Radiative-Z_{2}-breaking twin Higgs model

Jiang-Hao Yu

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The Radiative $\mathbb{Z}_2$ Breaking Twin Higgs

Jiang-Hao Yu[1]

[1] *Amherst Center for Fundamental Interactions, Department of Physics, University of Massachusetts-Amherst, Amherst, MA 01003, U.S.A.*

In twin Higgs model, the Higgs boson mass is protected by a $\mathbb{Z}_2$ symmetry. The $\mathbb{Z}_2$ symmetry needs to be broken either explicitly or spontaneously to obtain misalignment between electroweak and new physics vacua. We propose a novel $\mathbb{Z}_2$ breaking mechanism, in which the $\mathbb{Z}_2$ is spontaneously broken by radiative corrections to the Higgs potential. Two twin Higgses with different vacua are needed, and vacuum misalignment is realized by opposite but comparable contributions from gauge and Yukawa interactions to the potential. Because of fully radiative symmetry breaking, the Higgs sector is completely determined by twin Higgs vacuum, Yukawa and gauge couplings. There are eight pseudo-Goldstone bosons: the Higgs boson, inert doublet Higgs, and three twin scalars. We show the 125 GeV Higgs mass and constraints from Higgs coupling measurements could be satisfied.

The discovery of a 125 GeV Higgs boson at the LHC [1] sharpens existing naturalness problem in the Standard Model (SM): quadratically divergent quantum corrections to the Higgs boson mass destabilize the electroweak scale. This suggest the existence of new physics (NP) with a new symmetry which protects the Higgs mass against large radiative corrections. In supersymmetry and composite Higgs [2, 3] models, SM partners from new symmetry play the role of stabilizing the Higgs mass. Unfortunately, null results on new physics searches at the LHC put tight lower bounds on them. This leads to a sub-percent level of tuning between electroweak and NP cutoff scales, which is the little hierarchy problem [4].

The twin Higgs model [5] [see also [6–9]] is introduced to address the little hierarchy problem. It introduces a mirror copy of the SM: the twin sector, which is completely neutral under the SM gauge group. Since twin partners are colorless, they could have sub-TeV masses and thus soften the little hierarchy. The approximate global symmetry breaking $U(4)/U(3)$ at scale f produces a pseudo Goldstone boson (PGB), identified as the Higgs boson. Imposing a discrete $\mathbb{Z}_2$ symmetry between SM and twin sector ensures that there is no quadratically divergent quantum corrections to the Higgs mass term. The $\mathbb{Z}_2$ symmetry needs to be broken to realize vacuum misalignment mechanism: how to generate asymmetric vacua $v < f$ for the Higgs boson and twin Higgs boson. In original twin Higgs model, the $\mathbb{Z}_2$ symmetry is broken explicitly by introducing soft or hard $\mathbb{Z}_2$ breaking terms in scalar potential. This minimal model has been extended to incorporate two twin Higgses in non-supersymmetric [10] and supersymmetric [11] frameworks. The advantage of two twin Higgses setup is that it could accommodate a different $\mathbb{Z}_2$ breaking mechanism [12, 13]: the $\mathbb{Z}_2$ symmetry is spontaneously broken by a bilinear Higgs mass term between two twin Higgses. The vacuum expectation value (VEV) of one twin Higgs preserves $\mathbb{Z}_2$, while the other breaks $\mathbb{Z}_2$ completely and spontaneously. As the effective tadpole, this bilinear term transmits the $\mathbb{Z}_2$ breaking from the broken one to the unbroken one, and thus obtain vacuum misalignment.

In this work, we propose a novel approach to spontaneously break the $\mathbb{Z}_2$ symmetry: the radiative $\mathbb{Z}_2$ breaking mechanism. The Higgs potential is fully generated from gauge and Yukawa corrections, and the $\mathbb{Z}_2$ symmetry is broken spontaneously and radiatively. As a benefit, this model could explain the origin of the Higgs mass, which is fully determined by gauge and Yukawa contributions. On the other hand, in the original or tadpole induced twin Higgs model, one needs to introduce soft or bilinear mass terms by hand to break $\mathbb{Z}_2$, but the origin of these mass terms is unknown. Thus this model could be viewed as a UV completion of the original twin Higgs model: the soft mass term origins from the radiative corrections in gauge and Yukawa sector. The radiatively generated Higgs potential is parametrized as

