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Towards a beyond-the-Standard-Model model with elementary particle non-perturbative mass generation

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We show that a recently discovered non-perturbative field-theoretical mechanism giving mass to elementary fermions, is also capable of generating a mass for the electro-weak bosons, if weak interactions are switched on, and can thus be used as a viable alternative to the Higgs scenario. A detailed analysis of this remarkable feature shows that the non-perturbatively generated fermion and W masses have the parametric form $m_f \sim C_f(\alpha)\Lambda_{\rm RGI}$ and $M_W \sim g_w c_w(\alpha)\Lambda_{\rm RGI}$, respectively, where the coefficients $C_f(\alpha)$ and $c_w(\alpha)$ are functions of the gauge couplings, g_w is the weak coupling and $\Lambda_{\rm RGI}$ is the RGI scale of the theory. In view of these expressions, we see that to match the experimental value of the top quark and W masses, we need to conjecture the existence of a yet unobserved sector of massive fermions (that we denote Tera-fermions) subjected, besides ordinary Standard Model interactions, to some kind of super-strong gauge interactions (Tera-interactions), so that the full theory (SM plus Tera-particles) will have an RGI scale $\Lambda_{\rm RGI} \equiv \Lambda_T \gg \Lambda_{QCD}$ in the TeV region. This approach offers a solution of the Higgs mass naturalness problem (as there is no fundamental Higgs), an understanding of the fermion mass hierarchy and a physical interpretation of the electro-weak scale as a fraction of Λ_T.

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1. Introduction

In ref. [1] a field-theoretical renormalizable model was introduced where an SU(2) fermion doublet, subjected to non-abelian gauge interactions of the QCD type, is coupled to a complex scalar field via a $d = 4$ Yukawa term and an "irrelevant" $d > 4$ Wilson-like operator. A key property of the model is that an exact symmetry protects elementary particle masses against quantum power divergencies, unlike what happens in Wilson lattice QCD (WLQCD). Despite the fact that both Yukawa and Wilson-like terms break chiral invariance, it was shown in [1] (and numerically confirmed in [2]) that there exists a critical value of the Yukawa coupling where chiral symmetry is recovered, up to small effects vanishing as the UV cut-off is removed.

The interesting observation is that in the Nambu–Goldstone (NG) phase of this critical theory non-perturbative (NP) $\mathrm{O}(\Lambda_{\mathrm{RGI}})$ fermion masses get dynamically generated. They are a consequence of a sort of "interference" between residual UV chiral breaking terms and NP effects coming from the spontaneous breaking of the (recovered) chiral symmetry occurring in the NG phase of the theory. A detailed analysis of this remarkable field-theoretical feature shows that NP-ly generated fermion masses have the parametric form

$$m_f \sim C_f(\alpha)\Lambda_{\mathrm{RGI}}\,, \tag{1}$$

where Λ_{RGI} is the RGI scale of the theory and $C_f(\alpha)$ a function of the gauge coupling constants $\{\alpha\}$ of the theory. If we take the irrelevant chiral breaking Wilson-like term to be a $d = 6$ operator, one finds at the lowest loop order $C_f(\alpha) = \mathrm{O}(\alpha_f^2)$, where α_f is the coupling constant of the strongest among the gauge interactions the particle is subjected to.

A far reaching consequence of the formula (1), if applied to the top quark, is that since one gets $m_{top} = \mathrm{O}(\Lambda_{\mathrm{RGI}})$, to match its experimental value a new sector of super-strongly interacting particles, gauge-invariantly coupled to standard matter, needs to exist so that the complete theory, including the new sector and Standard Model (SM) particles, will have an RGI scale $\Lambda_{\mathrm{RGI}} \equiv \Lambda_T \gg \Lambda_{QCD}$ and of the order of a few TeV's. We immediately notice that the reason for assuming the existence of a super-strongly interacting sector here is very different from the reason why Technicolor was introduced [3, 4]. Technicolor was invoked to give mass to the EW bosons in the first place, while in the present approach super-strong interactions are introduced to give the right order of magnitude to the top quark and W (see eq. (2) below) masses. Anyway to avoid confusion and following Glashow [5], we will refer to this new set of particles as Tera-particles.

