\( W \) boson mass tension caused by its right-handed gauge coupling at high energies?

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Abstract. The CDF collaboration’s recent high-precision measurement of the \( W \) mass is in 7.0 \( \sigma \) disagreement with the Standard Model expectation. This tension will be relieved if the \( W \) boson has a non-trivial right-handed gauge coupling at high energies. At TeV scales, the SM gauge symmetric four-fermion interactions induce a right-handed gauge coupling, and SM fermions compose massive composite particles. We investigate the top-quark mass produced by spontaneous symmetry breaking and compute the \( W \) and \( Z \) boson propagators and decays. The right-handed coupling corrections to their masses and widths are consistent with experimental measurements. We discuss how SM gauge bosons and composite particles can restore parity-preserving gauge symmetries at TeV scales.
1 Introduction

The parity-violating gauge symmetries and spontaneous/explicit breaking of these symmetries for the $W$ and $Z$ gauge boson masses and hierarchy pattern of fermion masses have been at the centre of a conceptual elaboration that has played a major role in donating to Mankind the beauty of the Standard Model (SM) and possible scenarios beyond SM for fundamental particle physics. A simple description is provided on the one hand by the composite Higgs-boson model or the Nambu-Jona-Lasinio (NJL) model [1] with effective four-fermion operators, and on the other by the phenomenological model of the elementary Higgs boson [2–4]. These two models are effectively equivalent for the SM at low energies. The ATLAS [5] and CMS [6] collaborations have shown the first observations of a 125 GeV scalar particle in the search for the SM Higgs boson. The interpretation of the discovered new state as the SM Higgs boson implies that there is no anymore unknown parameter in the prediction for the $W$ boson mass. It is completely determined based on the internal symmetries of the theory and a set of high-precision measurements of observables, including the $Z$ and Higgs boson masses, the top-quark mass, etc. The recent high-precision measurement of the $W$ boson mass [7] shows a difference with a significance of 7.0 $\sigma$ level in comparison with the SM expectation, including the full SM high-order corrections [8, 9]. This tension suggests the possibility of theoretical extensions to the SM.

Several hypotheses with the additional symmetries, interactions and fields have been promulgated to provide a deeper explanation of the Higgs field, its potential, and the Higgs boson. These include supersymmetry models [10] and nonsupersymmetric models, e.g., compositeness, in which additional strong confining interactions produce the Higgs boson as a bound state [11]. These extensions to the SM would modify [9,12–18] the estimated mass of the $W$ boson relative to the SM expectation [19].
To address the $W$-boson mass tension, we investigate the SM gauge symmetric four-fermion interaction of NJL-type. It not only renders the effective parity-violating SM at its infrared (IR) fixed point of $v = 246$ GeV but also a parity-preserving theory of composite particles at its ultraviolet (UV) fixed point of $\Lambda \sim$ TeV scales. The $W$ boson gauge coupling is no longer purely left-handed but has a non-trivial right-handed coupling from the four-fermion interaction. As a result, the $W$ boson mass receives a positive correction to its SM expectation, and the tension is relieved. We made an early attempt to study this issue [20,21].

In Sec. 2, we briefly recall the four-fermion interaction at a natural cutoff and its UV fixed point of strong coupling for composite particles. In Secs. 3 and 4, we describe its IR fixed point of weak coupling for the SM via top-quark mass generation by spontaneous gauge symmetry breaking. We present in Sec. 5 the discussion on the right-handed gauge couplings at high energies. Their corrections to the masses and widths of the $W$ and $Z$ bosons and comparisons with experimental measurements are in Secs. 6 and 7. The examining parity-symmetry restoration by the left and right asymmetry is emphasised in the last section.

## 2 Theoretical ultraviolet completion

A well-defined quantum field theory for the SM Lagrangian requires a natural regularisation at the UV cutoff $\Lambda_{\text{cut}}$ fully preserving the SM chiral-gauge symmetry. The UV cutoff could be the Planck or the grand unified theory scale. Quantum gravity or another new physics naturally provides such regularisation. However, the No-Go theorem [22, 23] demonstrates the existence of right-handed neutrinos and the lack of consistent regularisation for the SM bilinear fermion Lagrangian to precisely preserve the SM chiral-gauge symmetries. It implies four-fermion operators for SM fermions and right-handed neutrinos at the UV cutoff. As a theoretical model, we adopt the four-fermion operators of the torsion-free Einstein-Cartan Lagrangian with all SM fermions and three right-handed neutrinos [24–26]. Among four-fermion operators, we consider here one for the third quark family

$$G_{\text{cut}} \left[ (\bar{\psi}_a^R t_R a)(\bar{t}_R^b \psi_L^b) + (\bar{\psi}_a^L b_R)(\bar{b}_R^b \psi_L^b) \right], \tag{2.1}$$

where $a$ and $b$ are the colour indexes of the top and bottom quarks. The $SU_L(2)$ singlets $\psi_R^a = t_R^a, b_R^a$ and doublet $\psi_L^{a} = (t_L^a, b_L^a)$ are the eigenstates of SM electroweak interactions. The effective four-fermion coupling $G_{\text{cut}} \propto \mathcal{O}(\Lambda_{\text{cut}}^{-2})$ and the dimensionless coupling $G_{\text{cut}} \Lambda_{\text{cut}}^2$ depends on the running scale $\mu$.

