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A new dynamics of electroweak symmetry breaking with classically scale invariance

Naoyuki Haba [a], Hiroyuki Ishida [a,*], Noriaki Kitazawa [b], Yuya Yamaguchi [a,c]

[a] Graduate School of Science and Engineering, Shimane University, Matsue 690-8504, Japan
[b] Department of Physics, Tokyo Metropolitan University, Hachioji, Tokyo 192-0397, Japan
[c] Department of Physics, Faculty of Science, Hokkaido University, Sapporo 060-0810, Japan

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ABSTRACT

We propose a new dynamics of the electroweak symmetry breaking in a classically scale invariant version of the standard model. The scale invariance is broken by the condensations of additional fermions under a strong coupling dynamics. The electroweak symmetry breaking is triggered by negative mass squared of the elementary Higgs doublet, which is dynamically generated through the bosonic seesaw mechanism. We introduce a real pseudo-scalar singlet field interacting with additional fermions and Higgs doublet in order to avoid massless Nambu–Goldstone bosons from the chiral symmetry breaking in a strong coupling sector. We investigate the mass spectra and decay rates of these pseudo-Nambu–Goldstone bosons, and show they can decay fast enough without cosmological problems. We further show that our model can make the electroweak vacuum stable.

1. Introduction

The origin of the electroweak symmetry breaking (EWSB) remains a mystery. In the standard model (SM), the EWSB requires a negative mass squared for the Higgs doublet scalar field, whose magnitude is set by hand. We expect a fundamental theory which naturally gives the negative mass squared with the suitable value. In a model of supersymmetric extension of the SM, the EWSB can be realized by so-called radiative breaking [2]. However, the supersymmetry breaking scale must be high because of no signal of super-particle at any experiments so far. In technicolor (TC) model [1], the Higgs doublet field is no longer an elementary scalar field, and the EWSB is triggered by the techni-fermion condensation under strongly coupled TC gauge interaction. However, the naive TC model, which is just scale up of QCD, has already been excluded by the electroweak precision measurements.

Recently, there are a lot of studies of other possibilities to solve the gauge hierarchy problem by imposing a classically scale invariance with an additional $U(1)$ gauge symmetry [3–27]. From the viewpoint of Bardeen's argument [28], we can only focus on logarithmic divergences, and the scale invariance protects large Higgs mass corrections. Under the classically scale invariance in terms of the cutoff regularization, the quadratic divergence itself can be subtracted by a boundary condition of the UV complete theory [8]. Once we subtract the quadratic divergence from the theory, it never appears in the observables. In the model with an additional $U(1)$ gauge symmetry, the scale invariance is broken by the Coleman–Weinberg mechanism [29], and if the breaking scale is not so far from the electroweak (EW) scale, there is no gauge hierarchy problem. On the other hand, a strong coupling dynamics can also realize such an EWSB with classically scale invariance [30, 31], where an additional singlet scalar mediates dimensional transmutation in the strong coupling sector to the SM sector. However, the sign of the coupling between the Higgs doublet and the additional scalar is assumed to be negative, so that the negative mass squared of the Higgs doublet is realized. Therefore, the origin of the EWSB is not necessary and inevitable in this scenario, and we are going to try the dynamical realization of negative mass squared by the bosonic seesaw mechanism [32].

In this paper, we expand the SM gauge group by $SU(N_{\mathrm{TC}})$ technicolor gauge symmetry with the classically scale invariant framework. The techni-fermions, which belong to vector-like representations under TC gauge symmetry as well as electroweak gauge symmetry, are introduced. Though the chiral symmetry breaking happens by techni-fermion condensations, the EWSB does not happen by this strong coupling TC dynamics itself. We show that the EWSB dynamically occurs in an inevitable way by the bosonic seesaw mechanism between the elementary Higgs scalar field and a

* Corresponding author.
E-mail address: ishida@riko.shimane-u.ac.jp (H. Ishida).

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Table 1
Charge assignments of techni-fermions and the Higgs doublet.

