

Chapter 2

Dynamical Electroweak Symmetry Breaking

Our work belongs to the category of beyond-Standard models where the electroweak symmetry is broken dynamically. It is closely related to the main representatives of this category, to the (Extended-) Technicolor models and to the Top-quark Condensation models. That is why we dedicate the whole chapter to present main ideas upon which they are built.

2.1 Technicolor Models

The most natural and the oldest solution to the gauge hierarchy problem follows the scenario already realized in Nature: The scale of QCD and consequent masses of hadrons are arbitrarily small relative to the Planck scale. The clue lies in the asymptotic freedom of a non-Abelian gauge theory, whose example is the QCD. The scale Λ of an asymptotically free theory, above which the gauge coupling parameter g tends to the zero UV fixed point as (1.6)

$$g^2(q^2 > \Lambda^2) \propto \frac{1}{\ln \frac{q^2}{\Lambda^2}}, \quad (2.1)$$

represents a natural cutoff for quantum corrections. The quantum corrections are kept under control and their effect decreases above the cutoff Λ without any reference to the huge Planck scale. Turning the logic around and keeping the reference to the Planck scale, because of the tiny value of the coupling constant at the Planck scale $g(\Lambda_{\text{Pl}}^2) \ll 1$ the huge exponential suppression of scales is in work

$$\Lambda^2 \sim e^{-1/g^2(\Lambda_{\text{Pl}}^2)} \Lambda_{\text{Pl}}^2. \quad (2.2)$$

2.1.1 The Technicolor Dynamics

Technicolor (TC) models [1, 2] assume the existence of a new strong asymptotically free gauge technicolor dynamics $SU(N_{TC})$ characterized by its scale Λ_{TC} and the coupling constant g_{TC} . The technicolor dynamics is acting among new fermions, the techniquark fields $Q(x)$. The techniquarks are electroweakly charged similarly to the usual quarks so that their hard mass terms are forbidden. Instead, the Lagrangian possesses even larger chiral symmetry $G_{\chi TC}$ out of which $SU(2)_L \times U(1)_Y$ electroweak subgroup is gauged. Due to the electroweak dynamics the chiral symmetry $G_{\chi TC}$ is only approximate except for the gauged subgroup. The electroweak symmetry breaking is achieved by a generation of techniquark self-energy $\Sigma_Q(p^2)$ spontaneously breaking the chiral symmetry

$$G_{\chi TC} \longrightarrow H_{\chi TC} \supset U(1)_{em} (\times SU(2)_{custodial}). \quad (2.3)$$

It gives rise to a set of composite pseudo-Nambu–Goldstone fields and three true composite Nambu–Goldstone fields, the technipions $\pi_{TC}^a(x)$, corresponding to the electroweak symmetry breaking. The three technipions couple to the corresponding broken electroweak currents

$$\langle 0 | j_{\pm}^{\mu}(x) | \pi_{TC}^{\pm}(q) \rangle = -iq^{\mu} F_{\pm} e^{-ix \cdot q}, \quad (2.4a)$$

$$\langle 0 | j_0^{\mu}(x) | \pi_{TC}^0(q) \rangle = -iq^{\mu} F_0 e^{-ix \cdot q}. \quad (2.4b)$$

The conservation of the electromagnetic charge $U(1)_{em}$ implies $F_+ = F_-$. The custodial symmetry $SU(2)_{custodial}$ protects the relation $F_{\pm} = F_0 \equiv F_{TC}$.

The technipions $\pi_{TC}^a(x)$ have the same quantum numbers as the QCD pions $\pi_{QCD}^a(x)$, thus they mix. The mixing results in the fact that the observed pion states have an admixture of technicolor states, roughly given by [3, 4]

$$|\pi^a\rangle = \cos \theta_{\pi} |\pi_{QCD}^a\rangle + \sin \theta_{\pi} |\pi_{TC}^a\rangle, \quad (2.5)$$

where $\tan \theta_{\pi} \equiv \frac{f_{\pi}}{F_{TC}}$, where f_{π} is the QCD pion decay constant. The orthogonal states are then the ‘would-be’ Nambu–Goldstone states