Apart from what is possible new physics at the UV cutoff $\Lambda_{\text{cut}}$ explaining the origin of these effective four-fermion operators (2.1), it is essential to study the following aspects of the interaction and ground state. (i) What the dynamics and ground state of these operators undergo in terms of their couplings as functions of running energy scale $\mu$; (ii) Associating to these dynamics and ground states where the IR or UV stable fixed point of physical couplings locates; (iii) In the domains (scaling regions) of these stable fixed points, which physically relevant operators become effectively
dimensional-4 \((d = 4)\) renormalizable operators following renormalization group (RG) equations (i.e. scaling laws); (iv) Which \((d > 4)\) irrelevant operators though suppressed by the UV cutoff scale have corrections to the relevant operators.

In the strong coupling \(G_{\text{cut}} \Lambda_{\text{cut}}^2 \gg 1\), it is a symmetric phase (strong-coupling ground state) where massive composite bosons and fermions form \([26–28]\)

\[ \Phi = \tilde{Z}_\Phi \tilde{\psi}_L \psi_R; \quad \Psi_{L,R} = Z_\Psi \Phi \psi_{L,R}. \]  

where the colour index sums and \(SU_L(2)\) isospin index is omitted. The \(Z_\Phi\) and \(Z_\Psi\) are energy-dependent form factors (wave function renormalisations). As long as their form factors do not vanish, composite particles behave as elementary particles after wave function renormalisations. An effective field theory for composite particles \((2.2)\) at the energy scale \(\Lambda < \Lambda_{\text{cut}}^2\) is realised in the scaling domain of the stable UV fixed point \(G_{\text{cut}}^c \Lambda_{\text{cut}}^2\), which is the critical coupling of the second-order phase transition from the strong-coupling symmetric phase to the weak-coupling symmetry breaking phase \([26, 31, 32]\). The phenomenology at the LHC of these composite particles has been initiated \([33]\).

We point out that the effective field theory is SM gauge symmetric because composite particles carry SM quantum numbers and couple to the SM gauge bosons. They are massive \(M_\Phi \Phi^\dagger \Phi\) and \(M_\Psi (\bar{\Psi}_L \psi_R + \bar{\psi}_L \Psi_R)\), but exactly preserving the SM chiral (parity-violating) gauge symmetries. For example, the \(W\) boson couples not only to the left-handed field \(\psi_L\), but also to the right-hand composite field \(\bar{\Psi}_R \propto \Phi^\dagger \psi_R \propto (\bar{\psi}_R \psi_L)\tilde{\psi}_R\). Namely, the \(W\) boson of the chiral \(SU_L(2)\) gauge symmetry has a vector-like (parity-preserving) coupling to composite fermions. It implies the parity symmetry restoration at the scale \(\Lambda\) \([34]\).

When the decreasing energy scale \(\mu\) is smaller than the energy scale \(\Lambda\), the coupling \(G_{\text{cut}} \Lambda_{\text{cut}}^2 < G_{\text{cut}}^c \Lambda_{\text{cut}}^2\) runs into the weak-coupling phase. Via the contact interactions \((\bar{\psi}_L \psi_R)\Phi\) and \([\bar{\psi}_L \psi_R] \Psi_{L,R}\), the composite bosons \(\Phi\) and fermions \(\Psi_{L,R}\) dissolve into SM fermions, as their negative binding energies \(B\), form factors \(Z_\Phi\) and \(Z_\Psi\) vanish. The dissolving dynamics is similar to composite particles (poles) dissolving into their constituents (cuts) in the energy-momentum plane, e.g. deuteron dissolves into a proton, and a neutron \([35–38]\). The four-fermion interacting dynamics run into the weak-coupling phase (SM ground state), the effective operators of elementary SM fermions at the energy scale \(\Lambda\) are

\[ G \left[ (\bar{\psi}_{Li} t_{Ra})(\bar{\psi}_{Ra} \psi_{Li}) + (\bar{\psi}_{Li} b_{Ra})(\bar{\psi}_{Ra} \psi_{Li}) \right]; \]  

and the four-fermion coupling \(G \propto \mathcal{O}(\Lambda^{-2})\) at the scale \(\Lambda\). The spontaneous SM symmetry breaking (SSB) dynamics proceeds. It is shown \([39, 40]\) that for an energetically favourable SSB ground state of the least numbers of Goldstone bosons, only one massive quark \(t\) and composite Higgs boson \(\tilde{t}t\) realise. Therefore, only the first term in Eq. (2.3) accounts for the SSB dynamics. It gives rise to the Bardeen, C. Hill and
Lindner (BHL) top-quark condensate model [41, 42], where the weak-coupling IR fixed point \( G_c \Lambda^2 = 8\pi^2/N_c \) (3.1) realises an effective SM theory of the massive top quark, composite Higgs, \( W^\pm \) and \( Z^0 \) at the electroweak scale \( v < \Lambda \). The approach has been generalised to the strongly-coupled Fermi liquid for the Bose-Einstein condensate [43] and the right-handed neutrino sector for discussing dark matter particles [44].

In summary, the ultraviolet completion (2.1) possesses: (i) the UV fixed point for an SM gauge symmetric theory of composite particles at the scale \( \Lambda \); (ii) the IR fixed point for an SM gauge symmetry breaking theory of elementary particles at the electroweak scale \( v \).

