(gg \to h)_{SM}$ from Ref. \[26\] with the 8 TeV energy. We then obtain the excluded ranges of $\tan \beta$ for the given values of $m_H$ and $m_A$ by requiring

$$\sigma(pp \to h \to \mu^+ \mu^-)_{95\% \text{ C.L.}} > \sigma(gg \to H/A \to \mu^+ \mu^-).$$

(III.3)

In Table III, excluded ranges of $\tan \beta$ with the 95% C.L. are listed by using Eq. (III.3) for each $\kappa$ value. In this table, the values written in the third, fourth and last columns respectively show the excluded range of $\tan \beta$ only by taking into account the $H$, $A$ contribution and both $H$ and $A$ contributions with $m_H = m_A$ to the dimuon process. We find that the region of $3 \lesssim \tan \beta \lesssim 30$ is excluded with the 95% C.L. in the mass range from 110 GeV to 150 GeV in the case of $m_H = m_A$.

Apart from the gluon fusion process, we discuss the pair production processes. In Table IV, the cross sections for the pair productions are listed with the collision energy to be 7 TeV and 14 TeV in the case of $m_H = m_A = m_{H^+}$. From these processes, we can obtain the same-sign dilepton events as follows

$$pp \to HA \to \ell^+ \ell^- \ell^+ \ell^-, \quad pp \to H^\pm H/H^\pm A \to \ell^\pm \nu \ell^+ \ell^-.$$  

(III.4)
| $m_A$ [GeV] | 100 110 120 130 140 150 160 170 180 190 200 |
|---|---|---|---|---|---|---|---|---|---|---|
| $\mu^+\mu^-\mu^+\mu^-$ [fb] | 59.5 42.8 31.4 23.4 17.7 13.6 10.1 8.35 6.68 5.37 4.32 |
| $\mu^+\mu^-\mu^+\nu$ [fb] | 67.8 49.9 37.3 28.5 21.8 17.1 13.4 10.8 9.05 7.06 5.74 |
| $\mu^+\mu^+X$ [fb] | 195 143 106 80.3 61.4 47.9 37.3 29.9 24.0 19.5 15.8 |

TABLE V: Cross sections for the $pp \rightarrow HA \rightarrow \mu^+\mu^-\mu^+\mu^-$ and $pp \rightarrow H^+H \rightarrow \mu^+\mu^-\mu^+\nu$ processes with the collision energy to be 7 TeV after taking the kinematic cuts given in Eqs. (III.5) and (III.6) for $\ell = \mu^+$. The total cross section of the $\mu^+\mu^+X$ final states are also shown in the last row. The masses of $H$ and $H^\pm$ are taken to be the same as $m_A$. The branching fractions of $H/A \rightarrow \mu^+\mu^-$ and $H^+ \rightarrow \mu^+\nu$ are taken to be 100%.

There are three (four) possible final states; i.e., $e^+e^-e^+e^-$, $\mu^+\mu^-\mu^+\mu^-$ and $e^+e^-\mu^+\mu^-$ ($e^+\nu e^+e^-$, $\mu^+\nu\mu^+\mu^-$, $\mu^+\nu e^+e^-$ and $e^+\nu\mu^+\mu^-$) for the $HA$ ($H^\pm H/H^\pm A$) production mode. The same-sign dilepton event search has been reported by the ATLAS Collaboration with the collision energy to be 7 TeV and the integrated luminosity to be 4.7 fb$^{-1}$ in [27]. The strongest constraint can be obtained from the $\mu^+\mu^+$ event whose 95% C.L. upper limit for the cross section is given by 15.2 fb. According to Ref. [27], we impose the following kinematic cuts which are used to obtain the above upper bound as

$$|\eta^\ell| < 2.5, \quad p_T^\ell > 20 \text{ GeV}, \quad (III.5)$$
$$M_{\ell\ell} > 15 \text{ GeV}, \quad (III.6)$$

where $\eta^\ell$, $p_T^\ell$ and $M_{\ell\ell}$ are the pseudorapidity, the transverse momentum for a charged-lepton and the invariant mass for a dilepton system, respectively. In order to compare the upper limit for the cross section of the $\mu^+\mu^+$ channel, the above cuts should be imposed for $\ell = \mu^+$. The signal cross sections are calculated by using CalcHEP [28] and Cteq6l for the parton distribution function (PDF).

