Top Quark Condensate in Grand Unified Theories

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We propose a top quark condensate scenario embedded in grand unified theories (GUTs), stressing that the gauged Nambu-Jona-Lasinio model has a nontrivial continuum limit (“renormalizability”) under certain condition which is actually satisfied in all sensible GUTs with simple group. The top quark mass prediction in this scenario is shown to be insensitive to the ultraviolet cutoff \(\Lambda\) thanks to the “renormalizability”. We also discuss a possibility to reduce the top mass prediction in this scenario.

0 Introduction

The large top quark mass \(m_t \approx 175\text{GeV}\) observed by CDF and D0 experiments\(^1\) is an important clue to explore physics beyond the standard model. The top quark is coupled to the electroweak symmetry breaking sector with coupling strength proportional to its large mass. It therefore inevitably influences the dynamics of the spontaneous electroweak symmetry breaking. Top quark condensate scenario\(^3\) is one of the most exciting possibilities in this direction. In this scenario the elementary Higgs field in the standard model is replaced by newly introduced Nambu-Jona-Lasinio (NJL)-type four-fermion interactions between ordinary quarks and leptons. Among these four-fermion interactions, one NJL coupling is assumed to be supercritical, causing dynamical electroweak symmetry breaking. The quark associated with the supercritical NJL coupling is identified as the heaviest quark, i.e., the top quark. This scenario is also referred to as “Top-Mode” Standard Model in contrast to the conventional “Higgs-Mode” scenario of electroweak symmetry breaking.

The top condensate scenario suffers from theoretical and phenomenological difficulties, however. Theoretical one arises from non-renormalizability of the

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NJL interaction\(^7\), which forces us to introduce an ultraviolet cutoff \(\Lambda\) to regularize loop integrals. The cutoff \(\Lambda\) is the scale where non-renormalizable NJL interaction is resolved and new physics takes the place of the NJL interaction above \(\Lambda\). The precise prediction from the top condensate scenario, therefore, depends on assumptions made for new physics behind the NJL interaction. Many references are assuming a sharp momentum cutoff and absence of higher dimensional interactions other than NJL. There is no physical reason for these assumptions, however. Moreover, the predictions from pure NJL model are sensitive to these assumption\(^7\),\(^8\),\(^9\).

This scenario also tends to predict too heavy top quark which is not acceptable phenomenologically. Since the top quark mass prediction is a decreasing function of the cutoff \(\Lambda\), we need to assume very large \(\Lambda\) to avoid too heavy top quark. Even for an extremely large cutoff \(\Lambda = 10^{19}\text{GeV}\), however, the prediction with the previous assumptions gives too large top quark mass\(^7\), \(m_t \approx 220\text{GeV}\), still incompatible with the observed one.

Topcolor\(^10\), which replaces the non-renormalizable NJL interaction with a broken topcolor gauge interaction, is a possible candidate to remedy these difficulties. Since the underlying dynamics behind NJL is specified in this model, there is no conceptual problem arising from the non-renormalizability. It is also claimed that the topcolor model can avoid serious fine tuning problem by introducing topcolor dynamics at the electroweak scale. Since the topcolor gauge group can be incorporated into technicolor model of dynamical electroweak symmetry breaking (topcolor-assisted technicolor\(^{11}\),\(^{12}\),\(^{13}\)), the existence of techni-fermion condensate can reduce the top quark mass prediction significantly. However, there is no dynamics which can naturally explain why the scale of the topcolor breaking coincides with the electroweak scale. The naturalness problem is therefore solved in a rather incomplete manner in the topcolor model. Moreover, the topcolor-assisted technicolor would have the notorious problems of technicolor models, i.e., the large positive \(S\)-parameter, \(\Delta \rho\) and the excess of the flavor-changing-neutral-currents.

In this talk, we propose a new class of top quark condensate scenario in which the top condensate is embedded in grand unified gauge theories (GUTs), i.e., top-mode GUTs.\(^\star\) The point is that the gauged NJL model (NJL plus gauge interactions) under certain condition is renormalizable in the sense that it has a nontrivial continuum theory in the \(\Lambda \to \infty\) limit\(^{16}\),\(^{19}\). Such a condition is pointed out\(^{16}\),\(^{19}\) to be equivalent to existence of the Pendleton-Ross-type infrared fixed point\(^{20}\) in the low energy effective (gauged) Yukawa theory, which is actually realized in all sensible GUTs with simple groups.\(^\star\) Thus we observe that the NJL model becomes “renormalizable” when cou-
pled to GUTs. Hence in the class of models we propose, the ambiguities of prediction caused by the non-renormalizability are washed out by a strongly attractive infrared fixed point of the renormalization group equation of GUT-Yukawa coupling. We have therefore less ambiguity than the conventional top-mode scenarios. We also point out a potential possibility to decrease the top quark mass prediction in this model.