The model can be naturally extended to incorporate EW interactions and leptons (see refs. [6–9]). EW bosons as well as leptons will acquire NP masses $\mathrm{O}(\Lambda_T)$ via the same mechanism that leads to eq. (1) [1], but with coefficient functions that scale like powers of the EW gauge coupling. For instance, at the leading loop order one gets for the W mass (see sect. 3)

$$M_W \sim \sqrt{\alpha_w}c_w(\alpha)\Lambda_{\mathrm{RGI}}\,, \quad c_w(\alpha) = \mathrm{O}(\alpha)\,, \tag{2}$$

where g_w is the weak coupling and α is a short-hand for the set of couplings $[\alpha_s, \alpha_T]$ with α_s and α_T the strong and Tera-strong gauge couplings, respectively.

Eqs. (1) and (2) are somewhat similar to the expression of the Higgs-masses of fermions and W's with, however, two fundamental differences. The first is that the scale of the masses is not

[1] With the SM hypercharge assignment, in our model neutrinos are massless because ν_R is sterile.

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the vev of the Higgs field, but a dynamical parameter related to a new interaction. The second is that the modulation of the Yukawa couplings that in the SM is introduced by hand to fit the values of fermion masses, here, as we said before, is controlled at leading order by the magnitude of the gauge coupling of the strongest among the interactions the particle feels.

We conclude from this analysis that the NP scenario for mass generation we are advocating here can be considered as a valid alternative to the Higgs mechanism, with the extra advantage that we will not have to deal with any Higgs mass related fine-tuning as there is no Higgs around. An extra conceptual bonus is that we have a natural interpretation of the magnitude of the EW scale, as (a fraction of) the dynamical physical parameter, Λ_T.

There is a number of further interesting features of the approach we are describing that are worth mentioning, but which we are not going to develop in this contribution for lack of space.

First of all, the dependence of the NP fermion masses upon the gauge couplings (at leading order one finds $C(\alpha_f)=\mathrm{O}(\alpha_f^2)$ for a $d=6$ Wilson-like operator, see eq. (7) below) offers a hint to understand the fermion mass hierarchy $m_\tau \ll m_t \ll m_T$, as due to the ranking among weak, strong and Tera-strong gauge couplings, $\alpha_Y \ll \alpha_s \ll \alpha_T$. We expect the gauge coupling dependence of the NP fermion mass estimate (1) to be more compelling for the heaviest of the SM fermion families which then we assume the following analysis refers to.

Secondly, lacking the need for a Higgs boson, we need an explanation for the existence of the 125 GeV resonance, recently identified at LHC. We propose to interpret it as a W^+W^-/ZZ composite state, bound by Tera-particles exchanges. Since this bound state is light on the Λ_T scale, it should be incorporated in the low energy effective Lagrangian (LEEL) obtained by integrating out the "heavy" Tera-degrees of freedom. One can show [8] that at (momenta)$^2 \ll \Lambda_T^2$ the $d=4$ piece of the LEEL of the model resembles very much to the Lagrangian of the SM (see also [10]).

Finally it was shown in ref. [11] that with a reasonable choice of the elementary particle content, a theory extending the SM with the inclusion of the new Tera-strong sector leads to gauge coupling unification at a $\sim 10^{18}$ GeV scale.

2. A simple toy-model

The simplest model enjoying the NP mass generation mechanism we outlined before is described by a Lagrangian where an SU(2) fermion doublet, subjected to non-abelian gauge interactions (of the QCD type), is coupled to a complex scalar via $d=4$ Yukawa and "irrelevant" $d=6$ Wilson-like chiral breaking terms. The Lagrangian of this toy-model reads [1]

$$\mathcal{L}_{\rm toy}(q,A;\Phi) = \mathcal{L}_{kin}(q,A;\Phi)+\mathcal{V}(\Phi)+\mathcal{L}_{Yuk}(q;\Phi)+\mathcal{L}_{Wil}(q,A;\Phi) \tag{3}$$

$$\bullet\ \mathcal{L}_{kin}(q,A;\Phi) = \frac{1}{4}(F^A\cdot F^A) + \bar q_L \mathcal{D}\!\!\!/^A q_L + \bar q_R \mathcal{D}\!\!\!/^A q_R + \frac{1}{2}\mathrm{Tr}\,[\partial_\mu\Phi^\dagger\partial_\mu\Phi] \tag{4}$$

