Experimentally distinguishable origin for electroweak symmetry breaking

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Abstract

We consider a classically conformal $U(1)$ gauge extension of the Standard Model (SM), in which the $U(1)$ gauge symmetry is radiatively broken by the Coleman-Weinberg mechanism. This breaking triggers the electroweak (EW) symmetry breaking through a mixed quartic coupling between the $U(1)$ Higgs field and the SM Higgs doublet. For two Higgs boson mass eigenstates after the symmetry breaking, $h_1$ (SM-like Higgs boson) and $h_2$ (SM singlet-like Higgs boson), we calculate the Higgs boson trilinear coupling ($g_{h_1 h_2 h_2}$) in the model by setting the Higgs boson mass spectrum to be $M_{h_1} > 2M_{h_2}$. For a common Higgs mass spectrum and mixing angle between two Higgs fields, we find that $g_{h_1 h_2 h_2}$ in the classically conformal model is highly suppressed compared to that calculated for the conventional Higgs potential, where the $U(1)$ and EW symmetry breaking originate from the negative squared masses for the Higgs fields at the tree-level. Thus, this coupling suppression is a striking nature of the radiative origin of EW symmetry breaking. We then consider how to distinguish this origin at the proposed International Linear Collider (ILC) via precise measurements of anomalous SM Higgs boson couplings and the search for anomalous SM Higgs boson decay $h_1 \rightarrow h_2 h_2$ followed by $h_2 \rightarrow b\bar{b}$. We conclude that once the anomalous couplings are measured at the ILC, the observation of the anomalous Higgs boson decay is promising in the conventional Higgs potential, while this decay process is highly suppressed and undetectable for the classically conformal model.

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I. INTRODUCTION

The existence of problems not addressed or answered by the Standard Model (SM) has motivated many searches for physics beyond the SM. To that end, many extensions of the SM have been proposed, involving various additional particles and/or interactions providing possible solutions to some or all of these problems. Among these problems is the origin of the electroweak (EW) symmetry breaking. While successful, there is no compelling explanation for how it arises, since the negative mass-squared term for the Higgs field is generally included by hand.

One interesting way to break a gauge symmetry is the so-called radiative symmetry breaking first proposed by Coleman and Weinberg [1], where classically conformal symmetry is imposed at the tree-level scalar potential, and radiative corrections to the potential from the gauge interaction trigger symmetry breaking. The origin of the symmetry breaking is thus due to quantum corrections, so no negative mass-squared term is added by hand. However, the Coleman-Weinberg mechanism cannot be naively applied to the SM Higgs sector, as the SM Higgs potential is unbounded from below by quantum corrections from the large top quark Yukawa coupling (see for example Ref. [2]). For this scheme to be implemented, extending the SM is necessary. A simple way to do so is to extend the SM with a new $U(1)$ gauge symmetry and to impose classical conformality to forbid explicit mass-squared terms [3, 4]. The new $U(1)$ gauge symmetry is then broken at some scale $v_\phi$ by the Coleman-Weinberg mechanism, where $v_\phi$ is the vacuum expectation value (vev) for a Higgs field $\Phi$ for the $U(1)$ gauge symmetry. A negative mixing quartic coupling between the SM Higgs doublet and $\Phi$ generates the negative mass squared term for the SM Higgs doublet, driving EW symmetry breaking. Therefore, the symmetry breaking of the new $U(1)$ gauge symmetry is the sole origin of the mass scale in this model. This simple mechanism can easily be incorporated into more complex SM extensions.

As the extension includes a new scalar, it is interesting to investigate Higgs boson phenomenology. We specifically consider the case where the extra Higgs boson is light enough to be produced by SM-like Higgs boson decay. Such a process originates from the mixing quartic coupling that generates the EW symmetry breaking. The same coupling also generates a mass mixing between the SM Higgs and $U(1)$ Higgs fields. Due to this mixing, the model predicts anomalous SM Higgs boson couplings.

