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The Role of a Standard Model Decoupling Limit in Electro-Weak Symmetry Breaking[a]

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We argue that a suitable Standard Model decoupling limit is generically the necessary ingredient which makes scenarios of electro-weak symmetry breaking viable. Additional requirements are only that the mass predictions of a given model (e.g. predictions or theoretical limits on the Higgs or top mass) are consistent with existing data. It is discussed how such a Standard Model limit can be realized in dynamical symmetry breaking models.

The Standard Model describes existing experimental data very well, including many processes and precision tests of radiative corrections. Despite this successes all experiments mainly test the gauge structure. For the Higgs sector we know essentially only that the gauge symmetry is broken by some suitable vacuum expectation value. In the Standard Model the elementary Higgs sector leads to the famous hierarchy problem[1], due to the quadratic divergences associated with scalar particles. A solution to this problem involves the embedding of the Standard Model in new physics at the TeV-scale, and there are two main approaches: Either the quadratic corrections are cancelled due to restored supersymmetry (SUSY) above the SUSY breaking scale $\Delta \simeq TeV$, or some strongly interacting dynamical electro-weak symmetry breaking scenario provides form-factors (i.e. un-binding) and eliminates elementary Higgs bosons in the underlying theory. There are genuine *technical* differences between low energy supersymmetry and models of dynamical electro-weak symmetry breaking which should not be used as an argument for or against either direction. While the MSSM is mostly perturbative and, therefore, its phenomenology relatively easy to analyse, dynamical symmetry breaking generically arises non-perturbatively. Therefore many quantities in such models are harder to calculate and predictions often very rough. There is also no or little guidance from a greater picture compared to a supersymmetric framework. Important are ultimately the *conceptual* differences between the approaches and future experiments will of course have to decide if one of the scenarios is realized in nature.

Both routes have attractive features, but what can we learn from the current experimental infromation? Within the framework of SUSY-GUTs one can obtain a remarkably accurate prediction of the weak mixing angle from gauge coupling unification. However, although being a nice feature, this is hardly direct evidence for supersymmetry. Apart from that, the so-called Minimal Supersymmetric

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Standard Model (MSSM) is unchallenged by current experimental data. This is however largely because in the MSSM all non-Standard Model effects can decouple if the soft supersymmetry breaking mass parameters are raised in the TeV-region and flavour changing effects are tuned to be small. In other words, the success of the MSSM does not depend crucially on SUSY but on the fact that it can hide sufficiently behind the Standard Model, at least as long as the Higgs mass is below the MSSM-upper bound.

The situation may appear more difficult for dynamical electro-weak symmetry breaking with a strongly interacting sector in the TeV-region. This view is however strongly influenced by *naive* Technicolor models, based on a rescaled version of QCD, where one finds severe constraints from the existing data, and most of the original models are now ruled out by experiment. Note, however, that this is not a generic problem of dynamical symmetry breaking, but essentially the failure of *any* model where the effective Lagrangian does not have a Standard Model-limit. Both perturbative or non-perturbative scenarios of physics beyond the Standard Model with a decoupling limit are, on the other hand, phenomenologically acceptable, unless Higgs or other mass predictions disagree with experimental constraints. The question is thus if models of dynamical symmetry breaking with a Standard Model limit can be built.

Conceptually the two routes have different attractive features. Low energy supersymmetry may radiatively induce electro-weak symmetry breaking and leads to a rich phenomenology which can be studied perturbatively. At the same time, the hierarchy problem can be solved, provided that one finds a compelling (dynamical) model of supersymmetry breaking that leads to the desired soft-breaking terms in the TeV-region. As an alternative to models of supergravity-mediated SUSY-breaking, a variety of models have been proposed where SUSY breaking is mediated at low energies by standard gauge interactions [2]. Building such models has been helped by a significant improvement in understanding the non-perturbative dynamics of N=1 supersymmetric gauge theories [3]. On the other hand, "immediate" dynamical electro-weak symmetry breaking is attractive from a simplistic point of view. Dynamical symmetry breaking is a natural effect in strongly interacting theories when an asymmetric ground state lowers the total energy of the system. Well known examples in other fields of physics are ferro magnetism and superconductivity, where in the latter case one finds a dynamical Higgs mechanism. Unfortunately, far less is know about the non-perturbative behaviour of *non*-supersymmetric field theories, compared to the supersymmetric case. Nevertheless it is worth to consider scenarios where the electro-weak gauge symmetry is dynamically broken by strong dynamics in the TeV-region. This might lead to interesting scenarios of new physics, which perhaps also involves supersymmetry at scales beyond the electro-weak scale.

The purpose of this article is to stress that the *current* experimental situation is almost insensitive to details of electro-weak symmetry breaking and essentially favours *any* framework with a suitable decoupling limit which can "hide" behind the Standard Model. We will argue that SUSY models have such a decoupling limit and that it is possible to systematically build models of dynamical symmetry breaking with such a decoupling limit which will be defined in the following paragraphs.