	$SU(N_{TC})$	$SU(2)_L$	$U(1)_Y$
H	1	2	1/2
χ	N_{TC}	2	1/2
ψ	N_{TC}	1	0

composite scalar field. To avoid massless Nambu–Goldstone (NG) bosons by the chiral symmetry breaking in strong coupling sector, we introduce a real pseudo-scalar singlet field and its interactions with techni-fermions and Higgs doublet. We analyze the mass spectrum of the pseudo-NG (pNG) bosons and estimate their decay rates. We show that the pNG bosons can decay fast enough to avoid cosmological problems. We will further show that our model can make the electroweak vacuum stable due to the contributions from additional vector-like fermion charged under the SM gauge group to the running of gauge coupling constants.

2. Bosonic seesaw mechanism

By imposing classically scale invariance, the mass term of the Higgs potential is forbidden and the Higgs potential becomes

$$V = \lambda \left(H^\dagger H\right)^2 . \tag{1}$$

The EWSB does not occur by this potential, and we try to use the dimensional transmutation in the strong coupling sector, where there are two vector-like techni-fermions as shown in Table 1. Due to the classically scale invariance, vector-like fermion masses are also forbidden. In the model, the chiral symmetry in the strong coupling sector $SU(3)_L \times SU(3)_R \times U(1)_A$ is explicitly broken by the SM gauge symmetry, $SU(2)_L \times U(1)_Y$, and the remaining symmetry is $SU(2)_{\chi_L} \times SU(2)_{\chi_R} \times U(1)_{\chi_A} \times U(1)_{\psi_A}$. There is also $U(1)_{\chi_V} \times U(1)_{\psi_V}$, which is similar to the baryon number symmetry.[1] This vector-like symmetry is expected to be unbroken by the strong-coupling technicolor dynamics due to the Vafa–Witten's theorem [34]. The chiral symmetry should be broken as preserving $SU(2)_L \times U(1)_Y$ symmetry, then we expect $\langle\bar{\chi}\psi\rangle = \langle\bar{\psi}\chi\rangle = 0$ and $\langle\bar{\chi}\chi\rangle \neq 0$, $\langle\bar{\psi}\psi\rangle \neq 0$. They cause chiral symmetry breaking $SU(2)_{\chi_L} \times SU(2)_{\chi_R} \times U(1)_{\chi_A} \times U(1)_{\psi_A} \to SU(2)_{\chi_V}$. There are five NG bosons; two massive pNG bosons of anomalous $U(1)_{\chi_A}$ and $U(1)_{\psi_A}$ breakings, and three massless NG bosons corresponding to the breaking $(SU(2)_{\chi_L} \times SU(2)_{\chi_R})/SU(2)_{\chi_V}$ symmetry. If we neglect $SU(2)_L \times U(1)_Y$, the chiral symmetry breaking of $SU(3)_L \times SU(3)_R \times U(1)_A \to SU(3)_V$ occurs. There are nine NG bosons; one massive pNG boson of $U(1)_A$ breaking, and eight massless NG bosons.

The techni-fermions interact with Higgs doublet H through the Yukawa interactions,

$$-\mathcal{L}_{\text{Yukawa}} = y_L \bar{\chi}_L H \psi_R + y_R \bar{\chi}_R H \psi_L + \text{h.c.} \tag{2}$$

After the techni-fermion condensation, $\chi_{L,R}$ and $\psi_{L,R}$ are confined by non-perturbative effects, and $\bar{\chi}_L\psi_R$ and $\bar{\chi}_R\psi_L$ couple to a "meson" state, that is just a composite Higgs doublet, $\Theta \sim \bar{\chi}\psi/\Lambda_{\text{TC}}^2$. When y_L and y_R are real, there is the charge conjugation invariance. Here, we assume $y_L = y_R = y$ for simplicity. The Yukawa interactions in Eq. (2) are CP invariant in this case. The composite Higgs doublet mixes with the elementary Higgs doublet, and the mass matrix becomes

$$-\mathcal{L}_{\text{mass}} = \begin{pmatrix} H^\dagger & \Theta^\dagger \end{pmatrix} \begin{pmatrix} 0 & y\Lambda_{\text{TC}}^2 \\ y\Lambda_{\text{TC}}^2 & \alpha\Lambda_{\text{TC}}^2 \end{pmatrix} \begin{pmatrix} H \\ \Theta \end{pmatrix} \tag{3}$$