In this paper, we consider a classically conformal extension to the SM by a (hidden) $U(1)$ gauge group with a Higgs scalar field $\Phi$. The Coleman-Weinberg mechanism induces radiative symmetry breaking of the new $U(1)$ gauge group, generating the negative mass-squared term for the SM Higgs doublet via a mixed quartic coupling between the SM Higgs doublet and $\Phi$. This in turn drives the EW symmetry breaking for the SM. We analyze the Higgs potential and extract a Higgs trilinear coupling $g_{h_1 h_2 h_2}$ between two mass eigenstates, the SM-like Higgs boson $h_1$ and the SM singlet-like Higgs boson $h_2$. For comparison, we also analyze the conventional Higgs potential, where the $U(1)$ and EW gauge symmetry breaking are triggered by negative squared masses for the two Higgs fields introduced by hand and calculate $g_{h_1 h_2 h_2}$. Very interestingly, we find a significant suppression in $|g_{h_1 h_2 h_2}|$ for our classically conformal potential compared to a naively expected value from the conventional Higgs potential while anomalous SM Higgs boson couplings are taken to be the same for both cases. We then consider how to experimentally distinguish the radiative breaking origin of the EW symmetry breaking from that in the conventional Higgs potential, namely, a negative mass squared introduced by hand. We point
out that the future $e^+e^-$ colliders will be a powerful tool for this job by combining a precise measurement of the anomalous SM Higgs boson coupling with the search for an anomalous Higgs boson decay to $h_1 \rightarrow h_2 h_2$ followed by $h_2 \rightarrow b\bar{b}$. If the anomalous SM Higgs boson couplings have been observed, the observation of the anomalous Higgs boson decay is promising for the conventional Higgs potential, while in the classically conformal case, the branching ratio of $h_1 \rightarrow h_2 h_2$ is highly suppressed and undetectable.

II. CLASSICALLY CONFORMAL $U(1)$ EXTENDED STANDARD MODEL

In this section, we first consider a minimal extension to the SM involving only an additional $U(1)$ gauge group, which we will denote by $U(1)_H$. In order to focus solely on Higgs phenomenology, we assume in this section that all SM fields to be neutral with respect to $U(1)_H$ and no additional chiral fermions are involved in the model. We now introduce the $U(1)_H$ Higgs field $\Phi$ to have charge $Q_\Phi = +2$ while it is singlet under the SM gauge group.

Imposing classically conformal invariance, explicit mass terms are forbidden, and the scalar potential at the tree-level is given by

$$V_{\text{tree}} = \lambda_h (H^\dagger H)^2 + \lambda_\phi (\Phi^\dagger \Phi)^2 - \lambda_{\text{mix}} (H^\dagger H)(\Phi^\dagger \Phi),$$

where $H$ is the SM Higgs doublet, and $\lambda_{\text{mix}}$ is chosen to be positive. We make the choice to take $\lambda_{\text{mix}} \ll 1$, effectively separating the SM and “hidden” $U(1)_H$ Higgs sectors from each other. The vacuum of the scalar potential at the tree-level appears at the origin, and EW and $U(1)_H$ symmetries are unbroken. For the $U(1)_H$ Higgs sector, we account for the radiative corrections at one-loop level $^{[1]}$,

$$V_{1\text{-loop}} = \frac{\beta_\phi}{8} \left( \ln \left[ \frac{\phi^2}{v^2} \right] - \frac{25}{6} \right) \phi^4,$$

where $\phi = \sqrt{2} \text{Re}[\Phi]$, $\beta_\phi$ is given by

$$\beta_\phi = \frac{1}{16\pi^2} (20\lambda_\phi^2 + 96g_H^4) \simeq \frac{1}{16\pi^2} (96g_H^4),$$

and $g_H$ is the gauge coupling of the hidden $U(1)_H$. With these corrections, the $U(1)_H$ symmetry is radiatively broken at $\langle \phi \rangle = v_\phi$.

