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Heavier fermions and fine-tuning problem in top-quark condensate scheme

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The analysis in the bubble approximation indicates that a heavier quark-lepton generation with the degenerate mass m_U in the range 163–353 GeV added in the top-quark condensate scheme could make the momentum cutoff Λ come down to 10^6–5×10^3 GeV when $m_t=160$ GeV is taken. This could greatly alleviate the fine-tuning problem. The Higgs boson ϕ_S^0 will obey the mass constraint $2m_t < m_{\phi_S^0} < 2m_U$. The maximal number of the allowed heavy fermion generations is estimated and the possible compositeness origin of the effective four-fermion interactions is discussed.

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In the top-quark condensate scheme [1,2], where the electroweak gauge group $SU_L(2)\times U_Y(1)$ is dynamically broken through the Nambu–Jona-Lasinio (NJL) mechanism [3], an accepted unsatisfactory situation is that the coupling constant must be finely tuned when the gap equation is resolved. Actually, the fine-tuning problem originates from the need for a very large (much larger than the t-quark mass m_t) momentum cutoff Λ of the loop integration and this need again comes from the limitation imposed by the basic relation in the scheme [2]:

$$\frac{4\sqrt{2}\pi^2}{3G_F}=m_t^2\ln\frac{\Lambda^2}{m_t^2}\ , \qquad (1)$$

where G_F is the Fermi constant. As indicated in Ref. [2], when $m_t=165$ GeV, the momentum cutoff Λ must be 10^{15} GeV, up to the energy scale of grand unification. The fact that so high of a Λ is needed reflects the fact that the top quark mass m_t is still a little small for the weak interaction scale $G_F^{-1/2}$. A possible way to decrease Λ is by adding to the right-hand side of Eq. (1) the contribution of some extra heavier fermions other than the top quarks. However, in doing so we must keep the theory in the minimal dynamical breaking scheme of the electroweak gauge group so as to assure that the great success of the standard model could be maintained mostly. Therefore, it is necessary to extend the NJL mechanism of minimal electroweak gauge symmetry breaking from one top flavor to many fermion generations, where the term "one generation" will be specified to the fermions in an $SU_L(2)\times U_Y(1)$ flavor doublet [a left-handed $SU_L(2)$ doublet and two right-handed $SU_L(2)$ singlets] and in a definite representation of some colorlike group G_c.

Fortunately such extensions have been proven to be successful in the bubble approximation first for one generation of fermions [4] and then for n generations of fermions [5]. This makes it possible to put the above idea into effect. Let us consider a model with two generations of fermions without bare masses, i.e., the simplest many generation extension of the NJL mechanism. These fermions are denoted by

$$Q_{\alpha L}=\begin{bmatrix} U_\alpha \\ D_\alpha \end{bmatrix}_L,\quad Q_{\alpha R}=U_{\alpha R},D_{\alpha R}\quad (\alpha=1,2)\ , \qquad (2)$$

with the identifications $(U_1,D_1)\equiv(t,b)$ and $(U_2,D_2)\equiv(U,D)$; i.e., one of the two generations is identified with the ordinary (t,b) quarks and the other one with some exotic (U,D) fermions. The (U,D) fermions could be the fourth generation of quarks and leptons which is assumed not to mix in the Cabibbo-like pattern with the (t,b) generation (e.g., forbidden by some horizontal symmetry) or they are the exotic fermions in the $SU_c(3)$ 3-plet or 6-plet or 8-plet [6] representation or the technifermions [7,8]. Their Y charges can be assigned as the ones of the ordinary quarks and leptons if they are the fourth generation of quarks and leptons, or as

$$Y_{U_L}=Y_{D_L}=0,\quad Y_{U_R}=1=-Y_{D_R} \qquad (3)$$

if they are the exotic fermions in the $SU_c(3)$ triplet. In this way we will avoid the $SU_L(2)\times U_Y(1)$ gauge anomaly [8] and the complexity caused by a possible Cabibbo-like mixture between the (t,b) and the (U,D) generation as well [5]. Our discussions will still be conducted in the bubble approximation and the results will be able to reflect the essential feature of such a kind of model.