3 Top-quark channel and effective SM Lagrangian

In the IR fixed point domain \( G_\Lambda \gtrsim G_c \Lambda^2 \) for the SSB dynamics, we use the BHL approach [41] to study the top-quark mass \( m_t \), \( W \) and \( Z \) mass generations, and the effective SM Lagrangian at \( v \). The mass gap-equation reads

\[
1 - \left( \frac{g_0}{g_0} \right)^2 = \left( \frac{m_t}{\Lambda} \right)^2 \ln \left( \frac{\Lambda}{m_t} \right)^2 ,
\]

where \( g_0 \equiv G_\Lambda^2, g_0 > g_0^c \) and critical value \( g_0^c = G_c \Lambda^2 = 8\pi^2/N_c \) (the colour number \( N_c = 3 \)). It appears the composite Higgs scalar \( \langle \bar{t}t \rangle \) and Nambu-Goldstone bosons, i.e. \( \langle \bar{t}(x)\gamma_5t(x) \rangle, \langle \bar{b}(x)\gamma_5t(x) \rangle \) and \( \langle \bar{t}(x)\gamma_5b(x) \rangle \). The latter becomes the longitudinal modes of the massive \( Z^0 \) and \( W^\pm \) gauge bosons.

The effective SM Lagrangian with the bilinear top quark kinetic term and Yukawa coupling to the composite Higgs boson \( H \) at the low-energy scale \( \mu \) is given by

\[
L = L_{\text{kinetic}} + \bar{\psi}_L \gamma_\mu t_R H + \text{h.c.} + \Delta L_{\text{gauge}} + \Delta L_{\text{irr}},
\]

\[
+ |D_\mu H|^2 - m_h^2 H^\dagger H - \frac{\tilde{\lambda}}{2} (H^\dagger H)^2 .
\]

The renormalized Yukawa coupling \( \bar{g}_t \), Higgs mass \( m_h \) and quartic coupling \( \tilde{\lambda} \) represent \( d = 4 \) relevant operators in the IR scaling domain. The \( \Delta L_{\text{gauge}} \) and \( L_{\text{kinetic}} \) are the usual SM renormalized Lagrangians of gauge bosons, top and bottom quarks. All renormalized quantities receive fermion-loop contributions and define for the low-energy scale \( \mu \). We add the \( \Delta L_{\text{irr}} \) to represent the \( d > 4 \) irrelevant operators suppressed by at least \( (v/\Lambda) \).

The conventional renormalization \( Z_v = 1 \) for fundamental fermions and the unconventional wave-function renormalization (form factor) \( \tilde{Z}_H \) for the composite Higgs boson are adopted

\[
\tilde{Z}_H(\mu) = \frac{1}{\bar{g}_t(\mu)^2} \bar{g}_t(\mu) = \frac{Z_{HY}}{Z_H^{1/2} g_0}; \quad \tilde{\lambda}(\mu) = \frac{\tilde{\lambda}(\mu)}{\bar{g}_t(\mu)^2}; \quad \tilde{\lambda}(\mu) = \frac{Z_{4H}}{Z_H^2} \lambda_0 ;
\]

where \( Z_{HY} \) and \( Z_{4H} \) are proper renormalization constants of the Yukawa and quartic couplings in the Lagrangian (3.2). The full one-loop RG equations for running
Figure 1. In the top-quark channel, the effective (renormalised) Higgs Yukawa coupling \( \bar{g}_t(\mu) \) (form factor \( \tilde{Z}_H(\mu) = 1/\bar{g}_t^2(\mu) \)) and quartic coupling \( \tilde{\lambda}(\mu) \) as functions of energy scale \( \mu \) are determined by RG equations (3.4,3.5), mass shell condition (4.2) of the experimentally measured top quark and Higgs mass. The effective Higgs quartic coupling \( \tilde{\lambda}(\mu) \) becomes negative at the energy scale \( \approx 5.1 \text{ TeV} \). These figures are reproduced from Refs. [32,45].

couplings \( \bar{g}_t(\mu^2) \) and \( \tilde{\lambda}(\mu^2) \) read

\[
16\pi^2 \frac{d\bar{g}_t}{dt} = \left( \frac{9}{2} g_t^2 - 8 \bar{g}_t^2 - \frac{9}{4} \bar{g}_2^2 - \frac{17}{12} \bar{g}_1^2 \right) \bar{g}_t, \tag{3.4}
\]

\[
16\pi^2 \frac{d\tilde{\lambda}}{dt} = 12 \left[ \tilde{\lambda}^2 + (\bar{g}_t^2 - A) \tilde{\lambda} + B - \bar{g}_1^2 \right], \quad t = \ln \mu \tag{3.5}
\]

where one can find \( A, B \) and RG equations for SM \( SU_c(3) \times SU_L(2) \times U_Y(1) \) running renormalised gauge couplings \( g_{1,2,3}^2(\mu^2) \) in Eqs. (4.7), (4.8) of Ref. [41].

The SSB-generated top-quark mass gives \( m_t(\mu) = \bar{g}_t^2(\mu) v/\sqrt{2} \). The pole-mass \( m_H^2(\mu) = 2 \tilde{\lambda}(\mu) v^2 \), form-factor \( \tilde{Z}_H(\mu) = 1/\bar{g}_t^2(\mu) \) and effective quartic coupling \( \tilde{\lambda}(\mu) \) describe a composite Higgs-boson, provided that \( \tilde{Z}_H(\mu) > 0 \) and \( \tilde{\lambda}(\mu) > 0 \) are obeyed.