The full potential is

$$V = \frac{\lambda_h}{4} h^4 + \left[ \frac{\lambda_\phi}{4} + \frac{\beta_\phi}{8} \left( \ln \left[ \frac{\phi^2}{v^2} \right] - \frac{25}{6} \right) \right] \phi^4 - \frac{\lambda_{\text{mix}} h^2 \phi^2}{4}$$

where $H = \frac{1}{\sqrt{2}} \begin{pmatrix} h \\ 0 \end{pmatrix}$ with a real scalar $h$ in the unitary gauge. In the presence of nonzero $\langle \phi \rangle$, the scalar $h$ develops a negative mass-squared term from the mixing quartic coupling. Subsequently, EW symmetry is broken at $\langle h \rangle = v_h = 246$ GeV.
Using the stationary conditions, \( \frac{\partial V}{\partial h} \bigg|_{h=\phi} = 0 \) and \( \frac{\partial V}{\partial \phi} \bigg|_{h=\phi} = 0 \), to eliminate \( \lambda_\phi \) and \( \lambda_{mix} \), we express the potential in the following form:

\[
V(h, \phi) = -\frac{1}{4} \left( \frac{m_h^2}{2v_h^2} \right) h^4 + \frac{1}{4} \left( \frac{11}{6} \beta_\phi + \frac{m_\phi^2 v_h^2}{2v_\phi^4} \right) \phi^4 + \frac{\beta_\phi}{8} \left( \ln \left[ \frac{\phi^2}{v_\phi^2} \right] - \frac{25}{6} \right) \phi^4 - \frac{1}{4} \left( \frac{m_h^2}{v_\phi^2} \right) h^2 \phi^2. \tag{5}
\]

where \( m_h^2 \equiv 2\lambda_h v_h^2 \). By shifting the fields \( h \rightarrow h + v_h \) and \( \phi \rightarrow \phi + v_\phi \), we can obtain the mass-squared matrix for the physical states,

\[
M_{sq} = \begin{pmatrix}
m_h^2 - M^2 & m_\phi^2 \\
-m_\phi^2 & m_\phi^2
\end{pmatrix}, \tag{6}
\]

where \( M^2 = m_h^2 \left( \frac{v_h}{v_\phi} \right) \) and \( m_\phi^2 \equiv m_h^2 \left( \frac{v_h}{v_\phi} \right)^2 + v_\phi^2 \beta_\phi \). We diagonalize \( M_{sq} \) by

\[
\begin{pmatrix}
h \\
\phi
\end{pmatrix} = \begin{pmatrix}
\cos(\theta) & \sin(\theta) \\
-\sin(\theta) & \cos(\theta)
\end{pmatrix} \begin{pmatrix}
h_1 \\
h_2
\end{pmatrix}, \tag{7}
\]

with a mixing angle \( \theta \) defined by \( \tan(2\theta) = \frac{2M^2}{m_h^2 - m_\phi^2} \), where \( h_1, h_2 \) are the mass eigenstates with mass eigenvalues \( M_{h_1} \) and \( M_{h_2} \), respectively. We see that for a small \( \theta \) (or equivalently, a small \( \lambda_{mix} \)), the SM Higgs is dominated by the \( h_1 \) mass eigenstate and \( \phi \) by \( h_2 \). In terms of observables (\( M_{h_1}, M_{h_2}, \theta \)), the parameters in Eq. (5) can be expressed as follows:

\[
m_h^2 = \frac{1}{2} \left( M_{h_1}^2 + M_{h_2}^2 + (M_{h_1}^2 - M_{h_2}^2) \cos(2\theta) \right), \tag{8}
\]

\[
m_\phi^2 = \frac{1}{2} \left( M_{h_1}^2 + M_{h_2}^2 - (M_{h_1}^2 - M_{h_2}^2) \cos(2\theta) \right), \tag{9}
\]

\[
v_\phi = \frac{m_\phi^2}{M^2} v_h, \tag{10}
\]

\[
\beta_\phi = \frac{m_\phi^2}{v_\phi^2} - \frac{m_h^2}{v_\phi^2} v_h. \tag{11}
\]

where \( M^2 = \frac{1}{2} \left( M_{h_1}^2 - M_{h_2}^2 \right) \sin(2\theta) \).