In Table V the cross sections for the $pp \rightarrow HA \rightarrow \mu^+\mu^-\mu^+\mu^-$ and $pp \rightarrow H^+H \rightarrow \mu^+\mu^-\mu^+\nu$ are listed after taking the cuts given in Eqs. (III.5) and (III.6) for $\ell = \mu^+$ for each fixed value of $m_A$ with the collision energy to be 7 TeV. We take $m_H$ and $m_{H^+}$ to be the same as $m_A$. The total cross section of $\mu^+\mu^+X$ final states are also shown, which is the sum of the contributions from $HA$, $H^+H$ and $H^+A$ productions. The values of the cross sections in this table are displayed by assuming 100% branching fractions of $H/A \rightarrow \mu^+\mu^-$ and $H^+ \rightarrow \mu^+\nu$, so that the actual cross sections are obtained by multiplying the branching fractions of the above modes.

In Fig. 4 the excluded regions are shown on the tan $\beta$-$m_A$ plane in the case of $m_H = m_A = m_{H^+}$. 

FIG. 4: Excluded regions with 95% C.L. on the \( \tan \beta \)-\( m_A \) plane from the gluon fusion process and the same-sign dimuon processes at the LHC.

\[
\begin{array}{cccccccccc}
\hline
m_A \text{ [GeV]} & 100 & 120 & 140 & 160 & 180 & 200 & 250 & 300 & 400 & 500 \\
\hline
\mu^+\mu^- e^+e^- \text{ [fb]} & 205 & 123 & 77.8 & 51.4 & 35.3 & 24.8 & 11.5 & 5.858 & 1.88 & 0.72 \\
\hline
\end{array}
\]

TABLE VI: Cross sections for the \( pp \rightarrow HA \rightarrow \mu^+\mu^- e^+e^- \) process after taking the basic kinematic cuts given in Eq. (III.5) with the collision energy to be 14 TeV.

The black and red shaded regions are respectively excluded with the 95% C.L. from the opposite-sign dimuon signal from the gluon fusion process and the same-sign dimuon signal from the pair production processes. We note that the region with \( \tan \beta > 100 \) is not so changed from that with \( \tan \beta \gtrsim 30 \) in this plot, because the branching fraction of \( H/A \rightarrow \mu^+\mu^- \) and \( H^+ \rightarrow \mu^+\nu \) are reached to be the maximal value, \textit{i.e.}, 50\%. Thus, when \( m_A \) is smaller than about 140 GeV, \( \tan \beta \gtrsim 3 \) is excluded with the 95% C.L. from the both constraints.

Finally, we discuss the discovery potential of \( H \) and \( A \) with the collision energy to be 14 TeV. We focus on the pair production process, especially for the \( pp \rightarrow HA \rightarrow e^+e^- \mu^+\mu^- \) event, because we can clearly see the electron and muon specific nature of \( H \) and \( A \). To estimate the background cross section, we use the \texttt{MadGraph5} \cite{29} and \texttt{Cteq6l} for the PDF. After we impose the basic kinematic cuts as given in Eq. (III.5) in which \( \ell \) is all the charged-leptons in the final state, we obtain the background cross section to be about 8.1 fb. The signal cross section is calculated by using \texttt{CalcHEP} and \texttt{Cteq6l}. In Table VI the cross section for the \( pp \rightarrow HA \rightarrow \mu^+\mu^- e^+e^- \) process after taking the kinematic cut is shown for each fixed value of \( m_A \). We here introduce the signal
FIG. 5: The 5σ discovery potential at the LHC with the collision energy to be 14 TeV. The black and red contours respectively show the parameter region giving $S = 5$ by assuming the integrated luminosity to be 300 fb$^{-1}$ and 3000 fb$^{-1}$.