4 Experimental measurements vs BHL composite conditions

To definitely solve the RG equations (3.4) and (3.5) for \( \bar{g}_t \) and \( \tilde{\lambda} \), one needs the boundary conditions at a physical energy scale. BHL naturally introduced the theoretical compositeness conditions,

\[
\tilde{Z}_H(\Lambda_{\text{cut}}) = 1/\bar{g}_t^2(\Lambda_{\text{cut}}) = 0; \quad \tilde{\lambda}(\Lambda_{\text{cut}}) = 0, \tag{4.1}
\]

at the composite scale \( \sim \Lambda_{\text{cut}} \), where the effective Lagrangian (3.2) is sewed together with the underlying four-fermion Lagrangian (2.1). It is the UV completion of the BHL approach at the cutoff \( \Lambda_{\text{cut}} \). However, their solutions cannot reproduce simultaneously correct experimental values of the electroweak scale \( v \), the top-quark mass \( m_t \), and the Higgs boson mass \( m_H \).
Instead, we obtained \[45, 46\] the solution to the RG equations (3.4) and (3.5) by using the boundary conditions based on the experimental values of top-quark and Higgs-boson masses, \( m_t \approx 173 \text{ GeV} \) and \( m_H \approx 126 \text{ GeV} \), via the mass-shell conditions

\[
m_t(m_t) = \frac{g_t^2(m_t)v}{\sqrt{2}} \approx 173 \text{GeV}, \quad m_H(m_H) = [2\lambda(m_H)]^{1/2}v \approx 126 \text{GeV},
\]

as well as the electroweak scale \( v = 246 \text{ GeV} \) determined by the measurement of the Fermi constant \( G_F \). As a result, we find the solutions for \( \tilde{Z}_H(\mu) \) and \( \tilde{\lambda}(\mu) \), as shown in Fig. 1. In low energies \( \mu \gtrsim M_z \), the effective Lagrangian and RG equations (3.2-3.5) with experimental boundary conditions (4.2) are equivalent to the SM Lagrangian and RG equations of elementary top-quark and Higgs fields. Extrapolating solutions to high energies \( \mu \gg M_z \), we find that (i) \( \tilde{Z}_H(\mu) \neq 0 \) is finite, the composite Higgs boson is a tightly bound state and behaves as an elementary particle after the wave-function renormalisation \( \tilde{Z}_H(\mu) \); (ii) the effective quartic coupling \( \tilde{\lambda}(\mu) \) becomes negative at the energy scale 5.1 TeV. These extrapolating results imply the aforementioned new physics beyond SM that a composite boson \( \Phi \propto Z \Phi(L_L L_R) \) (2.2) should exist at the energy scale \( \Lambda \approx 5.1 \text{ TeV} \) approximately.

Following the BHL approach to the leading order, we obtain the decay constants of the charged and neutral Goldstone bosons,

\[
f^2(0) \approx \bar{f}^2(0) \approx \frac{1}{2} N_c (4\pi)^2 m_t^2 \ln(\Lambda^2/m_t^2)
\]

at zero-momentum transfer \( (q^2 = 0) \). They will be explained by Eq. (6.4) in due course.

From the tree-level \( W \)-boson mass \( M_W = (1/2)\bar{g}_2(0)v \), the gauge-boson decay constant yields

\[
\frac{G_F}{\sqrt{2}} = \frac{1}{8\bar{f}^2(0)}; \quad \bar{f}(0) = \frac{v}{2}.
\]

On the other hand, from Eqs. (4.3) and (4.4), the gap equation (3.1) becomes,

\[
1 - \left( \frac{g_{t0}}{\bar{g}_{t0}} \right) = \frac{(4\pi)^2}{2N_c} \left( \frac{v}{\Lambda} \right)^2.
\]

These results are consistent with those given in Ref. [41] for the SM at low energies, apart from the energy scale \( \Lambda \) being much smaller than the cutoff \( \Lambda_{\text{cut}} \) of the four-fermion interaction (2.1). The energy scale \( \Lambda \) is an order of magnitude larger than the electroweak scale \( v \). Therefore, in Eq. (3.1) or (4.5) we do not need the fine-tuning \( g_{t0} \gtrsim \bar{g}_{t0} + \mathcal{O}(v^2/\Lambda^2) \) to achieve the electroweak scale \( v \). While, in the BHL model of the scale \( \Lambda_{\text{cut}} \), the extremely fine-tuning \( g_{t0} \gtrsim \bar{g}_{t0} \) is required to remove \( \Lambda^2_{\text{cut}} \) and achieve \( v^2 \ll \Lambda^2_{\text{cut}} \). It is also the reason why one introduces super-symmetry theories.

It is necessary to explain our solution in contrast with the BHL solution. The similarities between both solutions are that they approach the scaling domain of IR fixed point at low energies \( \mu \gtrsim M_z \), \( \tilde{Z}_H(\mu) \neq 0 \) and \( \tilde{\lambda}(\mu) \neq 0 \) are finite, the composite Higgs boson behaves as an elementary and interacting particle. The main differences are at high energies \( \mu \gg M_z \). The BNL solution \( \tilde{Z}_H(\mu) \) and \( \tilde{\lambda}(\mu) \) decrease to zero at
the scale $\Lambda_{\text{cut}}$, as required by the theoretical composite conditions. Instead, based on experimental measurements (4.2), our solution $\hat{Z}_H(\mu) \neq 0$ up to the scale $\Lambda \approx 5.1$ TeV, where the effective quartic coupling $\hat{\lambda}(\mu)$ becomes negative. It indicates the possible new physics (2.2) and (2.3) discussed in Sec. 2.

We give some discussions how our solutions (Fig. 1) match the effective theory of composite particles (2.2) at the scale $\Lambda$. As the energy scale $\mu$ increases, the four-fermion interaction becomes strong. Its dynamics run away from the weak-coupling asymmetric phase, where the IR ground state for the SM is, and enter the strong-coupling symmetric phase, where the UV ground state for composite particles is. In such a transition from the IR ground state to the UV one, the composite Higgs $H \sim \bar{t}t$ of the SM becomes a more massive and tightly bound state of composite boson $\Phi \sim (\bar{t}t)$ beyond the SM. The $\Phi$ can also combine with $t$ and $b$ quarks to form composite fermions $\Psi$ of $(\bar{t}t)t$ and $(\bar{t}t)b$. The transition should be described by RG equations for effective couplings and form factors, running from the scale $\nu$ (IR domain) to the scale $\Lambda$ (UV domain). The BHL effective Lagrangian (3.2) of the composite Higgs boson for the SM is sewed together with the effective Lagrangian of the composite bosons $\Phi$ and fermions $\Psi$ by matching at the scale $\Lambda$ their form factors $Z_{\Phi}$ and $Z_{\Psi}$ to those in Eq. (3.3). These are non-perturbative issues and will be subjects for future numerical studies.

