Maximally Symmetric Composite Higgs Models

Csaba Csáki,[1] Teng Ma,[2] and Jing Shu[2,3,4,*]
[1]Department of Physics, LEPP, Cornell University, Ithaca, New York 14853, USA
[2]CAS Key Laboratory of Theoretical Physics, Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100190, China
[3]School of Physical Sciences, University of Chinese Academy of Sciences, Beijing 100190, People's Republic of China
[4]CAS Center for Excellence in Particle Physics, Beijing 100049, China

(Received 15 February 2017; revised manuscript received 30 June 2017; published 27 September 2017)

Maximal symmetry is a novel tool for composite pseudo Goldstone boson Higgs models: it is a remnant of an enhanced global symmetry of the composite fermion sector involving a twisting with the Higgs field. Maximal symmetry has far-reaching consequences: it ensures that the Higgs potential is finite and fully calculable, and also minimizes the tuning. We present a detailed analysis of the maximally symmetric $SO(5)/SO(4)$ model and comment on its observational consequences.

DOI: 10.1103/PhysRevLett.119.131803

Electroweak symmetry breaking (EWSB) is a key ingredient of the standard model (SM) responsible for all elementary particle masses. While the discovery of the Higgs boson [1,2] was a major milestone towards understanding the mechanism of EWSB, several important issues remain unexplained. In analogy with superconductivity the Higgs potential is assumed to be of the simplest Landau-Ginzburg type [3]. In the condensed matter systems we understand that the potential describes the condensation of emergent collective modes; however, in particle physics we do not even know if the Higgs boson is elementary or composite, and what the true Higgs potential is. We would also like to understand whether the Higgs potential is stable under large quantum corrections in the ultraviolet and whether a small or large fine-tuning is needed to maintain the hierarchy.

The idea that the Higgs boson is a pseudo Nambu-Goldstone boson (pNGB) [4–6] of spontaneously broken approximate global symmetry of some strong dynamics gives intriguing answers to the above mysteries. In this scenario, the Higgs boson could be a bound state of some strongly coupled constituents, while the entire Higgs potential is radiatively generated via loops from the top and gauge sectors, which will trigger vacuum misalignment and EWSB. Because of its pNGB nature, the Higgs mass remains naturally light. The cutoff scale is reduced to the confinement scale $\Lambda \sim 4\pi f$ (where f is the scale of global symmetry breaking). The sensitivity of the Higgs potential to this confinement scale can be further reduced by different mechanisms [7–11]. However, the parameters of the existing models have to be tuned to achieve a realistic Higgs potential and particle spectrum. The origin of this tuning is to ensure that the EWSB vacuum expectation value (VEV) v is small compared to the global symmetry beaking scale f, $v/f \ll 1$, to evade electroweak precision [12] and direct detection bounds for the top partners [13,14].

In this Letter, we propose a novel type of composite Higgs model that can address the above issues and requires only the minimal structure of the general low-energy Goldstone boson (GB) Lagrangian. We present models with a "maximal symmetry": a remnant of the chiral global symmetries of the fermion sector which has wide-ranging consequences for the properties of the Higgs potential. This symmetry involves a twisting with the Higgs VEV, and such symmetry has not been previously considered. Because of this twist the maximal symmetry will not contain the original shift symmetry of the GBs; hence, a nonvanishing Higgs potential will be generated. However, the symmetry will be powerful enough to render the contribution of the top sector to the Coleman-Weinberg Higgs potential [15] automatically finite. Maximal symmetry will require the vanishing of one kind of the form factors $\Pi_1^{q,t}(p)$, and as a consequence imply a specific form of the Higgs potential, and hence, the tuning of the model reduces to the minimal universal tuning: a single parameter $\xi = \sin^2(v/f)$ must be fine-tuned to obtain a realistic electroweak symmetry breaking sector. There have been few truly new concepts for composite Higgs model building since the introduction of the ideas of collective symmetry breaking [7,8], partial compositeness, or the twin Higgs model [9]. Maximal symmetry is such a novel direction that, as we will demonstrate here, will lead to a much simpler and predictive model, and it is expected to have a wide range of applications for many types of composite Higgs models and related topics.

As usual in composite Higgs models, we will consider a strongly coupled system which dynamically breaks its global symmetry G to H, and the Higgs fields are identified with the pNGBs which lie in the coset space G/H. The additional assumption we will make is that the coset space is a "symmetric space" [16], which implies the existence of a parity operator V (called Higgs parity) and the automorphism of the form [17]

$$VT^aV^\dagger = T^a, \qquad VT^{\hat{a}}V^\dagger = -T^{\hat{a}}, \tag{1}$$

where $(T^a)T^{\hat{a}}$ are the (un)broken generators. Many of the most commonly used moduli spaces satisfy this requirement, including the $SO(5)/SO(4)$ of the minimal composite Higgs model, which we will study in detail below.