$$V(h) \simeq \frac{g_{\rm SM}^2 m_*^2}{16\pi^2}\left(-a\sin^2\frac{h}{f} + b\sin^4\frac{h}{f}\right), \qquad (1)$$

where $g_{\rm SM}$ is a typical SM coupling and m_* is the mass scale of twin partner. In the original twin Higgs model, the Z_2 symmetry ensures $a = b$, and thus induces symmetric vacua $v = f$. The Z_2 symmetry can only be broken if $a \neq b$. In this work, we introduce a second twin Higgs H_2 with $U(4)/U(3)$ symmetry breaking in additional to the twin Higgs H_1. If H_1 and H_2 receive opposite radiative corrections, a new radiative term $|H_1|^2|H_2|^2$ contributes to a and b in Eq. 1 differently, and thus triggers the Z_2 breaking between $H_{1A} \leftrightarrow H_{1B}$. However, the Z_2 breaking is not enough to obtain the 125 GeV Higgs mass. Usually a and b generated from radiative corrections are at the same order, which only induce symmetric vacua $v \sim f$. To obtain the realistic vacuum $v \ll f$, we need either a is suppressed or b is enhanced. For example, in littlest Higgs [3] the quartic term b is enhanced via adding tree-level quartic terms by hand. Without adding terms by hand, we could utilize possibly large cancellation among radiative corrections to suppress quadratic term a. In the original twin Higgs, we note that gauge and Yukawa corrections to the quadratic term a have opposite sign. However, a large cancellation cannot happen because gauge corrections are much smaller than Yukawa ones. Interestingly, in our setup, gauge corrections can be enhanced if the global symmetry breaking scale f' of the second twin Higgs is much greater than f. This causes comparable but opposite gauge and Yukawa corrections, and leads to vacuum misalignment with a moderate tuning between v and f. Thus, a different but more minimal spontaneous $\mathbb{Z}_2$ breaking mechanism is naturally realized without introducing either a soft breaking term or a bilinear tadpole term.

Two $U(4)$ invariant Higgs fields are introduced:

$$H_1 \equiv \begin{pmatrix} H_{1A} \\ H_{1B} \end{pmatrix}, \qquad H_2 \equiv \begin{pmatrix} H_{2A} \\ H_{2B} \end{pmatrix}. \tag{2}$$

where two twin Higgs doublets H_{1B} and H_{2B} are in twin sector. The $\mathbb{Z}_2$ symmetry maps the twin Higgses into visible Higgses: $H_{1B} \xrightarrow{\mathbb{Z}_2} H_{1A}$, $H_{2B} \xrightarrow{\mathbb{Z}_2} H_{2A}$. The scalar pontential, which respects both $\mathbb{Z}_2$ and global $U(4)_1 \times U(4)_2$ symmetries, reads

$$\begin{aligned} V(H_1,H_2) = &-\mu_1^2|H_1|^2 - \mu_2^2|H_2|^2 \\ &+\lambda_1|H_1|^4 + \lambda_2|H_2|^4 + \lambda_3|H_1|^2|H_2|^2. \end{aligned} \tag{3}$$

The Higgs sector is weakly gauged under both the SM and the mirror SM gauge symmetries. After symmetry breaking $\langle H_i\rangle \equiv f_i$ $(i = 1, 2)$, the symmetries of the Lagrangian have

$$\begin{aligned} &\text{global symmetry:} \quad U(4)\times U(4) \to U(3)\times U(3), \\ &\text{gauge symmetry:} \quad [SU(2)\times U(1)]_{A,B} \to [SU(2)\times U(1)]_A. \end{aligned} \tag{4}$$

In nonlinear σ Lagrangian, assuming radial modes in H_i are decoupled, the fields H_i are parametrized as

$$H_i = \exp\left[\frac{i}{f_i}\left(\begin{array}{c|cc} \mathbf{0}_{2\times 2} & \mathbf{0}_{1\times 2} & \boldsymbol{h}_1 \\ \hline \mathbf{0}_{2\times 1} & 0 & C_i \\ \boldsymbol{h}_i^* & C_i^* & N_i \end{array}\right)\right]\begin{pmatrix} \mathbf{0}_{1\times 2} \\ \hline 0 \\ f_i \end{pmatrix}, \tag{5}$$

where 14 GBs $\boldsymbol{h}_i, C_i, N_i (i = 1, 2)$ are generated.