$$\bullet\ \mathcal{V}(\Phi) = \frac{\mu_0^2}{2}\mathrm{Tr}\,[\Phi^\dagger\Phi] + \frac{\lambda_0}{4}(\mathrm{Tr}\,[\Phi^\dagger\Phi])^2 \tag{5}$$

$$\bullet\ \mathcal{L}_{Yuk}(q;\Phi) = \eta\,(\bar q_L\Phi q_R + \bar q_R\Phi^\dagger q_L) \tag{6}$$

$$\bullet\ \mathcal{L}_{Wil}(q,A;\Phi) = \frac{b^2}{2}\rho\,(\bar q_L \overleftarrow{\mathcal{D}}_\mu^A\Phi\mathcal{D}_\mu^A q_R + \bar q_R\overleftarrow{\mathcal{D}}_\mu^A\Phi^\dagger\mathcal{D}_\mu^A q_L)\,, \tag{7}$$

where $\Phi = \varphi_0 1\!\!1 + i\varphi_j\tau^j = [-i\tau_2\varphi^\star|\,\varphi]$ is a 2×2 matrix with $\varphi=(\varphi_2 - i\varphi_1, \varphi_0 - i\varphi_3)^T$ a complex scalar doublet, singlet under the color gauge SU(N_c), $b^{-1}\sim\Lambda_{UV}$ is the UV cutoff, η is the Yukawa coupling, ρ is to keep track of $\mathcal{L}_{Wil}$ and $\mathcal{D}_\mu^A$ the covariant derivative. Despite appearances, Φ

is not the Higgs field, but only an effective way to model the UV completion of the theory. The Lagrangian (3) is power counting renormalizable (like WLQCD is, notwithstanding the presence of the $d=5$ Wilson term).

Among other obvious symmetries, $\mathcal{L}_{\rm toy}$ is invariant under the (global) transformations involving fermions and scalars ($\Omega_{L/R} \in {\rm SU}(2)$)

$$\chi_L \times \chi_R = [\tilde\chi_L \times (\Phi \to \Omega_L \Phi)] \times [\tilde\chi_R \times (\Phi \to \Phi \Omega_R^\dagger)]\,, \tag{8}$$

$$\tilde\chi_{L/R} : q_{L/R} \to \Omega_{L/R} q_{L/R}\,, \quad \bar q_{L/R} \to \bar q_{L/R} \Omega_{L/R}^\dagger\,. \tag{9}$$

The exact $\chi_L \times \chi_R$ symmetry can be realized either *á la* Wigner or *á la* NG depending on the shape of the scalar potential. In any case no power divergent fermion mass can be generated by quantum corrections as the mass operator $(\bar q_L q_R + \bar q_R q_L)$ is not invariant under $\chi_L \times \chi_R$.

The operators $\mathcal{L}_{Wil}$ and $\mathcal{L}_{Yuk}$ break $\tilde\chi_L \times \tilde\chi_R$ and mix. Thus for generic values of η, $\mathcal{L}_{\rm toy}$ is not invariant under the fermionic chiral transformations $\tilde\chi_L \times \tilde\chi_R$. However, one can show [1, 12] that chiral invariance can be recovered (up to vanishingly small O(b^2) terms) at a critical value $\eta = \eta_{cr}$, where the Wilson-like term and the Yukawa terms "compensate", much like chiral symmetry is recovered (up to O(a) cutoff effects) in WLQCD by tuning the bare quark mass to a critical value, m_{cr}, where the Wilson and the mass term "compensate". The condition determining η_{cr} thus corresponds to the vanishing of the effective Yukawa term in the Quantum Effective Lagrangian (QEL) of the theory [2].

In the Wigner phase (where $\langle|\Phi|^2\rangle = 0$) this condition implies the cancellation diagrammatically represented at lowest order in the top panel of fig. 1.