$$|\pi_{W,Z}^a\rangle = -\sin \theta_{\pi} |\pi_{QCD}^a\rangle + \cos \theta_{\pi} |\pi_{TC}^a\rangle, \quad (2.6)$$

which are combined with the electroweak gauge bosons to give them masses

$$M_W = \frac{1}{2} g \sqrt{F_{TC}^2 + f_{\pi}^2}, \quad (2.7)$$

$$M_Z = \frac{1}{2} \sqrt{g^2 + g'^2} \sqrt{F_{TC}^2 + f_{\pi}^2}. \quad (2.8)$$

The QCD pion decay constant $f_\pi \simeq 93 \text{ MeV}$ is very small compared to the electroweak scale v , thus it contributes negligibly to the W and Z boson masses. In order to saturate their values, the technipion decay constant has to be

$$F_{\text{TC}} \simeq v \doteq 246 \text{ GeV}. \quad (2.9)$$

Consequently, as $F_{\text{TC}} \gg f_\pi$, the mixing (2.5) and (2.6) of QCD pions and technipions is truly negligible.

The technicolor dynamics can be defined simply by scaling-up of the QCD dynamics. It defines the technicolor scale in terms of the QCD scale $\Lambda_{\text{QCD}} \sim 300 \text{ MeV}$ as

$$\Lambda_{\text{TC}} \sim \sqrt{\frac{3}{N_{\text{TC}}}} \frac{F_{\text{TC}}}{f_\pi} \Lambda_{\text{QCD}} \simeq \sqrt{\frac{3}{N_{\text{TC}}}} 794 \text{ GeV}. \quad (2.10)$$

The scale Λ_{TC} sets the typical magnitude of masses of various bound-states, technibaryons. These states saturate the unitarity of scattering amplitudes of the longitudinally polarized electroweak gauge bosons. In the formula for F_{TC} , (2.15) below, the integral is dominated by low momenta, therefore

$$\Sigma_Q(0) \sim \Lambda_{\text{TC}}. \quad (2.11)$$

The technicolor is a beautiful idea and, like QCD, it has the potential to be the fundamental theory. In a natural way, it stabilizes the electroweak scale with respect to the Planck scale. However, it turns on its own inner scale tensions when applied to the fermion mass generation. In the rest of this section, we will briefly demonstrate that the technicolor dynamics cannot be simply QCD-like otherwise a conflict between magnitude of fermion masses and sufficient suppression of flavor-changing neutral currents (FCNC) pops up.

2.1.2 The Extended Technicolor

Even though the technicolor breaks the electroweak symmetry, it does not provide the mass generation for usual fermions, because it is not coupled to them. An additional interaction which connects usual fermions with the techniquarks is needed. This is provided by embedding $N_F = 3$ flavor copies (families) of ordinary fermions together with corresponding techniquarks in a representation of the extended-technicolor (ETC) gauge dynamics $\text{SU}(N_{\text{ETC}})$ [5, 6], where $N_{\text{ETC}} \equiv N_F + N_{\text{TC}}$.

Masses of flavors differ, hence the ETC gauge symmetry is not a property of the fermion mass spectrum. It has to be spontaneously broken down to its vector-like confining technicolor subgroup according to

$$\text{SU}(N_{\text{ETC}}) \rightarrow \text{SU}(N_{\text{TC}}). \quad (2.12)$$

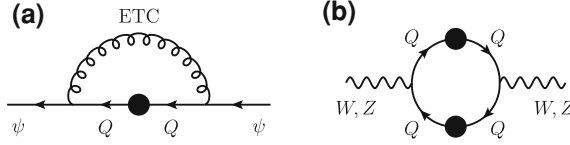


Fig. 2.1 **a** The loop expression for usual fermion mass (2.13). The formation of the techniquark condensate is communicated to the usual fermion sector via the exchange of massive ETC gauge bosons. **b** The loop expression for masses of W and Z . This is the pictorial representation of the Pagels–Stokar formula (2.15)

The ETC gauge symmetry has to be *chiral* otherwise it could not be broken [6]. At this point the technicolor construction loses its beauty as it revives the issue of chiral gauge symmetry breaking. Some mechanism of the spontaneous ETC symmetry breaking has to be added. Clearly, invoking condensing scalars is a possibility but it would negate the main original idea of dynamical symmetry breaking.

