Electroweak Symmetry Breaking via QCD

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We propose a new mechanism to generate the electroweak scale within the framework of QCD, which is extended to include conformally invariant scalar degrees of freedom belonging to a larger irreducible representation of $SU(3)_c$. The electroweak symmetry breaking is triggered dynamically via the Higgs portal by the condensation of the colored scalar field around 1 TeV. The mass of the colored boson is restricted to be 350 GeV $\lesssim m_S \lesssim$ 3 TeV, with the upper bound obtained from perturbative renormalization group evolution. This implies that the colored boson can be produced at the LHC. If the colored boson is electrically charged, the branching fraction of the Higgs boson decaying into two photons can slightly increase, and moreover, it can be produced at future linear colliders. Our idea of nonperturbative electroweak scale generation can serve as a new starting point for more realistic model building in solving the hierarchy problem.

Introduction.—With only the standard model (SM) Higgs-boson-like particle discovered and no new particle beyond the SM being found, there is no evidence for any of the generally proposed solutions to the hierarchy problem. With the currently measured Higgs boson mass and top quark mass, the SM could even survive up to the Planck scale [1]. However, one has to face the puzzle of why the electroweak (EW) scale is many orders of magnitude smaller than the Planck scale. A possible solution for the hierarchy problem is based on scale invariance, which is violated at the quantum level, and hence, a scale is introduced: The EW scale is generated dynamically by either the Coleman-Weinberg mechanism [2] or dimensional transmutation of a nonperturbatively created scale in a strongly coupled hidden sector [3]. Many of these attempts to generate the EW scale radiatively rely on the Higgs portal $\lambda_{HS} S^\dagger S H^\dagger H$, where the additional scalar field S (charged or neutral under a certain gauge group) obtains a vacuum expectation value (VEV) either directly or indirectly. In this Letter, we propose a new nonperturbative mechanism to generate the EW symmetry-breaking (EWSB) scale. Though the hierarchy problem between the EW scale and the Planck scale is not completely solved in our proposed minimal model, which is a least extension of the SM, our mechanism can be applied to more realistic model building scenario in solving the hierarchy problem. Specifically, EWSB is triggered by the condensation of an additional scalar field S, which belongs to a larger representation of $SU(3)_c$. In general, the condensation of S, i.e., $\langle S^\dagger S \rangle \neq 0$, takes place when

$$C_2(S)\alpha(\Lambda) \gtrsim 1, \qquad (1)$$

with C_2 representing the quadratic Casimir operator of a certain representation $\boldsymbol{R}$ of S and α the gauge coupling of the chosen non-Abelian gauge group. The crucial point to notice here is that confinement (we throughout assume that the confinement scale is the same as the condensation scale) can take place even if α is relatively small, provided that the representation of S is large enough. QCD is a part of the SM and generates dynamically an energy scale of $\mathcal{O}(1 \text{ GeV})$ by the gluon and quark condensates. However, we emphasize that these scales are closely related to the fact that the quarks belong to the fundamental representation of $SU(3)_c$. Therefore, according to Eq. (1), if there exist colored degrees of freedom belonging to a larger representation of $SU(3)_c$, QCD can generate a much higher energy scale in principle. In fact, exotic quarks that are confined at a higher energy scale have been considered in Ref. [4]. However, most of these exotic fermions with EW charges cannot generate the correct EW scale without large deviations from EW precision tests. This situation will change if we consider a colored EW singlet scalar field, as we will see in the next sections.

Electroweak symmetry breaking by scalar QCD.—We assume that the SM with the new scalar QCD extension is classically scale invariant and the EW scale is generated via the condensation scale of S. In fact, the $SU(3)_c$ sector of the SM itself before EWSB is scale invariant, contrary to ordinary QCD with explicit massive quarks. The full Lagrangian is given as

$$\mathcal{L} = \mathcal{L}_{\text{SM},m^2\to 0} + (D_{\mu,ij}S_j)^\dagger(D_{ik}^\mu S_k) + \lambda_{HS}H^\dagger H S^\dagger S - \lambda_{\mathbf{1}_i}[\bar{\boldsymbol{S}} \times \boldsymbol{S} \times \bar{\boldsymbol{S}} \times \boldsymbol{S}]_{\mathbf{1}_i}, \qquad (2)$$

