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A composite twin Higgs model

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Twin Higgs models are economical extensions of the standard model that stabilize the electroweak scale. In these theories the Higgs field is a pseudo Nambu-Goldstone boson that is protected against radiative corrections up to scales of order 5 TeV by a discrete parity symmetry. We construct, for the first time, a class of composite twin Higgs models based on confining QCD-like dynamics. These theories naturally incorporate a custodial isospin symmetry and predict a rich spectrum of particles with masses of order a TeV that will be accessible at the LHC.

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I. INTRODUCTION

Quantum corrections to the Higgs mass parameter in the standard model (SM) are quadratically divergent. Stabilizing the weak scale against these divergences generally requires a phenomenologically rich spectrum of new particles near a TeV, associated with the existence of a new symmetry of nature. One appealing idea is that the Higgs sector of the SM might actually be the nonlinear sigma model of some larger, dynamically generated pattern of symmetry breaking [1,2]. In such theories, the Higgs behaves much like a pion in QCD; the Higgs is a composite pseudo Nambu-Goldstone boson (pNGB), which is protected from the worst class of quadratic divergences that usually afflict scalar bosons. Precision electroweak measurements currently constrain the compositeness scale to lie above 5 TeV. The fact that the SM gauge couplings, top Yukawa coupling, and Higgs self-coupling necessarily break any global symmetry with order one strength then implies that the mass parameter of the composite Higgs needs additional protection, if the theory is to be natural.

Little Higgs theories [3–5] (for reviews see [6]), are a class of nonlinear sigma models, which realize the Higgs as a protected pNGB. The underlying concept behind little Higgs theories is the idea of ‘‘collective symmetry breaking’’—the global symmetry of which the Higgs is the pNGB is broken only when two or more couplings in the Lagrangian are nonvanishing. This is a significant restriction on the form of the quantum corrections to the pNGB potential, which can be used to engineer natural electroweak symmetry breaking. These theories stabilize the weak scale to about 5–10 TeV.

Twin Higgs theories [7,8] are an alternative class of nonlinear sigma models that also realize the Higgs as a protected pNGB. These theories possess a discrete Z_2 interchange symmetry, in addition to the approximate global symmetry of which the Higgs is the pNGB. In the existing twin Higgs models this Z_2 symmetry is identified either with mirror symmetry, or with left-right symmetry. This discrete symmetry is enough to ensure that any quadratically divergent contribution to the scalar potential accidentally respects the global symmetry, and therefore cannot contribute to the mass of the pNGB. These theories also stabilize the weak scale to about 5–10 TeV.

Since in general little Higgs and twin Higgs theories have been formulated only as nonlinear sigma models, above 5–10 TeV these theories require ultraviolet completions to maintain unitarity. Weakly coupled ultraviolet completions using supersymmetry have been constructed for both the little Higgs [9,10] and the twin Higgs [11], in the context of the supersymmetric little hierarchy problem. In the little Higgs case, nonsupersymmetric ultraviolet completions have also been constructed [12–14]. There has also been some work on the difficult problem of realizing the little Higgs as a strongly coupled composite [15–19] furthering the analogy to QCD. However, in the twin Higgs case the corresponding problem has not been addressed.

A significant challenge in dynamically realizing a composite twin Higgs is to ensure that the strong dynamics respects a custodial SU(2) symmetry. For this to happen, the custodial symmetry must be contained in the nonlinearly realized global symmetry of the Higgs sector. In the twin Higgs models currently in the literature this global symmetry is either SU(4), which is spontaneously broken to SU(3), or O(8), spontaneously broken to O(7). While the breaking of SU(4) to SU(3) is fairly straightforward to realize through QCD-like strong dynamics [16], this pattern does not admit a custodial SU(2). On the other hand, while the O(8) → O(7) pattern does preserve a custodial symmetry, this pattern is significantly more complicated to realize through strong dynamics.

In this paper we identify an alternative pattern of symmetry breaking for twin Higgs models that naturally incorporates a custodial isospin symmetry. We then show how this pattern can be realized through QCD-like dynamics, and apply these ideas to construct a class of composite twin Higgs models with left-right symmetry. These theories predict a rich spectrum of new particles at the TeV scale that will be accessible to the LHC.

