Dynamical electroweak symmetry breaking due to strong Yukawa interactions

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We present a new mechanism for electroweak symmetry breaking (EWSB) based on a strong Yukawa dynamics. We consider an SU(2)L × U(1)Y gauge invariant model endowed with the usual Standard model fermion multiplets and with two massive scalar doublets. We show that, unlike in the Standard model, EWSB is possible even with vanishing vacuum expectation values of the scalars. Such EWSB is achieved dynamically by means of the (presumably strong) Yukawa couplings and manifests itself by the emergence of fermion and gauge boson masses and scalar mass-splittings, which are expressed in a closed form in terms of the fermion and scalar proper self-energies. The ’would-be’ Nambu–Goldstone bosons are shown to be composites of both the fermions and the scalars. We demonstrate that the simplest version of the model is compatible with basic experimental constraints.

PACS numbers: 11.15.Ex, 12.15.Ff, 12.60.Fr
Keywords: Electroweak symmetry breaking, Dynamical mass generation, Strong Yukawa interactions.

I. INTRODUCTION

The majority of current models of electroweak symmetry breaking (EWSB) can be divided into two large classes, according to the mechanism the symmetry breaking is achieved (omitting exotic scenarios such as extra dimensions): (i) The electroweak symmetry is broken at tree level by non-vanishing vacuum expectation value(s) (VEV(s)) of some elementary scalar field(s). This class includes the Standard model (SM) [1] and its various extensions like the two-Higgs-doublet models [2], the minimal supersymmetric Standard model [3], or the little Higgs models [4, 5]. (ii) The electroweak symmetry is broken non-perturbatively by strong dynamics of some new, yet unobserved, interactions. This class includes theories with new strong gauge interactions such as (extended) technicolor [6, 7, 8, 9], top condensate [10, 11], or topcolor [12, 13], as well as effective models of the Nambu–Jona-Lasinio type [14, 15].

In this paper we study a different scenario, which combines features of both classes mentioned above (elementary scalars and strong dynamics): The electroweak symmetry is broken non-perturbatively by strong Yukawa interactions (a similar approach can be found in [16, 17] and, in the context of non-relativistic systems, in [18, 19, 20]). We believe that this option has certain advantages. For example, it allows to distinguish otherwise identical fermions at the level of Lagrangian by different Yukawa couplings and hence avoid dangerous (because unobserved) Nambu–Goldstone (NG) bosons resulting from the spontaneous breaking of the large flavor global symmetries. The idea has already been studied in our previous paper [21] using a prototype model with Abelian gauge symmetry. We have shown that the dynamically generated masses of different fermions may differ by orders of magnitude even with Yukawa coupling ratios of order one. We thus suppose that the present mechanism could shed some new light on the problem of the fermion mass hierarchy in SM.

Encouraged by the previous model study, we investigate in this paper the possibility of EWSB by the same mechanism, considering for simplicity a single family of SM fermions. At this exploratory stage we do not want to make any kind of phenomenological predictions, our aim is merely to demonstrate that the extension of the scenario to the electroweak symmetry brings no new principal problems. In Sec. II we first introduce a model with global SU(2)L × U(1)Y symmetry, which is broken dynamically by the strong Yukawa couplings. In Sec. III we then gauge the symmetry, which induces dynamical EWSB. In Sec. IV we present numerical results and perform a basic check of consistency with precision electroweak measurements. The outlook for future work as well as potential applications in other fields of physics are discussed in the conclusions.

II. THE MODEL

We consider an SU(2)L × U(1)Y invariant Lagrangian, equipped with one generation of the SM fermion multiplets: The left-handed isospin doublets, denoted as ℓL = (νL, eL)T, qL = (uL, dL)T, and the right-handed singlets eR, uR, dR. In addition, we introduce the right-handed neutrino singlet νR in order to generate the neutrino mass. However, we will consider only the Dirac mass here, since the Majorana mass would require a more elaborate formalism. We also introduce two complex scalar doublets, S = (S1, S0)T and N = (N0, N−1)T. Unlike in SM, they have ordinary masses M_S, M_N (i.e.,
\( M_{S,N}^2 > 0 \), which implies they do not condense at tree
level. Hence there is no need for the scalar self-couplings
to stabilize the vacuum and they will for simplicity be
neglected.