In a better way, it is sometimes assumed that the chiral gauge ETC dynamics is strongly coupled and that it *self-breaks*. That is why our model of flavor gauge dynamics described in Chap. 5 is relevant for ETC models. The flavor gauge model can be rephrased as “the ETC model without TC” as its symmetry breaking pattern of completely broken flavor symmetry can be understood in terms of (2.12) as $SU(N_{ETC} = N_F) \rightarrow \emptyset$.

After the spontaneous symmetry breaking at some scale $\Lambda_{ETC} > \Lambda_{TC}$, corresponding ETC gauge bosons acquire masses of order Λ_{ETC} . Because the exchange of the massive ETC gauge bosons provides the FCNC [6] among ordinary fermions, their masses have to be adequately large, setting $\Lambda_{ETC} > 10^3$ TeV. In order to keep the idea of suppression of hierarchy problem, the ETC dynamics is canonically assumed to be again asymptotically free above Λ_{ETC} .

Ordinary fermion masses are generated via the diagram Fig. 2.1a being the pictorial representation of the formula in the Landau gauge

$$\Sigma_\psi(q^2) = -3ig_{ETC}^2 \int \frac{d^4k}{(2\pi)^4} \frac{1}{(k-q)^2 - \Lambda_{ETC}^2} \frac{\Sigma_Q(k^2)}{k^2 - \Sigma_Q^2(k^2)}. \quad (2.13)$$

From this equation the fermion mass can be approximated after the Wick rotation as

$$m_\psi \approx \Sigma_\psi(0) = \frac{g_{ETC}^2}{\Lambda_{ETC}^2} \frac{3}{8\pi^2} \int_0^{\Lambda_{ETC}^2} dx \frac{x \Sigma_Q(x)}{x + \Sigma_Q^2(x)}. \quad (2.14)$$

The technipion decay constant F_{TC} is given by the Pagels–Stokar formula [7] pictorially represented by the diagram in Fig. 2.1b for which we derived the approximate formula in appendix (B.33),

$$F_{\text{TC}}^2 = \frac{N_{\text{TC}}}{8\pi^2} \int_0^{\Lambda_{\text{ETC}}^2} dx \, x \frac{\Sigma_Q(x)^2}{(x + \Sigma_Q^2(x))^2}. \quad (2.15)$$

The scale Λ_{ETC} is used to cutoff the integrals referring to the asymptotic freedom of the ETC dynamics. It is clear that the techniquark self-energy $\Sigma_Q(p^2)$ is instrumental for the electroweak symmetry breaking and mass generation. At first sight the two integrals differ by their sensitivity to the high-momentum details of $\Sigma_Q(p^2)$. Naively the integral for m_ψ in (2.14) is quadratically sensitive, while the integral for F_{TC} in (2.15) is only logarithmically sensitive. That is why $\Sigma_Q(0) \sim \Lambda_{\text{TC}}$ (2.11) in order to get the value of F_{TC} appropriate for the electroweak boson masses.

2.1.3 Technicolor Governed by Infrared Fixed Point

It is therefore of key importance to know the behavior of $\Sigma_Q(p^2)$ in the region of momenta $\Lambda_{\text{TC}}^2 < p^2 < \Lambda_{\text{ETC}}^2$. If dynamically generated, the fermion self-energy has a general Euclidean high-momentum dependence (to be found, e.g., in [8])

$$\Sigma_Q(p^2 \gg \Lambda_{\text{TC}}^2) \sim \frac{\Lambda_{\text{TC}}^3}{p^2} \exp \left[\int_0^t \gamma_m(t') dt' \right], \quad t \equiv \frac{1}{2} \ln \frac{p^2}{\Lambda_{\text{TC}}^2}, \quad (2.16)$$

where $\gamma_m(t)$ is the anomalous dimension of the mass operator. The dependence of the anomalous dimension on t is provided merely through the t dependence of the coupling constant $g_{\text{TC}}(t)$

$$\gamma_m(t) = \gamma_m(g_{\text{TC}}(t)). \quad (2.17)$$

Thus, through the anomalous dimension $\gamma_m(t)$, the evolution of the coupling constant determines the high-momentum dependence of the $\Sigma_Q(p^2)$.