We begin by constructing an alternative realization of the twin Higgs model, in its left-right symmetric incarnation. Consider a scalar field H that transforms as a fundamental under an Sp(4) global symmetry, and which is also charged under a global U(1). If H acquires a vacuum expectation value (VEV) such that $\langle H\rangle = (0, 0, 0, f)$, the $\mathrm{Sp}(4)\times\mathrm{U}(1)$ global symmetry is spontaneously broken to $\mathrm{SU}(2)\times\mathrm{U}(1)$, and there are 7 Goldstone bosons. We now break the global Sp(4) explicitly by gauging an $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ subgroup. The overall U(1), which is to be identified with $\mathrm{U}(1)_{B-L}$, is also gauged. The overall gauge structure is therefore that of a left-right symmetric model [20].

Under gauge transformations the field H decomposes into (H_L, H_R), where H_L is a doublet under $\mathrm{SU}(2)_L$ and H_R is a doublet under $\mathrm{SU}(2)_R$. If the VEV of $\langle H\rangle$ points along a direction that breaks $\mathrm{SU}(2)_R$ but preserves $\mathrm{SU}(2)_L$, the surviving gauge symmetry is the familiar $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ of the SM. Of the 7 Goldstone bosons, 3 are eaten. The remaining 4 Goldstone bosons, which are contained in H_L, are to be identified with the SM Higgs doublet. If the discrete parity symmetry, which interchanges $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$, is exact, H_L is protected against quadratic divergences by the twin Higgs mechanism. The key observation is that the discrete symmetry ensures that any quadratically divergent contribution to the scalar potential has an Sp(4) invariant form, and therefore cannot contribute to the mass of the Goldstones.

Yukawa interactions can take the same form as in the original left-right twin Higgs model, since they are only required to respect the gauge and parity symmetries, which are identical in both models. Although these couplings violate the global symmetry with order one strength, the discrete parity symmetry again ensures that quadratic divergences are absent. From this we infer that $[\mathrm{Sp}(4)\times\mathrm{U}(1)]/[\mathrm{SU}(2)\times\mathrm{U}(1)]$ constitutes an alternative symmetry breaking pattern that allows the realization of a twin Higgs model with left-right symmetry.

Although this construction is extremely simple, it does not admit a custodial SU(2) symmetry. Furthermore, it is not clear whether such a pattern of symmetry breaking can arise from strong dynamics. In the next section we show that a natural generalization of this model exists that addresses the first problem. We then go on to discuss how the required symmetry breaking pattern can be realized through the condensation of strongly coupled fermions, in analogy with QCD.

II. A CUSTODIAL SYMMETRY FOR THE TWIN HIGGS

Consider a theory with an $\mathrm{Sp}(4)\times\mathrm{Sp}(4)$ global symmetry, which is spontaneously broken down to the diagonal Sp (4) at approximately the scale f. We label the 10 resulting Nambu-Goldstone bosons (NGBs) that are produced by π^A, and define

$$X = f\exp(2i\pi^A T^A/f). \tag{1}$$

Here, the matrices T^A are the generators of Sp(4), and correspond to the matrices

$$\begin{pmatrix} \sigma^a & 0 \\ 0 & 0\end{pmatrix}\begin{pmatrix} 0 & 0 \\ 0 & \sigma^a\end{pmatrix}\begin{pmatrix} 0 & iI \\ -iI & 0\end{pmatrix}\begin{pmatrix} 0 & \sigma^a \\ \sigma^a & 0\end{pmatrix}, \tag{2}$$

where $a \in \{1, 2, 3\}$ and the σ^a are the three Pauli matrices. We now gauge an $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ subgroup of the first Sp(4), and an $\mathrm{SU}(2)_{L'} \times \mathrm{U}(1)_{R'}$ subgroup of the second Sp (4). Here, $\mathrm{U}(1)_{R'}$ is the diagonal generator of the $\mathrm{SU}(2)_{R'}$ contained in the second Sp(4). We label the gauge coupling constants of these four groups as g_L, g_R, g'_L, and g'_R, respectively. The unbroken gauge symmetry is then $\mathrm{SU}(2)\times\mathrm{U}(1)$, which is identified with the electroweak gauge group of the SM. Note that this symmetry breaking pattern is similar to that of the little Higgs model of Chang and Wacker [5]. Of the original 10 Goldstone bosons, 6 are eaten. The remaining 4 pseudo-Goldstone bosons are identified with the SM Higgs doublet, and correspond to the generators T^a,

$$\{T^a\} = \left\{\begin{pmatrix} 0 & iI \\ -iI & 0\end{pmatrix}\begin{pmatrix} 0 & \sigma^a \\ \sigma^a & 0\end{pmatrix}\right\}. \tag{3}$$