As usual, the pNGB fields $h^{\hat{a}}$ can be described by the Goldstone matrix

$$U = \exp(ih^{\hat{a}}T^{\hat{a}}/f). \tag{2}$$

The main consequence of the existence of V is that one can define a modified pNGB matrix which transforms linearly under the full set of symmetries G. The original pNGB matrix U has the nonlinear transformation properties [18,19] $U \to gUh(h^{\hat{a}}, g)^\dagger$, where $g \in G$ and $h \in H$, and h depends nonlinearly on the pNGB field $h^{\hat{a}}$ and the transformation element g. However, the parity transformed pNGB matrix $\tilde{U} = VUV^\dagger = U^\dagger$ transforms as $\tilde{U} \to VgV^\dagger \tilde{U} h^\dagger$. We can then define the modified pNGB matrix:

$$\Sigma' \equiv U\tilde{U}^\dagger V = U^2 V, \tag{3}$$

which transforms linearly under the full global symmetries $\Sigma' \to g\Sigma' g^\dagger$.

The linearly realized global symmetry can be used to fully fix the structure of the low-energy effective Lagrangian of the theory. The SM fermions are charged under the $SU(2)_L \times U(1)_Y$, which is a subgroup of the full global symmetries G; thus, they can always be embedded into the full symmetry group G, though this embedding may break the full global symmetry itself (hence, it is a spurionic embedding). We will assume that the left-handed (LH) top doublet q_L and right-handed (RH) top singlet t_R are embedded into Ψ_{q_L} and Ψ_{t_R}, which are in some representation of the full global symmetry group G. Thus imposing the original G symmetry will completely fix the most general effective action for the SM fermion fields coupled to the pseudo Goldstone Higgs bosons:

$$\begin{aligned}\mathcal{L}_{\text{eff}} = {} & \bar{\Psi}_{q_L} M_1^t(p)\Sigma'\Psi_{t_R} + \bar{\Psi}_{q_L}\not{p}[\Pi_0^q(p) + \Pi_1^q(p)\Sigma']\Psi_{q_L} \\ & + \bar{\Psi}_{t_R}\not{p}[\Pi_0^t(p) + \Pi_1^t(p)\Sigma']\Psi_{t_R} + \text{H.c.}\end{aligned} \tag{4}$$

The form factors $\Pi_{0,1}^{q,t}$ and M_1^t encode the effect of the strong dynamics, and for simplicity we have assumed that Ψ_{q_L,t_R} are in the fundamental of G. Note that since $\Sigma'^2 = 1$, we only need to expand Eq. (4) to linear order in Σ'. Once the embeddings of the SM fermions into G are fixed, this expression can be expanded to find the effective SM Lagrangian in terms of the form factors.

A careful examination of the symmetries of the effective Lagrangian Eq. (4) will allow us to identify an enlarged global symmetry in certain limits, which we will call the maximal symmetry. This maximal symmetry is the key new ingredient of composite Higgs models, which will be the focus of discussions for the rest of this Letter. Let us first start with the massless limit of Eq. (4) when $M_1^t = 0$ and also assume $\Pi_1^{q,t} = 0$. In this case the global symmetry G is enlarged to a chiral $G_L \times G_R$ symmetry acting on the left- or right-handed fermions Ψ_{q_L}/Ψ_{t_R} [20]. Now turning on the top mass term $\bar{\Psi}_{q_L} M_1^t(p)\Sigma'\Psi_{t_R}$ (while still keeping $\Pi_1^{q,t} = 0$), we observe that we do not break the enlarged global symmetry completely, but rather leave a subgroup $G_{V'}$ of $G_L \times G_R$ unbroken. We call this $G_{V'}$, the maximal symmetry, which is identified with the subgroup that keeps the pNGB field invariant, $g_L\Sigma' g_R^\dagger = \Sigma'$. A more convenient form of the defining relation for maximal symmetry can be obtained by absorbing the Higgs field into the definition of the transformation $g'_{L,R} = U^\dagger g_{L,R} U$, such that $g'_L V g'^\dagger_R = V$, forming the maximal symmetry group $G_{V'}$. Shortly we will see that the origin of this maximal symmetry can be traced back to a true global symmetry of the heavy composites. This symmetry will have far-reaching consequences for the structure of the Higgs potential: it will render the potential finite, and imply that the model has the minimal universal amount of tuning. One immediate consequence of maximal symmetry is that the Higgs potential is directly proportional to the top mass square, since in the limit when $M_1^t = 0$ the Lagrangian Eq. (4) does not break the global shift symmetry explicitly. Thus, if the top mass is collective, then maximal symmetry will automatically imply that the Higgs potential from the fermion sector is finite.