Because the new scale $\Lambda$ is not much larger than the electroweak scale $\nu$. Therefore, it deserves to study the $d > 4$ irrelevant operators $\Delta L_{\text{irr}}$ (3.2) suppressed at least $(\nu/\Lambda)$ to find experimentally sizable corrections to the SM observables, for instance, the $W$ boson mass.

5 $W$ boson right-handed coupling at high energies

Among $\Lambda$-suppressed operators $\Delta L_{\text{irr}}$ (3.2), we show [47] as an example that the Feynman diagram of Fig. 2 induces an effective one-particle-irreducible (1PI) vertex function of $W$-boson right-handed gauge coupling,

$$\Gamma^W_\mu(p', p) = i\frac{\hat{g}_2}{\sqrt{2}}\gamma_\mu P_L G^W_R(p', p), \quad q = p' - p, \quad (5.1)$$

and the dimensionless $G^W_R(p', p)$ is a Lorentz scalar. The external four momenta $q$, $p'$ and $p$ are respectively for the $W$ boson, top and bottom quarks. Figure 2 is a complicated two-loop diagram, i.e., a sun-set diagram interacting with an external $W$ boson wave line. The $G^W_R(p', p)$ receives contributions from two internal four-momenta integrated up to $\Lambda$, $G^2 \sim \Lambda^{-4}$ and $G^2 \Lambda^4 \sim (8\pi^2/N_c)^3$. From the Lorentz invariance and gauge symmetry (the Ward identity), the 1PI vertex function (5.1) can be expressed as the difference between two sun-set diagrams $\Sigma(p') \propto p'^2$ and $\Sigma(p) \propto p^2$, namely $G^W_R(p', p) \propto (GA^2)^2(p'^2 - p^2)/\Lambda^2$. Here, we treat the effective $W$ boson right-handed coupling (5.1) as a model for phenomenological studies.

\textsuperscript{3}More detailed calculations can be found there. As an example, the left-handed projector $P_L$ moves clockwise or anticlockwise to the interacting vertex $G$, giving the right-handed operator $G^W_R(p', p)$. Here we do not use the old notation $\Gamma^W(p', p)$ to avoid confusion with the $W$-boson decay width.
The dependence of the vertex function and/or (bosons are realized. For the case same for exchanging
Based on the Lorentz invariance, the 1PI vertex can be written as
matrix represents self-energy functions of Dirac fermions, which are the eigenstates of a mass operator. The CKM
are external momenta,
FIG. 3: We adopt the third quark family (t, b) as an example to illustrate the 1PI vertex function of W-boson coupling to right-handed Dirac fermions induced by four-fermion operators (2.3). \(\ell, p\) and \(p'\) are external momenta, \(q\) and \(k\) are internal momenta integrated up to the energy scale \(\Lambda\). The cross “\(\times\)” represents the self-energy functions of Dirac fermions, which are the eigenstates of a mass operator. The \(g_2g_\mu p_L\) is the W boson left-handed coupling vertex of the SM. This figure and caption are reproduced from figure 3 of Ref. [47].

To find the finite terms at the leading order \((\mu/\Lambda)^\delta\) for the least \(\delta\), we estimate in three cases the nontrivial \(G^W_R(p', p)\) for given external particles’ momenta and masses \(\mu = p', p, q, m_b, \Lambda, M_W\).

(i) W boson and two fermions are at low energies \(\mu = p', p, q \ll \Lambda\),

\[
G^W_R(p', p) \sim (G\Lambda^2)^2 (\frac{m^2}{\Lambda^2}) \ll 1, \quad p', p, q < m, \quad (5.2)
\]

\[
G^W_R(p', p) \sim (G\Lambda^2)^2 (\frac{\mu^2}{\Lambda^2}) \ll 1, \quad p', p, q > m. \quad (5.3)
\]

Even for the top and bottom quark case \(m^2 = m_t m_b\), Eq. (5.2) is negligible for \(\Lambda \gtrsim 1\) TeV, and thus irrelevant for low energies. We have used Eq. (5.3) as a momentum independent parameter \(G_R\) for studying the right-handed sterile neutrino radiative decay [48] and experiment XENON1T [49].

(ii) W boson and one fermion are at high energies, while another fermion is at low energies, e.g., \(\mu = q \approx p' \lesssim \Lambda\) and \(p \ll \Lambda\),

\[
G^W_R(p', p) \sim (G\Lambda^2)^2 (\frac{\mu^2}{\Lambda^2}) \lesssim 1. \quad (5.4)
\]

It is the case for studying the mass relation of top and bottom quarks [47].

(iii) W boson is at low energies \(q = p' - p \ll \Lambda\), while two fermions are at high energies \(p' \approx p \lesssim \Lambda\),

\[
G^W_R(p', p) \sim (G\Lambda^2)^2 (\frac{|q|}{\Lambda}) < 1. \quad (5.5)
\]

It gives the contribution to the effective low-energy W-boson coupling \(\tilde{g}_2(q)\), if two fermions \(p'\) and \(p\) are internal momenta of the bubble diagrams for the W-boson vacuum polarisation tensor.
The effective right-handed vertex function $G_R^W(p', p)$ has the following properties. For low energies $p', p, q \ll \Lambda$, $G_R^W(p', p) \to 0$ recovering the SM parity-violating symmetry. While $G_R^W(p', p) \to O(1)$ for high energies $p', p, q \lesssim \Lambda$, implying the parity-symmetry restoration at high energies [45]. It is consistent with the vectorlike spectra of composite particles (2.2).