If the technicolor were QCD-like, i.e., its coupling constant were running according to (2.1) all the way above Λ_{TC} , the $\Sigma_Q(p^2)$ would be damping quickly above Λ_{TC} , for illustration see Fig. 2.2 (dotted line). While this would not affect too much the value of F_{TC} , it would, together with FCNC constrain $\Lambda_{\text{ETC}} > 10^3$ TeV, allow only unacceptably small ordinary fermion masses. Therefore the high-momentum enhancement of the $\Sigma_Q(p^2)$ is needed in order to enhance the ordinary fermion masses while keeping the proper value of F_{TC} . This can be achieved if over the range $\Lambda_{\text{TC}}^2 < p^2 < \Lambda_{\text{ETC}}^2$ the technicolor dynamics is governed by a non-trivial fixed point g_{TC}^* . It is the idea pioneered in [9] and further elaborated in [10–14].

Within the range $\Lambda_{\text{TC}}^2 < p^2 < \Lambda_{\text{ETC}}^2$, the QCD-like technicolor is compared to the fixed-point-governed technicolor using the formula (2.16):

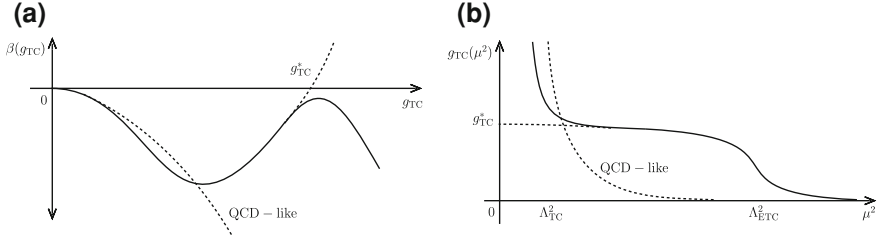


Fig. 2.2 Cartoon **a** of β -function $\beta(g_{TC})$ and **b** of running coupling constant g_{TC} in the ETC theory governed by an infrared fixed point g_{TC}^* (solid line) in comparison with a QCD-like theory (dotted line)

$$\begin{aligned}
 \text{QCD-like: } & g_{TC}^2(t) \sim t^{-1}, \quad \gamma_m(t) \sim g_{TC}^2(t) \sim t^{-1}, \quad \Sigma_Q(p^2) \sim \frac{\Lambda_{TC}^3}{p^2} \left(\ln \frac{p^2}{\Lambda_{TC}^2} \right)^a, \\
 \text{fixed point: } & g_{TC}(t) \sim g_{TC}^*, \quad \gamma_m(t) \sim \gamma_m(g_{TC}^*) = \text{const.}, \quad \Sigma_Q(p^2) \sim \frac{\Lambda_{TC}^3}{p^2} \left(\frac{p^2}{\Lambda_{TC}^2} \right)^{\gamma_m/2},
 \end{aligned}$$

where the exponent a is given by details of the technicolor dynamics and its order of magnitude is $|a| \sim 1$.

The most popular improved ETC models are based on the walking technicolor dynamics [9]. The walking technicolor dynamics is asymptotically free above Λ_{ETC} but, contrary to the QCD, below Λ_{ETC} it is attracted by an approximate infrared fixed point. That slows down the evolution of the coupling constant. Around the scale Λ_{TC} the dynamics finally confines. Over the relevant range of momenta $\Lambda_{TC}^2 < p^2 < \Lambda_{ETC}^2$ the anomalous dimension is approximately constant, $\gamma_m(t) \sim \gamma_m(g_{TC}^*)$. Such gauge theories do exist [15, 16], for illustration see Fig. 2.2 (solid line). The comfortable splitting of Λ_{TC} and Λ_{ETC} is achieved, e.g., within the Minimal Walking Technicolor model [17]. For completeness notice that the same enhancement lifts up masses of pseudo-Nambu–Goldstone bosons, which would be otherwise unacceptably light as well.