Until now, we have simply assumed certain properties of the low-energy effective action Eq. (4) and found the important consequences of the emerging maximally symmetric structure. Next, we would like to explain the conditions for the actual existence of this symmetry by examining the simplest realistic example of the $SO(5)/SO(4)$ MCHM$_5$ [21]: we will see that the existence of maximal symmetry will impose a condition on the spectrum of the composites as well as relations among the mixing terms between the elementary and the composite sectors. The explicit form of the Higgs parity operator is $V = \text{diag}(1,1,1,1,-1)$ with the properties $V = V^\dagger$ and $V^2 = 1$. The SM fermions are embedded into Ψ_{q_L} and Ψ_{t_R}, which are in the 5 of $SO(5)$ while we assume that the composite fermions (top partners) Ψ_Q and Ψ_S transform as a **4** (for Ψ_Q) and **1** (for Ψ_S) under $SO(4)$, which can be combined into a full **5** of $SO(5)$ [22]. The general fermionic Lagrangian can then be parametrized as [23]

$$\begin{aligned}\mathcal{L}_f = {} & \bar{\Psi}_Q(i\not{\nabla} - M_Q)\Psi_Q + \bar{\Psi}_S(i\not{\nabla} - M_S)\Psi_S \\ & + \frac{f}{\sqrt{2}}\bar{\Psi}_{t_R} P_L(\epsilon_{tS} U\Psi_S + \epsilon_{tQ} U\Psi_Q) \\ & + f\bar{\Psi}_{q_L} P_R(\epsilon_{qS} U\Psi_S + \epsilon_{qQ} U\Psi_Q) + \text{H.c.},\end{aligned} \tag{5}$$

where ∇ is a covariant derivative with respect to the SM gauge fields.

In order to understand the symmetry properties of this Lagrangian more easily, it is useful to combine Ψ_Q and Ψ_S back to complete representations **5** of the global symmetry $SO(5)$ (and assume for simplicity that CP is conserved): $\Psi_+ = (\Psi_Q + \Psi_S)/\sqrt{2}$, $\Psi_- = (\Psi_Q - \Psi_S)/\sqrt{2}$. Thus, Ψ_+ and Ψ_- are related by the Higgs parity operator: $\Psi_+ = V\Psi_-$, and are not independent fields. Our original fermion Lagrangian Eq. (5) in terms of $\Psi_\pm$ is

$$\begin{aligned}\mathcal{L}_f = {} & 2\bar{\Psi}_+ i\nabla\!\!\!/ \Psi_+ + f(c_{-R}\bar{\Psi}_{t_R} UV\Psi_{+L} + c_{+R}\bar{\Psi}_{t_R} U\Psi_{+L}) \\ & - (M_Q + M_S)\bar{\Psi}_{+L}\Psi_{+R} - (M_Q - M_S)\bar{\Psi}_{+L} V\Psi_{+R} \\ & + f(c_{-L}\bar{\Psi}_{q_L} UV\Psi_{+R} + c_{+L}\bar{\Psi}_{q_L} U\Psi_{+R}) + \text{H.c.}, \qquad (6)\end{aligned}$$

where the Yukawa couplings are $c_{\pm R} = (\epsilon_{tQ} \pm \epsilon_{tS})/2$, $c_{\pm L} = (\epsilon_{qQ} \pm \epsilon_{qS})/\sqrt{2}$.