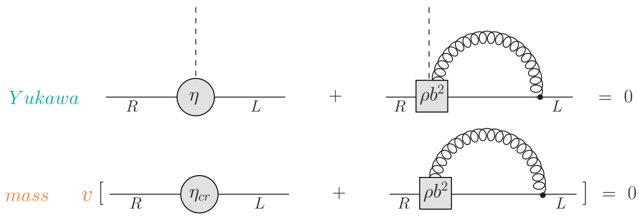

Figure 1: Top panel: the cancellation implied by the condition determining η_{cr} in the Wigner phase. The disc represents the Yukawa vertex, the box the insertion of the Wilson-like vertex. Bottom panel: the mechanism behind the cancellation of the Higgs-like mass of the quark occurring in the NG phase at $\eta = \eta_{cr}$.

In the NG phase (where at tree level $\langle|\Phi|^2\rangle = v^2 = \mu_0^2/\lambda_0$) the same condition implies that the Higgs-like mass of the fermion is cancelled out, as schematically shown in the bottom panel of fig 1.

It is a peculiar feature of the model (3) that the (quadratic) divergencies in the loop integrals in fig. 1 are exactly compensated by the b^2 factors coming from the insertion of the Wilson-like vertex. In the following we will encounter other instances of this UV vs. IR compensation.

2.1 NP mass generation

Now the key observation is that at $\eta = \eta_{cr}$ where in the NG phase the fermion Higgs-mass is cancelled, the quark acquires a non-vanishing NP finite mass via a UV vs. IR compensation mechanism, reminiscent of the one that makes finite the 1-loop diagrams of fig. 1. Because of this

[2]By Quantum Effective Lagrangian we mean the generating functional of the 1PI vertices of the theory.

kind of subtle UV vs. IR compensations, it is crucial to analyze and determine the structure of the formally O(b^2) terms of the regularized theory. This can be properly done by making recourse to the Symanzik expansion technique [13].

This analysis was carried out in ref. [1] where it was shown that, in order to properly describe O(b^2) effects in the presence of the spontaneous breaking of the (restored) chiral symmetry, one must include in the Symanzik expansion also NP-ly generated operators of the form

$$O_{6,\bar{q}q} \propto b^2\Lambda_s\alpha_s|\Phi|\left[\bar{q}\,\not{\mathcal{D}}^A q\right], \qquad O_{6,FF} \propto b^2\Lambda_s\alpha_s|\Phi|\left[F^A\!\cdot F^A\right]. \tag{10}$$

The expression of these operators is completely fixed by symmetries (in particular $\chi_L \times \chi_R$) and dimensional considerations. The presence of the Λ_s factor signals their NP origin. The existence of NP effects are taken into account by woking with the "augmented" Lagrangian

$$\mathcal{L}_{\rm toy} \to \mathcal{L}_{\rm toy} + \Delta\mathcal{L}_{NP}\,, \qquad \Delta\mathcal{L}_{NP} = b^2\Lambda_s\alpha_s|\Phi|\left[c_{FF}F^A\!\cdot F^A + c_{\bar{q}q}\bar{q}\,\not{\mathcal{D}}^A q\right] + \ldots\,. \tag{11}$$

The remarkable fact is that new diagrams are generated that contribute, among other correlators, to the quark self-energy. A couple of them are drawn in fig. 2. A straightforward analysis gives for the amputated zero momentum diagram in the right panel of fig. 2

Figure 2: Two lowest loop order quark self-energy diagrams. The blobs represent insertions of $\Delta\mathcal{L}^{Sym}_{NP}$.

$$m_q^{NP} \propto \alpha_s^2 \int^{1/b}\frac{d^4k}{k^2}\frac{\gamma_\mu k_\mu}{k^2}\int^{1/b}\frac{d^4\ell}{\ell^2+m^2_{\zeta_0}}\frac{\gamma_\nu(k+\ell)_\nu}{(k+\ell)^2}\cdot b^2\gamma_\rho(k+\ell)_\rho\, b^2\Lambda_s\gamma_\lambda(2k+\ell)_\lambda \sim \alpha_s^2\Lambda_s\,. \tag{12}$$

Here again we get a finite result, owing to an exact compensation between the UV power divergencies of the 2-loop integrals and the IR behaviour determined by the expression of the Wilson-like vertex (7) and the physics of the spontaneous breaking of the chiral $\tilde{\chi}_L \times \tilde{\chi}_R$ symmetry encoded in the form of the Symanzik operators (10). A similar compensation occurs in the calculation of the diagram in the left panel of fig. 2. Masses look like a sort of "anomaly" preventing the full recovery of chiral symmetry.