It is known from Ref. [5] that when $n=2$ the low energy effective four-fermion Lagrangian corresponding to the minimal dynamical breaking of the electroweak gauge group may be written as

$$\mathcal{L}_{4F}=\sum_{\alpha,\beta=1}^{2}\Big\{g_{U_\beta U_\alpha}(\bar{Q}_{\beta L}U_{\beta R})(\bar{U}_{\alpha R}Q_{\alpha L})+g_{D_\beta D_\alpha}(\bar{Q}_{\beta L}D_{\beta R})(\bar{D}_{\alpha R}Q_{\alpha L})+[g_{U_\beta D_\alpha}(\bar{U}_{\beta R}Q_{\beta L}^T i\sigma_2)(\bar{D}_{\alpha R}Q_{\alpha L})+\text{H.c.}]\Big\} \qquad (4)$$

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where σ_2 is the Pauli matrix and T means the transposition of an $SU_L(2)$ spinor. The four-fermion coupling constants satisfy the minimal Higgs condition

$$g_{Q'Q}=g_{Q'Q'}^{1/2}g_{QQ}^{1/2}, \quad Q,Q'=U_\alpha,D_\alpha \ (\alpha=1,2) . \tag{5}$$

Hence we have only the four independent real and nonnegative coupling constants $g_{QQ}[Q=U_\alpha,D_\alpha \ (\alpha=1,2)]$ altogether.

Suppose that $\mathcal{L}_{4F}$ will lead to the formation of the G_c-invariant vacuum condensates $\langle \bar{Q}Q \rangle$ and the generation of the dynamical masses m_Q; we will obtain the gap equation

$$\sum_Q g_{QQ}I_Q=1 , \tag{6}$$

and the three independent relations between the ratios of the masses and the coupling constants,

$$m_Q/m_{Q'}=g_{QQ}^{1/2}/g_{Q'Q'}^{1/2}, \quad Q,Q'=U_\alpha,D_\alpha \ (\alpha=1,2) , \tag{7}$$

where the sum of Q is always understood as that $Q=U_\alpha$, $D_\alpha(\alpha=1,2)$ and I_Q is expressed as

$$I_Q=2d_Q(R)\int \frac{i\,d^4l}{(2\pi)^4}\frac{1}{l^2-m_Q^2} = \frac{d_Q(R)\Lambda^2}{8\pi^2}\left[1-\frac{m_Q^2}{\Lambda^2}\ln\frac{\Lambda^2+m_Q^2}{m_Q^2}\right] \tag{8}$$

and $d_Q(R)$ is the dimension of the G_c representation of the Q fermions. It is clear that we may have the relation

$$I_Q/I_{Q'}=d_Q(R)m_Q^2/d_{Q'}(R)m_{Q'}^2 . \tag{9}$$

When the minimal Higgs condition (5) and the gap equation (6) are satisfied we will obtain that only a massive neutral Higgs scalar boson ϕ_S^0, a massless neutral pseudoscalar Goldstone boson ϕ_P^0 and two massless charged Goldstone bosons $\psi^{\mp}$ as the real physical modes, emerge from the theory and they are precisely the products of the minimal dynamical breaking of the global group $SU_L(2)\times U_Y(1)$. The propagators and configurations for these bosons may be expressed as

$$\Gamma^{\phi_S^0}(p^2)=iG\Big/\sum_Q g_{QQ}K_Q(p^2)(p^2-4m_Q^2) \quad \text{for } \phi_S^0=\sum_Q G^{-1/2}g_{QQ}^{1/2}(\bar{Q}Q), \tag{10a}$$