We rewrite the potential in terms of the mass eigenstates \( h_1 \) and \( h_2 \), and extract the trilinear coupling \( g_{h_1 h_2 h_3} \equiv \frac{1}{2} \frac{\partial^2 V}{\partial h_1 \partial h_2 \partial h_3} \bigg|_{h_1=h_2=0} \):

\[
g_{h_1 h_2 h_3} = -\frac{M_{h_2}^2 \cos(\theta) \left( M_{h_1}^4 - 5M_{h_1}^2 M_{h_2}^2 + 2M_{h_2}^4 \right.}{v_h \left( M_{h_1}^2 + M_{h_2}^2 + \left( M_{h_1}^2 - M_{h_2}^2 \right) \cos(2\theta) \right)^2} \left. \left( M_{h_1}^4 - 3M_{h_1}^2 M_{h_2}^2 + 2M_{h_2}^4 \right) \cos(2\theta) \right) \middle|_{h_1=h_2=0}. \tag{12}
\]

Assuming \( M_{h_1} > 2M_{h_2} \) and \( \theta \ll 1 \), Eq. (12) reduces to

\[
g_{h_1 h_2 h_3} \approx -\frac{M_{h_2}^2}{2v_h} \left( 1 - 4 \frac{M_{h_2}^2}{M_{h_1}^2} \right) \theta^2. \tag{13}
\]
Note that this trilinear coupling is extremely suppressed compared to the naively expected value derived from the conventional Higgs potential without classical conformal symmetry. To see this, we now look at the conventional Higgs potential.

Let us consider the conventional Higgs potential for \( h \) and \( \phi \) in the unitary gauge:

\[
V = \frac{1}{4} \lambda_h (h^2 - v_h^2)^2 + \frac{1}{4} \lambda_\phi (\phi^2 - v_\phi^2)^2 + \frac{1}{4} \lambda_{mix} (h^2 - v_h^2)(\phi^2 - v_\phi^2). \tag{14}
\]

The potential minimum appears at \( \langle h \rangle = v_h \) and \( \langle \phi \rangle = v_\phi \). Using the stationary conditions, we rewrite the potential in terms of physical states by shifting \( h \to h + v_h \) and \( \phi \to \phi + v_\phi \):

\[
V = \frac{1}{4} \left( \frac{m_h^2}{2v_h^2} \right) (h^2 + 2hv_h)^2 + \frac{1}{4} \left( \frac{m_\phi^2}{2v_\phi^2} \right) (\phi^2 + 2\phi v_\phi)^2 + \frac{1}{4} \lambda_{mix} (h^2 + 2hv_h)(\phi^2 + 2\phi v_\phi). \tag{15}
\]

From the above scalar potential we obtain the mass-squared matrix for the conventional case of the form:

\[
M_{sq} = \begin{pmatrix} m_h^2 & M^2 \\ M^2 & m_\phi^2 \end{pmatrix}, \tag{16}
\]

where \( M^2 = \lambda_{mix} v_h v_\phi \). We diagonalize this matrix with Eq. (7) with the mixing angle defined by \( \tan(2\theta) = \frac{-2M^2}{m_h^2 - m_\phi^2} \). In terms of observables \( M_{h_1}, M_{h_2}, \theta \), and \( v_\phi \), we can express the parameters in Eq. (15) as

\[
m_h^2 = \frac{1}{2} \left( M_{h_1}^2 + M_{h_2}^2 + (M_{h_1}^2 - M_{h_2}^2) \cos(2\theta) \right), \tag{17}
\]

\[
m_\phi^2 = \frac{1}{2} \left( M_{h_1}^2 + M_{h_2}^2 - (M_{h_1}^2 - M_{h_2}^2) \cos(2\theta) \right), \tag{18}
\]

\[
\lambda_{mix} = M^2 \frac{v_h}{v_\phi}, \tag{19}
\]

where \( M^2 = \frac{1}{2} \left( M_{h_1}^2 - M_{h_2}^2 \right) \sin(2\theta) \). Note that the conventional Higgs potential of Eq. (15) is controlled by four free parameters. In contrast to our conformal model, both \( \theta \) and \( v_\phi \) remain free after \( M_{h_1} \) and \( M_{h_2} \) are fixed.