By now there exists direct evidence for all fermions and gauge bosons of the Standard Model and we know that the electro-weak gauge group is broken. There are experimental lower limits and theoretical (indirect and model dependent) upper limits for the Higgs mass, but there is no direct evidence for the existence of a Higgs particle. Existing data provides, however, additionally numerous restrictions for modifications of the Higgs sector. Examples are rare decays, FCNC effects, contributions to R_b and other indirect effects via radiative corrections. Among those the so-called oblique corrections, commonly parametrized by the S-, T- and U- parameters, are particularly important. Thus, while a model of new physics attempting to explain electro-weak symmetry breaking need not feature an elementary Higgs boson, it certainly must be in accordance with all the above mentioned indirect constraints which are consistent with the Standard Model.

A model of new physics has a "Standard Model limit", if the parameters of the model can be chosen such that (within current experimental errors) the model becomes in a certain limit indistinguishable from the Standard Model, e.g., all effects from additional particles and interactions can be decoupled. If this is possible, it will allow a certain range of those parameters which are so far only poorly determined within the Standard Model, such as the mass of the Standard Model Higgs boson, certain mixing angles, details of CP violation or neutrino masses. Thus a Standard Model limit - if it exists - does not necessarily lead to every possible choice of Standard Model parameters, and examples for 'relic' mass restrictions (e.g., for the Higgs mass) will be seen below.

The minimal supersymmetric extension of the Standard Model (MSSM) exhibits such a Standard Model limit. To be phenomenologically viable, one must ensure sufficient separate conservation of the three lepton numbers, the suppression of flavour changing neutral currents (FCNC) as well as the smallness of the neutron and electron electric dipole moments. As emphasized in a recent review [4], this confines the MSSM to special, "exceptional" points in parameter space, and there are several attempts to explain this situation [b]. Once within this parameter regime, the virtual effects of supersymmetric particles can be decoupled from electro-weak observables at or below the Z-mass [5], when the masses of all supersymmetric particles are raised above ~ 200 GeV. Especially, all the non-minimal Higgs bosons of the MSSM (H^0, $H^\pm$ and A^0) can be decoupled, when the mass of the CP-odd Higgs boson is raised ($m_{A^0} \gg m_{Z^0}$). In this limit the CP-even Higgs boson h^0 remains light, with Standard Model equivalent couplings to all Standard Model particles. This is because supersymmetry relates the quartic Higgs coupling (and mass) to gauge couplings, and after taking radiative corrections into account the CP-even Higgs bosons is always lighter than [6] ~ 125 GeV. Thus, due to this Standard Model limit, the MSSM can be consistent with electro-weak precision data, provided the additional "relic" mass restriction is met. I.e., the global fit for the Standard Model Higgs mass must be consistent with the upper bound on h^0, $m_{h^0} \lesssim 125$ GeV.

Let us now see how the Standard Model limit can be realized in alternative scenarios of electro-weak symmetry breaking. Besides the issue of flavour changing neutral currents, it is very important to consider the so-called "oblique radiative corrections", which can be parametrized in the precision variables S, T, and U. These corrections are defined from the vacuum polarization amplitudes of γ, W, and Z and the γZ mixing which have the form

$$\Pi_{\gamma\gamma} = e^2 \Pi_{QQ} \ , \tag{1}$$

$$\Pi_{WW} = \frac{e^2}{s^2} \Pi_{11} \ , \tag{2}$$

$$\Pi_{ZZ} = \frac{e^2}{c^2 s^2} (\Pi_{33} - 2s^2 \Pi_{3Q} + s^4 \Pi_{QQ}) \ , \tag{3}$$

$$\Pi_{Z\gamma} = \frac{e^2}{cs} (\Pi_{3Q} - s^2 \Pi_{QQ}) \ . \tag{4}$$

Here the relation $J_Z = J_3 - s^2 J_Q$ has been used. J_Q is the electro-magnetic current, $s^2 = \sin^2\theta_w$, $c^2 = \cos^2\theta_w$, and the weak coupling constants have been expressed in terms of e, s^2 and c^2. The indices i, j of Π_{ij} on the right hand side indicate the relevant currents. The Standard Model contributions to

[b] In the MSSM with minimal particle content and R-parity conservation assumed, the number of parameters of the model [4] equals 124. One possibility to largely reduce this set of parameters, and to naturally solve the SUSY-flavour problem in the process, opens up in the context of gauge mediated supersymmetry breaking. The flavour blindness of gauge interactions leads to flavour blind soft breaking squark and slepton mass terms, which ensure the desired FCNC suppression.