We can write an effective field theory for the pNGBs, which is valid at low momenta. This takes the form of a nonlinear sigma model. In general the Lagrangian for this theory will contain all operators involving the field X consistent with the nonlinearly realized $\mathrm{Sp}(4)\times\mathrm{Sp}(4)$ global symmetry. Nonrenormalizable operators are suppressed by the cutoff Λ of the nonlinear sigma model, and their coefficients are determined by the specific ultraviolet completion. The cutoff Λ must be less than about $4\pi f$, where the upper bound corresponds to strong coupling.

In this low-energy theory, the masses of the pseudo Goldstones are protected against one loop quadratic divergences from gauge interactions. This can be understood as a consequence of the little Higgs mechanism. The theory has an exact Sp(4) global symmetry in the limit that g_L and g_R are zero, and also in the limit that g'_L and g'_R are zero. Any diagram that results in a quadratic divergence must therefore involve both these sets of couplings. The leading contributions to the pseudo-Goldstone masses arise at order g^4, and are therefore necessarily suppressed by at least two loop factors.

If the low-energy effective theory is weakly coupled at the scale Λ, in the special case that the $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$ of the first Sp(4) are related by a discrete interchange symmetry, so that $g_L = g_R$, there is an alternative way of understanding this cancellation based on the twin Higgs mechanism. This discrete symmetry ensures that at quadratic order in X all radiative corrections to the pseudo-Goldstone potential are invariant under the first Sp(4), and therefore must simply vanish. To see this let us consider all possible operators consistent with the gauge symmetry at quadratic order in X in the nonlinear sigma model. At one loop these terms are the only ones generated with a quadratically divergent coefficient in the effective potential. Schematically these operators include

$$\begin{aligned} &X_{LL'}X^{\dagger L'L}, \quad X_{RL'}X^{\dagger L'R}, \quad X_{L3}X^{\dagger 3L}, \\ &X_{R3}X^{\dagger 3R}, \quad X_{L4}X^{\dagger 4L}, \quad X_{R4}X^{\dagger 4R}, \end{aligned} \tag{4}$$

and also (suppressing Hermitian conjugates)

$$\begin{aligned} &\epsilon_{LL}\epsilon_{L'L'}X_{LL'}X_{LL'}, \quad \epsilon_{RR}\epsilon_{L'L'}X_{RL'}X_{RL'}, \\ &\epsilon_{LL}X_{L3}X_{L4}, \quad \epsilon_{RR}X_{R3}X_{R4}. \end{aligned} \tag{5}$$

Here, L and L' take values 1 and 2, while R takes values 3 and 4. The discrete $L \leftrightarrow R$ symmetry ensures that in the Lagrangian operators on the same line above necessarily have the same coefficient. Then it is clear that at quadratic order in X the Lagrangian is actually invariant under the first global Sp(4) symmetry. This symmetry is only broken at quartic order in X, and therefore corrections to the pNGB mass are loop suppressed, and at most logarithmically divergent.

The argument above does not carry over to the case where the low-energy theory is strongly coupled at the cutoff Λ, because now the quartic terms in X, though still loop suppressed, need not be small. The reason is that the quartic terms can now be generated at order g^2 by loops involving operators that are strongly coupled at the cutoff, and this could potentially compensate for the loop suppression.[1] However, the little Higgs mechanism still ensures that any such term is invariant under the second Sp(4), and so does not contribute to the pNGB potential. Therefore, the leading terms that contribute to the mass of the pNGB only arise at order $g^2 g'^2$, and are suppressed by an additional loop factor.

This construction ensures that the strong dynamics does not violate the custodial SU(2) symmetry. To see this explicitly, note that we can write

$$2\pi^a T^a \equiv \begin{pmatrix} 0 & \phi \\ \phi^\dagger & 0 \end{pmatrix}, \tag{6}$$

where $\phi = (i\sigma_2 h_L^*, h_L)$. The full expression for X is

$$\cos\left(\frac{|h_L|}{f}\right)f + \frac{if}{|h_L|}\sin\left(\frac{|h_L|}{f}\right)\begin{pmatrix} 0 & \phi \\ \phi^\dagger & 0 \end{pmatrix}. \tag{7}$$

The Lagrangian written as a function of X preserves the $SU(2)_L \times SU(2)_R$ subgroup of the diagonal Sp(4), under which $\phi \to U_L \phi U_R^\dagger$, and $|h_L| \to |h_L|$. After electroweak symmetry breaking, $\langle h_L \rangle = (0, v)$, and the diagonal SU(2) symmetry is preserved. This is precisely the custodial symmetry we are looking for.