Of key importance are the Yukawa couplings of the
fermions to \( S, N \), by assumption responsible for EWSB:

\[
\mathcal{L}_{\text{Yukawa}} = y_{\nu} \ell_L e_R S + y_{\nu} \ell_L \nu_R N + y_d \bar{q}_L d_R S + y_u \bar{q}_L u_R N \\
+ \text{h.c.}
\]  

(1)

The Yukawa couplings can always be made real by a suit-
able change of phases of the fields. For simplicity we do
not consider interactions of the charge conjugated scalar
doublets, \( \bar{S} \equiv i \sigma_2 S^* \) and \( \bar{N} \equiv i \sigma_2 N^* \). Our ‘minimal’
Yukawa interaction Lagrangian in fact turns out to be
sufficient for generating fermion and gauge boson masses.

In the present paper we deliberately neglect renormal-
ization and running of the Yukawa couplings. Within our
model this is decently justified by the fact that thanks to
the structure of the interaction Lagrangian (1), the one-
loop perturbative corrections to the Yukawa couplings
are absent. Throughout the calculation, the couplings
are therefore implicitly assumed to be fixed and renor-
malized at the scale of the scalar masses \( M_{S,N} \).

Of course, once the couplings become large and more-
over the chiral symmetry is spontaneously broken, this
naive argument breaks down. We do not attempt to de-
termine the full nonperturbative flow of the couplings
but merely note that it can be analyzed using the
renormalization-group techniques, as was done for simi-
lar models in Refs. [22, 23]. Our results, in particular the
existence of a critical coupling for spontaneous symmetry
breaking, are in qualitative agreement.

With the above remarks in mind, we assume the full
fermion propagators to have the form (\( \nu, e, u, d \))

\[
\langle \bar{f} \bar{f} \rangle = i \left( \rho - \Sigma_{f,1} - i \gamma_5 \Sigma_{f,2} \right)^{-1},
\]

(2)

where the \( \Sigma \)'s are real functions of \( p^2 \) (the argument will
be often omitted throughout the paper). Notice that the
\( \Sigma \)'s are scalar functions, which means that we also neglect
the perturbative wave function renormalization.

Likewise, for the full scalar propagators we adopt the
Ansatz (\( X = S, N \))

\[
\langle \Phi_X \Phi_X \rangle = i \left( p^2 - M_X^2 \right)^{-1},
\]

(3a)

\[
\langle \Phi_{SN} \Phi_{SN} \rangle = i \left( p^2 - M_{SN}^2 \right)^{-1},
\]

(3b)

written using the Nambu doublet notation \( \Phi_X = (X^{(0)}, X^{(0)*})^T \), \( \Phi_{SN} = (S^{(+)}, N^{(-)*})^T \). The \( \Pi \)'s are
complex functions of \( p^2 \).

The spectrum can now be revealed by looking for
the poles of the full propagators. In the fermionic
case, provided there are non-vanishing self-energies \( \Sigma_f \equiv \Sigma_{f,1} + \Sigma_{f,2} \), one arrives at the simple pole equations
\( p^2 - |\Sigma_f(p^2)|^2 = 0 \). For the neutral scalar doublets \( \Phi_X \)
the pole equation reads \( (p^2 - M_X^2)^2 - |\Pi_X(p^2)|^2 = 0 \).

Its two solutions \( M_{X1,2} \) are to be interpreted as that in-
stead of one electrically neutral, yet complex scalar field
\( X^{(0)} \) with the mass \( M_X \) there are \( \text{two real scalar fields} \)
\( X_{1,2}^{(0)} \) with distinct masses \( M_{X1,2} \) in the spectrum. These
mass eigenstates are apparently linear combinations of
\( \text{Re} X^{(0)} \) and \( \text{Im} X^{(0)} \). Similarly considering the charged
doublet \( \Phi_{SN} \), its pole equation is
\( (p^2 - M_{SN}^2)^2 - |\Pi_{SN}(p^2)|^2 = 0 \). In this case the two solutions \( M_{SN1,2} \) indicate that the charged fields \( S^{(+)} \) and \( N^{(-)*} \) need to be
mixed in order to get the mass eigenstates. In prin-
ciple, charge conservation also allows mixing of \( S^{(0)} \) and
\( N^{(0)} \), but this only appears at the two-loop level.