As a result, the $W$-boson coupling is no longer purely left-handed. We generally assume that the charged $W$-boson $SU_L(2)$ coupling $g_2$ takes the form 4.

$$i \frac{g_2}{\sqrt{2}} \gamma^\mu [V_L P_L + V_R P_R G_R^W(p', p)].$$

(5.6)

Since all SM fermions process four-fermions interactions $G \bar{\psi}_L \psi_R \psi_R \psi_L$ [25, 26] in the same structure as the top and bottom channel (2.3), the interaction (5.6) is generalised to all families of SM leptons and quarks. The family mixing matrix $V_L$ is the CKM matrix $U_{e_{i}}^d U_{e_{i}}^d$ for quarks or the PMNS matrix $U_{e_{i}}^d U_{e_{i}}^d$ for leptons. Whereas, their counterparts $V_R$ in the right-handed sector of $U_{e_{i}}^d U_{e_{i}}^d$ for quarks or the $U_{e_{i}}^d U_{e_{i}}^d$ for leptons [25]. In family space, all unitary matrices $U_{L,R}$ transform quark and lepton fields from gauge eigenstates to mass eigenstates.

We generalize the above discussions to the neutral gauge boson channels: the $SU_L(2)$ coupling $\bar{g}_3 t_3^i \bar{\psi}_L^i \gamma^\mu \psi_L^i W_3^\mu$; the $U_Y(1)$ coupling $\bar{g}_1 t_3^i \bar{\psi}_L^i \gamma^\mu \psi_L^i B_3^t$; and $\bar{g}_q t_3^i \bar{\psi}_R^i \gamma^\mu \psi_R^i B_3^q$.

Here $t_3^i = (1/2, -1/2)$ is the third isospin component, $q_i$ is the electric charge, and the isospin index $i = t, b$ for the top and bottom quarks. In the Feynman diagram 2, replacing $t (b)$ by $b (t)$ and $W_3^\mu$ by $W_3^\mu$ or $B_3^t$, we obtain the induced 1PI vertex functions of $W_3^\mu$ and $B_3^t$ right-handed gauge couplings. The neutral $SU_L(2)$ field $W_3^\mu$ and $U_Y(1)$ field $B_3^t$ vertices coupling to the left-handed fermions are modified by

$$i \frac{g_2}{\sqrt{2}} t_3^i \gamma^\mu [P_L + P_R G_R^W(p', p)]; \quad i \frac{g_1}{\sqrt{2}} t_3^i \gamma^\mu [P_L + P_R G_R^W(p', p)].$$

(5.7)

The 1PI function $G_R^W(p', p)$ has the same properties (i), (ii) and (iii) of Eqs.(5.2-5.5).

6 $W$ and $Z$ boson mass corrections

Using the modified couplings (5.6) and (5.7), we compute the radiative corrections (bubble diagrams) to the $W$ and $Z$ boson propagators and find the right-handed
corrections to their masses, namely the corrections from the right-handed couplings $G_R^{W,Z}(p', p)$. For the (iii) case (5.5) of low energy $W$ and $Z$ bosons $q = p' - p \approx M_{W,Z} \ll \Lambda$, while two fermions are at high energies $p' \approx p \lesssim \Lambda$, we parametrise the right-handed 1PI functions as
\[
G_R^W(q) = c_w |q|/\Lambda; \quad G_R^Z(q) = c_z |q|/\Lambda,
\]
where the scale $\Lambda \approx 5.1$ TeV and coefficients $c_w$ and $c_z$ should be the order of unity $O(1)$. Here the scale $\Lambda$ reflects the characteristic scale for the parity symmetry restoration. Namely, when $q \to \Lambda$ the $W$ and $Z$ boson couplings tend to become vectorlike.

In the top and bottom quark channel, following the BHL approach [41] and neglecting the bottom quark mass $m_b$, we arrive at the renormalised and gauge-invariant $W$ and $Z$ propagators
\[
D_{\mu\nu}^W(q) = \frac{q_{\mu}q_{\nu}/q^2 - g_{\mu\nu}}{q^2 - \tilde{g}_W^2(q) [1 + c_w^2(|q|/\Lambda)^2] f^2(q)}, \quad (6.2)
\]
\[
D_{\mu\nu}^Z(q) = \frac{q_{\mu}q_{\nu}/q^2 - g_{\mu\nu}}{q^2 - [\tilde{g}_1^2(q) + \tilde{g}_2^2(q)] [1 + c_z^2(1/2)(|q|/\Lambda)^2] f^2(q)}. \quad (6.3)
\]
The $\tilde{f}^2(q)$ and $f^2(q)$ are the decay constants of the charged and neutral Goldstone bosons, which become the longitudinal modes of massive $W$ and $Z$ bosons. The energy running $W^\pm$ and $Z^0$ masses are
\[
M_{W^\pm}^2(q) = \tilde{g}_2^2(q) [1 + c_w^2(|q|/\Lambda)^2] \tilde{f}^2(q), \quad (6.4)
\]
\[
M_Z^2(q) = [\tilde{g}_1^2(q) + \tilde{g}_2^2(q)] [1 + 2^{-1} c_z^2(|q|/\Lambda)^2] f^2(q), \quad (6.5)
\]
and the energy scale $q^2 = \mu^2$. These are similar to the running top-quark mass $m_t(\mu) = \tilde{g}_2(\mu)v/\sqrt{2}$ and Higgs mass $m_h(\mu) = [2\lambda(\mu)]^{1/2} v$, see Eq. (4.2). At the tree-level SM Lagrangian, the right-handed couplings (6.1) vanish, the gauge couplings are $\tilde{g}_{1,2}(0)$ and the SM gauge-boson masses $M_W = (1/2)\tilde{g}_2(0)v$ and $M_Z = (1/2)(\tilde{g}_2^2(0) + \tilde{g}_1^2(0))^{1/2} v$ in terms of the electroweak scale $v = (\sqrt{2}G_F)^{-1/2} \approx 246\text{GeV}$.