1 “Renormalizability” of the gauged NJL model

We here consider the “renormalizability” of the gauged NJL model using the original formulation of the top condensate, i.e., the Schwinger-Dyson gap equation and the Pagels-Stokar formula for the decay constant of the NG boson. The renormalization group formulation will also be used to discuss the “renormalizability.”

Let us start with the pure NJL-type interaction in the top-mode Lagrangian:

$$\mathcal{L} = \frac{G_t}{2} (\bar{q}_L t_R) (\bar{t}_R q_L).$$

(1)

The gap equation can be written as

$$m_t = \frac{N_c}{8\pi^2} G_t m_t \left( \Lambda^2 - m_t^2 \ln \frac{\Lambda^2}{m_t^2} \right),$$

(2)

with $\Lambda$ being ultraviolet cutoff. We obtain a nontrivial solution $m_t \neq 0$ only if the NJL coupling strength exceeds the critical value:

$$G_t > G_{\text{crit}} = \frac{8\pi^2}{N_c} \frac{1}{\Lambda^2}.$$  

(3)

Although we need a fine-tuning of $G_t$ to obtain finite $m_t$, the second order phase transition property of Eq. (2) enables us to take such a fine-tuning. The decay constant of Nambu-Goldstone boson can be evaluated in terms of the top quark mass and the ultraviolet cutoff (Pagels-Stokar formula):

$$v^2 = \frac{N_c}{8\pi^2} m_t^2 \ln \frac{\Lambda^2}{m_t^2}.$$  

(4)

We note here that it is impossible to take $\Lambda \to \infty$ limit keeping $v$ and $m_t$ finite. This is nothing but a consequence of the non-renormalizability of the pure NJL model. In other words the top quark mass $m_t$ vanishes in $\Lambda \to \infty$ limit if we keep the decay constant $v$ finite. This behavior indicates triviality of the effective Yukawa interaction $m_t/v$. Prediction of the pure NJL model is
therefore very sensitive to the value of Λ, the exact definition of the ultraviolet cutoff Λ and the existence of higher dimensional operators, i.e., the details of the physics at the ultraviolet cutoff region.

This situation drastically changes if there exists gauge interaction in addition to the NJL (gauged NJL model): The theory becomes “renormalizable” in the sense that it has a finite continuum theory in the limit Λ → ∞. To illustrate this phenomenon, we first consider QCD effects neglecting its running of coupling strength. Unlike the case of pure NJL model, the top quark mass as a solution of the gap equation becomes a function of momentum scale in the gauged NJL model. The asymptotic solution of the gap equation is then given by:

\[ \Sigma_t(p^2) \simeq m_t \left( \frac{p^2}{m_t^2} \right)^{-\frac{1}{2}} \frac{1 - \sqrt{1 - \alpha/\alpha_c}}{2}, \]  

(5)

with \( m_t \) being the on-shell mass. By using this high energy behavior, we may write the decay constant of the NG boson as

\[ v^2 \simeq \frac{N_c}{8\pi^2} \int_{m_t^2}^{\Lambda^2} dp^2 \frac{\Sigma_t(p^2)}{p^2} = \frac{N_c}{8\pi^2} \frac{m_t^2}{1 - \sqrt{1 - \alpha/\alpha_c}} \left[ 1 - \left( \frac{m_t^2}{\Lambda^2} \right)^{1 - \sqrt{1 - \alpha/\alpha_c}} \right]. \]  

(6)

Unlike the pure NJL model we can take the Λ → ∞ limit keeping both \( v \) and \( m_t \) finite and non-zero. The prediction of the gauged NJL model is therefore insensitive to the ultraviolet cutoff, i.e., the detail of physics at the cutoff region. In other words, the gauged NJL model can be “renormalized” in the sense of its ultraviolet cutoff insensitivity. We can show the “renormalizability” in this sense even in the symmetric vacuum.