$$\simeq \begin{pmatrix} H_1^\dagger & H_2^\dagger \end{pmatrix} \begin{pmatrix} -\frac{y^2}{\alpha}\Lambda_{\text{TC}}^2 & 0 \\ 0 & \alpha\Lambda_{\text{TC}}^2 \end{pmatrix} \begin{pmatrix} H_1 \\ H_2 \end{pmatrix} , \tag{4}$$

where α is a dimensionless positive coefficient of $\mathcal{O}(1)$. Here, $y \ll \alpha$ is assumed, since the chiral symmetry breaking terms should be small to be treated perturbatively. As we will see later, the small Yukawa coupling y is also necessary for the hierarchy between the EW and TC condensation scales. As a result, the lighter (heavier) mass eigenstate H_1 (H_2) is almost H (Θ). The field H_1 is regarded as the SM-like Higgs doublet, and the negative mass squared is dynamically obtained through the bosonic seesaw mechanism (a similar mechanism has been discussed in [35]). The field H_2 has mass of $\mathcal{O}(\Lambda_{\text{TC}})$.

There are massless NG bosons in the present stage. To avoid the massless NG bosons, we introduce real pseudo-scalar field, S, which has interactions,

$$-\mathcal{L}_S = g_S S \bar{\chi} i\gamma_5 \chi + g'_S S \bar{\psi} i\gamma_5 \psi , \tag{5}$$

where g_S and g'_S are taken to be real to keep the CP invariance. Since we can expect $\langle\bar{\chi}i\gamma_5\chi\rangle = 0$ and $\langle\bar{\psi}i\gamma_5\psi\rangle = 0$ in vector-like technicolor dynamics, the no tadpole term of S is not generated.

Now the potential in Eq. (1) is modified as

$$V_{\text{eff}} = \lambda\left(H^\dagger H\right)^2 + \kappa S^2 H^\dagger H + \lambda_S S^4 + y\Lambda_{\text{TC}}^2\left(H^\dagger\Theta + \Theta^\dagger H\right) + \alpha\Lambda_{\text{TC}}^2\Theta^\dagger\Theta , \tag{6}$$

where fourth and fifth terms are obtained from Eq. (3). Since we have assumed the hierarchy between the light and heavy mass eigenstates, the heavier mass eigenstate H_2 is decoupled at low energies. Therefore, the effective potential at low energy is

$$V_{\text{eff}} \simeq \lambda\left(H_1^\dagger H_1\right)^2 + \kappa S^2 H_1^\dagger H_1 + \lambda_S S^4 - \frac{y^2}{\alpha}\Lambda_{\text{TC}}^2 H_1^\dagger H_1 - \frac{1}{2}m_S^2 S^2 , \tag{7}$$

where we include the mass term of S which is generated by bosonic seesaw mechanism again. We will give an analysis about this issue shortly in the next section.

The vacuum expectation values (VEVs) of H_1 and S can be evaluated by the effective potential Eq. (7). The stationary conditions are

$$\left(\lambda v_H^2 + \kappa v_S^2 - \frac{y^2}{\alpha}\Lambda_{\text{TC}}^2\right) v_H = 0 , \tag{8}$$

$$\left(\kappa v_H^2 + 4\lambda_S v_S^2 - m_S^2\right) v_S = 0 , \tag{9}$$

where $\langle H_1\rangle = (0, v_H/\sqrt{2})^{\text{T}}$ and $\langle S\rangle = v_S$. Note that v_H should corresponds to the EW scale ($v_H = 246$ GeV), and nonzero v_S causes spontaneous CP violation. Except for a trivial solution $v_H = v_S = 0$, there are three possibilities of solutions as follows.

- $v_H = 0$ and $v_S \neq 0$
  In this case the EW symmetry is not broken, while v_S can be estimated as

  $$v_S^2 = \frac{m_S^2}{4\lambda_S} , \tag{10}$$

  where m_S^2 must be positive. To satisfy $v_H = 0$, i.e., to realize the positive mass squared of H_1, the following condition must be satisfied:

  $$\kappa v_S^2 - \frac{y^2}{\alpha}\Lambda_{\text{TC}}^2 > 0 . \tag{11}$$

[1] They guarantee the stability of the lightest techni-baryon which can be a candidate of the dark matter. (For instance, see Ref. [33].)