Expressing the potential in terms of the mass eigenstates \( h_{1,2} \), we extract the trilinear coupling \( g_{h_1 h_2 h_2} \):

\[
g_{h_1 h_2 h_2} = \left( M_{h_1}^2 + 2M_{h_2}^2 \right) (-v_h \cos(\theta) + v_\phi \sin(\theta)) \sin(2\theta) \over 2v_h v_\phi \tag{20}
\]

This coupling has different behavior for \( \theta \ll \frac{v_h}{v_\phi} \ll 1 \) and \( \frac{v_h}{v_\phi} \ll \theta \ll 1 \). For \( \theta \ll \frac{v_h}{v_\phi} \), the trilinear coupling \( g_{h_1 h_2 h_2} \) reduces to

\[
g_{h_1 h_2 h_2} \simeq -\frac{M_{h_1}^2}{2v_\phi} \left( 1 + 2 \frac{M_{h_2}^2}{M_{h_1}^2} \right) \theta, \tag{21}
\]

while for \( \frac{v_h}{v_\phi} \ll \theta \), the coupling is roughly given by

\[
g_{h_1 h_2 h_2} \simeq \frac{M_{h_1}^2}{2v_h} \left( 1 + 2 \frac{M_{h_2}^2}{M_{h_1}^2} \right) \theta^2. \tag{22}
\]
Comparing Eqs. (21) and (22) with Eq. (13), we see a remarkable difference. For \( \theta \ll \frac{v_h}{v_t} \), the coupling from the conformal potential is proportional to \( \theta^2 \), while the conventional coupling is proportional to \( \theta \), as we would naively expect: \( g_{h_1h_2h_2} \sim \lambda_{\text{mix}} v_h \). This is because in the very small \( \theta \) expansion for the trilinear coupling, the lowest order in \( \theta \) is cancelled out in the conformal potential. On the other hand, for \( \frac{v_h}{v_t} \gtrsim \theta \), the couplings from both potentials have the same \( \theta \) dependence, but that of the conformal model still exhibits relative suppression. The ratio of the trilinear couplings in this region is found to be

\[
R \equiv \frac{|g_{h_1h_2h_2}|_{\text{conf}}}{|g_{h_1h_2h_2}|_{\text{conv}}} = \left( \frac{M_{h_2}^2}{M_{h_1}^2} \right) \frac{1 - 4 \frac{M_{h_2}^2}{M_{h_1}^2}}{1 + 2 \frac{M_{h_2}^2}{M_{h_1}^2}} < 0.0505, \tag{23}
\]

for \( M_{h_2} < \frac{1}{2} M_{h_1} \), with the maximum value obtained when \( M_{h_2} \simeq 0.33 M_{h_1} \).

Note that in both cases, the Higgs anomalous coupling is controlled by \( \cos(\theta) \). This result has interesting implications for Higgs phenomenology, namely, even if the anomalous SM Higgs coupling is detectable in size, the SM-like Higgs boson \( h_1 \) decay to \( h_2 h_2 \) can be much harder to detect in our model. This suppression is a characteristic feature of our extended conformally symmetric model, in which EW symmetry breaking is triggered by the radiative \( U(1)_H \) symmetry breaking.