This simple form of the Lagrangian allows us to identify the possible symmetry breaking patterns and the conditions for the emergence of the maximal symmetry. We have assumed here that the composite fermions Ψ_Q and Ψ_S fill out a full $SO(5)$ representation. This does not generically have to be the case, but it will be a necessary condition on the spectrum of composites in order to obtain maximal symmetry. Once the composites do fill out a complete $SO(5)$ representation, its kinetic term will have the enlarged $SO(5)_L \times SO(5)_R$ chiral global flavor symmetry, which can have various symmetry breaking patterns depending on the structure of the Yukawa couplings and composite mass terms. These symmetry breaking patterns will determine the form of the radiatively induced Higgs potential and its degree of divergence. If $c_{-L/R}$ and $c_{+L/R}$ were to appear simultaneously in the Lagrangian, one would not be able to maintain an entire $SO(5)$ global symmetry as needed for maximal symmetry. This requirement for maximal symmetry is equivalent to the assumption that the elementary-composite mixing terms are fully $SO(5)$ invariant. Since our goal is to find an implementation of the maximal symmetry, we will set $c_{-L} = c_{-R} = 0$ in the general Lagrangian. In this case, the Lagrangian is

$$\begin{aligned}\mathcal{L}_f = {} & 2\bar{\Psi}_+ i\nabla\!\!\!/ \Psi_+ + f c_{+R}\bar{\Psi}_{t_R} U\Psi_{+L} + f c_{+L}\bar{\Psi}_{q_L} U\Psi_{+R} \\ & - \bar{\Psi}_{+L}[(M_Q + M_S) + (M_Q - M_S)V]\Psi_{+R} + \text{H.c.} \qquad (7)\end{aligned}$$

Once we impose $c_{-L,R} = 0$, the mixing terms will have the full $SO(5)_L \times SO(5)_R$ chiral global symmetry, and the breaking pattern depends on the relation of the mass terms $M_{Q,S}$, giving rise to the following possible breaking patterns:

$$\begin{aligned} M_Q - M_S = 0 &\Rightarrow SO(5)_L \times SO(5)_R/SO(5)_V, \\ M_Q + M_S = 0 &\Rightarrow SO(5)_L \times SO(5)_R/SO(5)_{V'}, \\ |M_Q| \neq |M_S| &\Rightarrow SO(5)_L \times SO(5)_R/SO(4)_V. \qquad (8)\end{aligned}$$

Clearly, the second case $M_Q + M_S = 0$ corresponds to the maximally symmetric scenario. There is still a remaining global symmetry, the maximal symmetry $SO(5)_{V'}$, under which the composite fields $\Psi_{+L,R}$ transform as $\Psi_{+L} \to g'_L\Psi_{+L}$, $\Psi_{+R} \to g'_R\Psi_{+R}$, with $g'_L V g'^{\dagger}_R = V$, shedding light on the true origin of the maximal symmetry. Maximal symmetry which, however, does *not* contain the entire Goldstone shift symmetry [since the Goldstone shift symmetry is a subgroup of $SO(5)_V$ and not of $SO(5)_{V'}$]; thus, a nonvanishing Higgs potential will be generated.

We can easily understand the structure of the radiatively induced Higgs potential by examining the collective symmetry breaking properties of the maximally symmetric case corresponding to $M_Q + M_S = 0$ in Eq. (7). The Goldstone shift symmetry corresponds to the broken generators of the original global symmetry G. Every time we have a particular set of terms turned on in the Lagrangian, we need to check whether the global symmetries that are not explicitly broken by turning on these new terms do contain the broken generators of the original global symmetry. For example, when c_{+L}, c_{+R} and the twisted mass term $M_Q - M_S$ are turned on, the unbroken symmetry is the maximal symmetry $SO(5)_{V'}$, which does not coincide with the original global symmetry $SO(5)_V$. Thus, the shift symmetry in this case is explicitly broken, and one expects a corresponding term generated in the potential. However, if we turn off one of the c's (for example, c_{+L}), then the situation changes. In this case, we can define a global symmetry transformation under which $\Psi_{+L} \to g'_L\Psi_{+L}$, and identify this with the original $SO(5)$ global symmetry which contains the Higgs shift symmetry $h^{\hat{a}} \to h^{\hat{a}} + \alpha^{\hat{a}}$ [$g'_L = \exp(i\alpha^{\hat{a}}T^{\hat{a}})$] so that $f c_{+R}\bar{\Psi}_{t_R} U\Psi_{+L}$ is invariant. Since Ψ_{+R} does not couple to U, we can just do an additional rotation on $\Psi_{+R} \to V g'_L \Psi_{+R}$ to keep the twisted mass term invariant, which keeps the whole Lagrangian invariant. Thus, we need all three terms to be simultaneously turned on, making this a "triply collective" breaking. In fact, to be able to actually generate a potential all three breaking terms have to show up twice, giving rise to a finite contribution of the form $|c_{+L}|^2|c_{+R}|^2 f^4(M_Q - M_S)^2/\Lambda^2$. The maximal symmetry automatically implies a highly collective symmetry breaking pattern, ensuring the finiteness of the generated Higgs potential.