$$\Gamma^{\phi_P^0}(p^2)=iG\Big/\sum_Q g_{QQ}K_Q(p^2)p^2 \quad \text{for } \phi_P^0=\sum_Q iG^{-1/2}(-1)^{I_Q^3}g_{QQ}^{1/2}(\bar{Q}i\gamma_5Q) , \tag{10b}$$

$$\Gamma^{\phi^-}(p^2)=iG/2\sum_{\alpha=1}^{2}[g_{U_\alpha U_\alpha}J_{U_\alpha D_\alpha}(p^2)+g_{D_\alpha D_\alpha}J_{D_\alpha U_\alpha}(p^2)]p^2 \quad \text{for } \phi^-=\sum_{i=1,5}\sum_{\alpha=1}^{2}G^{-1/2}\eta_i^\alpha(-1)^{1+\delta_{i5}}(\bar{U}_\alpha\Gamma_iD_\alpha) , \tag{10c}$$

$$\Gamma^{\phi^+}(p^2)=\Gamma^{\phi^-}(p^2) \quad \text{for } \phi^+=(\phi^-)^\dagger \tag{10d}$$

where

$$G=\sum_Q g_{QQ} , \tag{11}$$

$$K_Q(p^2)=-\frac{d_Q(R)}{8\pi^2}\int_0^1 dx\left[\ln\frac{\Lambda^2+M_Q^2}{M_Q^2}-\frac{\Lambda^2}{\Lambda^2+M_Q^2}\right], \quad M_Q^2=m_Q^2-p^2x(1-x) , \tag{12}$$

$$J_{U_\alpha D_\alpha}(p^2)=-\frac{d_Q(R)}{8\pi^2}\int_0^1 dx(1-x)\left[\ln\frac{\Lambda^2+M_{U_\alpha D_\alpha}^2}{M_{U_\alpha D_\alpha}^2}-\frac{\Lambda^2}{\Lambda^2+M_{U_\alpha D_\alpha}^2}\right], \quad M_{U_\alpha D_\alpha}^2=(m_{U_\alpha}^2-p^2x)(1-x)+m_{D_\alpha}^2x , \tag{13}$$

$$\Gamma_1\equiv 1, \quad \Gamma_5\equiv\gamma_5, \quad \eta_1^\alpha=\tfrac{1}{2}(g_{U_\alpha U_\alpha}^{1/2}-g_{D_\alpha D_\alpha}^{1/2}), \quad \eta_5^\alpha=\tfrac{1}{2}(g_{U_\alpha U_\alpha}^{1/2}+g_{D_\alpha D_\alpha}^{1/2}) \tag{14}$$

and I_Q^3 is the third component of the weak isospin of the Q fermions. It can be derived from Eq. (10a) that the mass of the Higgs boson ϕ_S^0 must obey the constraints

$$2(m_Q)_{\min}\le m_{\phi_S^0}\le 2(m_Q)_{\max} \tag{15}$$

where $(m_Q)_{\min}$ and $(m_Q)_{\max}$ are respectively the smallest and the largest masses among the Q fermions. However, $m_{\phi_S^0}$ will be closer to $2(m_Q)_{\max}$ [5].

Once the electroweak gauge interactions are opened the three massless Goldstone bosons will enter the vacuum polarizations of the electroweak gauge bosons and lead to the generation of the masses of the $W^{\mp}$ and Z^0 bosons, i.e., realization of the composite Higgs mechanism. For instance, the inverse propagator for the W boson may be written as

$$\frac{1}{g_2^2}D_{\mu\nu}^{W}(p)^{-1}=i\left[\frac{p_\mu p_\nu}{p^2}-g_{\mu\nu}\right][p^2/\bar{g}_2^2(p^2)-\bar{f}^2(p^2)] \tag{16}$$

where g_2 is the classical $SU_L(2)$ gauge coupling constant and $\bar{g}_2^2(p^2)$ and $\bar{f}^2(p^2)$ are defined by