2.2 The QEL of the critical theory in the NG phase

The resulting non-vanishing quark mass (12), the occurrence of which was checked in the explicit lattice simulations of the model (3) in ref. [2], can be incorporated in the QEL, Γ^{NG}, of the theory, upon introducing the non-analytic field U in the polar decomposition

$$\Phi = (v+\zeta_0)U\,, \quad U = \exp[i\vec{\tau}\vec{\zeta}/v]\,. \tag{13}$$

On the basis of symmetries and observing that U transforms like Φ, new NP $\chi_L\times\chi_R$ invariant operators can be constructed in terms of U and one obtains [1, 9] [3]

$$\Gamma^{NG}_{d=4} = \frac{1}{4}(F^A\!\cdot F^A) + \bar{q}_L\not{\mathcal{D}}^A q_L + \bar{q}_R\not{\mathcal{D}}^A q_R + c_q\Lambda_s[\bar{q}_L U q_R + \bar{q}_R U^\dagger q_L] + \frac{c^2\Lambda_s^2}{2}{\rm Tr}\,[\partial_\mu U^\dagger\partial_\mu U]\,. \tag{14}$$

Naturally, of special interest is the fourth term in the r.h.s. of eq. (14) because, upon expanding $U = 1\!\!1 + \vec{\tau}\,\vec{\zeta}/v + \ldots$, it gives rise to a fermion mass plus a wealth of NG boson non-linear interactions.

[3] There are subtleties in the derivation of eq. (14) that we cannot detail here for lack of space (see [8, 9].

3. Introducing weak and Tera-interactions

To proceed to the construction of a possibly realistic beyond-the-SM-model we clearly need to switch on EW interactions. At the same time, as mentioned in the Introduction, it is also necessary to extend the model by incorporating a super-strongly interacting sector in order for the whole theory to have an RGI scale, Λ_T, much larger than Λ_{QCD} and of the order of a few TeVs. Only in this way there is the chance that eq. (1) can yield the correct order of magnitude of the top quark mass. The extended Lagrangian is obtained by doubling the gauge structure of quarks in order to encompass Tera-particles (Q=Tera-quarks and G=Tera-gluons) and gauging the exact χ_L symmetry to introduce weak interactions. The whole Lagrangian will thus read [8]

$$\mathcal{L}(q,Q;\Phi;A,G,W)=\mathcal{L}_{kin}(q,Q;\Phi;A,G,W)+\mathcal{V}(\Phi)++\mathcal{L}_{Yuk}(q,Q;\Phi)+\mathcal{L}_{Wil}(q,Q;\Phi;A,G,W) \quad (15)$$

$$\bullet\, \mathcal{L}_{kin}(q,Q;\Phi;A,W)=\frac{1}{4}\Big(F^A\cdot F^A+F^G\cdot F^G+F^W\cdot F^W\Big)+$$

$$+\Big[\bar q_L \mathcal{D}\!\!\!/^{AW} q_L+\bar q_R \mathcal{D}\!\!\!/^{A} q_R\Big]+\Big[\bar Q_L \mathcal{D}\!\!\!/^{AGW} Q_L+\bar Q_R \mathcal{D}\!\!\!/^{AG} Q_R\Big]+\frac{k_b}{2}\mathrm{Tr}\,[(\mathcal{D}^W_\mu\Phi)^\dagger\mathcal{D}^W_\mu\Phi] \quad (16)$$

$$\bullet\, \mathcal{V}(\Phi)=\frac{\mu_0^2}{2}k_b\mathrm{Tr}\,[\Phi^\dagger\Phi]+\frac{\lambda_0}{4}(k_b\mathrm{Tr}\,[\Phi^\dagger\Phi])^2 \quad (17)$$

$$\bullet\, \mathcal{L}_{Yuk}(q,Q;\Phi)=\eta_q\,(\bar q_L\Phi\, q_R+\bar q_R\Phi^\dagger q_L)+\eta_Q\,(\bar Q_L\Phi\, Q_R+\bar Q_R\Phi^\dagger Q_L) \quad (18)$$

