Higgs as a Top-Mode Pseudo

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In this talk, in the spirit of the top quark condensation, we introduce a model which has a naturally light composite Higgs boson, “tHiggs”, to be identified with the 126 GeV Higgs discovered at the LHC. The tHiggs emerges as a pseudo Nambu-Goldstone boson (NGB), “Top-Mode Pseudo”, together with the exact NGBs (eaten by the W and Z bosons) as well as another Top-Mode Pseudo (CP-odd composite scalar). Those five NGBs are dynamically produced simultaneously by a four-fermion dynamics.

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*This talk is based on [1].
1. Introduction

A key clue to access a dynamical origin of the 126 GeV Higgs boson at the LHC \cite{2, 3} would be deduced from an observed coincidence among masses of top quark, Higgs boson and electroweak gauge bosons. This coincidence may imply that the top quark plays a crucial role for both the generation of the electroweak symmetry breaking (EWSB) scale and the generation of the mass of the Higgs boson. Top quark condensation \cite{4, 5, 6} naturally provides such a close relation between those mass scales. However, the original top quark condensate model is somewhat far from a realistic situation, e.g. a Higgs boson predicted as a $t\bar{t}$ bound state has the mass in a range of $m_t < m_H < 2m_t$, which cannot be identified with the 126 GeV Higgs boson at the LHC.

Based on \cite{1} \footnote{At almost the same time as \cite{1} a similar model was proposed in a slightly different context \cite{7}.}, we introduce a new class of the top quark condensate model, where a composite Higgs boson emerges as a pseudo Nambu–Goldstone boson (PNGB) associated with the spontaneous breaking of a global symmetry, which can be as light as the 126 GeV Higgs boson at the LHC.

2. Model

Let us consider a Nambu–Jona-Lasinio (NJL)-like model constructed from the third generation quarks in the SM, $q = (t, b)$, and an $SU(2)_L$ singlet quark ($\chi$). The left-handed quarks $q_L$ and $\chi_L$ form a triplet $\psi^i_L \equiv (t_L, b_L, \chi^i_L)^T$, ($i = 1, 2, 3$) under the flavor $U(3)_{\psi_L}$ group, while the right-handed top and bottom quarks $q^i_R \equiv (t^i_R, b^i_R)$, ($i = 1, 2$) and $\chi_R$ are a doublet and singlet under the $U(2)_{\chi_R}$ group, respectively. The electroweak gauge symmetry is embedded as a subgroup of the global symmetry. We thus write the global $U(3)_{\psi_L} \times U(2)_{\chi_R} \times U(1)_{\chi_L}$-invariant Lagrangian: $\mathcal{L}_{\text{kin}} + \mathcal{L}_{\text{4f}}^{\text{A}}$ where

$$\mathcal{L}_{\text{4f}}^{\text{A}} = G(\psi^i_L, \chi^i_R)(\bar{\chi}^i_R \psi^i_L),$$

and $G$ denotes the four-fermion coupling strength. We can derive the gap equations for fermion dynamical masses $m_\chi$ and $m_{\chi \chi}$ through the mean field relations $m_\chi = -G(\langle \bar{\chi}_{\chi L} \chi_L \rangle)$ and $m_{\chi \chi} = -G(\langle \chi_{\chi L} \chi_L \rangle)$ in the large $N_c$ limit: $m_{\chi \chi \chi} = m_{\chi \chi \chi}[N_cG/(8\pi^2)]\left[\Lambda^2 - (m^2_\chi + m^2_{\chi \chi}) \ln \Lambda^2/(m^2_\chi + m^2_{\chi \chi})\right]$ where $\Lambda$ stands for the cutoff of the model. There exist nontrivial solutions $m_\chi \neq 0$ and $m_{\chi \chi} \neq 0$ when the criticality condition is satisfied: $G > G_{\text{crit}} = 8\pi^2/(N_c\Lambda^2)$ under which we have the nonzero dynamical masses as well as the nonzero condensates, $\langle \chi_{\chi L} \chi_L \rangle \neq 0$ and $\langle \chi_{\chi L} \chi_L \rangle \neq 0$.