$$\frac{1}{\bar{g}_2^2(p^2)}=\frac{1}{g_2^2}+\sum_{\alpha=1}^{2}\frac{d_{Q_\alpha}(R)}{8\pi^2}\int_0^1 dx\, x(1-x)\left[\ln\frac{\Lambda^2+M_{U_\alpha D_\alpha}^2}{M_{U_\alpha D_\alpha}^2}-\frac{\Lambda^2}{\Lambda^2+M_{U_\alpha D_\alpha}^2}\right] \tag{17}$$

and

$$\bar{f}^2(p^2)=\sum_{\alpha=1}^{2}\frac{d_{Q_\alpha}(R)}{16\pi^2}\int_0^1 dx[m_{U_\alpha}^2(1-x)+m_{D_\alpha}^2x]\left[\ln\frac{\Lambda^2+M_{U_\alpha D_\alpha}^2}{M_{U_\alpha D_\alpha}^2}-\frac{\Lambda^2}{\Lambda^2+M_{U_\alpha D_\alpha}^2}\right]. \tag{18}$$

The inverse propagator for the neutral gauge bosons will be expressed by a 2×2 matrix; however, we will not put it down here since it will not be used in the following discussions. From Eq. (16) one can derive the equation determining the W-boson mass m_W and then the basic relation in the model [2],

$$\frac{G_F}{\sqrt{2}}=\frac{\bar{g}_2^2(0)}{8m_W^2}=\frac{1}{8\bar{f}^2(0)} \tag{19}$$

which will be our starting point to tackling the fine-tuning problem.

For the sake of simplicity and without loss of essentiality we will assume that

$$m_b=0 \text{ and } m_U=m_D \tag{20a}$$

or equivalently by the relation (7) that

$$g_{bb}=0 \text{ and } g_{UU}=g_{DD}\,. \tag{20b}$$

In this case, when Eqs. (7) and (9) are used the gap equation (6) will take the form

$$\frac{g_{UU}}{g_c}\left[1-\frac{m_U^2}{\Lambda^2}\ln\left[1+\frac{\Lambda^2}{m_U^2}\right]\right]=1; \tag{21}$$

with the denotation

$$g_c=8\pi^2\Big/\left[2+\frac{3}{d_U(R)}\frac{m_t^4}{m_U^4}\right]d_U(R)\Lambda^2\,, \tag{22}$$

the mass constraints (15) will become

$$2m_t\le m_{\phi_S^0}\le 2m_U \tag{23}$$

and the basic relation (19) may be expressed as

$$\frac{4\sqrt{2}\pi^2}{G_F}=2d_U(R)m_U^2\left\{\ln\left[1+\frac{\Lambda^2}{m_U^2}\right]-\frac{\Lambda^2}{\Lambda^2+m_U^2}\right\}+3m_t^2\left\{\ln\left[1+\frac{\Lambda^2}{m_t^2}\right]+\frac{\Lambda^2}{m_t^2}\left[\frac{\Lambda^2}{m_t^2}\ln\left[1+\frac{\Lambda^2}{m_t^2}\right]-1\right]\right\} \tag{24}$$

which is just a modified version of Eq. (1). Here we do not assume that $\Lambda\gg m_U$ and $\Lambda\gg m_t$ so the complete form of Eq. (24) has been maintained. Equation (24) establishes a connection among the parameters m_t, m_U, and Λ. However, m_t must be subject to the constraints imposed by experiment. Such constraints are now $m_t>89$ GeV [9], $m_t=120_{-28}^{+27}$ GeV [10], $m_t\gtrsim150$ GeV [11], and $m_t=140\pm35$ GeV [12], and in general m_t should be greater for a greater Higgs boson mass. Considering these constraints and that in our model the Higgs boson ϕ_S^0 would have a rather large mass owing to the limitation (23), we will take $m_t=160$ GeV. As a result Eq. (24) is left merely as a relation between m_U and Λ. It is found that when m_t is fixed the allowed momentum cutoff Λ is of a definite lower bound. To see this let us rewrite Eq. (24) as