Table 2
Summary of the pNG bosons. σ_i are Pauli matrices.

	Operators	$SU(2)_L$	$U(1)_Y$	Masses
η_χ	$\bar{\chi} i\gamma_5 \chi$	1	0	$\beta\Lambda_{\rm TC}^2$
η_ψ	$\bar{\psi} i\gamma_5 \psi$	1	0	$\beta\Lambda_{\rm TC}^2$
Π_i $(i=1,2,3)$	$\bar{\chi} i\gamma_5 \sigma_i \chi$	3	0	$\frac{8\pi^2 g_S^4}{\lambda_S\beta^2}\Lambda_{\rm TC}^2$
$\Sigma = \begin{pmatrix}\Sigma^0\\ \Sigma^-\end{pmatrix}$	$\bar{\psi} i\gamma_5 \chi$	2	−1/2	$\frac{16\pi^2 g_S^4}{\lambda_S\beta^2}\Lambda_{\rm TC}^2$
$\bar{\Sigma} = \begin{pmatrix}\bar{\Sigma}^0\\ \bar{\Sigma}^+\end{pmatrix}$	$\bar{\chi} i\gamma_5 \psi$	2	1/2	$\frac{16\pi^2 g_S^4}{\lambda_S\beta^2}\Lambda_{\rm TC}^2$

Thus, a certain large value of κ is required when v_S is $\mathcal{O}(\Lambda_{\rm TC})$. Anyway, we do not consider this case, since the EW symmetry is unbroken.

- $v_H \neq 0$ and $v_S = 0$
  In this case the EWSB occurs and its scale is given by

$$v_H^2 = \frac{y^2}{\lambda\alpha}\Lambda_{\rm TC}^2\,. \tag{12}$$

This is really a solution if the mass squared of S is positive, that is,

$$\frac{\kappa y^2}{\lambda\alpha}\Lambda_{\rm TC}^2 - m_S^2 > 0\,. \tag{13}$$

This condition is always satisfied for a sufficiently large κ. Since we would like to treat κ perturbatively in good approximation, we do not adopt this case also.

- $v_H \neq 0$ and $v_S \neq 0$
  This case leads a suitable result. The stationary conditions give

$$v_H^2 = \frac{1}{4\lambda\lambda_S - \kappa^2}\left(-\kappa m_S^2 + 4\lambda_S\frac{y^2}{\alpha}\Lambda_{\rm TC}^2\right), \tag{14}$$

$$v_S^2 = \frac{1}{4\lambda\lambda_S - \kappa^2}\left(\lambda m_S^2 - \kappa\frac{y^2}{\alpha}\Lambda_{\rm TC}^2\right). \tag{15}$$

Since the squared VEVs must be positive, a certain small value of κ are required. In the limit of $\kappa \to 0$, the VEVs are approximately given by

$$v_H^2 \simeq \frac{y^2}{\lambda\alpha}\Lambda_{\rm TC}^2\,, \qquad v_S^2 \simeq \frac{1}{4\lambda_S}m_S^2\,, \tag{16}$$

where m_S^2 must be positive. Since S obtains a nonzero VEV, a mixing term with Higgs doublet affects the Higgs mass through $\kappa|H_1|^2S^2$. However, it is negligible because κ is assumed to be sufficiently small. (In the case of $\kappa \simeq 0$, we can treat H_1 and S independently.) From now on, we adopt this case with taking sufficiently small value of κ.

3. Mass spectra and decay rates of pNG bosons

Now let us investigate the mass spectra and decay rates of the pNG bosons. Actually, all nine NG bosons should become massive due to the introduction of S, since the chiral symmetries $SU(2)_{\chi_L} \times SU(2)_{\chi_R} \times U(1)_{\chi_A} \times U(1)_{\psi_A}$ is explicitly broken down into $SU(2)_{\chi_V}$ by the interactions of Eq. (5). The results of the mass spectra are summarized in Table 2.