### III. NUMERICAL ANALYSIS OF THE HIGGS BOSON TRILINEAR COUPLING

For an example numerical value, we use the known SM Higgs vev and fix the Higgs mass eigenvalues \( M_{h_1}, M_{h_2} \) as follows, alongside a mixing angle \( |\theta| = 0.1 \),

\[
M_{h_1} = 125 \text{ GeV}, \quad M_{h_2} = 25 \text{ GeV}. \tag{24}
\]

For the classically conformal model, \( v_\phi \) is totally fixed by the above choices of model parameters to be \( v_\phi = 2555 \) GeV. We find the triple coupling to be

\[
g_{h_1h_2h_2} = -0.0107, \tag{25}
\]

while for the conventional Higgs potential, \( v_\phi \) is a free parameter chosen to be \( v_\phi = 10^4 \) GeV, and the coupling becomes

\[
g_{h_1h_2h_2} = 0.424. \tag{26}
\]

As explained in the previous section (see Eqs. (13), (21), and (22)), for \( |\theta| = 0.1 \gg \frac{v_h}{v_\phi} \), the trilinear coupling for the conformal case is suppressed by a factor of \( R \simeq 0.04 \). We use \( v_\phi = 10^4 \) GeV as a benchmark value for all calculations for the conventional Higgs potential.

In the left panel of Fig. 1, we plot \( |g_{h_1h_2h_2}| \) for the two cases as a function of \( |\theta| \). For \( |\theta| \gtrsim \frac{v_h}{v_\phi} \simeq 0.025 \), we can see that \( |g_{h_1h_2h_2}| \) in the conformal case is suppressed by a factor \( R \simeq 0.04 \). As discussed in the previous section (see Eqs. (21) and (22)), we can see the transition in the \( \sin(\theta) \) dependence from \( \sin(\theta) \) to \( \sin^2(\theta) \) at \( |\theta| \simeq \frac{v_h}{v_\phi} \simeq 0.025 \). For \( |\theta| \ll \frac{v_h}{v_\phi} \), we can see stronger relative suppression between the couplings, with a suppression factor,

\[
\left( \frac{M_{h_2}^2}{M_{h_1}^2} \right) \frac{1 - 4 \frac{M_{h_2}^2}{M_{h_1}^2}}{1 + 2 \frac{M_{h_2}^2}{M_{h_1}^2}} \left( \frac{v_\phi}{v_h} |\theta| \right) \ll R. \tag{27}
\]
FIG. 1. Left Panel: $|g_{h_1h_2h_2}|$ as a function of $|\theta|$ in our model (solid line) and the conventional Higgs potential (dotted line). Right Panel: $|g_{h_1h_2}|$ as a function of $|\theta|$ in our model (solid line) and the conventional Higgs potential (dotted line). We set $M_{h_1} = 125 \text{ GeV}$ and $M_{h_2} = 25 \text{ GeV}$ in our analysis. The grey shaded region is excluded by the LEP-II results for the Higgs boson search [5], and the red shaded region represents the prospective search reach of the proposed International Linear Collider (ILC) [6–8].

The grey shaded region is excluded by the LEP-II results for the Higgs boson search [5]. Assuming the Higgs boson mass were 25 GeV, the Higgs boson coupling with the $Z$ boson must be suppressed by a factor of $\sin^2(\theta) \leq 0.02$. The red shaded region represents the prospective search reach of the proposed International Linear Collider (ILC), as low as $\sin^2(\theta) \approx 0.002$ [6–8]. We also calculate the trilinear coupling $g_{h_1h_1h_2}$ with the same parameter choice. The result is shown in the right panel of Fig. 1. Now, the resultant coupling from our conformal model overlaps well with the same coupling from the conventional Higgs potential, and there is no relative suppression between the two.