**significance $S$ defined as**

$$S = \frac{N_{\text{sig}}}{\sqrt{N_{\text{sig}} + N_{\text{bg}}}},$$

**(III.7)**

where $N_{\text{sig}}$ and $N_{\text{bg}}$ denote the event number of the signal and background processes, respectively.

In Fig. 5, we show the discovery potential of the $e^+e^-\mu^+\mu^-$ signal from the $pp \rightarrow HA$ production. The signal significance $S$ is larger than 5 in the regions inside the black and red curves, where the integrated luminosity is assumed to be 300 fb$^{-1}$ and 3000 fb$^{-1}$. Because the top quark pair decay of $H$ and $A$ opens, the discovery reach is saturated at about 350 GeV. We find that $H$ and $A$ with their masses up to 350 GeV can be discovered by 5σ in the case of $\tan \beta \gtrsim 30$ with 300 fb$^{-1}$. In the 3000 fb$^{-1}$ luminosity, the discovery reach can be above 350 GeV when $\tan \beta \gtrsim 30$.

**IV. SUMMARY AND CONCLUSION**

We have studied the THDMs in the framework based on the $S_3$ flavor symmetry. Assigning the first and second generation lepton fields (two Higgs doublet fields) to be the doublet (singlet) under $S_3$, the mass matrix for the charged-leptons is obtained to be the diagonal form in the weak eigenbasis. The quark masses and mixings are explained as the same way in the SM by assuming the $S_3$ charge for quarks to be the singlet. The $S_3$ charge assignment to the Higgs doublet fields in our model, which is different from the previous studies for $S_3$ models where the Higgs fields are usually taken to be the $S_3$ doublet, is suitable to explain the SM-like Higgs boson with the mass of 126 GeV discovered at the LHC.
The tree level FCNC appearing in the general THDMs is forbidden by the $S_3$ symmetry in our model set up in which four types of the Yukawa interaction are allowed depending on the $S_3$ charge assignments for fermions named as Type-I, Type-II, Type-X and Type-Y $S_3$ models. We have found that the extra Higgs bosons $H$, $A$ and $H^\pm$ can be electron and muon specific in the Type-I $S_3$ model in the large tan $\beta$ regions. Namely, the decay modes of $H/A \rightarrow \mu\mu$, $H/A \rightarrow e e$ and $H^\pm \rightarrow \mu^\pm\nu/\mu^\pm\nu$ are dominant, and the branching fraction for the muon final state is almost the same as that for the electron final state. This property does not appear in any other THDMs without the tree level FCNC such as a $Z_2$ symmetric version of the THDMs. Therefore, measuring signatures of the $\mu/e$ specific extra Higgs bosons can be a direct probe of our model.

We have explored excluded regions on the tan $\beta$-$m_A$ plane has been evaluated as shown in Fig. 4 by using the Higgs boson search data of the dimuon decay mode data and the same-sign dimuon event. We also have estimated the $5\sigma$ discovery potential of the $pp \rightarrow HA \rightarrow e^+e^-\mu^+\mu^-$ signal assuming the center of mass energy to be 14 TeV and the integrated luminosity to be 300 fb$^{-1}$ and 3000 fb$^{-1}$.

Acknowledgments

H.O. thanks to Prof. Eung-Jin Chun for fruitful discussion. Y.K. thanks Korea Institute for Advanced Study for the travel support and local hospitality during some parts of this work. K.Y. was supported in part by the National Science Council of R.O.C. under Grant No. NSC-101-2811-M-008-014.