There are two ways to incorporate fermions. In the "mirror fermion" assignment [5, 6], the SM fermions have mirror fermions: $q_A(3, 2; 1, 1) \xrightarrow{\mathbb{Z}_2} q_B(1, 1; 3, 2)$ and $t_A(3, 1; 1, 1) \xrightarrow{\mathbb{Z}_2} t_B(1, 1; 3, 1)$, with quantum number assignment $[SU(3), SU(2)]_{A,B}$. The general top-Yukawa Lagrangian reads

$$-\mathcal{L}_{\text{Yuk}} = y_1\left(H_{1A}^\dagger q_A \bar{t}_A + H_{1B}^\dagger q_B \bar{t}_B\right) + (1 \leftrightarrow 2) + h.c. \tag{6}$$

To avoid Higgs mediated flavor changing neutral current in A sector, similar to the two Higgs doublet model (2HDM), either the discrete Z_2' symmetry or aligned Yukawa structure [14] are imposed. We will discuss the following two Yukawa structures: Type-I Yukawa structure ($y_2 = 0$), and a special aligned Yukawa structure $y_2 = \epsilon y_1$ ($\epsilon \ll 1$) [15]. In the "$U(4)$ fermion" assignment [5], the following $U(4)$ fermions are introduced:

$$\begin{aligned} Q &= q_A + q_B + \tilde{q}_A\,(3, 1; 1, 2) + \tilde{q}_B\,(1, 2; 3, 1), \\ U &= t_A\,(3, 1; 1, 1) + t_B\,(1, 1; 3, 1). \end{aligned} \tag{7}$$

The top Yukawa Lagrangian is

$$-\mathcal{L}_{\text{Yuk}} = y_1 H_1^\dagger Q\bar{U} + y_2 H_2^\dagger Q\bar{U} + M\bar{\tilde{q}}_{A,B}\tilde{q}_{A,B} + h.c. \tag{8}$$

Here either Type-I or aligned Yukawa structure is used.

The global $U(4) \times U(4)$ symmetry is weakly broken by the radiative corrections from the gauge and Yukawa interactions. The dominant radiative corrections to the scalar potential are written as

$$\begin{aligned} V_{\text{loop}} = &\,\delta_1|H_{1A}|^4 + \delta_2|H_{2A}|^4 + \delta_3|H_{1A}|^2|H_{2A}|^2 \\ &+\delta_4|H_{1A}^\dagger H_{2A}|^2 + \frac{\delta_5}{2}\left[(H_{1A}^\dagger H_{2A})^2 + h.c.\right] \\ &+\left[(\delta_6|H_{1A}|^2 + \delta_7|H_{2A}|^2)H_{1A}^\dagger H_{2A} + h.c.\right] \\ &+(A \leftrightarrow B). \end{aligned} \tag{9}$$

Here for gauge bosons $W_{A,B}$ and $Z_{A,B}$, the one-loop corrections are

$$\begin{aligned} \delta^B_{1,2} &\simeq -\frac{1}{16\pi^2}\left(\frac{9}{8}g^4 + \frac{3}{4}g^2g'^2 + \frac{3}{8}g'^4\right)\log\frac{\Lambda}{f}, \\ \delta^B_3 &\simeq -\frac{1}{16\pi^2}\left(\frac{9}{4}g^4 - \frac{3}{2}g^2g'^2 + \frac{3}{4}g'^4\right)\log\frac{\Lambda}{f}, \\ \delta^B_4 &\simeq -\frac{1}{16\pi^2}\left(3g^2g'^2\right)\log\frac{\Lambda}{f}, \quad \delta^B_{5-7} \equiv 0, \end{aligned} \tag{10}$$

where $f \equiv \sqrt{f_1^2 + f_2^2}$ and $\Lambda \equiv 4\pi f$. For fermions, radiative corrections depend on the fermion assignment and Yukawa structure. In the Type-I Yukawa structure, we obtain [5]