We close this discussion by summarizing the conditions for the presence of maximal symmetry in the following: the elementary-composite mixing must respect the full $SO(5)$ symmetry, while the only source for the composite masses should be twisted by the Higgs parity.

Integrating out the heavy top partner Ψ_+ from the Lagrangian in Eq. (6), we obtain the form factors $\Pi_0^{q,t}$, $\Pi_1^{q,t}$, and M_1^t for the effective Lagrangian of the elementary quarks as in Eq. (4). As expected, maximal symmetry will automatically enforce $\Pi_1^{q,t} = 0$, while the expression for the top mass in this case simplifies to $m_t = c_{+L}c_{+R}(M_Q - M_S)f^2(\sin 2\langle h\rangle/f)/(2M_T M_{T_1})$, where M_T,

M_{T_1} are the top partner masses. We can see that the contribution to the Higgs potential is proportional to the top mass square $V \sim (M_1^t \Sigma')^2 \sim |c_{+L}|^2 |c_{+R}|^2 f^4 (M_Q - M_S)^2 (\sin 2\langle h \rangle / f)^2 / \Lambda^2$, which is finite and has the form as expected from the general symmetry arguments.

Next, we consider the tuning in models with maximal symmetry. We parametrize the potential as usual as

$$V(h) = -\gamma s_h^2 + \beta s_h^4, \tag{9}$$

with the minimum at $\xi \equiv s_h^2 = \gamma / 2\beta$. In generic composite Higgs models the fermionic (f) contribution to γ_f always dominates over that to β_f due to the presence of a nonzero $\Pi_1^{q,t}$ form factor. The reason behind this can be seen by counting the powers of couplings appearing in the form factors yielding $\Pi_1^{q/t} \sim |\epsilon_{L/R}|^2 / g_f^2$, $M_1^t \sim \epsilon_L \epsilon_R f / g_f$, where the characteristic top partner mass is $M_f = g_f f$, and $\epsilon \sim \epsilon_{L/R}$ is a characteristic Yukawa coupling. As a result, the contributions to $\gamma_f = [2N_c/(4\pi)^2] \int dp^4 (\Pi_1^q - \Pi_1^t + |M_1^t|^2/p^2)$ are $\mathcal{O}(\epsilon^2/g_f^2)$ while β_f are $\mathcal{O}(\epsilon^4/g_f^4)$. This statement holds for both UV divergent and finite contributions, and leads to the so-called double tuning of the Higgs potential: the tuning needed to obtain a realistic electroweak symmetry breaking sector is $\Delta \simeq g_f^2/\xi\epsilon^2$ [24], a factor of g_f^2/ϵ^2 larger than the irreducible $1/\xi$ minimal tuning. Double tuning is a generic problem even for holographic composite Higgs models [25] based on a warped extra dimension and their deconstructed versions [26–29]: these models do regulate the UV behavior of $\Pi_1^{q/t}$ to yield a log divergent or finite Higgs potential; however, double tuning is present. Maximal symmetry offers an elegant and completely novel solution for eliminating double tuning: maximal symmetry implies the vanishing of the entire $\Pi_1^{q,t}$ form factor, thereby eliminating all $\mathcal{O}(\epsilon^2/g_f^2)$ contributions to the Higgs potential. As a result, both $\gamma_f, \beta_f = \mathcal{O}(\epsilon^4/g_f^4)$, yielding a potential with $\beta_f = \gamma_f$ and a minimum at $\xi \simeq 0.5$, implying that the tuning will be $\Delta \simeq 1/2\xi$, which is the minimal universal tuning necessary for a small ξ. In order to actually achieve this tuning and reduce ξ to experimentally allowed values, one needs to include gauge contributions and impose a cancellation between the fermionic and gauge contributions of the γ terms $\gamma_f \simeq -\gamma_g$ [while β_g is at order $\mathcal{O}(g^4/g_\rho^4)$, which is always negligible compared to β_f].

Another simple way to see that maximal symmetry implies minimal tuning is to realize that maximal symmetry will imply the existence of an additional Z_2 symmetry in the Higgs potential corresponding to the $s_h \Leftrightarrow -c_h$ exchange, analogous to the case of twin Higgs models [30]. This symmetry will forbid the $\epsilon^2 s_h^2$ term (similar to composite twin Higgs models [31–33]) and eliminate the double tuning.