First, we investigate pNG boson mass spectra. The SM singlet pNG bosons (η_χ and η_ψ) mix with S, and the mass matrix is written by

$$-\mathcal{L}_{S\text{-}\eta_\chi\text{-}\eta_\psi} = \frac{1}{2}\begin{pmatrix} S & \eta_\chi^\dagger & \eta_\psi^\dagger\end{pmatrix}\begin{pmatrix} 0 & g_S\Lambda_{\rm TC}^2 & g'_S\Lambda_{\rm TC}^2 \\ g_S\Lambda_{\rm TC}^2 & \beta_\chi\Lambda_{\rm TC}^2 & 0 \\ g'_S\Lambda_{\rm TC}^2 & 0 & \beta_\psi\Lambda_{\rm TC}^2\end{pmatrix}\begin{pmatrix} S \\ \eta_\chi \\ \eta_\psi\end{pmatrix}, \tag{17}$$

where β_χ and β_ψ are dimensionless positive coefficients of $\mathcal{O}(1)$. All off-diagonal elements are induced from Eq. (5). The determinant of this mass matrix is $-(g_S^2\beta_\psi + g_S'^2\beta_\chi)\Lambda_{\rm TC}^6 < 0$, thus S has a negative mass term. Taking $g_S = g'_S \ll \beta_\chi = \beta_\psi = \beta$, for simplicity, mass eigenvalues of S, η_χ, and η_ψ can be estimated as

$$-m_S^2 \simeq -\frac{2g_S^2}{\beta}\Lambda_{\rm TC}^2\,, \qquad m_{\eta_\chi}^2 = m_{\eta_\psi}^2 \simeq \beta\Lambda_{\rm TC}^2\,, \tag{18}$$

respectively. The smallness of g_S is natural, since it is expected to break the chiral symmetry perturbatively. Note that S has the negative mass term by the bosonic seesaw mechanism again (See Eq. (7)).

Using Dashen's formula [36] and VEVs of H_1 and S in Eq. (16), the masses of Π and Σ are estimated as

$$m_\Pi^2 f_\Pi^2 = \langle 0|\,[Q\,,\,[Q\,,\,\mathcal{H}_S]]\,|0\rangle \simeq \frac{g_S^4}{2\lambda_S\beta^2}\Lambda_{\rm TC}^4\,, \tag{19}$$

$$m_\Sigma^2 f_\Sigma^2 = \langle 0|\,[Q\,,\,[Q\,,\,\mathcal{H}_S]]\,|0\rangle \simeq \frac{g_S^4}{\lambda_S\beta^2}\Lambda_{\rm TC}^4\,, \tag{20}$$

where $\mathcal{H}_S = g_S S\bar{\chi}i\gamma_5\chi + g_S S\bar{\psi}i\gamma_5\psi$ from Eq. (5) and, f_Π and f_Σ are decay constants of Π and Σ, respectively. Both decay constants are evaluated by naive dimensional analysis [37,38] as $\Lambda_{\rm TC} \simeq 4\pi f_{\Pi,\Sigma}$ by analogy with QCD. Therefore, the masses of Π and Σ are estimated as

$$m_\Pi^2 \simeq \frac{8\pi^2 g_S^4}{\lambda_S\beta^2}\Lambda_{\rm TC}^2\,, \qquad m_\Sigma^2 \simeq \frac{16\pi^2 g_S^4}{\lambda_S\beta^2}\Lambda_{\rm TC}^2\,. \tag{21}$$

In the following, we take $\Lambda_{\rm TC} = 10$ TeV and $\alpha = \beta = 1$ for an explicit example. Then, the coupling y is evaluated as $y \simeq 0.068$ from Eq. (16). If we also take $\lambda_S = 10^{-3}$ and $g_S = 0.05$, v_S and the mass of Π and Σ are evaluated as $v_S \simeq 11$ TeV, $m_\Pi \simeq 7$ TeV and $m_\Sigma \simeq 10$ TeV, respectively.