In order to make the structure of the symmetry breaking clearer, we may change the flavor basis of fermions $\psi_L \rightarrow \tilde{\psi}_L$ by an orthogonal rotation. The above gap equations are then reduced to a single gap equation, $1 = [N_cG/(8\pi^2)]\left[\Lambda^2 - (m^2_\chi + m^2_{\chi \chi}) \ln \Lambda^2/m^2_{\chi \chi}\right]$ with $m^2_{\chi \chi} = m^2_\chi + m^2_{\chi \chi} \neq 0$. Accordingly, the associated two condensates are reduced to a single nonzero condensate on the basis of $\tilde{\psi}^i_L$: $\langle \chi_{\chi L} \chi_L \rangle \neq 0$. We thus see that, with the criticality condition satisfied, the four-fermion dynamics triggers the global symmetry breaking pattern $U(3)_{\psi_L} \times U(1)_{\chi_R} \rightarrow U(2)_{\chi_R} \times U(1)_{\chi_L}$.

The broken currents associated with this symmetry breaking are found to be $J_{\chi_L}^{a\mu} = \bar{\psi}_L \gamma^\mu \lambda^a \psi_L$ and $J_A^{a\mu} \equiv (1/4)(\chi_{\chi L} \gamma^\mu \chi_R - \bar{\psi}_L \lambda^A \psi_L)$ where $\lambda^a (a = 4, 5, 6, 7, A)$ are the Gell-Mann matrices normalized as $\text{tr}[\lambda^a \lambda^b] = 2\delta^{ab}$ and $\lambda^A = \text{diag}(0, 0, \sqrt{2})$. The associated five NGBs emerge with the
decay constant $f$ as $\langle 0 | J^3_\mu (x) | \pi^i (p) \rangle = - if \delta^{ab} p_\mu e^{-ip.x}$ where the decay constant $f$ is calculated through the Pagels-Stokar formula [8]: $f^2 = (N_c/8\pi^2)m^2_{\pi} \ln(\Lambda^2 / m^2_{\pi})$. The five NGBs ($\pi^i$) can be expressed as composite fields (interpolating fields) made of the fermion bilinears on the basis of $(\bar{\psi}_L, \chi_R)$. Besides these composite NGBs, there exists a composite scalar $(H^0)$ corresponding to the $\sigma$ mode in the usual NJL model, $H^0 \sim \bar{\chi} R \chi L + \bar{\chi} L \chi R$ with the mass $m^2_{H^0} = 4m^2_{\pi}$. The $H^0$ will be regarded as a heavy Higgs boson with the mass of $\mathcal{O}(1)$ TeV, not the light Higgs boson at around 126 GeV.

We incorporate explicit breaking terms into the Lagrangian to give masses to some NGBs: $\mathcal{L}_{\text{kin.}} + \mathcal{L}^{4f} + \mathcal{L}^{\text{h}}$ where $(\Delta XX, G' > 0)$

$$\mathcal{L}^{\text{h}} = -[\Delta XX \bar{\chi} R \chi L + \text{h.c.}] - G'(\bar{\chi} L \chi R)(\bar{\chi} R \chi L). \quad (2.2)$$

The gap equations for fermion dynamical masses $m_{XX}$ are given by $m_{XX} = m_{XX}[N_cG/(8\pi^2)][\Lambda^2 - m^2_{\pi} \ln(\Lambda^2 / m^2_{\pi})], m_{XX} = \Delta XX + m_{XX}[N_c(G - G')/(8\pi^2)][\Lambda^2 - m^2_{\pi} \ln(\Lambda^2 / m^2_{\pi})]$ where $m^2_{XX} = m^2_{\pi} + m^2_{\pi}$. In this case, the nonzero $\Delta XX$ and $G'$ allow to determine the ratio of two dynamical masses $m_{XX}$ and $m_{XX}$, i.e. $\tan \theta = m_{XX}/m_{XX}$, in contrast to the previous gap equations which only determine the squared-sum of two, $m^2_{XX} + m^2_{XX}$. It turns out that Eq.(2.2) does not affect the criticality of the four-fermion dynamics at all. Eq.(2.2) forces the vacuum to choose a specific direction, $(\bar{\chi} R \chi L) \neq 0$, and gives masses to some of the NGBs. In fact, Eq.(2.2) is invariant under the chiral transformation associated with the broken currents $(f^6 \mu \pm i f^7 \mu)$ and $(f^4 \mu \cos \theta + f^3 \mu \sin \theta)$, but not for $f^5 \mu$, and $f^4 \mu \sin \theta + f^3 \mu \cos \theta$ $(\chi^0, w^\pm)$ eaten by the $Z$ and $W$ bosons are found to be $z^0 \equiv \pi^0 \cos \theta + \pi^0 \sin \theta, w^\pm \equiv (1/\sqrt{2})(\pi^0 \pm i\pi^7)$. Other NGBs remain as physical states: $h^0 \equiv \pi^0, A^0 \equiv -\pi^0 \sin \theta + \pi^0 \cos \theta$ and these NGBs become pseudo NGBs, called “Top-Mode Pseudos”, obtaining their masses once explicit breaking effects are introduced as Eq.(2.2).