In order to write down Yukawa couplings, first make the identification

$$X_{i4} = H_i = (H_L, H_R). \tag{8}$$

Yukawa couplings can be written down exactly as in the original left-right twin Higgs model, in terms of H_L and H_R, so that the discrete $L \leftrightarrow R$ symmetry is preserved. The twin Higgs mechanism then ensures that quadratic divergences from the fermion sector preserve the first global Sp (4) symmetry and vanish from the pseudo-Goldstone potential, just as in the gauge sector.

The fermionic content of the theory then contains three generations of

$$\begin{aligned} Q_L &= (u, d)_L = [2, 1, 1/3], \\ L_L &= (\nu, e)_L = [2, 1, -1], \\ Q_R &= (u, d)_R = [1, 2, 1/3], \\ L_R &= (\nu, e)_R = [1, 2, -1], \end{aligned} \tag{9}$$

where the square brackets indicate the quantum numbers of the corresponding field under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. We identify $U(1)_{R'}$ with $U(1)_{(B-L)/2}$. As dictated by left-right symmetry the theory includes right-handed neutrinos in addition to the SM fermions.

The Higgs fields have quantum numbers

$$H_L = [2, 1, 1], \qquad H_R = [1, 2, 1] \tag{10}$$

under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The down-type Yukawa couplings of the SM arise from nonrenormalizable couplings of the form

$$\left\{\frac{\bar{Q}_R H_R H_L^\dagger Q_L + \bar{L}_R H_R H_L^\dagger L_L}{\Lambda}\right\} + \text{H.c.} \tag{11}$$

Here, Λ is an ultraviolet cutoff, which we take to be about 10 TeV, the limit of validity of the nonlinear sigma model. Similarly, the up-type Yukawa couplings of the SM emerge from

$$\left\{\frac{\bar{Q}_R H_R^\dagger H_L Q_L + \text{H.c.}}{\Lambda}\right\}. \tag{12}$$

The top Yukawa coupling is too large to be naturally obtained from a nonrenormalizable operator. As in the original left-right twin Higgs model, we therefore introduce a pair of vectorlike quarks T_L and T_R, which have the quantum numbers

$$T_L = [1, 1, 4/3], \qquad T_R = [1, 1, 4/3] \tag{13}$$

under $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. We can then write the Yukawa coupling

$$(y\bar{Q}_R H_R^\dagger T_L + y\bar{Q}_L H_L^\dagger T_R + M\bar{T}_L T_R) + \text{H.c.} \tag{14}$$

[1]Whether a quartic term is generated at order g^2 in a general twin Higgs model in the limit of strong coupling depends on the pattern of symmetry breaking. For example, if the symmetry breaking pattern is $O(8) \to O(7)$, a quartic term is only generated at order g^4, and is therefore always loop suppressed [8].

Here, Q_L and Q_R are the usual left and right-handed third generation quark doublets of the left-right model.

Since the top Yukawa gives the largest contribution to the Higgs potential, let us understand the cancellation of quadratic divergences in this case. As in the gauge case, the discrete $L \leftrightarrow R$ symmetry ensures that terms quadratic in X are invariant under the first Sp(4), and do not contribute to the potential for the pNGB. Terms quartic and higher order in X that violate Sp(4) are only generated at order y^4, and not at order y^2, and are therefore suppressed by one loop factor, even in the limit that the nonlinear sigma model is strongly coupled at the cutoff.

It is also possible to generate the smaller Yukawa couplings from renormalizable interactions [21], (see also [22]). To do this we introduce three generations of vector-like fermions with the following charge assignments:

$$\begin{aligned} U_L &= [1,1,4/3], \qquad U_R = [1,1,4/3],\\ D_L &= [1,1,-2/3], \qquad D_R = [1,1,-2/3],\\ E_L &= [1,1,-2], \qquad E_R = [1,1,-2]. \end{aligned} \tag{15}$$

Then the Yukawa couplings for the lighter fermions can be written down in analogy with that for the top. For example, the charged lepton Yukawa couplings arise from the interactions

$$\{\bar{L}_R H_R E_L + \bar{L}_L H_L E_R + M\bar{E}_L E_R\} + \mathrm{H.c.} \tag{16}$$

We choose the mass parameter M to be of order several TeV. On integrating out E_L and E_R we get back exactly the same nonrenormalizable operator that earlier generated the charged lepton masses.