This observation, however, might depend on our over-simplification on the gauge interaction, i.e., the non-running behavior of the gauge coupling strength. We therefore need to study the dynamics of the gauged NJL model including the running effects. We parametrize one-loop running of QCD gauge coupling:

\[ \alpha(p^2) = \frac{\pi}{2} \frac{A}{\ln(p^2/\Lambda^2_{QCD})}, \quad A = \frac{24}{33 - 2N_f}, \]  

(7)

Here large \( A \) implies slowly running gauge coupling. The asymptotic solution of the gap equation is given by:

\[ \Sigma_t(p^2) \simeq m_t \left( \frac{\alpha(p^2)}{\alpha(m_t^2)} \right)^{A/2} \]  

(8)
which leads to
\[ v^2 \approx \frac{N_c}{8\pi^2} \int_{m_t^2}^{\Lambda^2} \frac{d^2 \Sigma_d(p^2)}{p^2} = \frac{N_c}{16\pi} \frac{m_t^2}{\alpha(m_t^2)} \frac{A}{A-1} \left[ 1 - \left( \frac{\alpha(\Lambda^2)}{\alpha(m_t^2)} \right)^{A-1} \right]. \quad (9) \]

We note here that we can take the \( \Lambda \to \infty \) limit keeping both \( m_t \) and \( v \) finite if \( A > 1 \), i.e., sufficiently slow running coupling. Unlike the analysis within non-running gauge coupling, however, the cutoff scale physics decouples only through logarithmic suppression.

The dynamics of the gauged NJL model can also be analyzed in terms of a low energy induced gauged Yukawa model by imposing a “compositeness condition” on its renormalization group equation. We can further confirm our observation on the “renormalizability” of the gauged NJL model by using this technique.

The one-loop renormalization group equation of the top quark Yukawa coupling is given by
\[ (4\pi)^2 \mu \frac{d}{d\mu} g_t = (N_c + 3/2) g_t^3 - 8g_t g_3^2, \quad (4\pi)^2 \mu \frac{d}{d\mu} g_3 = -\frac{8}{A} g_3^3, \quad (10) \]
with \( N_c = 3 \). Here we have neglected effects of the \( SU(2)_W \) and \( U(1)_Y \) gauge interactions for simplicity. The compositeness condition, i.e., the absence of the Higgs kinetic term at \( \Lambda \) is described as
\[ 1/g_t^2(\mu = \Lambda) = 0. \quad (11) \]
It is easy to solve Eq. (11) and Eq. (11):
\[ g_t^2(\mu) = \frac{8(A - 1)}{A(N_c + 3/2) g_3^2(\mu) - g_3^2(A - 1)(\Lambda)}. \quad (12) \]
which leads to
\[ v^2 \approx \frac{2m_t^2}{g_t^2(\mu = m_t)} = \frac{N_c + 3/2}{16\pi} \frac{m_t^2}{\alpha(m_t^2)} \frac{A}{A-1} \left[ 1 - \left( \frac{\alpha(\Lambda^2)}{\alpha(m_t^2)} \right)^{A-1} \right]. \quad (13) \]
This expression almost coincides with the expression Eq. (9) except for a small difference coming from subleading effects in the large \( N_c \) expansion.

2 Top quark mass prediction

In the minimal version of the top condensate scenario, \( A \) is given by
\[ A = \frac{24}{33 - 2N_f} = \frac{8}{7} > 1, \quad N_f = 6. \quad (14) \]
We can therefore naively expect decoupling of the cutoff physics. Unfortunately this is not true quantitatively in the minimal scenario. The decoupling is controlled by the logarithmic suppression

\[
\frac{\ln(m_t^2/\Lambda_{QCD}^2)}{\ln(\Lambda^2/\Lambda_{QCD}^2)} A^{-1} \to 0 \quad \text{as} \quad \Lambda \to \infty,
\]

which is not enough due to its very slow convergence. Actually even if we take \(\Lambda \sim M_{pl}\) we obtain \(\ln(m_t^2/\Lambda_{QCD}^2)/\ln(\Lambda^2/\Lambda_{QCD}^2)\) \(A^{-1} \simeq 0.8\) due to small \(A = 1/7 \ll 1\).

We therefore conclude that decoupling of the cutoff physics due to the “renormalizability” of the gauged NJL model, is insufficient in the minimal scenario. The prediction of the minimal top condensate therefore relies on the assumptions made for the cutoff physics.