$$By-Dy\left\{\ln\left[1+\frac{1}{y}\right]+\frac{1}{y}\left[\frac{1}{y}\ln(1+y)-1\right]\right\}$$
$$=x\left[\ln\left[1+\frac{1}{x}\right]-\frac{1}{1+x}\right]\equiv f(x) \tag{24a}$$

with the denotations

$$x=m_U^2/\Lambda^2,\quad y=m_t^2/\Lambda^2,\quad 0<x<1,\quad 0<y<1 \tag{25}$$

TABLE I. The relation between the allowed values of the momentum cutoff Λ and the U fermion mass m_U. Four cases with $d_U(R)=3$, $3+1$, 6, and 8 are distinguished. The t quark mass m_t is taken to be 160 GeV. At the critical value Λ_c the relation $m_U=0.68\Lambda_c$ is valid.

	Λ (GeV) / $d_U(R)$	10^{15}	10^{10}	10^8	10^6	10^5	10^4	5×10^3	
m_U (GeV)	3	29	98	133	190	240	350	426	1295
	3+1	25	85	115	163	205	294	353	1092
	6	20	69	93	131	164	231	272	895
	8	17	59	80	113	141	195	228	777

TABLE II. The values of $1-g_c/g_{UU}$ for the acceptable values of Λ.

	$d_U(R)$ \ Λ (GeV)	10^6	10^5	10^4	5×10^3
$1-g_c/g_{UU}$	3	6.19×10^{-7}	6.95×10^{-5}	8.21×10^{-3}	3.58×10^{-2}
	3+1	4.63×10^{-7}	5.20×10^{-5}	6.10×10^{-3}	2.64×10^{-2}
	6		3.45×10^{-5}	4.02×10^{-3}	1.72×10^{-2}
	8			2.99×10^{-3}	1.28×10^{-2}

and

$$B=2\sqrt{2}\pi^2/d_U(R)G_F m_t^2,\quad D=3/2d_U(R)\ . \tag{26}$$

The function $f(x)$ in the right-hand side of Eq. (24a) will be up to its maximum 0.216 in the region $0<x<1$ when $x=0.462$. From this result one can prove that, in the region $0<y<1$, Eq. (24a) is solvable only if $y\geq y_c$ or correspondingly $\Lambda\geq\Lambda_c$. This implies that Λ indeed has the lower bound Λ_c which may be called the critical momentum cutoff. At the scale Λ_c by the definition of x in Eq. (25) the U fermions will have the mass $m_U=0.68\Lambda_c$. Then we may solve Eq. (24) and obtain the values of m_U for variant allowed values of $\Lambda\geq\Lambda_c$. The Fermi constant G_F is taken to be 1.166×10^{-5} GeV^{-2} [13]. In addition, four different cases with $d_U(R)=3$, $3+1$, 6, and 8 are distinguished. This means that we may identify the (U,D) fermions with the exotic fermions respectively in the $SU_c(3)$ 3-plet, $(3+1)$-plet (the fourth generation of quark leptons), 6-plet, and 8-plet representation. The results are listed in Table I.

It can be seen from Table I that as m_U increases the corresponding Λ will decrease. Since $m_U>m_t$ has been demanded and $\Lambda\gg m_U$ is just physically plausible the acceptable values of Λ will be $5\times10^3-10^6$ GeV for $d_U(R)=3$ and $3+1$, $5\times10^3-10^5$ GeV for $d_U(R)=6$ and $5\times10^3-10^4$ GeV for $d_U(R)=8$. In any case the addition of the heavier (U,D) fermions will inevitably lead to a great descent of the momentum cutoff Λ. The corresponding values of m_U in the three cases have been marked in Table I by the underlines. The above results also indicate that the bigger is $d_U(R)$, the smaller is the range of the acceptable values of Λ. In other words, the maximal number of the allowed exotic fermion generations is limited. This number is approximately equal to three if the exotic fermions are in the $SU_c(3)$ 3-plets, two if in the $SU_c(3)$ $(3+1)$-plets (exotic quark-lepton generations), one if in the $SU_c(3)$ $(3+6)$-plet and one if in the $SU_c(3)$ 8-plet.