III. A TWIN HIGGS MODEL FROM STRONG DYNAMICS

We now explain how the symmetry breaking pattern $\mathrm{Sp}(4)\times\mathrm{Sp}(4)\rightarrow\mathrm{Sp}(4)$ may be obtained from QCD-like strong dynamics. Our discussion will closely follow that of [18], where the same problem was considered in the context of a dynamical realization of the little Higgs model of Chang and Wacker [5]. Consider an $\mathrm{SU(N}_c)$ gauge group, with a set of four fermions, $\chi_{\alpha i}$, in the fundamental representation. Here, α represents an $\mathrm{SU(N}_c)$ gauge index and i labels the fermions from 1 through 4. We also add a set of four right-handed fermions $\psi_{\alpha i}$. When the $\mathrm{SU(N}_c)$ theory gets strong, a condensate $\langle \chi_i \bar{\psi}_j\rangle \propto \delta_{ij}$ forms and breaks the $\mathrm{SU}(4)^2$ flavor symmetry to the diagonal SU(4). We label the 15 resulting NGBs that are produced by π^A, and define $X = f\exp(2i\pi^A T^A/f)$, where the matrices T^A are generators of SU(4). We also add to the theory a nonrenormalizable term

$$\frac{m^2}{(4\pi f^2)^2}\,\mathrm{Tr}[(\chi\bar{\psi})J(\chi\bar{\psi})^T J] \sim m^2\,\mathrm{Tr}[XJX^TJ], \tag{17}$$

where J is the matrix

$$J = \begin{pmatrix} i\sigma^2 & 0 \\ 0 & i\sigma^2 \end{pmatrix}. \tag{18}$$

The effect of this term is to explicitly break the global $\mathrm{SU}(4)^2$ symmetry to $\mathrm{Sp}(4)^2$, thereby giving a mass of order m to 5 of the 15 NGBs. With the addition of this term the pattern of global symmetry breaking is in fact $\mathrm{Sp}(4)^2\rightarrow\mathrm{Sp}(4)$, which accounts for the 10 surviving NGBs. The unbroken global symmetry, the diagonal Sp(4), contains the custodial SU(2) symmetry we desire.

In order to recreate the low-energy structure of the model of the previous section we gauge the subgroups

$$[\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R]\times[\mathrm{SU}(2)_{L'}\times\mathrm{U}(1)_{R'}]\subset\mathrm{Sp}(4)^2 \tag{19}$$

as shown in Fig. 1. After symmetry breaking, this gauge symmetry is broken down to the $\mathrm{SU}(2)\times\mathrm{U}(1)_Y$ gauge symmetry of the SM, where SU(2) is the diagonal subgroup of $\mathrm{SU}(2)_L\times\mathrm{SU}(2)_{L'}$, while $\mathrm{U}(1)_Y$ is the unbroken linear combination of the diagonal generator of $\mathrm{SU}(2)_R$ and $\mathrm{U}(1)_{R'}$. Of the 10 surviving NGBs, 6 are eaten by the broken gauge symmetries, while the 4 that are left over precisely constitute the SM Higgs doublet.

In order to write down the Higgs couplings to fermions we simply make the replacement

$$H_i \rightarrow \frac{\chi_i\bar{\psi}_4}{4\pi f^2} \tag{20}$$

in the Yukawa couplings of the previous section. For example, the left-right symmetric top Yukawa couplings become

$$\left\{y\bar{Q}_R\left(\frac{\chi_R\bar{\psi}_4}{4\pi f^2}\right)T_L + y\bar{Q}_L\left(\frac{\chi_L\bar{\psi}_4}{4\pi f^2}\right)^\dagger T_R\right\}. \tag{21}$$

These interactions are nonrenormalizable, and therefore require additional new physics to generate them. We leave the question of the ultraviolet origin of these operators for future work.