where $D_{ij}^\mu = \delta_{ij}\partial^\mu - ig_s(T_R)_{ij}^k G_k^\mu$ and T_R represents the generator for the representation $\boldsymbol{R}$ of $SU(3)_c$. The term $\lambda_{\mathbf{1}_i}$

denotes the quartic scalar coupling for the ith invariant formed by the four tensor products of the S representation. Because of classical scale invariance, the Lagrangian in Eq. (2) does not contain quadratic and cubic terms of S. Conventional scalar QCD would be quadratically sensitive to an embedding scale, and it would, therefore, not solve the hierarchy problem. Note, however, that our scenario is based on conformal QCD, which should have only logarithmic scale dependence. Note that accidental U(1) symmetry appears for the S sector due to the absence of a cubic term, and this has interesting phenomenology on its own if this U(1) symmetry is identified with the $U(1)_Y$ hypercharge of the SM, which we will discuss later. EWSB triggered by QCD is as follows: The strong coupling g_s runs as usual from a finite value set at high energy (Planck or GUT) scale to the condensation scale of S. The scalar condensate $\langle S^\dagger S \rangle$ forms when Eq. (1) is satisfied for $\mathcal{O}(1 \text{ TeV})$ where the small value of $\alpha_s(\Lambda = 1 \text{ TeV}) \approx 0.09$ is compensated by the large C_2 of S in higher representation. Note that the confinement scale is fixed once a representation for S is chosen; see Table I. The condensate generates a scale that enters the portal

$$\lambda_{HS}\langle S^\dagger S \rangle H^\dagger H \rightarrow \lambda_{HS}\Lambda^2 H^\dagger H \qquad (3)$$

and triggers the EWSB radiatively. The Higgs mass after EWSB is determined by

$$m_h^2 = 2\lambda_{HS}\Lambda^2, \qquad (4)$$

and this in turn determines the value of Higgs quartic coupling $\lambda_h = \lambda_{HS}\Lambda^2/v^2$, with the Higgs VEV $v = 246$ GeV. The coupling λ_{HS} is determined once the confinement scale is fixed to be any value higher than the EW scale, as we require that confinement happens before EWSB. In general, we have no upper bound on Λ, except that larger representation of S is required as α_s decreases with a higher value of Λ.

The low-energy QCD remains unaltered by our new additional field as the coupling of higher representation of field S with the quarks in fundamental representation to form a singlet requires typically higher-dimensional operators. It is important to remember that such condensation takes place despite the small coupling of α_s at scales of

TABLE I. Values of the quadratic Casimir operator and index for certain representations of QCD. The approximate confinement scale Λ for each representation is listed.

Rep ($\boldsymbol{R}$)	$C_2(\boldsymbol{R})$	$C(\boldsymbol{R})$	Λ (GeV)
8	3	3	1
10	6	15/2	20
15	16/3	10	10
15′	28/3	35/2	1000
21	40/3	35	10^5

$\mathcal{O}(\text{TeV})$ due to a large C_2 value for larger representation. As we can read off from Table I, **15′** is the unique representation for our purpose as it generates the desired condensation scale at $\mathcal{O}(1 \text{ TeV})$.

The phenomenology of this new scalar QCD extension with the representation of S being **15′** will now be discussed in detail. First, we can constrain the coupling λ_{1_i} and λ_{HS} from the requirement that all the scalar couplings do not hit a Landau pole or destabilize the vacuum. For the case of **15′**, we have three quartic couplings λ_{1_i} due to the existence of three invariants formed from the four tensor products of **15′**. The invariants formed by the tensor products can be calculated with proper Clebsch-Gordan coefficients, and subsequently, one-loop beta functions for the quartic couplings can be calculated [5]. To simplify our calculation further, we assume that the order of each λ_{1_i} is roughly the same, i.e., $\lambda_{1_i} \approx \lambda_S/3$ and normalized such that the mass term m_S of S can be extracted from the Lagrangian. Notice that the bare m_S of S does not exist in Eq. (2) due to scale invariance. The mass term can be approximately obtained from self-consistent mean field approximation [6] after confinement has taken place, where the mean field serves as a back-reaction to the field S and the mass is obtained from

$$\frac{\lambda_S}{2}(S^\dagger S)(S^\dagger S) \rightarrow \lambda_S\langle S^\dagger S \rangle S^\dagger S = \lambda_S\Lambda^2 S^\dagger S. \qquad (5)$$