[1] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) [arXiv:1207.7214 [hep-ex]].
[2] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) [arXiv:1207.7235 [hep-ex]].
[3] S. Pakvasa and H. Sugawara, Phys. Lett. B 73, 61 (1978); S. Pakvasa and H. Sugawara, Phys. Lett. B 82, 105 (1979).
[4] J. Kubo, A. Mondragon, M. Mondragon and E. Rodriguez-Jauregui, Prog. Theor. Phys. 109, 795 (2003) [Erratum-ibid. 114, 287 (2005)] [hep-ph/0302196]; S. -L. Chen, M. Frigerio and E. Ma, Phys. Rev. D 70, 073008 (2004) [Erratum-ibid. D 70, 079905 (2004)] [hep-ph/0404084]; E. Ma, [hep-ph/0409075]; F. Gonzalez Canales, A. Mondragon and M. Mondragon, Fortsch. Phys. 61, 546 (2013) [arXiv:1205.4753 [hep-ph]].
[5] E. Ma and G. Rajasekaran, Phys. Rev. D 64, 113012 (2001) [hep-ph/0106291]; G. Altarelli and F. Feruglio, Nucl. Phys. B 720, 64 (2005) [hep-ph/0504165].
[6] S. L. Glashow and S. Weinberg, Phys. Rev. D 15, 1958 (1977);
E. A. Paschos, Phys. Rev. D 15, 1966 (1977).
[7] V. D. Barger, J. L. Hewett and R. J. N. Phillips, Phys. Rev. D 41, 3421 (1990); Y. Grossman, Nucl.
Phys. B 426, 355 (1994).
[8] A. G. Akeroyd and W. J. Stirling, Nucl. Phys. B 447, 3 (1995); A. G. Akeroyd, Phys. Lett. B 377, 95
(1996).
[9] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher and J. P. Silva, Phys. Rept. 516, 1
(2012) [arXiv:1106.0654 [hep-ph]].
[10] E. Ma and B. Melic, Phys. Lett. B 725, 402 (2013) [arXiv:1303.6928 [hep-ph]].
[11] H. E. Logan and D. MacLennan, Phys. Rev. D 79, 115022 (2009) [arXiv:0903.2246 [hep-ph]].
[12] H. E. Logan and D. MacLennan, Phys. Rev. D 81, 075016 (2010) [arXiv:1002.4916 [hep-ph]].
[13] M. Aoki, S. Kanemura, K. Tsumura and K. Yagyu, Phys. Rev. D 80, 015017 (2009) [arXiv:0902.4663
[hep-ph]].
[14] C. -W. Chiang and K. Yagyu, JHEP 1307, 160 (2013) [arXiv:1303.0168 [hep-ph]].
[15] K. Kannike, M. Raidal, D. M. Straub and A. Strumia, JHEP 1202, 106 (2012) [Erratum-ibid. 1210,
136 (2012)] [arXiv:1111.2551 [hep-ph]].
[16] M. Krawczyk and J. Zochowski, Phys. Rev. D 55, 6968 (1997) [hep-ph/9608321].
[17] F. Jegerlehner and A. Nyffeler, Phys. Rept. 477, 1 (2009).
[18] M. Benayoun, P. David, L. Delbuono and F. Jegerlehner, Eur. Phys. J. C 72, 1848 (2012).
[19] M. Misiak, H. M. Asatrian, K. Bieri, M. Czakon, A. Czarnecki, T. Ewerth, A. Ferroglia and P. Gambino
et al., Phys. Rev. Lett. 98, 022002 (2007).
[20] F. Mahmoudi and O. Stal, Phys. Rev. D 81, 035016 (2010).
[21] U. Haisch, [arXiv:0805.2144 [hep-ph]].
[22] H. S. Cheon and S. K. Kang, JHEP 1309, 085 (2013) [arXiv:1207.1083 [hep-ph]]; N. Craig and S. Thomas, JHEP 1211, 083 (2012) [arXiv:1207.4835 [hep-ph]]; S. Chang, S. K. Kang, J. -P. Lee, K. Y. Lee, S. C. Park and J. Song, JHEP 1305, 075 (2013) [arXiv:1210.3439 [hep-ph]]; Y. Bai, V. Barger, L. L. Everett and G. Shaughnessy, Phys. Rev. D 87, no. 11, 115013 (2013) [arXiv:1210.4922 [hep-ph]]; P. M. Ferreira, R. Santos, H. E. Haber and J. P. Silva, Phys. Rev. D 87, no. 5, 055009 (2013) [arXiv:1211.3131 [hep-ph]]; A. Drozd, B. Grzadkowski, J. F. Gunion and Y. Jiang, JHEP 1305, 072 (2013) [arXiv:1211.3580 [hep-ph]]; J. Chang, K. Cheung, P. -Y. Tseng and T. -C. Yuan, Phys. Rev. D 87, no. 3, 035008 (2013) [arXiv:1211.3849 [hep-ph]]; C. -Y. Chen and S. Dawson, Phys. Rev. D 87, no. 5, 055016 (2013) [arXiv:1301.0309 [hep-ph]]; A. Celis, V. Ilisie and A. Pich, JHEP 1307, 053 (2013) [arXiv:1302.4022 [hep-ph]];
B. Grinstein and P. Uttayarat, JHEP 1306, 094 (2013) [Erratum-ibid. 1309, 110 (2013)] [arXiv:1304.0028 [hep-ph]];