$$\delta^F_1 \simeq \begin{cases} \frac{3y^4}{16\pi^2}\log\frac{\Lambda^2}{f^2} & \text{(mirror fermion)} \\ \frac{3y^4}{16\pi^2}\frac{y}{x(z-x)}\left[x\log\frac{z+x}{x} - (x \leftrightarrow z)\right] & (U(4)\text{ fermion}) \end{cases} \tag{11}$$

where $x = y_1^2 f^2$ and $z = M^2$, and $\delta^F_{2-7} = 0$. In the aligned Yukawa structure with $y_2 \ll y_1$, we have

$$\begin{aligned} \delta^F_{1,2} &\simeq \frac{3y_{1,2}^4}{16\pi^2}\log\frac{\Lambda^2}{f^2}, \quad \delta^F_{3-5} \simeq \frac{3y_1^2y_2^2}{16\pi^2}\log\frac{\Lambda^2}{f^2}, \\ \delta^F_6 &\simeq \frac{3y_1^3y_2}{16\pi^2}\log\frac{\Lambda^2}{f^2}, \quad \delta^F_7 \simeq \frac{3y_1y_2^3}{16\pi^2}\log\frac{\Lambda^2}{f^2}. \end{aligned} \tag{12}$$

The radiatively generated scalar potential in Eq. 9 further triggers electroweak symmetry breaking and induces VEVs for the GBs $h_{1,2}$ in visible sector. The VEVs of the fields $H_{1,2}$ are parametrized as

$$\langle H_1\rangle \equiv \begin{pmatrix} 0 \\ f_1\sin\theta_1 \\ 0 \\ f_1\cos\theta_1 \end{pmatrix}, \qquad \langle H_2\rangle \equiv \begin{pmatrix} 0 \\ f_2\sin\theta_2 \\ 0 \\ f_2\cos\theta_2 \end{pmatrix} \tag{13}$$

where $\theta_1 \equiv \frac{\langle h_1\rangle}{f_1}$, $\theta_2 \equiv \frac{\langle h_2\rangle}{f_2}$. Similar to 2HDM, $t_\beta \equiv \tan\beta = \frac{f_2}{f_1}$, $\delta_{45} = \delta_4 + \delta_5$ and $\delta_{345} \equiv \delta_3 + \delta_4 + \delta_5$ are used. Imposing tadpole conditions on Eq. 9 determine $\theta_{1,2}$. We will neglect $\delta_{6,7}$ terms, because either $\delta_{6,7} = 0$ in Type-I or $\delta_{6,7} \ll \delta_{1-5}$ ($y_2 \ll g \ll y_1$) in aligned Yukawa structure. The tadpole conditions are

$$\begin{aligned} \sin 4\theta_1 + \Omega_1\sin 4\theta_2 + \Omega_2\sin 2(\theta_1+\theta_2) = 0, \\ \sin 4\theta_1 - \Omega_1\sin 4\theta_2 + \Omega_2\sin 2(\theta_1-\theta_2) = 0, \end{aligned} \tag{14}$$

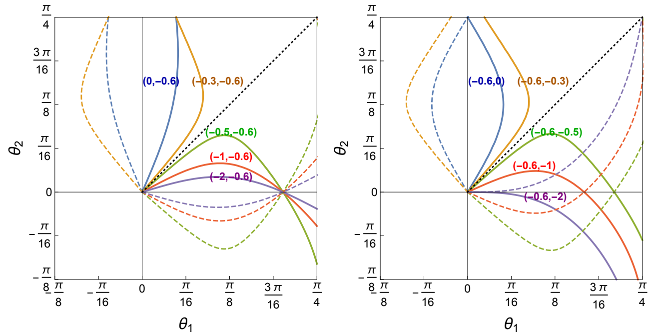

FIG. 1. The contour lines shows the relations between θ_1 and θ_2 imposed by two tadpole conditions, denoted by solid and dashed lines respectively, with fixed $\Omega_2 = -0.6$ (left) and fixed $\Omega_1 = -0.6$ (right).