The coupling λ_S dictates directly $m_S^2 = \lambda_S\Lambda^2$ while the mixing parameter λ_{HS} determines m_h. The large m_S value prevents the S field from obtaining a nonzero VEV; hence, color symmetry is not spontaneously broken. From the renormalization group equation (RGE) analysis, we obtain the running of scalar couplings once the confinement scale is set. The measured m_h fixes λ_{HS}, while the mass $m_S \sim \lambda_S$ cannot be pushed arbitrarily high due to the emergence of a Landau pole, yielding an upper bound on $m_S \lesssim 3$ TeV while the lower bound can be obtained from the collider phenomenology. The running of scalar mixing parameter λ_{HS} is relatively slow, and it will only hit the triviality bound when λ_S hits the Landau pole; this subsequently drives λ_h to a Landau pole. We would like to stress that other RGE scenarios maybe viable if the parameters λ_{1_i} and the confinement scale Λ are varied independently. In this Letter, we study only the simplest model to explain EWSB triggered by QCD. More realistic models should include dark matter and neutrino masses, and their coupling to our new scalar could alter the high UV behavior of the RGE of λ_S significantly. The Landau pole at 10 TeV may, therefore, be absent in a more realistic model or be a signal for nonperturbativity.

Collider phenomenology.—The scalar S can change the Higgs production rate in the gluon fusion channel due to λ_{HS}. We have calculated $\sigma(pp \rightarrow H)$ to the next-to-leading (NLO) order with this additional scalar. We followed the calculation of Ref. [7] and utilize the heavy scalar
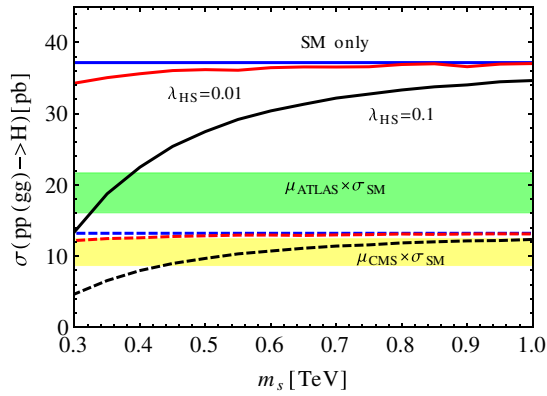
FIG. 1 (color online).   The Higgs production cross section from gluon fusion channel at NLO is calculated for different values of λ_{HS}. The solid (dashed) curves represent the prediction of $\sigma(gg \to H)$ at $\sqrt{s} = 14$ TeV (8 TeV). The combined signal strength μ for ATLAS [10] and CMS [11] is shown where we have assumed a SM-like BR.

approximation. The MSTW2008 parametrization of parton density function (PDF) [8] implemented in LHAPDF [9] has been used in our computation with the factorization scale μ_F and the renormalization scale μ_R set to be equal to m_h. We have also utilized the zero-width approximation for the Higgs boson to simplify the calculation, and the resulting production cross section is shown in Fig. 1. Since our model does not modify the branching ratio (BR) of the SM Higgs boson (the decay $H \to \gamma\gamma$ is modified with accidental symmetry, which we will discuss later on, but this loop-induced decay is very small compared to the tree-level decays), we can compare the signal strength μ times $\sigma(pp \to H)_{SM}$ measured by ATLAS [10] and CMS [11] to our model's prediction. The additional S field decreases the Higgs gluon fusion production rate, with almost half the rate for large λ_{HS} (small Λ) and small m_S. We obtain the suppression of ggH production rate as opposed to the enhancement due to the negative sign of λ_{HS}.