B. Coleppa, F. Kling and S. Su, JHEP 1401, 161 (2014) [arXiv:1305.0002 [hep-ph]];

C. -Y. Chen, S. Dawson and M. Sher, Phys. Rev. D 88, 015018 (2013) [arXiv:1305.1624 [hep-ph]];

O. Eberhardt, U. Nierste and M. Wiebusch, JHEP 1307, 118 (2013) [arXiv:1305.1649 [hep-ph]];

N. Craig, J. Galloway and S. Thomas, [arXiv:1305.2424 [hep-ph]];

R. V. Harlander, S. Liebler and T. Zirke, JHEP 1402, 023 (2014); [arXiv:1307.8122 [hep-ph]]

N. Chen, C. Du, Y. Fang and L.-C. L, Phys. Rev. D 89, 115006 (2014) [arXiv:1312.7212 [hep-ph]];

J. Baglio, O. Eberhardt, U. Nierste and M. Wiebusch, Phys. Rev. D 90, 015008 (2014) [arXiv:1403.1264 [hep-ph]];

P. M. Ferreira, J. F. Gunion, H. E. Haber and R. Santos, Phys. Rev. D 89, 115003 (2014) [arXiv:1403.4736 [hep-ph]];

B. Coleppa, F. Kling and S. Su, [arXiv:1404.1922 [hep-ph]];

L. Wang and X.-F. Han, [arXiv:1404.7437 [hep-ph]];

B. Dumont, J. F. Gunion, Y. Jiang and S. Kraml, [arXiv:1405.3584 [hep-ph]].

[23] A. Pich and P. Tuzon, Phys. Rev. D 80, 091702 (2009) [arXiv:0908.1554 [hep-ph]].

[24] T. P. Cheng and M. Sher, Phys. Rev. D 35, 3484 (1987). D. Atwood, L. Reina and A. Soni, Phys. Rev. D 55, 3156 (1997) [hep-ph/9609279]; P. Ball and R. Zwicky, Phys. Rev. D 71, 014015 (2005) [hep-ph/0406232]; J. L. Diaz-Cruz, J. Hernandez-Sanchez, S. Moretti, R. Noriega-Papaqui and A. Rosado, Phys. Rev. D 79, 095025 (2009) [arXiv:0902.4490 [hep-ph]].

[25] ATLAS Collaboraboration, ATLAS-CONF-2013-010.

[26] https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CERNYellowReportPageAt8TeV

[27] G. Aad et al., [ATLAS Collaboration], JHEP 1212, 007 (2012) [arXiv:1210.4538 [hep-ex]].

[28] A. Pukhov, [hep-ph/0412191].

[29] J. Alwall, M. Herquet, F. Maltoni, O. Mattelaer and T. Stelzer, JHEP 1106, 128 (2011) [arXiv:1106.0522 [hep-ph]].