For later purposes it is convenient to define the ‘anoma-
lous’ propagators,

\[
S_f(p) = \frac{\Sigma_f}{p^2 - |\Sigma_f|^2},
\]

(4a)

\[
D_X(p) = \frac{\Pi_X}{(p^2 - M_X^2)^2 - |\Pi_X|^2},
\]

(4b)

\[
D_{SN}(p) = \frac{\Pi_{SN}}{(p^2 - M_{SN}^2)(p^2 - M_{SN}^2) - |\Pi_{SN}|^2}.
\]

(4c)

It should be emphasized that, once the symmetry is bro-
ken dynamically by the above defined anomalous propa-
gators, nothing any longer prevents the scalars from ac-
quiring nonzero VEVs. In general, they will be induced
by perturbative loop effects. However, since it is our aim
to here to demonstrate the possibility of dynamical EWSB,
not driven by scalar VEVs, we will neglect these VEVs al-
together in our Schwinger–Dyson (SD) equations. Their
self-consistent incorporation will be studied in a future
work.

The self-energies \( \Sigma, \Pi \) are subject to the SD equations,
which we approximate by the following set (see Fig. 1):

\[
\Sigma_U = ig_U^2 \int \frac{d^4k}{(2\pi)^2} D_N(k)S_U^T(k+p) \\
+ig_Uy_D \int \frac{d^4k}{(2\pi)^2} D_{SN}(k)S_D^T(k+p),
\]

\[
\Sigma_D = ig_D^2 \int \frac{d^4k}{(2\pi)^2} D_S(k)S_D^T(k+p) \\
+ig_Dy_D \int \frac{d^4k}{(2\pi)^2} D_{SN}(k)S_D^T(k+p),
\]

\[
\Pi_S = -2i \sum_D N_Cg_D^2 \int \frac{d^4k}{(2\pi)^2} S_D(k)S_D(k+p),
\]

\[
\Pi_N = -2i \sum_U N_Cg_U^2 \int \frac{d^4k}{(2\pi)^2} S_U(k)S_U(k+p),
\]

\[
\Pi_{SN} = -2i \sum_D N_Cg_Uy_D \int \frac{d^4k}{(2\pi)^2} S_U(k)S_D(k+p),
\]

\((N_C\text{ denotes the number of colors). The sum in the last equation is over all isospin doublets, }D = (U, D)^T\). In our case, }D = \ell, q\). The numerical solution of these equations is discussed in Sec. [IV]

### III. GAUGE BOSONS

Since the }SU(2)_L \times U(1)_Y\text{ symmetry is broken down to electromagnetic }U(1)_{em},\text{ the question about the nature of the corresponding NG bosons (to become the longitudinal parts of the }W^\pm, Z\text{ bosons) arises. It turns out that, according to general principles, they can be }seen\text{ as massless poles in the proper vertices of the currents associated with the broken generators [24]. The residua of these poles are proportional to the symmetry-breaking proper self-energies }\Sigma \text{ and }\Pi.\text{ }

Without going too much into technical details (interested reader can find them in [21]), let us only sketch the way leading to computation of the gauge boson masses. In this context it is useful to consider the vertex functions of the broken currents,

\[
G_{\text{DA}}^a(x, y, z) = \langle 0 | T \{ j^a_\mu(x) D(y) D(z) \} | 0 \rangle,
\]

\[
G_{\Phi_a(x, y, z) = \langle 0 | T \{ j^a_\mu(x) \Phi(y) \Phi^\dagger(z) \} | 0 \rangle. \quad (6a)
\]