3 “Renormalizability” of gauged NJL model in GUTs

The top condensate scenario embedded in a grand unified gauge theory might be a viable possibility to solve these problems. In fact, it was shown by Vaughn\(^\text{21}\) that the condition \(A > 1\) is satisfied in all sensible GUTs with simple group, e.g., \(SU(5), SO(10), E_6\), etc.. Thus the “renormalizability” of the gauged NJL model is naturally realized in GUT scenarios.

Let us first consider the minimal \(SU(5)\) model as an example. The top quark Yukawa coupling above GUT scale obeys the renormalization group equation:

\[
(4\pi)^2 \mu \frac{d}{d\mu} g_t = 6 g_t^3 - \frac{108}{5} g_t g_5^2,
\]

with \(g_5\) being \(SU(5)\) gauge coupling strength. The renormalization group equation of \(g_5\) is given by

\[
(4\pi)^2 \mu \frac{d}{d\mu} g_5 = -\frac{108}{5A} g_5^3 = -\left(\frac{55}{3} - \frac{5}{6} N_H - \frac{4}{3} N_g \right) g_5^3,
\]

with \(N_g = 3\) and \(N_H = 1\) being the number of generations and the number of 5-dimensional Higgs field, respectively. Here we have assumed \(SU(5)\) breaks into the standard model gauge group by elementary Higgs field of 24 representation.\(^\text{24}\)

\(^c\) The predictions of the minimal top condensate is very stable for \(\Lambda \simeq 10^{15} \sim 10^{19}\) GeV due to small \(A = 1/7\), however. It can be understood by the “quasi”-infrared fixed point.\(^d\)

\(^d\) If we consider dynamical breaking of \(SU(5)\), the renormalization group equation of \(g_5\) can be modified.
We note here $A = 81/50 > 1$, suggesting the existence of the infrared fixed point of Yukawa coupling. We can therefore take the limit $\Lambda \to \infty$ with the “compositeness condition” keeping finite Yukawa coupling. We also note that the deviation of $A$ parameter from unity, $A - 1 = 31/50$, is reasonably large, which implies relatively fast decoupling of the cutoff physics above the GUT scale.

4 Top condensate in grand unified theories

The present measurement of Weinberg angle is not consistent with the minimal $SU(5)$ GUT. We thus need to introduce non-minimal unification models. Since we are dealing with the dynamical electroweak symmetry breaking, we restrict ourselves within models which do not contain elementary scalar particles below the GUT scale. Particularly, we do not consider SUSY extension of $SU(5)$ model.

Here, we discuss effects of extra fermions in the $SU(5)$ GUT. The minimal extension in this direction was given by Murayama and Yanagida, who introduced extra vector-like quark $Q$ with $(3, 2)_{1/6}$ representation of the standard model gauge group. The renormalization group equations of the standard model gauge couplings are modified above the mass of the vector like quark $Q$:

\[
\begin{align*}
(4\pi)^2 \mu \frac{d}{d\mu} g_1 &= - \left( -\frac{1}{10} N_H - \frac{4}{3} N_g - \frac{1}{15} N_Q \right) g_1^3, \\
(4\pi)^2 \mu \frac{d}{d\mu} g_2 &= - \left( \frac{22}{3} - \frac{1}{6} N_H - \frac{4}{3} N_g - N_Q \right) g_2^3, \\
(4\pi)^2 \mu \frac{d}{d\mu} g_3 &= - \left( \frac{33}{3} - \frac{4}{3} N_g - \frac{2}{3} N_Q \right) g_3^3,
\end{align*}
\]

with $N_H = 1$ and $N_Q = 2$ being the numbers of Higgs and extra quark species, respectively. Here $N_Q = 2$ implies a pair of extra vector-like quarks. We can achieve grand unification of gauge couplings by taking $M_Q \sim O(10^6) GeV$.

The $(3, 2)_{1/6}$ representation can be embedded in 10 representation of $SU(5)$. The renormalization group equation above the GUT scale can be written as

\[
(4\pi)^2 \mu \frac{d}{d\mu} g_5 = - \left( \frac{55}{3} - \frac{4}{3} N_g - \frac{5}{6} - \frac{1}{6} N_H - N_Q \right) g_5^3.
\]

It is now straightforward to evaluate the top quark mass prediction in the top quark condensate scenario embedded in this particular GUT model. We obtain $m_t \approx 224 GeV$ for $\Lambda = 10^{19} GeV$. Here we have assumed $\alpha_3(\mu = M_Z) = 0.118$. Unlike the minimal top condensate, we can take $\Lambda \to \infty$ limit and
obtain $m_t \simeq 201\text{GeV}$ for $\Lambda = \infty$. The prediction becomes cutoff insensitive thanks to the "renormalizability" of the gauged NJL model.