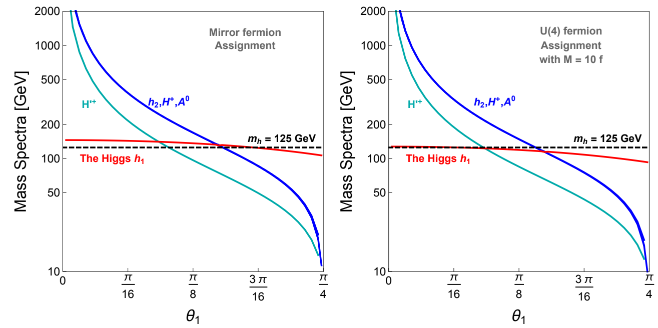

FIG. 2. The mass spectra of the PGBs as function of θ_1 for "mirror fermion" (left panel) and "$U(4)$ fermion" (right panel, $M = 10f$ is taken) assignments.

where $\Omega_1 = t_\beta^4 \delta_2/\delta_1$ and $\Omega_2 = t_\beta^2 \delta_{345}/\delta_1$. We are interested in the region $\Omega_{1,2} < 0$ because $\delta_1 > 0, \delta_{2-5} < 0$. If $|\Omega_1 + \Omega_2| > 1$ the two conditions are symmetric under $\theta_2 \leftrightarrow -\theta_2$. While if $|\Omega_1 + \Omega_2| < 1$ they are symmetric under $\theta_1 \leftrightarrow -\theta_1$. The solutions should be

$$\begin{cases} \theta_2 = 0,\ \theta_1 < \pi/4, & \text{for } |\Omega_1 + \Omega_2| > 1 \\ \theta_1 = 0,\ \theta_2 > \pi/4, & \text{for } |\Omega_1 + \Omega_2| < 1. \end{cases} \tag{15}$$

Thus only one H_i further generates a VEV after radiative symmetry breaking. We plot the (θ_1, θ_2) contours imposed by tadpole conditions for different (Ω_1, Ω_2) in Fig. 1. We note that Ω_2 alone could determine $\theta_{1,2}$ which is intersection point between solid and dashed curves, while Ω_1 only controls the convex behavior of the curves.

To obtain the electroweak vacuum $v \ll f$, $\theta_i < \pi/4$ is required, which spontaneously breaks the $\mathbb{Z}_2$ symmetry. This implies $|\Omega_1 + \Omega_2| > 1$ and $\theta_2 = 0$ in Eq. 15. The electroweak vacuum thus has $v \equiv f_1 \sin\theta_1 = 174$ GeV. The tadpole conditions reduce to one condition

$$\sin^2\theta_1 = \frac{v^2}{f_1^2} \equiv \frac{1}{2}\left(1 + t_\beta^2 \frac{\delta_{345}}{2\delta_1}\right). \tag{16}$$

Because of $\delta_1 > 0, \delta_{345} < 0$, we have $\theta_1 < \pi/4$. Furthermore, if t_β is large ($f_2 > f_1$), θ_1 could be much smaller than $\pi/4$. Therefore, t_β controls the tuning behavior between v and f_1, and it is natural to realize such tuning by setting $f_2 \gg f_1$. Let us understand physics behind this mechanism. The radiative corrections takes the form of Eq. 9. By taking approximation $|H_{iB}|^2 \simeq f_i^2 - |H_{iA}|^2$, we obtain

$$\begin{aligned} V_{\text{loop}} \supset & \left[-|\mu_{h_1}^2||H_{1A}|^2 + \lambda_{1A}|H_{1A}|^4\right] \\ & + \left[|\mu_{h_2}^2||H_{2A}|^2 + \lambda_{2A}|H_{2A}|^4 + \delta_{345}|H_{1A}|^2|H_{2A}|^2\right], \end{aligned} \tag{17}$$

where $\lambda_{iA} = \frac{8\delta_i}{3} + \frac{\delta_{345}}{3} t_\beta^{2(3-2i)}$ with $i = 1, 2$ and the quadratic terms

$$\mu_{h_1}^2 = f_1^2\left(2\delta_1 + \delta_{345} t_\beta^2\right), \mu_{h_2}^2 = f_2^2\left(2\delta_2 + \delta_{345} t_\beta^{-2}\right) \tag{18}$$