Next, we estimate decay rates of the pNG bosons by analogy with light mesons in QCD. A charged components of Σ and Π can decay into their neutral components and the SM fermions through the weak interactions. η_χ and the neutral component of Π (Π^0) decay into two photons by analogy with π^0 decay in the SM. The decay rate of η_χ is evaluated as

$$\Gamma(\eta_\chi \to \gamma\gamma) = \left(\frac{N_{\rm TC}e^2}{4\pi^2 f_{\eta_\chi}}\right)^2\frac{m_{\eta_\chi}^3}{64\pi} \simeq \frac{N_{\rm TC}^2\alpha_{\rm em}^2}{4\pi}\frac{m_{\eta_\chi}^3}{\Lambda_{\rm TC}^2}\,, \tag{22}$$

where we have used $f_{\eta_\chi} \simeq \Lambda_{\rm TC}/4\pi$ and $\alpha_{\rm em} = e^2/4\pi$. When we take $N_{\rm TC} = 3$ for example, the decay rate is estimated by $\Gamma(\eta_\chi \to \gamma\gamma) \simeq 400$ MeV.

The neutral component of Σ (Σ^0) also decays into two photons via a mixing with η_χ. The effective Σ-η_χ mixing is evaluated by Dashen's formula as

$$m_{\Sigma\text{-}\eta_\chi}^2 \simeq (4\pi)^2\sqrt{2}y v_H\Lambda_{\rm TC}\,, \tag{23}$$

and hence, the magnitude of Σ-η_χ mixing is given by

$$V_{\Sigma\text{-}\eta_\chi} \equiv \frac{m_{\Sigma\text{-}\eta_\chi}^2}{m_\Sigma^2} \simeq \frac{\sqrt{2}y\lambda_S\beta^2}{g_S^4}\frac{v_H}{\Lambda_{\rm TC}}\,. \tag{24}$$

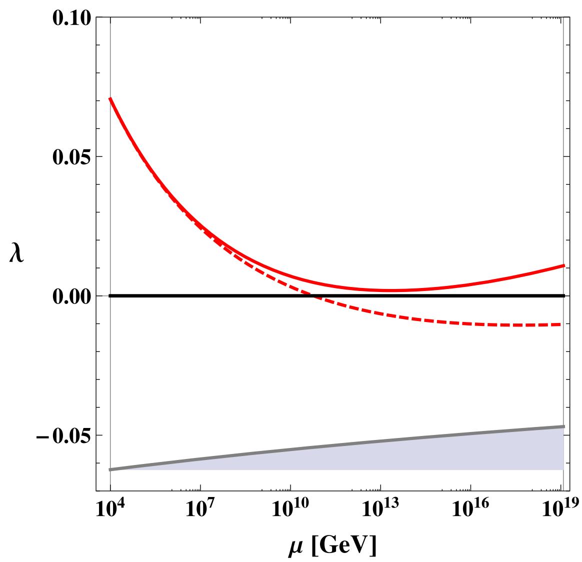

Fig. 1. Running of Higgs quartic couplings between $\Lambda_{\rm TC} = 10$ TeV and the Planck scale, which are denoted by the black vertical lines. The solid and dashed lines correspond to our model and the SM, respectively. The gray shaded region means the unstable vacuum.

Thus, we obtain $V_{\Sigma\text{-}\eta_\chi} \simeq 0.4$ by using the same numerical values as above. As a result, we find $\Gamma(\Sigma^0 \to \gamma\gamma) \simeq V^2_{\Sigma\text{-}\eta_\chi} \times \Gamma(\eta_\chi \to \gamma\gamma) \simeq 60$ MeV.

The decay mode of the lightest neutral pNG boson η_ψ is a little bit tricky. The decay process is $\eta_\psi \to S \to \eta_\chi \to \gamma\gamma$ through mass mixings. Therefore, the lifetime of η_ψ would be the longest among the pNG bosons. Since S-η_ψ and S-η_χ effective mixing couplings can be evaluated from Eq. (17) as g'_S/β_ψ and g_S/β_χ, respectively, the decay rate can be estimated as

$$\Gamma(\eta_\psi \to \gamma\gamma) \simeq \left(\frac{g'_S}{\beta_\psi}\frac{g_S}{\beta_\chi}\right)^2 \times \Gamma(\eta_\chi \to \gamma\gamma). \tag{25}$$

Thus, $\Gamma(\eta_\psi \to \gamma\gamma)$ is around 3 keV using the same numerical values as above. Even the lightest pNG boson can decay much faster than the QCD neutral pion ($\Gamma(\pi^0 \to \gamma\gamma) \simeq 7.7$ eV). As a result, we can expect that all the pNG bosons decay into the SM particles fast enough without cosmological problems.