If we allow fine tuning of NJL interaction of bottom quark, the low energy effective theory becomes a two-Higgs-doublet model. In this model, the renormalization group equation of Yukawa coupling is modified and the prediction of $m_t$ is reduced, $m_t \simeq 193\text{GeV}$ for $\Lambda = \infty$. We also note that the prediction of $m_t$ is sensitive to $\alpha_3(\mu = M_Z)$ as indicated in Eq.(13). For $\alpha_3(M_Z) = 0.110$ we obtain $m_t \simeq 196\text{GeV}$, $188\text{GeV}$ for $\Lambda = \infty$ in one- and two-Higgs doublet models, respectively. The $m_b/m_\tau$ ratio can also be calculated in our framework. We find $m_b \simeq 4.5\text{GeV}$ for $\alpha_3(M_Z) = 0.110$ and $\Lambda = \infty$ in the two-Higgs-doublet model, which agrees with the present measurement of $m_b$.

5 Discussion

The top quark mass prediction in the previous section was somewhat heavier than the present measurement. We therefore discuss how our model can be improved to predict lighter $m_t$.

The one-loop renormalization group equation of the effective gauged Yukawa interaction above the GUT scale can be parametrized by

$$
(4\pi)^2 \mu \frac{d}{d\mu} g_t = \gamma g_t \left[ B g_t^2 - g_b^2 \right], \quad (4\pi)^2 \mu \frac{d}{d\mu} g_5 = -\frac{\gamma}{A} g_5^2.
$$

The model described in the previous section corresponds to the set of parameters:

$$
A = \frac{81}{50}, \quad B = \frac{15}{54}, \quad \gamma = \frac{108}{5}.
$$

The solution of the compositeness condition is given by

$$
g_t^2(\mu = M_{\text{GUT}}) = \frac{2\pi}{B} \alpha_5(\mu = M_{\text{GUT}}) \frac{A - 1}{A} \left[ 1 - \left( \frac{\alpha_5(\mu = \Lambda)}{\alpha_5(\mu = M_{\text{GUT}})} \right)^{A-1} \right]^{-1}.
$$

Since we want to construct models which are free from the cutoff ambiguity, the coefficient $A$ should be large enough. We therefore obtain

$$
g_t^2(\mu = M_{\text{GUT}}) \simeq \frac{2\pi}{B} \alpha_5(\mu = M_{\text{GUT}}).
$$

The GUT scale gauge coupling is constrained by the low energy measurements of the gauge coupling strength. Combining the proton decay constraint for $M_{\text{GUT}}$ we obtain the lower bound of the GUT gauge coupling $\alpha_5^{-1}(M_{\text{GUT}}) < \alpha_{1,\text{SM}}^{-1}(\mu > 7 \times 10^{14}\text{GeV}) < 40$. We thus need to construct
models with sufficiently large coefficient $B$ to obtain the top quark mass prediction consistent with the present measurement. Actually such a large $B$ can be implemented in our scenario by a minor extension. The gauge symmetry allows Yukawa interaction between the extra vector-like quark $Q$, $\tilde{5}$ Higgs field and the $\tilde{5}$ fermion. If we admit the existence of such a Yukawa coupling, the coefficient $B$ can be enhanced, leading to a lighter top quark mass, where we have assumed that the newly introduced Yukawa coupling is proportional to the top quark Yukawa coupling.

It is also possible to enhance $B$ parameter by introducing additional scalar field and its Yukawa coupling above the GUT scale. To make the $m_t$ prediction consistent with the present measurement $m_t \approx 175$GeV, we need to construct a GUT model with large $B$ parameter, $B \approx 0.7$. Here we have assumed Eq.(20) and renormalization group equation of the $SU(5)$ GUT containing extra vector-like quarks below the GUT scale.

We have discussed a top quark condensate scenario embedded in an $SU(5)$ GUT, stressing the “renormalizability” of the gauged NJL interaction. The top quark mass prediction in this scenario is shown to be rather insensitive to the ultraviolet cutoff $\Lambda$. This result can be considered as a consequence of the “renormalizability” of the gauged NJL interaction.

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