We now consider outcomes for experimental probes of our model. The Higgs physics of interest here are the decay mode $h_1 \rightarrow h_2h_2$ and the Higgs anomalous coupling, determined by $\cos(\theta)$. The partial decay width for the process $h_1 \rightarrow h_2h_2$ is given by

$$\Gamma_{h_1 \rightarrow h_2h_2} = \frac{|g_{h_1h_2h_2}|^2}{8\pi M_{h_1}} \sqrt{1 - 4 \frac{M_{h_2}^2}{M_{h_1}^2}}. \quad (28)$$

Using the SM-like Higgs boson total decay width to SM particles $\Gamma_{h \rightarrow SM} = 4.07 \text{ MeV}$ [14], the branching ratio to a pair of $h_2$ is defined as

$$\text{Br } (h_1 \rightarrow h_2h_2) = \frac{\Gamma_{h_1 \rightarrow h_2h_2}}{\Gamma_{h \rightarrow SM} + \Gamma_{h_1 \rightarrow h_2h_2}}. \quad (29)$$
In Fig. 2 we show the results of $\text{Br}(h_1 \to h_2 h_2)$ for the values of $M_{h_2} =$ 10 (red), 25 (black), and 50 (blue) GeV, with solid lines corresponding to the conformal model and dashed lines to the conventional Higgs potential (for which all three lines are nearly overlapping). As expected from the difference in the couplings, the branching ratios in our conformal model are also extremely suppressed. The grey shaded region, as before, is excluded by ATLAS($h_1 \to h_2 h_2 \to b\bar{b} b\bar{b}$) [11, 12] and LEP-II for $M_{h_2} =$ 25 GeV [13].

The combination of anomalous Higgs decay and anomalous Higgs coupling results will provide a way to distinguish our conformal scenario from the conventional Higgs potential. For a benchmark value of $\sin(\theta) = 0.1$, for the conventional Higgs case, the anomalous branching ratio and anomalous coupling can both be within the ILC search region, so they can be measured simultaneously. On the other hand, for our conformal model the branching ratio is highly suppressed, and even if the anomalous coupling were measured, the anomalous decay mode will still evade detection. This is a way to determine the origin of EW symmetry breaking, i.e. does it occur via the conventional Higgs potential with the negative mass-squared introduced by hand or does it originate from radiative symmetry breaking.

In our conformal model, once a value for $M_{h_2}$ is fixed, we can determine the $U(1)_H$ gauge
coupling $g_H$ and $Z'$ boson mass $m_{Z'}$ as functions of $|\theta|$. In Fig. 3, we show the relation between $g_H$, $m_{Z'}$, and $\sin^2(\theta)$ for $M_{h_2} = 10$ GeV, 25 GeV, and 50 GeV, from bottom to top. If the $Z'$ boson couples to SM fermions, we may consider high-energy collider bounds on the $(m_{Z'}, g_H)$ parameter space. As a well known example of such a model, one may consider the minimal $B-L$ (baryon number minus lepton number) model, based on the $U(1)_{B-L}$ gauge group [15–21]. The final results from the LHC Run 2 for the $B-L Z'$ gauge boson resonance search provide us with an upper bound on $g_H = g_{B-L}$ as a function of $m_{Z'}$, which is presented in Ref. [22]. In Fig. 4, we show the relationship between $g_H$ and $m_{Z'}$, along with the excluded region (grey shaded) by the Large Hadron Collider (LHC) experiments from Ref. [22], if we identify our $U(1)_H$ with $U(1)_{B-L}$. The three lines show the relationship for $M_{h_2} = 10$ GeV (red), 25 GeV (black), and 50 GeV (blue) for the range of $0.001 \leq \sin^2(\theta) \leq 0.8$, corresponding to the range shown in Fig. 3. The lines will extend to the right as we take lower values of $\sin^2(\theta) < 0.001$. To be consistent with the LHC bounds, $\sin^2(\theta)$ must be much smaller than 0.001 (see Fig. 3). Thus, to explore interesting regions for Higgs phenomenology ($\sin^2(\theta) \gtrsim 0.002$), we do not consider the $B-L$ model identification of $U(1)_H$. However, note that the identification of $U(1)_H$ with well-known flavor-dependent $U(1)$ extended SMs, such as $U(1)_{B-L_3}$ [23–25] and $U(1)_{L_\nu-L_\tau}$ [26, 27], is still possible, as in these models the $Z'$ boson has no coupling with the first generation quarks and can easily evade the LHC constraints.