In the end of this section, we mention the EW vacuum stability. Due to the existence of techni-fermion, χ, which has an electroweak charge, the $U(1)_Y$ and $SU(2)_L$ gauge coupling constants become larger than the SM case. Then, a running of the Higgs quartic coupling is lifted up, and hence, the EW vacuum can be completely stable. Using the parameter values above, the running of the Higgs quartic coupling is shown in Fig. 1. We have solved renormalization group equations (RGEs) at one-loop level, and taken $m_t = 173$ GeV and $m_h = 126$ GeV as a reference values. When we use two-loop RGEs and/or the different values of Higgs and top-quark masses, the running can be modified, but it does not change our statement. As a result, we find that our model can realize the EW vacuum stability.

4. Discussions and conclusions

The origin of the EWSB is not established yet, although the SM-like Higgs boson has been discovered. In this paper, we have investigated the dynamical origin of the EWSB via the bosonic seesaw mechanism in a classically scale invariant version of the SM. We have introduced the $SU(N_{\rm TC})$ technicolor gauge symmetry for the dimensional transmutation by the techni-fermion condensations. In this model, the mixing between the elementary and composite Higgs doublets becomes the origin of EWSB. An extra real pseudo-scalar singlet field has also been introduced to avoid massless NG bosons. We have estimated mass spectra and decay rates of the pNG bosons. We have checked that all of the pNG bosons can decay fast enough without cosmological problems. The EW vacuum stability becomes better than the SM due to the stronger $U(1)_Y$ and $SU(2)_L$ gauge coupling constants by introducing vector-like fermion, χ.

Finally, we comment on the collider phenomenology. When the singlet pseudo-scalar is light enough to be produced at the collider, some vestiges could be searched in the future collider experiments. In addition, since there can exist light new mesons depending on the parameters, they might be detectable at collider experiments.