For general composite Higgs models one usually needs some additional tuning to get the Higgs mass down to 125 GeV. However, the model with maximal symmetry has the special property that the top mass is maximized: $m_t \sim \sin\theta_L \sin\theta_R\, |M_Q - M_S| s_h$, where θ_L and θ_R are the degrees of LH and RH top compositeness. Since maximal symmetry implies $M_Q = -M_S$, the $|M_Q - M_S|$ factor is maximized; hence, the degree of compositeness can be minimized while the top mass is held fixed at the physical value. This also implies that the mass of the lightest top partner $\min\{M_T, M_{T_1}\} = \min\{M_S/\cos\theta_L, M_Q/\cos\theta_R\}$ is also automatically reduced, which in turn cuts off the top contribution to the Higgs mass earlier, $m_H \propto \min\{M_T, M_{T_1}\} m_t/f$, and allows us to obtain a light 125 GeV Higgs mass in the maximally symmetric limit.

To explicitly verify our estimates, we have numerically evaluated the tuning in the model of Eq. (7) with maximal symmetry where we have used the measure of tuning from Ref. [24]. To obtain the contribution of the gauge sector one can extend the concept of maximal symmetry by imposing that the vector meson ρ_μ and the axial-vector meson a_μ form a full adjoint representation of $SO(5)$ (which again automatically renders the Higgs potential finite). The analytic expression for the maximal amount of tuning is $\Delta_m \simeq 1/\xi - 2$. The numerical values of the tuning are shown in Fig. 1. We can clearly see that the largest tuning is from m_ρ, which is from the requirement $\gamma_f \simeq -\gamma_g$ and is slightly smaller than 8 for $\xi = 0.1$ because corrections beyond those at $\mathcal{O}(c_{+L}^2 c_{+R}^2 / g_f^4)$ can also contribute to β_f, making it slightly smaller than γ_f. We also show scatter plots of the Higgs and top partner masses in Fig. 2 for the maximally symmetric MCHM.

The main consequence of maximal symmetry is the vanishing of the form factors $\Pi_1^{q,t} = 0$, which is what one would like to check experimentally by testing the properties of the top Yukawa coupling. For the MCHM the top Yukawa coupling is parametrized by

$$\mathcal{L}_Y \sim M_1^t \sin\frac{2h}{f}\left(1 + (\alpha_q \Pi_1^q + \alpha_t \Pi_1^t) \sin^2\frac{h}{2f}\right)\bar{t}t. \tag{10}$$

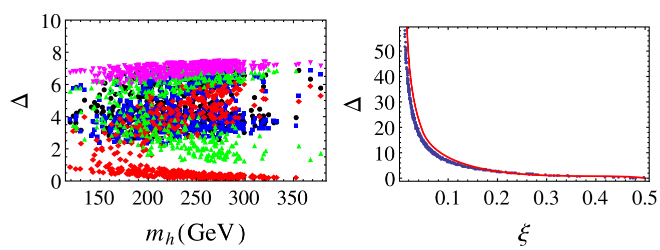

FIG. 1. Left: Scatter plot of tuning Δ_i for the various input parameters x_i, c_{+R} (black), c_{+L} (blue), f (red), M_S (green), and m_ρ (magenta), as a function of m_h with $\xi = 0.1$ held fixed. Right: The tuning Δ_m as a function of ξ for Higgs mass $m_h = 125$ GeV. The red solid line is the analytic expression of Δ_m.

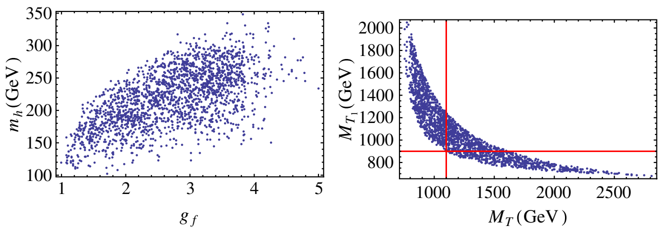

FIG. 2. Scatter plots for the Higgs mass as a function of g (left) and the top partner masses (right) for $\xi = 0.1$. The chosen range of the parameters is $m_t \in [150, 170]$ GeV, $m_\rho \geq 2$ TeV. The horizontal and vertical red lines, corresponding to 900 and 1100 GeV, respectively, are the lower bounds of the doublet and singlet top partners from the most recent 13 TeV LHC data [34–36].