2.2 Top-Quark Condensation Models

The mechanism of the electroweak symmetry breaking manifests itself by the fact that both the top-quark mass m_t and the Higgs boson mass M_h are of the same order as the electroweak scale v , which is the scale for masses of the weak bosons, i.e., $M_W = gv/2$ and $M_Z = \sqrt{g^2 + g'^2}v/2$. A simple contemplation leads to the suspicion that both the longitudinal components of the weak vector bosons and the Higgs boson are in fact bound states containing top-quark. This is the idea of Top-quark Condensation models introduced in [18–20].

The top-quark condensation is underlain by some new dynamics characterized by some scale Λ_t and coupling constant g_t . It can be for instance some asymptotically free gauge dynamics whose symmetry gets broken below Λ_t providing masses of the

order $\sim \Lambda_t$ for its gauge bosons. Then the infrared effects of the new dynamics can be parametrized by four-fermion interactions induced by the exchange of the massive gauge bosons, thus weighted by the scale Λ_t . For the top-quark, the four-fermion interaction upon the Fierz identity (C.17a) can be written as

$$(\bar{t}_R \gamma^\mu t_R) \frac{g_t^2}{\Lambda_t^2} (\bar{t}_L \gamma_\mu t_L) \equiv (\bar{t}_L t_R) \frac{g_t^2}{\Lambda_t^2} (\bar{t}_R t_L). \quad (2.18)$$

This interaction is part of the $SU(2)_L \times U(1)_Y$ invariant Lagrangian $\mathcal{L}_t$, which defines the Top-quark Condensation model. In a suggestive form, the four-top-quark interaction (2.18) can be introduced as

$$\mathcal{L}_t \supset \frac{\kappa_t}{\Lambda_t^2} (\bar{t}_L t_R) (\bar{t}_R t_L) \equiv \frac{\kappa_t}{4\Lambda_t^2} [(\bar{t}t)^2 - (\bar{t}\gamma_5 t)^2], \quad (2.19)$$

where we have introduced the coupling parameter $\kappa_t \propto g_t^2$. This is in the analogy with the old proposal of Nambu and Jona-Lasinio (NJL) [21] formulated as the model of a spontaneous chiral $U(1)$ global symmetry breaking.

2.2.1 The Top-Quark Mass

The four-top-quark interaction generates a top-quark mass m_t which can be approximated by a solution of the gap equation

$$m_t = i \frac{\kappa_t}{\Lambda_t^2} \int_{\Lambda_t} \frac{d^4 k}{(2\pi)^4} \frac{m_t}{k^2 - m_t^2} \quad (2.20)$$

$$\stackrel{\text{Wick}}{=} \frac{\kappa_t}{16\pi^2 \Lambda_t^2} \int_0^{\Lambda_t^2} dx \frac{x m_t}{x + m_t^2}. \quad (2.21)$$

The cutoff Λ_t expresses the fact that the underlying dynamics is assumed to become weaker above that scale.

The gap equation (2.20) should be compared with the ETC formula (2.14). The ETC formula is mere *expression* for the top-quark mass $m_t \sim \frac{g_{\text{ETC}}^2}{\Lambda_{\text{ETC}}^2} \langle \bar{Q} Q \rangle$ in terms of the techniquark condensate $\langle \bar{Q} Q \rangle$ given by the integral. On the other hand, (2.20) is the algebraic transcendent *equation* for the top-quark mass

$$m_t = \frac{\kappa_t m_t}{16\pi^2} \left(1 - \frac{m_t^2}{\Lambda_t^2} \ln \frac{\Lambda_t^2 + m_t^2}{m_t^2} \right). \quad (2.22)$$

Within this approach, the top-quark mass generation is a non-perturbative phenomenon which exhibits several characteristic features common to various dynamical mass generation models:

- The trivial solution $m_t = 0$ exists.
- There is a critical value of the coupling parameter $\kappa_t \equiv \kappa_{t,\text{crit.}} = 16\pi^2$ below which only the trivial solution exists and above which also the non-trivial solution $m_t \neq 0$ exists.
- The non-trivial solution for the top-quark mass $m_t \neq 0$ is proportional to the only mass parameter in the model Λ_t through the numerical scaling factor $f(\kappa_t)$ as

$$m_t = f(\kappa_t)\Lambda_t. \quad (2.23)$$

- At the critical value of the coupling constant the scaling factor is a non-analytic function $f(\kappa_t) \simeq \sqrt{1 - \frac{\kappa_{t,\text{crit.}}}{\kappa_t}}$ providing arbitrarily strong amplification $m_t \ll \Lambda_t$ when $\kappa_t \rightarrow \kappa_{t,\text{crit.}}$.

2.2.2 The Electroweak Scale

The original NJL model was immediately applied by authors as a model of pions [22] arising as pseudo-Nambu–Goldstone bosons from spontaneous breaking of chiral isospin symmetry explicitly violated by nucleon bare masses. These were found as massless poles in fermion scattering amplitudes in the pseudo-scalar channel. In the scalar channel a massive pole was found. It corresponds to a composite σ -particle with mass twice as big as the fermion mass. In the top-quark condensation model, the situation is analogous, but the spontaneously broken symmetry is the gauge electroweak symmetry and the Nambu–Goldstone modes are eaten by the W and Z bosons. The NJL treatment of the model leads to the composite Higgs boson with the mass $M_h = 2m_t$.

Once the top-quark mass is generated the electroweak symmetry is broken and the electroweak bosons acquire masses, $M_W^2 = g^2 F_W^2/4$ and $M_Z^2 = (g^2 + g'^2)F_Z^2/4$. The dimensionful factor F_Z is given by the Pagels–Stokar formula, for which we derived approximate formula in appendix (B.33)

$$F_Z^2 = \frac{N_C}{8\pi^2} \int_0^{\Lambda_t^2} dx \, x \frac{\Sigma_t^2(x)}{(x + \Sigma_t^2(x))^2}, \quad (2.24)$$

which is analogous to (2.15). For F_W there is a similar formula. Contrary to the TC models, here the loop integral is formed by top-quark propagators. It expresses the assumption that it is the top-quark condensation which stands behind the electroweak

symmetry breaking. Following the same level of approximation as in (2.20) we use the cutoff Λ_t and $\Sigma_t(p^2) = m_t$ and get the formulae

$$F_Z^2 = \frac{N_c}{8\pi^2} m_t^2 \left[\ln \frac{\Lambda_t^2}{m_t^2} - \frac{1}{2} \right], \quad (2.25a)$$

$$F_W^2 = \frac{N_c}{8\pi^2} m_t^2 \left[\ln \frac{\Lambda_t^2}{m_t^2} + 1 \right], \quad (2.25b)$$

notice that the bigger the Λ_t is, the closer to unity is the ρ -parameter, $\rho \approx 1$, which is defined as usual

$$\rho \equiv \frac{M_W^2}{M_Z^2 \cos^2 \theta_W} = \frac{F_W^2}{F_Z^2}. \quad (2.26)$$

This actually simulates the effect of the custodial symmetry which is not the property of the Lagrangian $\mathcal{L}_t$ (2.19). The Eqs. (2.22) and (2.25a) (neglecting the difference between F_W and F_Z) have to be satisfied simultaneously. By setting $F_Z = v \doteq 246$ GeV and $m_t \doteq 172$ GeV we fix the condensation scale Λ_t from the Eq. (2.25a),

$$\Lambda_t \sim 10^{14} \text{ GeV}. \quad (2.27)$$

The Eq. (2.22) then fixes the coupling constant κ_t

$$\kappa_t \approx (1 + m_t^2/\Lambda_t^2) \kappa_{t,\text{crit.}}. \quad (2.28)$$

This result means enormous fine-tuning of the coupling constant.

The lesson is the following. By assuming some dynamics characterized by the scale Λ_t leading effectively to the four-fermion interaction (2.19), the electroweak symmetry breaking can in principle take place. It turns out however that *the hierarchy problem is not solved*, because numerically Λ_t is pushed very high in order to get the correct values of $M_{W,Z}$ and m_t . The extreme fine tuning now reappears in the value of the coupling parameter κ_t (2.28).