Here the mass parameters $\mu_{h_1}^2 < 0$ but $\mu_{h_1}^2 > 0$, and thus H_{1A} obtains VEV but no VEV for H_{2A}. The Z_2 symmetry between $H_{2A} \leftrightarrow H_{2B}$ is spontaneously broken. Thus the term $|H_{1B}|^2|H_{2B}|^2$ generates additional mass term for H_{1A} radiatively [1]. This additional mass term by radiative corrections triggers the Z_2 symmetry breaking between $H_{1A} \leftrightarrow H_{1B}$ spontaneously. This can also be seen from the H_1 potential:

$$\begin{aligned} V(h_1) = & f_1^4\delta_1\left[\sin^4(\frac{h_1}{f_1}) + \cos^4(\frac{h_1}{f_1})\right] + f_1^4 t_\beta^2 \delta_{345}\cos^2(\frac{h_1}{f_1}) \\ & \simeq -(2 + \Omega_2)\delta_1 f_1^2 |h_1|^2 + \frac{8 + \Omega_2}{3}\delta_1 |h_1|^4, \end{aligned} \tag{19}$$

which recovers the first line of the Eq. 18. As mentioned in introduction, if the quadratic term is much smaller than the quartic term, one could obtain the electroweak VEV. Since the Yukawa and gauge corrections have $\delta_1 < 0$ and $\delta_{345} > 0$ respectively, the Higgs mass squared is suppressed by cancellation between Yukawa and gauge corrections. Note t_β plays an important role: only when t_β is not so small, cancellation in quadratic term is adequate. This implies a moderate tuning $f_1 < f_2$, which induces tuning between Yukawa and gauge corrections correspondingly. As a measure of the naturalness, the estimation of the fine-tuning is

$$\Delta = \left|\frac{2\delta m^2}{m_h^2}\right|^{-1} \simeq \left|\frac{3y_t^2 m_{t_B}^2}{4\pi^2 m_h^2}\right|^{-1} \sim \frac{2v^2}{f_1^2} \simeq 1 + \frac{\delta_{345} f_2^2}{2\delta_1 f_1^2} \tag{20}$$

Unlike the soft breaking or tadpole breaking mechanism, the tuning is realized via balancing between the gauge and Yukawa contributions. This is reflected in Eq. 20 via the hierarchy between f_1 and f_2. For example, for a level of tuning 10%, $t_\beta \simeq 3$ ($f_2 \simeq 3f_1$).

The purely radiative symmetry breaking only generate VEV $\langle H_{1A}\rangle$, but not $\langle H_{2A}\rangle$. Zero H_{2A} VEV ($\theta_2 = 0$) implies that the second Higgs H_{2A} in A sector is inert

[1] If the Z_2 symmetry between $H_{2A} \leftrightarrow H_{2B}$ is exact, the terms $|H_{1B}|^2|H_{2B}|^2$ and $|H_{1A}|^2|H_{2A}|^2$ generate opposite but equal mass terms for H_{1A}. Thus the Z_2 symmetry between $H_{1A} \leftrightarrow H_{1B}$ is still unbroken.

Higgs doublet [16], which does not mix with H_{1A}. In visible sector, particles in H_{iA} are identified as GBs $\boldsymbol{h}_i$. Among them, $(z^{0,\pm})$ in H_{1A} and $(H^\pm, A^0)$ in H_{2A} have

$$\begin{aligned} m_{z^0}^2 &= m_{z^\pm}^2 = 0, \quad \text{(exact GBs eaten by } W_A, Z_A), \\ m_{H^\pm}^2 &\simeq -2\delta_2 f_2^2 - \delta_{345} f_1^2 \cos 2\theta_1 - \delta_{45} f_1^2 \sin^2\theta_1, \\ m_{A^0}^2 &\simeq -2\delta_2 f_2^2 - \delta_{345} f_1^2 \cos 2\theta_1 - 2\delta_5 f_1^2 \sin^2\theta_1. \end{aligned} \tag{21}$$

Since the inert Higgs masses depends on $\delta_{2,3}$, their masses are not so light. And two CP-even GBs (h_1, h_2) in H_{1A}, H_{2A}, which do not mix due to zero H_{2A} VEV, have

$$\begin{aligned} m_{h_1}^2 &\simeq 2\delta_1 f_1^2 \sin^2 2\theta_1, \\ m_{h_2}^2 &\simeq -2\delta_2 f_2^2 - \delta_{345} f_1^2 \cos 2\theta_1. \end{aligned} \tag{22}$$