A test of maximal symmetry would be to precisely measure the various $t\bar{t}h^n$ couplings at future colliders and establish whether it can arise from a single trigonometric function or not. On the other hand, the ggh coupling from the top and top partner loops never gets corrected from $\Pi_1^{q,t}$, so its ratio over the SM value would be the same as the $t\bar{t}h$ coupling in the maximally symmetric limit [37]. Another way to test maximal symmetry is via the properties of the additional resonances if they are within the reach of the LHC (or future colliders). One can then derive sum rules for the conditions of the cancellation of the quadratic and log divergences in the Higgs potential, similar to those in little Higgs models [38]. For example, for the case of the top partners, we obtain the sum rules [39] $\mathrm{Tr}[Y_m M_D] = \mathrm{Tr}[Y_m M_D^3] = 0 + \mathcal{O}(v^2/M_f^2)$, where Y_m is the Yukawa coupling matrix of the top partners and the top quark while M_D is their mass matrix. Measuring the masses and couplings of all charge 2/3 top partner resonances, one can test these sum rules and thereby maximal symmetry, following the methods of Ref. [38].

In this Letter, we introduced the new concept of maximal symmetry. This symmetry is a leftover from the chiral symmetries of the fermion sector of composite Higgs models, involving a twisting by the nonlinear Higgs field. Maximal symmetry has powerful consequences: it renders the contributions to the Higgs potential finite, and also ensures that the tuning is minimal. We have presented a realistic model with maximal symmetry, evaluated its spectrum and the tuning, and commented on possible avenues to test the presence of the maximal symmetry experimentally. This novel symmetry breaking pattern has many other potential applications. For example, one can use it to build composite twin Higgs models with an unbroken Z_2 and no tuning of the Higgs potential at all [40], one can recast the UV Higgs parity operator to be an adjoint VEV providing a UV completion in terms of a "twisted moose" model [41], we can try to apply the maximal symmetry to the gauge sector to explain the vector meson spectra, or one can use maximal symmetry to obtain improved natural inflation models [42,43]. We expect these many possible applications to trigger interest for more explorations of the structure, formalism, and phenomenology of maximal symmetry.

We thank Brando Bellazzini, Tao Liu, Michael Geller, and Javi Serra for useful discussions and comments, and Tao Liu for sharing a preliminary version of Ref. [39] containing the sum rule. C. C. thanks the KITPC in Beijing for its hospitality while this project was initiated, and C. C. and J. S. thank the Aspen Center for Physics while this this work was in progress. C. C. is supported in part by the NSF Grant No. PHY-1316222. J. S. is supported by the NSFC under Grants No. 11647601, No. 11690022, and No. 11675243, and also supported by the Strategic Priority Research Program of the Chinese Academy of Sciences under Grants No. XDB21010200 and No. XDB23030100. T. M. is supported in part by Project No. Y6Y2581B11 supported by 2016 National Postdoctoral Program for Innovative Talents.

*Corresponding author.
jshu@itp.ac.cn

[1] G. Aad et al. (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] S. Chatrchyan et al. (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] V. L. Ginzburg and L. D. Landau, Zh. Eksp. Teor. Fiz. **20**, 1064 (1950).
[4] D. B. Kaplan and H. Georgi, Phys. Lett. **136B**, 183 (1984).
[5] H. Georgi and D. B. Kaplan, Phys. Lett. **145B**, 216 (1984).
[6] M. J. Dugan, H. Georgi, and D. B. Kaplan, Nucl. Phys. **B254**, 299 (1985).
[7] N. Arkani-Hamed, A. G. Cohen, and H. Georgi, Phys. Rev. Lett. **86**, 4757 (2001).
[8] N. Arkani-Hamed, A. G. Cohen, E. Katz, and A. E. Nelson, J. High Energy Phys. 07 (2002) 034.
[9] Z. Chacko, H. S. Goh, and R. Harnik, Phys. Rev. Lett. **96**, 231802 (2006).
[10] G. F. Giudice, C. Grojean, A. Pomarol, and R. Rattazzi, J. High Energy Phys. 06 (2007) 045.
[11] For a review, see B. Bellazzini, C. Csaki, and J. Serra, Eur. Phys. J. C **74**, 2766 (2014); G. Panico and A. Wulzer, Lect. Notes Phys. **913**, 1 (2016).
[12] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990); for recent constraints, see, for example, M. Baak, J. Cúth, J. Haller, A. Hoecker, R. Kogler, K. Mönig, M. Schott, and J. Stelzer (Gfitter Group), Eur. Phys. J. C **74**, 3046 (2014).
[13] G. Aad et al. (ATLAS Collaboration), J. High Energy Phys. 08 (2015) 105.
[14] V. Khachatryan et al. (CMS Collaboration), Phys. Rev. D **93**, 012003 (2016).
[15] S. R. Coleman and E. J. Weinberg, Phys. Rev. D **7**, 1888 (1973).
[16] See, for example, R. Contino, D. Marzocca, D. Pappadopulo, and R. Rattazzi, J. High Energy Phys. 10 (2011) 081; B. Bellazzini, C. Csaki, J. Hubisz, J. Serra, and J. Terning, J. High Energy Phys. 11 (2012) 003.