$$\bullet\, \mathcal{L}_{Wil}(q,Q;\Phi;A,G,W)=\frac{b^2}{2}\rho_q\,(\bar q_L\overleftarrow{\mathcal{D}}^{AW}_\mu\Phi\mathcal{D}^A_\mu q_R+\bar q_R\overleftarrow{\mathcal{D}}^{A}_\mu\Phi^\dagger\mathcal{D}^{AW}_\mu q_L)+$$

$$+\frac{b^2}{2}\rho_Q\,(\bar Q_L\overleftarrow{\mathcal{D}}^{AGW}_\mu\Phi\mathcal{D}^{AG}_\mu Q_R+\bar Q_R\overleftarrow{\mathcal{D}}^{AG}_\mu\Phi^\dagger\mathcal{D}^{AGW}_\mu Q_L) \quad (19)$$

with obvious notations for the covariant derivatives. Besides the Yukawa (eq. (18)) and the Wilson-like (eq. (19)) operators, now also the kinetic term of the scalar (last term in eq. (16)) breaks $\tilde\chi_L\times\tilde\chi_R$ and mixes with $\mathcal{L}_{Yuk}$ and $\mathcal{L}_{Wil}$. Thus to get the critical theory (invariant under $\tilde\chi_L\times\tilde\chi_R$), on top of η_q and η_Q, a further parameter, k_b, needs to be introduced and appropriately tuned. The conditions determining the critical theory correspond to have vanishing effective Yukawa interactions and vanishing scalar kinetic term in the QEL. Notice that the coefficient k_b appears also in the expression of the scalar potential (17). The reason for this is that in this way the (bare) vev^2 (in the NG phase) and the quartic coupling of the canonically normalized scalar field will have the standard definitions, $v^2=|\mu_0^2|/\lambda_0$ and λ_0.

The tuning conditions that determine $\eta_{q\,cr}$, $\eta_{Q\,cr}$ and $k_{b\,cr}$ imply in the Wigner phase the cancellations schematically indicated in fig. 3. At the lowest loop order one finds

$$\eta_{q\,cr}^{(1-loop)}=\rho_q\,\eta_{1q}\alpha_s\,,\qquad \eta_{Q\,cr}^{(1-loop)}=\rho_Q\,\eta_{1Q}\alpha_T\,,\qquad k_{b\,cr}^{(1-loop)}=[\rho_q^2N_c+\rho_Q^2N_cN_T]k_1\,, \quad (20)$$

with η_{1q}, η_{1Q} and k_1 computable coefficients and SU(N_T) the Tera-gauge group. As before, UV loop divergencies are exactly compensated by the IR behaviour of the inserted Wilson-like vertices. In the NG phase the tuning conditions entail the cancellation of the Higgs-like mass of quarks, Tera-quarks and W's, as schematically represented in the two panels of fig. 4. Elementary particles will, however, get a NP-ly generated mass. By extending the analysis of [1] of the O(b^2) operators necessary to describe the NP effects related to the spontaneous breaking of the (recovered) $\tilde\chi_L\times\tilde\chi_R$ symmetry, one identifies the following set of (leading) operators [8]

$$O^T_{6,\bar QQ}=b^2\alpha_T\,\rho_Q\Lambda_T|\Phi|\Big[\bar Q_L\mathcal{D}\!\!\!/^{AGW}Q_L+\bar Q_R\mathcal{D}\!\!\!/^{AG}Q_R\Big] \quad (21)$$