Here h_1 is identified as the SM Higgs boson. However, in twin sector B, the GBs in H_{1B} and H_{2B} are mixed due to the VEVs $f_{1,2}$. The rotation angle β_B between $C_1^\pm(N_1^0)$ and $C_2^\pm(N_2^0)$ is defined as $\tan\beta_B = t_\beta/\cos\theta_1$. Performing rotation to mass basis, we obtain

$$\begin{aligned} m_{N^0}^2 &= m_{C^\pm}^2 = 0, \quad \text{(exact GBs eaten by } W_B, Z_B), \\ m_{H'^\pm}^2 &= -\left(\delta_{45} + \delta_7 \tan\beta_B\right) f_1^2 (\cos^2\theta_1 + t_\beta^2), \\ m_{A'^0}^2 &= -\left(2\delta_5 + \delta_7 \tan\beta_B\right) f_1^2 \frac{(\cos^2\theta_1 + t_\beta^2)^2}{\cos^2\theta_1}. \end{aligned} \tag{23}$$

Fig. 2 shows mass spectra of the pGBs in two cases. In "mirror fermion" case, the pGB masses only depend on single parameter θ_1. Thus the requirement of a 125 GeV Higgs mass determines $\theta_1 = 0.57$, which corresponds to $t_\beta = 2$. In "$U(4)$ fermion" case, mass spectra depend on both θ_1 and vectorlike fermion $\tilde{q}_A$ mass M, which should have $M \leq 4\pi f$. As the M takes smaller value than $4\pi f$, the θ_1, obtained from the 125 GeV Higgs mass condition, gets smaller value. However, when $M = 8f$, θ_1 reaches zero, which put a lower cutoff for M. In the following, we take $M = 10f$ as the benchmark point.

The current limits on NP searches at the LHC put very strong constraints on new particles. New particles in A sector are the inert Higgses $H^\pm, h_2, A^0$, which is typically constrained by electroweak precision tests, Higgs data and dark matter candidates [16]. Approximately, the oblique correction ΔT due to inert Higgses is

$$\Delta T \simeq \frac{1}{24\pi^2 \alpha v^2}(m_{H^\pm} - m_A)(m_{H^\pm} - m_{h_2}). \tag{24}$$

The global fitting result on $\Delta T \simeq 0.07 \pm 0.08$ [16] implies that a mass difference between $H^\pm$ and a neutral inert scalar boson is typically less than 8 GeV. Furthermore, assuming the neutral inert scalar boson being candidate of dark matter, the current data of direct detection constrain the mass of dark matter candidate to be greater than 600 GeV [17]. On the other hand, if masses of inert Higgses are nearly degenerate, it is very difficult to probe the compressed parameter region at current LHC, which has been studied in Ref. [18]. And if we assume the dark matter candidate is the particle in the twin sector, the inert scalars could be as light as 100 GeV with degenerate masses. For particles in twin sector, it is harder to directly probe them due to zero SM charges. However, because of twin colorness, there are rich twin hadron phenomenology, which has been discussed in Ref. [19, 20]. For simplicity, in the following we adopt minimal twin matters: fraternal twin Higgs [8], in which only the third generation twin fermions are introduced, and typically twin lepton is identified as dark matter candidate. In this scenario, the twin photon and A'^0 could be either massless or massive depending on gauge and fermion assignments. For example, the aligned Yukawa structure could lift the A'^0 mass from zero value. If they are massless, they should contribute to dark radiation. Depending on temperature of thermal decoupling between visible and twin sector [20], the number of effective neutrino species $\Delta N_{\rm eff}$ could be adjusted to be within the range of recent Planck measurement 0.11 ± 0.23 [21]. We leave the detailed discussion in future study [22].

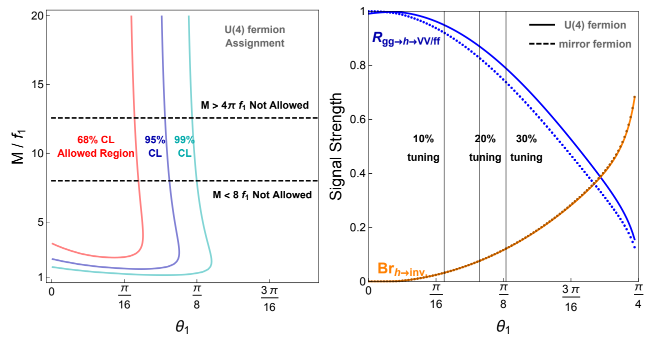

FIG. 3. (Left) the allowed contours on $(\theta_1, M/f_1)$ at 68%, 95%, 99% CLs in "$U(4)$ fermion" assignment. (Right) signal strength in gluon fusion channel (blue) and invisible branching ratio (orange) as function of θ_1 in "$U(4)$" (solid) and "mirror" (dotted) fermions.