IV. CONCLUSION AND DISCUSSION

We have considered a classically conformal $U(1)$ extension of the SM where the $U(1)$ Higgs field ($\Phi$) with a $U(1)$ charge $+2$ is introduced. The $U(1)$ symmetry is radiatively broken by
FIG. 4. The LHC bounds (grey shaded region) on \((m_{Z'}, g_H)\) parameter space in the case of \(U(1)_H = U(1)_{B-L}\), for \(M_{h_2} = 10\) GeV (red), 25 GeV (black), and 50 GeV (blue). The lines are shown for the range of \(0.001 \leq \sin^2(\theta) \leq 0.8\), corresponding to the range shown Fig. 2.

the Coleman-Weinberg mechanism, which generates a negative mass squared to the SM Higgs doublet through the mixed quartic coupling in the scalar potential \(V \supset -\lambda_{mix}(\Phi^\dagger\Phi)(H^\dagger H)\), and as a result, the EW symmetry is broken. Therefore, the radiative symmetry breaking of the new U(1) gauge symmetry is the origin of the EW symmetry breaking. This is a crucial difference from the conventional Higgs potential in which the EW symmetry breaking is triggered by a negative mass squared introduced by hand at the tree-level. In order to investigate the Higgs phenomenology of the classically conformal model, we have analyzed the effective Higgs potential and read off the trilinear coupling of the SM-like Higgs boson \((h_1)\) with a pair of the SM singlet-like Higgs bosons \((h_2)\) for the case of \(M_{h_1} > 2M_{h_2}\). For comparison, we have also calculated the trilinear coupling for the conventional Higgs potential. What we have found is very intriguing: for the same mixing angle between \(H\) and \(\Phi\) and the same Higgs boson mass spectrum, the trilinear coupling in the conformal model is highly suppressed compared to the one from the conventional Higgs potential, which is likely a striking nature of the classically conformal potential. We then have considered experimental signals for such conformal structure via anomalous SM Higgs boson couplings from the mixing between two Higgs fields and the anomalous (SM-like) Higgs boson decay \(h_1 \rightarrow h_2h_2\) followed by \(h_2 \rightarrow b\bar{b}\). Future \(e^+e^-\) colliders, in particular, the proposed ILC will operate as a Higgs factory which allows us to precisely measure the SM Higgs boson properties. The result of our analysis indicates that once the anomalous SM-like Higgs boson coupling is measured at the ILC, the observation of anomalous Higgs boson decay \(h_1 \rightarrow h_2h_2\) is promising in the conventional Higgs potential, while this decay process is highly suppressed and not detectable in the classically conformal model. This is a way to to distinguish the origin of EW symmetry breaking.
As previously discussed, we may extend our classically conformal $U(1)$ model by identifying $U(1)$ with $U(1)_{B-L}$ or $U(1)_{\mu-\tau}$. Another interesting possibility of such an extension is to consider the $U(1)$ sector as a “Dark Sector” which supplements the SM with a dark matter candidate. Note that although the $Z'$ gauge boson can generally have a kinetic mixing with the SM hyper-charge gauge boson, if the kinetic mixing is tuned to be zero, the $Z'$ gauge boson is stable and hence a dark matter candidate. Therefore, the classically conformal gauged $U(1)$ extension can not only provide the SM with the dynamical origin of the EW symmetry breaking but also supplement the SM with the $Z'$ boson dark matter. Based on a new $SU(2)$ gauge symmetry, such a model has already been proposed in Ref. [28], although the Higgs boson phenomenology that we have investigated in this paper is not addressed. In the present paper, we have considered the $U(1)$ gauge theory, but extending it to a non-Abelian gauge group is straightforward and what we have investigated in this paper remains essentially the same for such extension. It is worth investigating the Dark Sector in light of complementarity between dark matter physics and Higgs boson phenomenology at future collider experiments [29].

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