Here }j^a_\mu (a = 1, \ldots, 4)\text{ are the conserved currents associated with the generators of }SU(2)_L \times U(1)_Y\text{ and }\Phi = \Phi_S, \Phi_N, \Phi_{SN}.\text{ Taking the divergence of the vertex functions we arrive at the Ward–Takahashi (WT) identities for the proper vertices of the gauge bosons with the fermions and scalars. For instance, for the proper vertex }Z\ff\text{ we get (similarly for other gauge bosons and fermions/scalars)

\[
q_\mu T^{\nu}_{Z\ff}(p', p) = i(f \bar{f})^{-1} T_f - \bar{T}_f i(f \bar{f})^{-1}, \quad (7)
\]

\[
T_f = \frac{g}{2\cos \theta_W} \left[ T^3_f (1 - \gamma_5) - 2Q_f \sin^2 \theta_W \right], \quad (8)
\]

with }q = p' - p\text{ denoting the momentum of the incoming gauge boson and }T_f = \gamma_\ell T^3_f \gamma_\ell. (T_3 \text{ and } Q\text{ denote the isospin and charge matrices, diagonal in the flavor space.) These WT identities remain valid even though the corresponding symmetries are spontaneously broken. Consequently, it is easy to see that upon inserting the full propagators [2, 4] in the WT identities and taking the limit }q \to 0\text{, the proper vertices }\Gamma^\mu\text{ associated with the broken symmetries (i.e., }\Gamma^\mu_{W^z}, \Gamma^\mu_{Z})\text{ develop a massless pole. It is to be interpreted as the propagator of an intermediate scalar excitation, bilinearly coupled to the gauge boson – the NG boson. In the case of }\Gamma^\mu_{Z}\ff\text{ its pole part has the form }\langle \Sigma_f \equiv \Sigma_{f1} + i\gamma_5 \Sigma_{f2} \rangle

\[
\Gamma^\mu_{Z\ff\text{pole}}(p', p) = \frac{q^\mu}{q^2} \left[ -\Sigma_f(p')T_f + \bar{T}_f \Sigma_f(p) \right]. \quad (9)
\]

Since the broken symmetry is gauged, the NG bosons decouple from the spectrum and manifest themselves only as the longitudinal polarizations of the }W^\pm \text{ and } Z\text{ bosons. By the Schwinger mechanism [22], these poles in turn give rise to the gauge boson masses, which are determined by the corresponding residues [24]. Within our approximation, the polarization tensors are given by the graphs in Fig. 2 with the insertion of the pole parts of the proper vertices.

As a net result, the gauge boson masses can be expressed via the sum rules

\[
M_Z^2 = \frac{g^2}{\cos^2 \theta_W} \left( I_S + I_N + \sum_f I_f \right), \quad (10a)
\]

\[
M_W^2 = g^2 \left( I_{S,SN} + I_{N,SN} + \sum_D I_D \right), \quad (10b)
\]
with
\[
I_X = \frac{1}{2} \int \frac{d^4k}{(2\pi)^4} k^2 |D_X(k)|^2, \tag{11a}
\]
\[
I_{X,SN} = \frac{1}{4} \int \frac{d^4k}{(2\pi)^4} k^2 D_X(k) D_{SN}(k) \frac{|\Pi_X|^2 + |\Pi_{SN}|^2}{\Pi_X \Pi_{SN}}, \tag{11b}
\]
\[
I_f = -\frac{1}{2} N_C \int \frac{d^4k}{(2\pi)^4} |S_f(k)|^2, \tag{11c}
\]
\[
I_D = -\frac{1}{2} N_C \int \frac{d^4k}{(2\pi)^4} S_U(k) S_D(k) \frac{|\Sigma_U|^2 + |\Sigma_D|^2}{\Sigma_U \Sigma_D}. \tag{11d}
\]

Thanks to the fact that the dynamically generated self-energies are suppressed at high momentum, all integrals above are UV-finite.

IV. RESULTS AND DISCUSSION

The numerical analysis of the (Wick-rotated) SD equations \([5]\) shows that: (i) Non-trivial, UV-finite solutions do exist. (ii) The solutions are found only for relatively large values of the Yukawa coupling constants (of order of tens). (iii) Large ratios of fermionic masses can be accommodated while having the corresponding Yukawa coupling constants of the same order of magnitude.