The measured Higgs production and decay cross sections, and the upper limits on Higgs invisible decays at the LHC [23] also provide strong constraints on model parameters. The tree-level couplings of the Higgs boson to fermions and bosons in $A(B)$ sector are altered by a factor $\cos\theta_1$ $(\sin\theta_1)$ relative to SM. We assume masses of twin particles are altered by a factor $\cot\theta_1$ relative to SM. In "$U(4)$ fermion" case, there is an additional heavier vectorlike top T which mixes with the top quark through mixing angle $\cos\theta_R = y_t f/\sqrt{M^2 + y^2 f^2}$. This modifies the top-Higgs coupling to $y_t \cos\theta_1 \cos\theta_R$, and a new TTh coupling has $y_t \sin\theta_1 \sin\theta_R$. We calculate various Higgs signal strengths $\mu_{pp\to h_1 \to ii} = \sigma(pp \to h_1){\rm Br}_{h_1\to ii}/\sigma_{\rm SM}{\rm Br}_{\rm SM}$, and invisible decay width in the fraternal twin Higgs setup. Based on Higgs signal strengths at the 8 TeV LHC with 20.7 fb^{-1} data [23], we perform a global fit on model parameters [22]. Fig. 3 (left) shows the allowed contours on $(\theta_1, M/f_1)$ at 68%, 95%, 99% confidence levels (CLs) in "$U(4)$ fermion" case. Note that the Higgs data put strong limits on the new fermion masses $M > 1.5 f_1$, which is comparable with the expected exclusion limit from run-2 LHC [19]. Furthermore, the positive θ_1 requires $M > 8f_1$, which sets much tighter limit than the LHC searches. Fig. 3 (right) plots the signal strength of gluon fusion $gg \to h_1 \to VV/ff$, and invisible decay width in two assignments. We list our global χ^2-fitting results on Higgs signal strengths in "mirror fermion" ("$U(4)$ fermion" with $M = 10f$) assignments:

$$\theta_1 \equiv \frac{v}{f_1} < 0.25\ (0.31) \quad @\ 95\%\ \mathrm{CL}. \tag{25}$$

This limit rules out the whole parameter region of the "mirror fermion" case, but "$U(4)$ fermion" case is still viable. This limit corresponds to around 15% (18%) fine-tuning. In principle, electroweak precision test puts additional constraints on θ_1. Since contributions from the decoupled radial modes are negligible, the T parameter is dominated by the inert Higgses, which is given by Eq. 24. Due to degenerate masses of inert Higgses, contributions from inert Higgs are also negligible, and thus electroweak constraints on θ_1 are quite weak [24]. We also estimate the levels of tuning of about 10%, 20% and 30%, which are shown in Fig. 3 (right). The high luminosity LHC will improve sensitivity of signal strengths to around 5% assuming current uncertainty with 3 ab^{-1} luminosity [25]. The projected limits on θ_1 in "mirror fermion" ("$U(4)$ fermion" with $M = 10f$) assignments are $\theta_1 < 0.17$ (0.23) at 95% CL. This indicates that we could probe this model with about 9% (12%) tuning by the end of high luminosity LHC run.

In summary, we have investigated a minimal two twin Higgs model, which explains the origin of the $\mathbb{Z}_2$ breaking dynamically. The $\mathbb{Z}_2$ symmetry is typically broken by adding new terms, such as soft mass or bilinear mass terms. We discussed the possibility that the $\mathbb{Z}_2$ symmetry is spontaneously broken by purely radiatively generated Higgs potential. The vacuum misalignment $v < f$ is realized radiatively via cancellation of gauge and Yukawa corrections to the Higgs mass term. This theory can be viewed as a UV completion of a natural inert Higgs model. This minimal setup has no free parameter in the Higgs potential, and thus it has rich but predictive phenomenology.

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