In the absence of the right-handed correction $c_w^2(M_W/\Lambda)^2$, the top-quark mass $m_t = \tilde{g}_t(m_t)v/\sqrt{2}$ and $W$-boson mass $M_W = \tilde{g}_2(M_W)v/2$ yield
\[
M_W = [\tilde{g}_2(M_W)/\tilde{g}_t(m_t)] m_t/\sqrt{2}, \quad (6.6)
\]
where the $W$-boson gauge coupling $\tilde{g}_2(M_W)$ is the RG solution in the SM, and the Yukawa coupling $\tilde{g}_t(m_t)$ is the RG solution (Fig. 1) obtained by the experimentally measured top-quark and Higgs masses (4.2). From Eqs. (6.4) and (6.5), we define the experimentally measured $W^\pm$ and $Z^0$ masses on their mass-shell conditions,
\[
M_W^{\exp} = \tilde{g}_2(q) [1 + c_w^2(|q|/\Lambda)^2]^{1/2} f(q) \big|_{q^2=M_W^2} 
\approx M_W^{\text{SM}} [1 + c_w^2(M_W^{\text{SM}}/\Lambda)^2]^{1/2}, \quad (6.7)
\]
\[
M_Z^{\exp} = [\tilde{g}_1^2(q) + \tilde{g}_2^2(q)]^{1/2} [1 + 2^{-1} c_z^2(|q|/\Lambda)^2]^{1/2} f(q) \big|_{q^2=M_Z^2} 
\approx M_Z^{\text{SM}} [1 + 2^{-1} c_z^2(M_Z^{\text{SM}}/\Lambda)^2]^{1/2}, \quad (6.8)
\]

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Figure 3. This figure is a modified version of Fig. 1 of Refs. [7,9], and the detailed captions can be found there. The “Standard Model” line indicates the SM result with the Higgs mass $M_{H}^{\text{SM}} = 125.6 \pm 0.7$ GeV. The green shaded line shows the result from supersymmetric extensions to the SM. The blue shaded region represents the result (6.7) of the scale $\Lambda \approx 5.1$ TeV and the upper (lower) bound $c_w \approx 2.09 (1.68)$ from the right-handed corrections.

up to the leading order of the correction $c_w^2 (M_W^{\text{SM}}/\Lambda)^2$ or $c_z^2 (M_Z^{\text{SM}}/\Lambda)^2$. The $M_W^{\text{SM}}$ and $M_Z^{\text{SM}}$ stand for the gauge boson masses that receive the full high-order corrections in the SM with the measured Higgs mass $m_H = 126$ GeV [9]. Using the relation (6.6), we plot in Fig. 3 the corrected $W$-boson mass (6.7) in the blue area for the values $1.68 \leq c_w \leq 2.09$. It shows that for these $c_w$ values of the order of unity, the result is consistent with the recent experimental result $M_W^{\text{exp}}$ (the red elliptical circle) for $m_t \approx 173$ GeV.

The $Z$-boson mass (6.8) also receives the right-handed correction. The constraint $c_z < 3.79 \times 10^{-1}$ is given by the uncertainty in the reported measurement $M_Z^{\text{exp}} = 91.1876 \pm 0.0021$ GeV [53]. The correction to the on-shell $\cos^2 \theta_W$ reads

$$
\cos^2 \theta_W \equiv \frac{M_W^2}{M_Z^2} \approx \cos^2 \theta_W \bigg|_{\text{SM}} \left[ 1 + (c_w^2 - 2^{-1} c_z^2)(|q|/\Lambda)^2 \right] |q|=M_W
$$

(6.9)

The correction $\delta \cos^2 \theta_W \lesssim 5 \times 10^{-4}$ is consistent to the experimental uncertainty on the $\delta \cos^2 \theta_W = -\delta \sin^2 \theta_W < 2 \times 10^{-3}$ [53,54]. The parameter $\rho = M_W^2 [M_Z^2 \cos^2 \theta_W]^{-1} = 1$ does not receive the right-handed correction.