We briefly consider the precision electroweak constraints on this theory. In general, bounds from the S parameter on any composite Higgs force the compositeness scale $\Lambda \sim 4\pi f$ to be larger than or of order 5 TeV. Another source of corrections to precision electroweak observables arises from higher order operators in the ex-

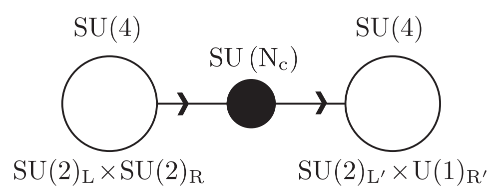

FIG. 1. A UV completion for the custodial Twin Higgs model. The theory has a global $\mathrm{SU}(4)^2$ flavor symmetry, with the indicated gauged subgroups. The two link fields represent the sets of $\mathrm{SU(N}_c)$ fundamentals, ψ_i and χ_j.

pansion of the kinetic term for X in terms of the π fields that contribute to the ρ parameter. Since the sum over the T^A in $X = f\exp(2i\pi^A T^A/f)$ now runs over all SU(4) generators, and not just the generators of Sp(4), there are fields with mass below the compositeness scale that correspond to the generators of SU(4)/Sp(4). These fields, which we denote by H_L', have exactly the same gauge quantum numbers as the light Higgs. The nonrenormalizable terms that arise in the expansion of X in terms of the π fields involve custodial SU(2) violating couplings of H_L' to the light Higgs H_L, and thereby contribute to the ρ parameter.

The effect of the nonrenormalizable term in Eq. (17) is to give a mass m to the fields in H_L', and to thereby decouple them from the low-energy spectrum. The precision electroweak constraints therefore translate into a lower bound on the parameter m. A quick estimate of the size of the correction to ρ yields

$$\frac{\delta m_Z^2}{m_Z^2} \sim \frac{\langle H_L' \rangle^2}{f^2}. \tag{22}$$

A VEV for H_L' arises from the radiatively generated mass term, which mixes H_L' with the light Higgs,

$$\langle H_L' \rangle \sim \left(\frac{f}{4\pi m}\right)^2 v. \tag{23}$$

Here, v is the electroweak VEV. From these formulas we estimate that the precision electroweak constraints on deviations of the ρ parameter from its SM value are comfortably satisfied provided that m is greater than or of order 500 GeV.

As in any general two Higgs doublet model, the presence of a second Higgs doublet can also lead to contributions to the ρ parameter at loop level. However, in this specific model, this contribution translates into a lower bound on m somewhat weaker than the one we have already found.

The additional $SU(2)_R$, $U(1)_{B-L}$ and $SU(2)_{L'}$ gauge bosons also contribute to the precision electroweak observables [5]. In general, these force f to be of order 1500 GeV or larger, reintroducing fine-tuning. However, f can be as low as 500 GeV if, as in the original left-right twin Higgs model, there is a second field $\hat{X}$ with exactly the same quantum numbers as X that exhibits exactly the same pattern of symmetry breaking, but where the decay constant $\hat{f}$ somewhat larger than f. Then, provided that $\hat{f}$ is greater than about 1500 GeV the precision electroweak constraints from the new gauge bosons are satisfied, and the fine-tuning is under control. The field $\hat{X}$ can also be used to generate neutrino masses [23] and dark matter [24], as in the original left-right twin Higgs model.

Much of the heavy spectrum of particles predicted by this theory will be accessible at the LHC. The new fields in the gauge and top sector must have mass of order the TeV scale if they are to be relevant for stabilizing the electroweak scale. Production and decay of the heavy top partner, as well as the massive gauge bosons associated with $SU(2)_R$ and $U(1)_{B-L}$, have been studied in the context of the left-right twin Higgs model [25]. The heavy electroweak doublet of scalars, from the explicit breaking of the global $SU(4)^2$ symmetry to $Sp(4)^2$ in Eq. (17), and the massive gauge bosons that constitute the linear combination of $SU(2)_L$ and $SU(2)_{L'}$ that is orthogonal to SU(2) of the SM, are key predictions of the underlying composite structure of this model. While the electroweak doublet decays primarily into third generation quarks and antiquarks, the new gauge bosons can decay either into SM fermions, or into electroweak gauge bosons.

In summary, we have identified a new class of left-right twin Higgs models that naturally incorporate a custodial SU(2) symmetry, and shown how the relevant pattern of symmetry breaking can be realized through QCD-like strong dynamics. This constitutes an important first step in the construction of completely realistic twin Higgs models from strong dynamics.

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