The point (iii) above is promising in the quest for realistic fermion mass hierarchy. Because of the large parameter space which needs to be scanned this has not been accomplished yet. However, some achievements, which suggest that it should be possible, have been made. First, we accommodated the hierarchy between the lepton and quark doublet. For \(y_e = 63, y_u = 84, y_d = 90\) (and \(M_2^2 = 2, M_3^2 = 1\)) we found \(m_\nu \approx m_c = O(10^{-4})\) and \(m_u \approx m_d = O(10^{-2})\). (Note that all masses are expressed in the units of \(M_N\).) Second, we managed to generate a large hierarchy within one doublet. Considering only the leptons and ‘neglecting’ the quarks (\(y_u = y_d = 0\)), we found \(m_e/m_\nu = O(10^3)\), calculated for \(y_e \approx 50, y_c = 80\) (and again \(M_2^2 = 2, M_3^2 = 1\)). Nevertheless, it should be emphasized that this lepton mass ratio will be significantly enhanced by the seesaw mechanism once the Majorana mass term is taken into account.

The numerical results above were based on the exact solution of the Wick-rotated SD equations. In this numerical procedure one chooses some values of the input parameters (the Yukawa coupling constants and the bare scalar masses), solves the equations and checks whether the resulting spectrum (fermion and gauge boson masses) corresponds to the real values. It is, however, difficult to find in this manner the combination of the parameters, giving realistic spectrum; not only because of the huge-ness of the parameter space, but also since solving the SD equations even for a single set of parameters is rather time-consuming.

Thus, in order to be able to invert the above procedure, i.e., to choose the spectrum first and then find the corresponding parameters of the model, we relaxed the exactness of the solution and solved the SD equations only approximately. For that purpose we chose the Ansatz for the fermion and scalar self-energies in the form of step-functions with a common position of the step \(\Lambda\) (i.e., \(\Sigma_f(p^2) \rightarrow m_f \theta(\Lambda^2 - p^2)\) and \(\Pi_X(p^2) \rightarrow \mu_X \theta(\Lambda^2 - p^2)\) and plugged them into the SD equations. This Ansatz is decently justified by the observation that the self-energies obtained by exact solution of the SD equations exhibit similar behavior: They are nearly constant in a wide range of momenta and then drop fast to zero around a scale common to all particles in the model.

With an Ansatz on the shape of the self-energies, the SD equations can of course no longer be satisfied identically. Because mass generation is a low-energy phenomenon, we demand instead that the left- and right-hand sides are equal at \(p^2 = 0\). This results in a set of several algebraic equations for the unknown parameters of the model, which can be solved much more easily than the full integral equations \([3]\).

According to the sum rules \([10]\), the heavy quark flavors contribute most to the gauge boson masses. Likewise, they also dominate the scalar self-energies \([6-8, 55]\), which tie the individual fermion masses together. We therefore neglected for the moment the lepton doublet and treated the quark one as the \(t\)- and \(b\)-quarks. We kept the gauge boson masses at the physical values. Due to numerical reasons, however, we could only fix the quark masses to the rather large values \(m_t = 405\) GeV, \(m_b = 10\) GeV. Nevertheless, note that the ratio of these masses is still realistic. We then found the following solution:

\[
y_t = 375, \tag{12a}
\]
\[
y_b = 235, \tag{12b}
\]
\[
M_N = 433\text{ TeV}, \tag{12c}
\]
\[
M_S = 283\text{ TeV}, \tag{12d}
\]
\[
\Lambda = 71\text{ TeV}. \tag{12e}
\]

This result shows again some important features. First, the Yukawa couplings are of the same order of magnitude, while keeping the realistic ratio of the quark masses in the doublet. Second, the scalar masses are much higher than the EW scale, which is gratifying, as explained above. Our numerical investigations suggest that if we managed to pull the quark masses down to the physical values, the scalar masses would even increase by a few orders of magnitude.