2.2.3 Including the QCD Effect

The situation becomes even much more inconvenient when more sophisticated approximations are used. In particular, the effect of QCD dynamics to the top-quark masses is significant and it should be taken into account. It leads to the momentum dependent self-energy $\Sigma_t(p^2)$ which is monotonically decreasing and given by the well-known formula [23, 24]

$$\Sigma_t(p^2) \simeq m_t \left[\frac{\alpha_s(p^2)}{\alpha_s(m_t^2)} \right]^{4/7}, \quad (2.29)$$

where $\alpha_s(q^2) = g_s^2(q^2)/4\pi$ is the QCD effective charge

$$\alpha_s^{-1}(q^2) = \alpha_s^{-1}(m_t^2) + \frac{7}{4\pi} \ln \frac{q^2}{m_t^2}. \quad (2.30)$$

Using the improved top-quark self-energy (2.29) within the safely simplified Pagels–Stokar formula (2.24)

$$F_Z^2 = \frac{3}{8\pi^2} \int_{m_t^2}^{\Lambda_t^2} \frac{dk^2}{k^2} \Sigma_t^2(k^2) \quad (2.31)$$

we get the improvement of the relation (2.25a)

$$F_Z^2 = \frac{3m_t^2}{2\pi\alpha_s(m_t^2)} \left(1 - \left[\frac{\alpha_s(\Lambda_t^2)}{\alpha_s(m_t^2)} \right]^{1/7} \right). \quad (2.32)$$

2.2.4 The Higgs Boson Mass

The electroweak symmetry breaking driven by the top-quark condensation is accompanied by a composite Higgs boson. The Higgs boson mass can be estimated by the formula [25, 26] similar to (2.31)

$$M_h^2 = \frac{3}{2\pi^2 F_Z^2} \int_{m_t^2}^{\Lambda_t^2} \frac{dk^2}{k^2} \Sigma_t^4(k^2). \quad (2.33)$$

It reflects the momentum dependence of the $\Sigma_t(p^2)$ and thus it is an improvement of the NJL relation $M_h = 2m_t$. Using (2.29), it leads to the expression for the Higgs boson mass

$$M_h^2 = \frac{2m_t^4}{3\pi\alpha_s(m_t^2)F_Z^2} \left(1 - \left[\frac{\alpha_s(\Lambda_t^2)}{\alpha_s(m_t^2)} \right]^{9/7} \right). \quad (2.34)$$

The two Eqs.(2.32) and (2.34) relate masses $M_{W,Z}$, M_h and m_t . The single free parameter is the scale of the new dynamics, Λ_t . As Table 2.1 shows such a simple top-quark condensation scenario is strictly incompatible with the experimental values of $M_{W,Z}$, m_t , and M_h . There are two main problems. First, the top-quark does not weight enough to saturate the electroweak scale. Second, even if it were heavy enough, the

Table 2.1 Results of the top-quark condensation model

	Λ_t [GeV]	10^{10}	10^{19}	10^{42}
(a)	m_t [GeV]	305	252	215
	M_h [GeV]	450	338	257
(b)	F_Z [GeV]	140	168	196
	M_h [GeV]	258	231	205

For different values of the top-quark condensation scale $\Lambda_t = 10^{10}$ GeV (axion), $\Lambda_t = 10^{19}$ GeV (Planck), $\Lambda_t = 10^{42}$ GeV (Landau), and taking $\alpha_s(m_t^2) \simeq 0.11$, using the Eqs.(2.32) and (2.34), we evaluate **a** the M_h and m_t masses while keeping the electroweak gauge boson masses at their experimental values, i.e., $F_Z = 246$ GeV, **b** the M_h and F_Z while keeping the top-quark mass at its experimental value, i.e., $m_t = 172$ GeV

mass of the Higgs boson as the top-quark bound-state would come out larger than the top-quark mass, which is in contradiction with the observation of the Higgs mass $M_h \simeq 125$ GeV.

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