[17] This is the form of the automorphism for $SU(N+M)/[SU(N)\times SU(M)\times U(1)]$ and $SO(N+M)/[SO(N)\times SO(M)]$.

[18] S. R. Coleman, J. Wess, and B. Zumino, Phys. Rev. **177**, 2239 (1969).

[19] C. G. Callan, Jr., S. R. Coleman, J. Wess, and B. Zumino, Phys. Rev. **177**, 2247 (1969).

[20] In fact, the global symmetry at this point is enhanced to $(G_L\times G_R)^2$, but since we later break it into a single copy of the chiral symmetry, this additional doubling has no important physical consequences.

[21] K. Agashe, R. Contino, and A. Pomarol, Nucl. Phys. **B719**, 165 (2005).

[22] A. De Simone, O. Matsedonskyi, R. Rattazzi, and A. Wulzer, J. High Energy Phys. 04 (2013) 004.

[23] D. Marzocca, M. Serone, and J. Shu, J. High Energy Phys. 08 (2012) 013.

[24] G. Panico, M. Redi, A. Tesi, and A. Wulzer, J. High Energy Phys. 03 (2013) 051.

[25] R. Contino, Y. Nomura, and A. Pomarol, Nucl. Phys. **B671**, 148 (2003).

[26] H. C. Cheng, J. Thaler, and L. T. Wang, J. High Energy Phys. 09 (2006) 003.

[27] R. Foadi, J. T. Laverty, C. R. Schmidt, and J. H. Yu, J. High Energy Phys. 06 (2010) 026.

[28] G. Panico and A. Wulzer, J. High Energy Phys. 09 (2011) 135.

[29] S. De Curtis, M. Redi, and A. Tesi, J. High Energy Phys. 04 (2012) 042.

[30] The additional Z_2 symmetry is defined as $\Psi_{+L}\to P_1\Psi_{+L}$, $\Psi_{+R}\to VP_1V\Psi_{+R}$, $U\to VUVP_1V$, $\Psi_{q_L}\to V\Psi_{q_L}=\Psi_{q_L}$, $\Psi_{t_R}\to P_2\Psi_{t_R}=\Psi_{t_R}$, where $P_1=\mathrm{diag}(1_{3\times3},\sigma_1)$, $P_2=\mathrm{diag}(1_{3\times3},-\sigma_3)$.

[31] M. Geller and O. Telem, Phys. Rev. Lett. **114**, 191801 (2015); C. Csaki, M. Geller, O. Telem, and A. Weiler, J. High Energy Phys. 09 (2016) 146.

[32] R. Barbieri, D. Greco, R. Rattazzi, and A. Wulzer, J. High Energy Phys. 08 (2015) 161.

[33] M. Low, A. Tesi, and L. T. Wang, Phys. Rev. D **91**, 095012 (2015).

[34] ATLAS Collaboration, Report No. ATLAS-CONF-2016-104.

[35] ATLAS Collaboration, Report No. ATLAS-CONF-2016-101.

[36] ATLAS Collaboration, Report No. ATLAS-CONF-2016-102.

[37] D. Liu, I. Low, and C. E. M. Wagner, Phys. Rev. D **96**, 035013 (2017).

[38] G. Burdman, M. Perelstein, and A. Pierce, Phys. Rev. Lett. **90**, 241802 (2003); M. Perelstein, M. E. Peskin, and A. Pierce, Phys. Rev. D **69**, 075002 (2004); T. Han, H. E. Logan, and L. T. Wang, J. High Energy Phys. 01 (2006) 099.

[39] C. R. Chen, J. Hajer, T. Liu, I. Low, and H. Zhang, arXiv:1705.07743.

[40] C. Csaki, T. Ma, and J. Shu, (to be published).

[41] C. Csaki, T. Ma, and J. Shu, (to be published).

[42] N. Arkani-Hamed, H. C. Cheng, P. Creminelli, and L. Randall, J. Cosmol. Astropart. Phys. 07 (2003) 003.

[43] M. Bastero-Gil, A. Berera, R. O. Ramos, and J. G. Rosa, Phys. Rev. Lett. **117**, 151301 (2016).