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References

[1] S. Weinberg, Phys. Rev. D 13 (1976) 974;
L. Susskind, Phys. Rev. D 20 (1979) 2619.
[2] K. Inoue, A. Kakuto, H. Komatsu, S. Takeshita, Prog. Theor. Phys. 68 (1982) 927, Prog. Theor. Phys. 70 (1983) 330.
[3] R. Hempfling, Phys. Lett. B 379 (1996) 153, arXiv:hep-ph/9604278.
[4] W.F. Chang, J.N. Ng, J.M.S. Wu, Phys. Rev. D 75 (2007) 115016, arXiv:hep-ph/0701254.
[5] S. Iso, N. Okada, Y. Orikasa, Phys. Rev. D 80 (2009) 115007, arXiv:0909.0128 [hep-ph].
[6] S. Iso, N. Okada, Y. Orikasa, Phys. Lett. B 676 (2009) 81, arXiv:0902.4050 [hep-ph].
[7] N. Okada, Y. Orikasa, Phys. Rev. D 85 (2012) 115006, arXiv:1202.1405 [hep-ph].
[8] S. Iso, Y. Orikasa, PTEP 2013 (2013) 023B08, arXiv:1210.2848 [hep-ph].
[9] C. Englert, J. Jaeckel, V.V. Khoze, M. Spannowsky, J. High Energy Phys. 1304 (2013) 060, arXiv:1301.4224 [hep-ph].
[10] E.J. Chun, S. Jung, H.M. Lee, Phys. Lett. B 725 (2013) 158, Phys. Lett. B 730 (2014) 357, arXiv:1304.5815 [hep-ph].
[11] I. Oda, Phys. Lett. B 724 (2013) 160, arXiv:1305.0884 [hep-ph].
[12] V.V. Khoze, G. Ro, J. High Energy Phys. 1310 (2013) 075, arXiv:1307.3764 [hep-ph].
[13] M. Hashimoto, S. Iso, Y. Orikasa, Phys. Rev. D 89 (1) (2014) 016019, arXiv:1310.4304 [hep-ph].
[14] M. Lindner, D. Schmidt, A. Watanabe, Phys. Rev. D 89 (1) (2014) 013007, arXiv:1310.6582 [hep-ph].
[15] M. Hashimoto, S. Iso, Y. Orikasa, Phys. Rev. D 89 (5) (2014) 056010, arXiv:1401.5944 [hep-ph].
[16] S. Benic, B. Radovcic, Phys. Lett. B 732 (2014) 91, arXiv:1401.8183 [hep-ph].
[17] V.V. Khoze, C. McCabe, G. Ro, J. High Energy Phys. 1408 (2014) 026, arXiv:1403.4953 [hep-ph].
[18] S. Benic, B. Radovcic, J. High Energy Phys. 1501 (2015) 143, arXiv:1409.5776 [hep-ph].
[19] H. Okada, Y. Orikasa, arXiv:1412.3616 [hep-ph].
[20] J. Guo, Z. Kang, P. Ko, Y. Orikasa, Phys. Rev. D 91 (11) (2015) 115017, arXiv:1502.00508 [hep-ph].
[21] P. Humbert, M. Lindner, J. Smirnov, J. High Energy Phys. 1506 (2015) 035, arXiv:1503.03066 [hep-ph].
[22] N. Haba, Y. Yamaguchi, arXiv:1504.05669 [hep-ph].
[23] S. Oda, N. Okada, D.s. Takahashi, Phys. Rev. D 92 (1) (2015) 015026, arXiv:1504.06291 [hep-ph].
[24] P. Humbert, M. Lindner, S. Patra, J. Smirnov, arXiv:1505.07453 [hep-ph].
[25] A.D. Plascencia, arXiv:1507.04996 [hep-ph].
[26] N. Haba, H. Ishida, N. Okada, Y. Yamaguchi, arXiv:1508.06828 [hep-ph].
[27] A. Das, N. Okada, N. Papapietro, arXiv:1509.01466 [hep-ph].
[28] W.A. Bardeen, FERMILAB-CONF-95-391-T, C95-08-27.3.

[29] S.R. Coleman, E.J. Weinberg, Phys. Rev. D 7 (1973) 1888.
[30] T. Hur, D.W. Jung, P. Ko, J.Y. Lee, Phys. Lett. B 696 (2011) 262, arXiv:0709.1218 [hep-ph];
T. Hur, P. Ko, Phys. Rev. Lett. 106 (2011) 141802, arXiv:1103.2571 [hep-ph];
M. Heikinheimo, A. Racioppi, M. Raidal, C. Spethmann, K. Tuominen, Mod. Phys. Lett. A 29 (2014) 1450077, arXiv:1304.7006 [hep-ph].
[31] M. Holthausen, J. Kubo, K.S. Lim, M. Lindner, J. High Energy Phys. 1312 (2013) 076, arXiv:1310.4423 [hep-ph].
[32] X. Calmet, Eur. Phys. J. C 28 (2003) 451, arXiv:hep-ph/0206091;
H.D. Kim, Phys. Rev. D 72 (2005) 055015, arXiv:hep-ph/0501059;
N. Haba, N. Kitazawa, N. Okada, Acta Phys. Pol. B 40 (2009) 67, arXiv:hep-ph/0504279.
[33] O. Antipin, M. Redi, A. Strumia, E. Vigiani, J. High Energy Phys. 1507 (2015) 039, arXiv:1503.08749 [hep-ph].
[34] C. Vafa, E. Witten, Phys. Rev. Lett. 53 (1984) 535.
[35] O. Antipin, M. Redi, A. Strumia, J. High Energy Phys. 1501 (2015) 157, http://dx.doi.org/10.1007/JHEP01(2015)157, arXiv:1410.1817 [hep-ph].
[36] R.F. Dashen, Phys. Rev. 183 (1969) 1245.
[37] A. Manohar, H. Georgi, Nucl. Phys. B 234 (1984) 189.
[38] H. Georgi, L. Randall, Nucl. Phys. B 276 (1986) 241.