The condensate $\langle S^\dagger S \rangle$ has to be heavier than the Higgs boson to trigger the EWSB; therefore, it will decay to Higgs particles or two gluons. The scalar S can be produced at the LHC, with the dominating production channel $gg \to S_i^* S_j$. The pair production of colored scalars with higher-dimensional representation at LO in the gluon fusion channel has been calculated in Ref. [12], and the result for our case is given in Fig. 2. The resulting particles $S_i^* S_j$ will form two bound state pairs, with each pair decaying predominantly to gg (2 jets) or to Higgs particles. Since the BR of $H \to b\bar{b}$ dominates, we would expect almost 70% for $S^* S \to jjjj$ in the total cross section. The width of the band in Fig. 2 represents the factorization and renormalization scale dependence and the α_s uncertainty from RGE with extra S contribution. In Fig. 2, we plot the ATLAS exclusion limit on pair production of new color scalar decaying to four jets [13], where we have assumed
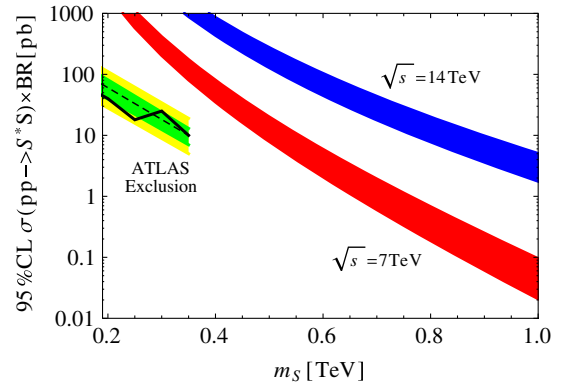

FIG. 2 (color online).   The S pair production cross section from gluon fusion channel is calculated for different values of m_S. The 95% confidence level exclusion limit on $\sigma \times$ BR for $\sqrt{s} = 7$ TeV by ATLAS is plotted. We assume 100% BR of $\langle S^\dagger S \rangle$ into two jets.

100% BR to four jets. $m_S \lesssim 350$ GeV is excluded at the 95% confidence level and serves as our lower bound on m_S. In combining this result with the upper bound due to the triviality constraint above, the mass parameter of this model is very constrained, i.e.,

$$350 \text{ GeV} \lesssim m_S \lesssim 3 \text{ TeV}. \tag{6}$$

The S field in Eq. (2) possesses an accidental U(1) symmetry due to the absence of the cubic term as we have imposed classical scale invariance in the Lagrangian. A priori, this U(1) is another global symmetry, but if it is identified with the local $U(1)_Y$ of the SM, we would obtain more interesting phenomenology. For instance, the $H \to \gamma\gamma$ channel is enhanced by the additional S running in the loop. Contrary to other scalar extension, enhancement of $H \to \gamma\gamma$ is obtained instead of suppression due to the minus sign of λ_{HS} [14]. Strong enhancement of signal strength $\mu_{\gamma\gamma}$ for different values of m_S can be obtained, with the result normalized to the SM prediction shown in Fig. 3. The signal strength $\mu_{\gamma\gamma}$ can be only enhanced by increasing the electric charge or λ_{HS} of S to compensate the suppression of production cross section. Compared to the $\mu_{\gamma\gamma} \approx 1.65$ (0.77) value reported by ATLAS [15] (CMS [11]) with the average $\mu_{\gamma\gamma} \approx 1.21$, our model would require large electric charge to explain the large $H \to \gamma\gamma$ anomaly. The large electric charge provides a possible alternative to study the S particle via Drell-Yan production in a linear collider.

Confinement of strongly coupled scalar field.—So far, we have discussed the perturbative sector of the colored scalar S. We restricted the nonperturbative aspect of the model to the upscaling of the gap equation in Eq. (1). Let us discuss a bit the physics in Eq. (1). An analytical way to understand confinement in the quarks sector of QCD is to calculate the scaling of the gap equation from the Dyson-Schwinger equation (DSE)
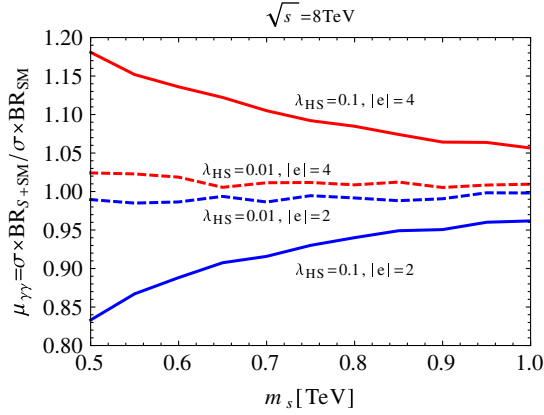
FIG. 3 (color online). The signal strength of the $H \to \gamma\gamma$ branching ratio with the additional S contribution relative to the SM prediction are plotted for different values of electric charge e and λ_{HS} of S. The large electric charge has to compensate the suppression of production cross section for $\mu_{\gamma\gamma}$ enhancement.