7 W and Z boson width corrections

We turn to examine the right-handed corrections to the $W$ and $Z$ boson decay widths $\Gamma_{W,Z}$. It will be a crosscheck of the present theoretical scenario. Analogously to the masses $M_{W,Z}$, the widths $\Gamma_{W,Z}$ are also the SM quantities fixed without any free parameter after the Higgs mass is measured. Calculating total width of $W$-boson
decaying to all leptons and quarks, except $t$ and $b$ quarks, see for example Refs. [55,56], we obtain in the SM

$$\Gamma_{W}^{SM} = \frac{3}{16\pi} g_w^2 \bar{\sigma}_W \approx 2.07 \text{ GeV},$$  \tag{7.1}$$

using the unitary conditions $\sum_{\ell=e,\mu,\tau} |V_{\ell L}|^2 = 1$ for $i = \nu_e, \nu_\mu, \nu_\tau$, and $\sum_{q=d,s,b} |V_{iq}^R|^2 = 1$ for $i = u, c$. The decay to the top and bottom quark channel ($i = t$) is kinetically forbidden. The superscript “SM” stands for full SM high-order corrections have been taken into account. In the same way as the SM calculations, considering the modified $W$-boson coupling (5.6) we obtain up to the leading order

$$\Gamma_W^{\text{exp}} \approx \frac{3}{16\pi} g_w^2 \bar{\sigma}_W \left[ 1 + (G_W^R)^2 \right]^{1/2} \left[ 1 + c_w^2 (M_W^{SM}/\Lambda)^2 \right]^{1/2},$$  \tag{7.2}$$

using the unitary conditions $\sum_{\ell=e,\mu,\tau} |V_{\ell L}|^2 = 1$ for $i = \nu_e, \nu_\mu, \nu_\tau$ and $\sum_{q=d,s,b} |V_{iq}^R|^2 = 1$ for $i = u, c$. Right-handed corrections come from two aspects. One is the mass $M_W^{SM}$ correction (6.7). Another is the correction factor $\left[ 1 + (G_W^R)^2 \right]^{1/2}$ to the $W$-boson vertex accounting for the decay. In the decay process, the $W$ boson and two fermions are at low energies $\mu \sim M_W \ll \Lambda$, the right-handed coupling $G_W^R \sim (M_W^{SM}/\Lambda)^2$, as discussed in the (i) case (5.3). Therefore, at the leading order correction $O[(M_W^{SM}/\Lambda)^2]$, the $W$-boson decay width receives the same correction (7.2) as its mass (6.7). Based on the experimental measurement $\Gamma_W^{\text{exp}} = 2.085 \pm 0.042 \text{ GeV}$ [53], and the values $1.68 \leq c_w \leq 2.09$ (6.7) fixed by the high-precision $W$ mass measurement [7], we obtain

$$\Gamma_W \approx \Gamma_W^{SM} \left[ 1 + (3.51 \sim 5.44) \times 10^{-4} \right].$$  \tag{7.3}$$

Using the experimental $\Gamma_W^{\text{exp}}$ centre value $2.085 \text{ GeV}$ for $\Gamma_W^{SM}$, we find the right-handed correction is within the error bar $\pm 0.042 \text{ GeV}$.

In the same way, considering the modified $Z$-boson coupling (5.7), we obtain the leading order right-handed correction to the $Z$ boson decay width

$$\Gamma_Z^{\text{exp}} \approx \Gamma_Z^{SM} \left[ 1 + c_z^2 (M_Z^{SM}/\Lambda)^2 \right]^{1/2},$$  \tag{7.4}$$

which is the same as the $Z$ boson mass correction. Based on the experimental measurement $\Gamma_Z^{\text{exp}} = 2.4952 \pm 0.0023 \text{ GeV}$ [53] and the constraint $c_z < 3.79 \times 10^{-1}$ (6.8) from the uncertainty in the $M_Z$ measurement, the positive correction to the $Z$ boson width (7.4) should be smaller than $5.73 \times 10^{-5} \text{ GeV}$. It is within the error bar $\pm 0.0023 \text{ GeV}$ of the measurement $\Gamma_Z^{\text{exp}}$. Therefore, the high-precision measurements of the $Z$ boson mass, the $W$ and $Z$ boson decay widths are important for cross-checks.
8 Conclusion and remarks

We discuss the possible new physics scenario beyond the SM. One of the features is that the $W$ and $Z$ boson gauge couplings are not maximally parity-violating, non-vanishing right-handed couplings depend on momenta, and the parity symmetry could restore around the energy scale $\Lambda \sim \text{TeV}$ scales. We adopt the value $\Lambda \approx 5.1 \text{ TeV}$ indicated by the composite Higgs high-energy behaviours, which are solutions to the RG equations demanding mass-shell conditions of experimentally measured v.e.v $v$, top-quark and Higgs masses. The $W$-boson right-handed coupling gives rise to the leading order right-handed correction $c_w^2 (M_{W}^{SM}/\Lambda)^2$ to the $W$ boson mass of the SM. For $c_w^2 \sim \mathcal{O}(1)$ be the order of unity, this correction consistently relieves the $W$-boson mass $7\sigma$ tension recently discovered [53]. The analogous corrections contribute to the $Z$ boson mass, the $W$ and $Z$ boson decay widths. They are tiny and within the error bars of existing experimental measurements. The high-precision measurements of these quantities will be important for cross-checks.

We mention that the asymmetry measurements can possibly examine the right-handed correction to the $W$ boson gauge coupling. In general, it was suggested [34,45] to measure the asymmetry

$$A_{L,R} = \frac{\sigma_L - \sigma_R}{\sigma_L + \sigma_R} \approx 1 - 2|G_W^R(\mu)|^2,$$

(8.1)

and $\sigma_{L,R}$ are cross-sections of left- and right-handed polarised particles in interactions. Which energy-dependent case of (5.2-5.5) for the right-handed coupling $G_W^R(\mu)$ depends on the process under consideration. It is expected that $\sigma_R = |G_W^R(\mu)|^2 \sigma_L \to 0$ and $A_{L,R} \to 1$ for $\mu \ll \Lambda$, whereas $A_{L,R} \to 0$ for $\mu \sim \Lambda$. The asymmetry signatures should be more evident in high energies, for example the CDF [57] and D0 [58] experiments that measured the forward-backwards asymmetry in top-quark pair production at the Fermilab Tevatron $p\bar{p}$ collisions.

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