$$O_{6,AA}=b^2\alpha_s\,\rho_Q\Lambda_T|\Phi|F^A\cdot F^A \quad (22)$$

$$O_{6,GG}=b^2\alpha_T\,\rho_Q\Lambda_T|\Phi|F^G\cdot F^G \quad (23)$$

$$O_{6,WW}=b^2\alpha_w\,\rho_Q\Lambda_T|\Phi|F^W\cdot F^W \quad (24)$$

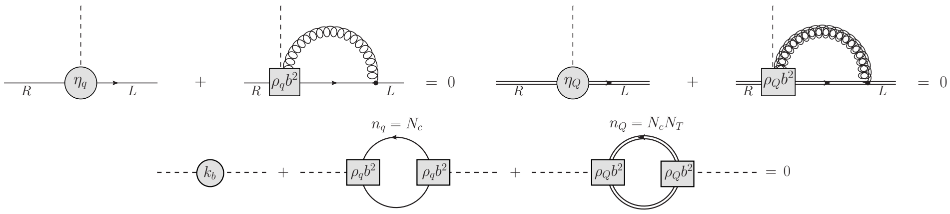

Figure 3: The lowest loop order cancellations of the Yukawa terms (top panel) and the scalar kinetic operator (bottom panel) implied in the Wigner phase by the tuning conditions determining $\eta_{q\,cr}$, $\eta_{Q\,cr}$ and $k_{b\,cr}$, respectively. Boxes represent the insertion of the quark and Tera-quark Wilson-like vertices, the disc the insertion of the scalar kinetic term. Double lines are Tera-particles.

These operators need to be added to the fundamental Lagrangian as in eq. (11). NP masses

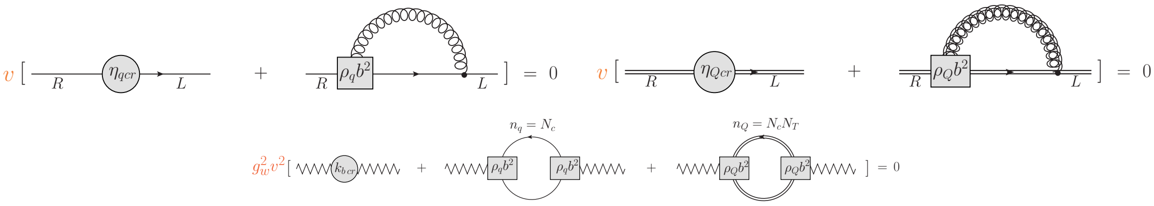

Figure 4: The cancellation mechanism of the Higgs-like mass of quark, Tera-quark (top panel) and W (bottom panel) occurring in the NG phase of the critical theory.

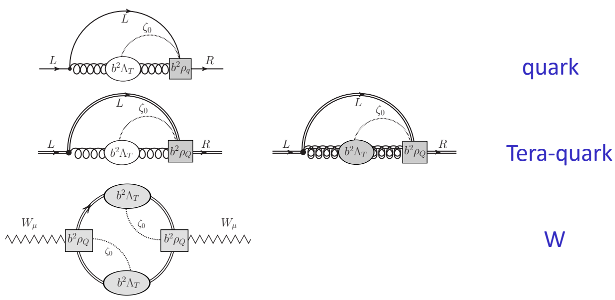

Figure 5: Self-energy diagrams giving NP masses to quarks, Tera-quarks and W. Blobs represent insertions of the NP Symanzik operators $O_{6,AA}$ (quark and Tera-quark), $O_{6,GG}$ (Tera-quark) and $O^T_{6,\bar{Q}Q}$ (W).

emerge from the type of self-energy diagrams shown in fig. 5 in which the operators (21)–(24) are combined with appropriate Wilson-like terms. These diagrams are finite owing to the exact UV-IR compensation and all of O(Λ_T) times gauge coupling dependent coefficients.

3.1 The critical QEL in the NG phase

With the same line of arguments we used before to derive eq. (14), one obtains for the $d = 4$ piece of the QEL of the critical theory in the NG phase the expression [4]

$$\Gamma^{NG}_{4\,cr}(q, Q; \Phi; A, G, W) = \frac{1}{4}\Big(F^A \cdot F^A + F^G \cdot F^G + F^W \cdot F^W\Big) + \\ +\Big[\bar{q}_L \not{\mathcal{D}}^{WA} q_L + \bar{q}_R \not{\mathcal{D}}^A q_R\Big] + C_q \Lambda_T \left(\bar{q}_L U q_R + \bar{q}_R U^\dagger q_L\right) + \\ +\Big[\bar{Q}_L \not{\mathcal{D}}^{WAG} Q_L + \bar{Q}_R \not{\mathcal{D}}^{AG} Q_R\Big] + C_Q \Lambda_T \left(\bar{Q}_L U Q_R + \bar{Q}_R U^\dagger Q_L\right) + \frac{1}{2} c_w^2 \Lambda_T^2 \mathrm{Tr}\,[(\mathcal{D}^W_\mu U)^\dagger \mathcal{D}^W_\mu U]\,. \tag{25}$$

[4]Actually getting eq. (25) is not completely trivial and requires some discussion [8, 9].