The ‘cut-off’ \(\Lambda\) also deserves a remark. It corresponds to the energy scale at which the self-energies fall rapidly to zero. The fact that \(\Lambda \ll M_S, M_N\) suggests that the scalar self-energies are practically vanishing for momenta corresponding to the scalar pole masses. Consequently, the scalar mass splitting is negligible and the \(S\)-parameter tends to be very close to zero. (For \(\Pi_S = \Pi_N = \Pi_{SN} = 0\) we would have exactly \(S = 0\).) On
the other hand, since $\Lambda \gg m_t, m_b$, the quark self-energies are almost constant (and non-zero) at the energy scale of fermion masses (under 1 TeV). If we considered all three fermion generations together with a mixing, this would consequently lead to the CKM matrix being significantly close to the unitary form [28].

A. Compatibility with electroweak observables

While the realistic fermion spectrum together with the Yukawa coupling constants can be presumably accommodated, it on the other hand brings the problem how to keep the $\rho$-parameter (defined as $\rho \equiv M_W^2/(M_Z^2 \cos^2 \theta_W)$) close to 1. Note that if we set $y_e = y_e$, $y_u = y_d$ and $M_S = M_N$, the SD equations yield $\Sigma_e = \Sigma_e$, $\Sigma_u = \Sigma_u$ and $\Pi_S = \Pi_N = \Pi_{SN}$ and consequently $\rho = 1$ exactly. (Which is not surprising, since in this case the Lagrangian possesses the custodial symmetry.) Of course, in reality $y_e \neq y_e$, $y_u \neq y_d$, so the custodial symmetry is broken by the fermions. However, for the values $m_e \leq m_e = O(10^{-4})$ and $m_u \leq m_d = O(10^{-2})$ mentioned above, we find $\rho \approx 1.012$. It is thus apparent that one can achieve realistic fermion mass ratios and still keep $\rho$ reasonably close to one. A detailed analysis reveals that key role is here played by the scalars: If their bare masses and the self-energies are of the same order of magnitude (i.e., $M_S \approx M_N$ and $\Pi_S \approx \Pi_N \approx \Pi_{SN}$) so that they approximately conserve the custodial symmetry, they can render $\rho$ close to 1, provided they are heavy enough so that they can overcome the effect of the custodial symmetry breaking in the fermion sector.

The fact that the scalars tend to be heavy is in fact reassuring. First, we do not have to deal with the usual hierarchy problem, that is, radiative instability of the masses of elementary scalars at the TeV scale. Of course, the problem is just postponed to a higher scale. In this sense we understand our model as an effective approach which is valid at energy scales around and below the heavy scalar masses. We do not pretend to have an ultraviolet complete theory.

Second, the new scalars must be heavy enough in order to avoid constraints from flavor-changing neutral currents (FCNC). Even though there are no FCNC in the present model, for we only consider a single family at the moment, anticipating a future extension to all three families of SM fermions, we can make at least a rough, order-of-magnitude estimate. With several fermion families there will be explicit flavor-changing Yukawa interactions, allowing, for instance, for the decay $\mu^- \to e^- S^{(0)}$. The virtual heavy scalar can subsequently decay as $S^{(0)} \to e^+ e^-$. Our Yukawa interactions will therefore induce the flavor-changing muon decay, $\mu^- \to e^- e^+ e^-$, with the amplitude being roughly given by $y^2/M_S^2$. (We assume that in the absence of fine tuning, all Yukawa couplings, including the flavor-changing ones, will be of the same order of magnitude.) The dominant muon decay channel, with branching ratio close to 100%, is $\mu^- \to e^- \bar{\nu}_e \nu_{\mu}$, whose amplitude is analogously proportional to $G_F$. From here we infer the estimate $BR(\mu^- \to e^- e^+ e^-) \sim (y^2/G_F M_S^2)^2$. Taking the current experimental limit [27], $BR(\mu^- \to e^- e^+ e^-) < 10^{-12}$, we find $M_S/y \gtrsim 10^{5.5}$ TeV. Our numerical calculations reported above suggest that this constraint might impose some tension on our model, but it is certainly possible to satisfy.