where we have utilized the rainbow-ladder approximation and only kept the leading-order contribution to our analysis. The diagram above resembles the DSE for quark propagator, which can be solved within a certain truncation scheme in order to obtain the critical value X in

$$C_2(S)\alpha(\Lambda) \gtrsim X, \qquad (8)$$

for confinement to take place. However, there are subtleties that one has to be careful with when trying to extract the exact bound of X. First, the value X is gauge and truncation scheme dependent. Different values ranging from 0.6 to $\pi/3$ have been obtained [16–18]. Lowering X will allow us to consider a lower representation of S, but in our analysis above we assume the conservative bound $X > 0.8$. Second, the DSE for quark is linearizable with its linear form a Fredholm integral equation [16,19,20] as the wave function renormalization part, and the self-energy part can be dealt separately in Landau gauge. However, such a privilege is not enjoyed by the scalar DSE, as the integral equation

$$F(p^2) = p^2 + \frac{3C_2\alpha_s}{4\pi}\left[\int_0^{p^2} dq^2 \frac{q^4}{p^2 F(q^2)} + \int_{p^2}^{\infty} dq^2 \frac{p^2}{F(q^2)}\right], \qquad (9)$$

is not linearizable, where we have denoted the function $F(k^2) = Z(k^2)k^2 + \Sigma^2(k^2)$. The main reason for such a difficulty is due to the lack of confinement order parameter

for scalar QCD. In comparison to fermionic QCD, the order parameter for confinement can be related to the degree of chiral symmetry breaking. From the perturbative calculation of the anomalous dimension of operator $\langle \bar\psi\psi \rangle$ and $\langle S^\dagger S \rangle$ in the same representation, it can be shown that

$$\gamma_{\langle \bar\psi\psi \rangle} = \gamma_{\langle S^\dagger S \rangle} + \mathcal{O}(\lambda_S). \qquad (10)$$

Hence, one can conjecture that the relevant order parameter $C_2\alpha_s$ at leading order for determining confinement should be the same for both fermionic QCD and scalar QCD, which we have assumed. In fact, it has been argued that the scaling property for scalar and quark propagator in the infrared range is identical [21]. This result can be verified in lattice QCD.

Note that the QCD coupling becomes nonperturbative in the TeV regime even though the coupling is pretty small. This stems from the large value of C_2 that is responsible for condensation. As a consequence, the exact evolution of α_s cannot be precisely calculated in the TeV regime. However, the coupling may become perturbative again at sufficiently small α_s or high energy. A similar conclusion was made in Ref. [4]. Measuring α_s at high energy will provide an independent test for our model.

Conclusion.—With no signature of any SM extension at the LHC and in other searches, the notion of naturalness deserves to be reexamined, and other ideas of explaining the EW scale should be considered. We discussed in this Letter a scenario where conformal symmetry plays an essential role and where the EW scale is a consequence of quantum effects. The idea of mass scale generation from a quantum effect, so-called dimensional transmutation, is already implemented in the QCD sector of the SM. We have shown that it is possible to extend the success of QCD and to explain the existence of the EW scale by including a new scalar particle belonging to $\mathbf{15}'$ of $SU(3)_c$. The extension is rather minimal and, moreover, unique if X in Eq. (8) is greater than 0.8. The mass of this new colored boson is constrained such that it can be explored or ruled out by the LHC. The signature of this colored scalar field is comparatively clean. The accidental U(1) symmetry can also be probed in the $H \to \gamma\gamma$ signal strength if it is identified with $U(1)_Y$ of the SM. Furthermore, with a nonzero hypercharge the new colored boson can be directly produced at linear colliders, which will be our next target to investigate. We leave the more detailed investigation of the nonperturbative aspect of this model and the implication of EW phase transition to future work.

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