Because of the great descent of Λ the fine-tuning problem appearing in the gap equation will become no longer so serious. This can be seen from Table II in which the values of $1-g_c/g_{UU}$ are listed for the acceptable values of Λ. When the acceptable momentum cutoff Λ and the corresponding mass m_U are fixed, by Eq. (15) the Higgs boson ϕ_S^0 will get the definite upper bound mass $2m_U$ (of course, it also has the lower bound mass $2m_t$). For example, when $d_U(R)=3+1$, i.e., the (U,D) fermions are the fourth generation of quark leptons, the Higgs boson mass will obey the limitations 320 GeV $\leq m_{\phi_S^0}\leq 588$ GeV for $\Lambda=10^4$ GeV and 320 GeV $\leq m_{\phi_S^0}\leq 410$ GeV for $\Lambda=10^5$ GeV and $m_{\phi_S^0}\simeq 320$ GeV for $\Lambda=10^6$ GeV. These predictions could provide a basis for the experimental test of such a kind of model. Certainly, more accurate estimates of m_t, m_U, and $m_{\phi_S^0}$ depend on the consideration of the full dynamical effects of gauge bosons and the composite Higgs boson and this is a problem to be researched further.

The low energy effective four-fermion interactions used in the above models could originate from some gauge interactions at the higher energy scales than the momentum cutoff Λ. However, noting the fact that the acceptable values of Λ in Table I, $5\times10^3-10^6$ GeV have been at or close to the conventionally assumed compositeness scales of quarks and leptons, a natural question will arise about whether the four-fermion interactions could have the origin of compositeness. This implies that the t quarks and the (U,D) fermions would be considered as composite particles and the four-fermion interactions would be assumed to be the residuals on the composite level of some new strong binding forces among the constituents of these composite particles. To answer this question we must examine if the four-fermion coupling strength, enough for satisfying the gap equation, could be attained by the residual interaction among the composite particles. For convenience of comparison, this strength will be defined by the parameter

$$g_U\equiv g_{UU}\Lambda^2 \tag{27}$$

where U is specified to the most heavy fermion in the model and the momentum cutoff Λ will also be identified with the compositeness scale. Hence g_U/Λ^2 will correspond to the four-fermion coupling constant among the composite fermions. The values of g_U, enough for satisfying the gap equation, can be calculated from Eq. (21)

TABLE III. The values of g_U sufficient for satisfying the gap Eq. (21) when Λ is at the acceptable values.

	$d_U(R)$ \ Λ (GeV)	10^6	10^5	10^4	5×10^3
g_U	3	10.5	12.0	13.0	13.5
	3+1	7.3	8.7	9.6	10.0
	6		5.4	6.2	6.5
	8			4.6	4.8

and the results for the acceptable values of Λ as shown in Table III.

As seen from Table III only if $d_U(R)=6$ or 8 could the values of g_U decrease to near the magnitude of order of unity and be achieved through the effective four-fermion interactions among the composite fermions. This implies that only if the (U,D) fermions are in the $SU_c(3)$ high dimension representations or equivalently there are two heavy quark-lepton generations or three (U,D) fermion flavor-doublets in the $SU_c(3)$ triplets, it is just possible that the four-fermion Lagrangian to realize the NJL mechanism of dynamical breaking of the electroweak gauge group comes from composite fermion models, though the range of the acceptable values of Λ will be more confined in this case.

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