Masses are not affected by quantum power divergencies. They are "naturally small" owing to the enhanced (chiral) symmetry, of the massless theory, in line with 't Hooft notion of naturalness [14].

From a detailed analysis of the diagrams in fig. 5 one can read-off the parametric dependence of the mass of quarks, Tera-quarks and Ws at the leading loop order. One finds

$$m_q^{NP} = C_q \Lambda_T, \quad C_q = \mathrm{O}(\alpha_s^2), \qquad m_Q^{NP} = C_Q \Lambda_T, \quad C_Q = \mathrm{O}(\alpha_T^2, \dots) \tag{26}$$

$$M_W^{NP} = C_w \Lambda_T, \quad C_w = \sqrt{\alpha_w}\, c_w, \quad c_w = k_w \mathrm{O}(\alpha_T, \dots). \tag{27}$$

As for the key question of the relation of this scheme with the SM, it can be proved that, as remarked in the Introduction, upon integrating out the (heavy) Tera-dof's in the critical model (15), the resulting LEEL, which incorporates the "light" (on the Λ_T scale) 125 GeV bound W^+W^-/ZZ state detected at LHC, displays new terms with respect to the ones that appear in eq. (25) making the resulting LEEL to closely resemble the SM Lagrangian. For lack of space we do not reproduce here the proof of this statement. It can be found in ref. [8] (see also [10]).

4. Universality

A key point that needs to be thoroughly discussed is to what extent the NP mass predictions we have derived above are "universal", or in other words to what extent their value depends on the exact form of the "irrelevant" $d > 4$ Wilson-like terms that one decides to introduce in the fundamental Lagrangian. Actually it turns out that the power of the gauge coupling multiplying the RGI at the leading loop order in the formula expressing the NP masses may depend on the dimensions of such chiral breaking terms. For instance, one can see that generically the higher is the dimension of the Wilson-like terms the larger will be the power of the gauge couplings in front of the RGI scale in the mass formulae.

This state of affairs may appear to be a blunt violation of universality, however, one could imagine exploiting this dependence to try to interpret family mass hierarchy (from heavy to light) as due to Wilson-like terms of increasing dimensions.

As for the ρ dependence of the NP masses, one can prove [9] that in the case of a theory with only quarks and W's (and not Tera-quarks or leptons) their mass is actually ρ independent. If other fermions are present, be them Tera-particles or leptons, NP masses will be functions of the ratios of the ρ parameters of the various fermions.

5. Conclusions and Outlook

We have shown that, as an alternative to the Higgs mechanism, elementary particle masses can be NP-ly generated in strongly interacting theories where chiral symmetry, broken at the UV cutoff level by irrelevant Wilson-like terms, is recovered at low energy owing to the tuning of certain Lagrangian parameters.

Since in the theory there is no Higgs, in the scheme we are advocating here, there is no Higgs mass naturalness problem. Furthermore, as no power divergencies can affect physical masses, they are "naturally small", i.e. $\mathrm{O}(\Lambda_{\mathrm{RGI}})$, in line with the 't Hooft idea of naturalness [14], because the massless theory enjoys an enhanced (chiral) symmetry.

We have seen that to cope with the magnitude of the W and top mass, a super-interacting sector, gauge invariantly coupled to standard matter, must exists in order to get a theory with an RGI scale

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$\Lambda_{\rm RGI} = \Lambda_T$ of the order of a few TeVs. The simplest model discussed in ref. [1] enjoying the NP mass generation mechanism can be straightforwardly extended to incorporate weak interactions and the new Tera-degrees of freedom [8]. In a similar way one could introduce leptons and hypercharge interactions, an issue that is deferred to a forthcoming paper [9]. Upon integrating out the heavy Tera-dof's the resulting LEEL, valid at momenta$^2 \ll \Lambda_T^2$, it is seen to closely resemble the SM Lagrangian [8] (see also [10]).

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