We identify the CP-even Top-Mode Pseudo, $h^0$, as the 126 GeV Higgs, called tHiggs. The mass of $h^0$ is proportional to $m_{XX}$ associated with the EWSB scale $v_{EW}$ as $\sin \theta = m_{XX}/m_{XX} = v_{EW}/f$, just like the case of the SM Higgs boson, while the mass of $A^0$ is not. We thus set the mass of $h^0$ to $\simeq 126$ GeV: $m_{h^0} = m_{A^0} \sin \theta = 126$ GeV.

3. Phenomenological constraints on Top-Mode Pseudos

After adding four-fermion interactions to give the SM fermion masses, we find the tHiggs $h^0_i$ couplings to the SM particles in the present model are described by (see [1, 10] for details)

$$g_{hVV} \frac{v_{EW}}{2} \left( g^2 h^0_i W^+ W^- + \frac{g^2 + g^2}{2} h^0_i Z^+ Z^- \right) - \sum_{f = 1, b, t} g_{hf} \frac{m_f}{v_{EW}} h^0_i t f f, \quad (3.1)$$

where $g_{hVV} = g_{hbb} = g_{h\tau \tau} = \cos \theta$ and $g_{hct} = \sqrt{(1 + \cos^2 \theta)/2}$. We see that the $h^0_i$ couplings to the $W$ and $Z$ bosons and to the SM fermions become the same as the SM Higgs ones when we take the limit $\cos \theta \rightarrow 1$, i.e., $g_{hVV} = g_{hbb} = g_{h\tau \tau} = g_{hct} = g_{hct}^{\text{SM}} (= 1)$ when $\sin \theta = v_{EW}/f \rightarrow 0$ by $f \rightarrow \infty$ with
\( v_{\text{ew}} = 246 \text{GeV} \) fixed. Examining Eq. (3.1), we see that the couplings of \( h^0_t \) to the \( W \) and \( Z \) bosons deviate from the SM Higgs ones by \( \kappa_V \equiv g_{hVV}/g_{hVV}^{\text{SM}} = \cos \theta \) where \( V = W \) and \( Z \). The current LHC data give the constraint on \( \kappa_V \) to be \( \kappa_V > 0.94 \) at 95\% C.L. for the 126 GeV Higgs boson \([9]\). Therefore, we obtain the following constraint on the angle \( \theta \): \( \sin \theta < 0.34 \). This bound combined with the mass relation in Eq. (2.3) constrains the \( A^0_t \) mass to be \( m_{A^0_t} \gtrsim 370 \text{GeV} \). The \( A^0_t \) does not couple to the \( W \) and \( Z \) bosons due to the CP-symmetry and couplings to other SM particles are generically suppressed by \( \sin \theta (\lesssim 0.34) \). Hence the \( A^0_t \) is distinguishable from that of the SM-like Higgs boson in the high-mass SM Higgs boson search at the LHC. More detailed LHC study is discussed in \([10]\).

4. Summary

We introduced a model which has a naturally light composite Higgs boson, \( t\text{Higgs} \), to be identified with the 126 GeV Higgs in the spirit of the top quark condensation based on \([1]\). The \( t\text{Higgs} \), a bound state of the top quark and its flavor (vector-like) partner, emerges as a pseudo NGB. The coupling properties of the \( t\text{Higgs} \) are shown to be consistent with the currently available data reported from the LHC.

References