The very introduction of new scalars also affects the Peskin–Takeuchi $S$-parameter [28]. In order to estimate the scalar contribution to it, we set for simplicity the scalar contribution to it, we set for simplicity the $S$-parameter can then be written as $S = S_S + S_N + S_{SN}$, where

$$S_X = \frac{1}{12\pi} \left[ \frac{5}{6} - \frac{2M_{X_1} M_{X_2}}{(M_{X_1} - M_{X_2})^2} - \frac{1}{2} \ln \frac{M_{X_1}^2 M_{X_2}^2}{\mu^4} \right],$$

$$S_{SN} = \frac{1}{12\pi} \ln \frac{M_{SN_1} M_{SN_2}}{\mu^4}.$$  

(13a)

(13b)

(The $\mu$ is just an arbitrary mass scale introduced for convenience; the total $S$-parameter is independent of it.) Taking into account the previous discussion of the $\rho$-parameter and the scalar masses, we plotted the $S$-parameter for the special case $M_{S1,2} = M_{N1,2} = M_{SN1,2} \equiv M_{1,2}$. The resulting $S$-parameter (see Fig. 3) meets the experimental bounds for any value of the mass ratio $M_1/M_2$ from 0.01 up to 100. When in this special case the ratio $M_1/M_2$ is far from one, the $S$-parameter is well approximated by the simple formula $S = \frac{1}{6\pi} \left( \frac{5}{6} - \ln \frac{M_1^2}{M_2^2} \right)$. On the other hand, for $M_1/M_2$ close to one the $S$-parameter behaves like $\frac{1}{6\pi} \left( 1 - \frac{M_1}{M_2} \right)^2$. 

![FIG. 3: The $S$-parameter plotted for the special case $M_{S1,2} = M_{N1,2} = M_{SN1,2} \equiv M_{1,2}$. Note that, according to the Particle Data Group [27], $S = -0.10 \pm 0.10$.](image-url)
V. CONCLUSIONS AND OUTLOOK

We have presented a new possible mechanism and the corresponding formalism for EWSB. We have introduced elementary scalars and argued that EWSB could be driven not by their VEVs, like in SM, but rather dynamically, by their Yukawa couplings to fermions. If this was the case, then the ‘would-be’ NG bosons should be, as shown by general arguments, composites of both the fermions and the scalars.

The presented model certainly has some advantages: (i) It distinguishes the fermion species already at the level of Lagrangian by their distinct Yukawa couplings and hence avoids the dangerous flavor symmetries. (ii) It ties the gauge boson masses up with the fermion and scalar masses by means of the sum rules (10). (iii) Since the resulting fermion masses are nonlinear functions of the Yukawa coupling constants, there is a possibility to accommodate large ratios of fermion masses, as seen in experiment, while keeping the Yukawa couplings of the same order of magnitude. This is supported by the numerical results shown above, as well as by the numerical analysis of our preceding model [21], dealing with a simpler toy version of the present model of strong Yukawa dynamics.

So far we have considered only one fermionic generation. Including all three generations will require to introduce the flavor mixing through the Yukawa coupling constant matrices. This opens a new way to study dynamical CP violation in both the quark and lepton sector, since even though the Yukawa coupling constants are real, the solutions to the SD equations are in general complex, hence leading to complex ‘mixing matrices’.

Upon necessary modifications, the present mechanism also has the potential to describe non-relativistic bosonic superfluidity without a single-field condensate. The idea is supported by Refs. [17] and [18], in which the Cooper-type pairing in the many-boson system of $^4$He is studied. The anomalous Green’s function of the type $\langle \Phi \Phi \rangle$, considered here, was also studied in Ref. [20], along with the conventional Bose condensate.

Acknowledgments

The authors gratefully acknowledge discussions with J. Hošek and H. Bíla. This work was supported in part by the Institutional Research Plan AV0Z10480505, and by the GACR Grant No. 202/06/0734. The work of T. B. was also supported by the Alexander von Humboldt Foundation.

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