the vacuum polarizations can be separated such that the remaining contributions of new physics are then functions of q^2/M^2, where M^2 is the mass of new, heavy particles. If M^2 is large enough one can expand in powers of q^2. Using QED Ward-Identities for $q^2 = 0$ the expansion leads at order q^2 to six coefficients:

$$\Pi_{QQ} = q^2\Pi'_{QQ}(0) + ... \tag{5}$$

$$\Pi_{11} = \Pi_{11}(0) + q^2\Pi'_{11}(0) + ... \tag{6}$$

$$\Pi_{3Q} = q^2\Pi'_{3Q}(0) + ... \tag{7}$$

$$\Pi_{33} = \Pi_{33}(0) + q^2\Pi'_{33}(0) + ... \tag{8}$$

When the three most precise measured observables, α_{em}, G_F and m_Z, are used as input there remain three independent variables which parametrize effects of new physics. These three variables can be defined as [7]

$$S = 16\pi\left[\Pi'_{33}(0) - \Pi'_{3Q}(0)\right] , \tag{9}$$

$$T = \frac{4\pi}{s^2c^2m_Z^2}\left[\Pi_{11}(0) - \Pi_{33}(0)\right] , \tag{10}$$

$$U = 16\pi\left[\Pi'_{33}(0) - \Pi'_{11}(0)\right] . \tag{11}$$

The last variable, U, is so small that it is currently irrelevant and will not be considered further. The variable T, measuring custodial SU(2) violating effects, and S, which is related to axial SU(2) and which is sensitive to the unitarity partner of Goldstone bosons, will both be discussed in more detail below.

For dynamical symmetry breaking it is instructive to consider the effects of naive Technicolor scenarios on oblique radiative corrections. The original models of this type, as well as more sophisticated versions like extended Technicolor (ETC) or "walking" Technicolor, have largely been ruled out in the past [8], and in the following we describe the basic difficulties. In the simplest version of these models of dynamical electro-weak symmetry breaking, the Goldstone boson decay constant of QCD is scaled up such that the correct W mass arises. This leads (up to $N_{TC} \neq 3 = N_c$ corrections) to a more or less fixed spectrum of QCD-like bound states in the TeV-region. Consequently the Techni-pions, the Goldstone bosons giving mass to W and Z in a dynamical Higgs mechanism, do not have a Higgs-like, *i.e., scalar*, unitarity partner [c]. Instead, the role of unitarity partners is played predominantly by low lying composite vector resonances, the Techni-rhos. Due to the QCD-like dynamics it is therefore not possible to obtain the Standard Model spectrum, where a suitable composite Higgs-like scalar mimics the Standard Model Higgs boson, and the remaining spectrum is decoupled. However, the main phenomenological problem of naive Technicolor is not per se the absence of the scalar partner of the Goldstone bosons, but the low mass of the Techni-rho resonances. We will see in a moment that low lying (composite or fundamental) vector-like states are in general a problem for the experimentally small precision variable S, even if there is a (composite or fundamental) Higgs particle. If a Higgs-like scalar is absent, like in Technicolor, then there exist of course upper bounds for the masses of vector states due to the unitarity of Goldstone boson scattering amplitudes. Such upper bounds for the masses of vector states do not exist in models where a Higgs particle exists in addition to extra vector states. An example is given by models which contain a Z', which can in principle become arbitrarily heavy.

The phenomenological problem with light vector states is that they can mix with the W and Z-bosons, which is severely constrained by precision electro-weak data, i.e. the smallness of the S parameter. This explains why in Technicolor the problem becomes even more severe as the number of

[c] The so-called sigma particle would – if it exists – be broad and too heavy.

Techni-colors N_{TC} is increased, since for large N_{TC} the ratio of Techni-rho mass and Goldstone boson decay constant becomes smaller, and the mixing with W and Z is increased. This effect can be seen best by expressing S with the help of dispersion relations [7] as

$$S = \frac{1}{3\pi} \int_0^\infty \frac{ds}{s} \left[R_V(s) - R_A(s) - H(s) \right] \ , \tag{12}$$

where R_V and R_A measure the contribution of vector and axial-vector states, respectively. R_V and R_A are defined as the ratios of the cross sections of a photon which couples to the isospin current $J^{\mu 3}$ divided by the Compton process, in analogy to the famous R ratio in QCD. The definition of S depends on a reference Standard Model Higgs mass value which was often chosen to be 300 GeV. The function $H(s)$ allows one to remove the Standard Model Higgs contribution, e.g., in Technicolor by sending the Higgs mass into the continuum of bound states (i.e., $m_H \simeq 1$ TeV). The function H can be written as

$$H(s) = \frac{1}{4} \left(1 - \frac{m_H^2}{s} \right)^3 \theta(s - m_H^2) \tag{13}$$

which leads essentially to a small logarithmic dependence on the Higgs mass value.

In Technicolor the quantities R_V and R_A measure the sum of charges involved in $I = 1$ vector resonances. One can estimate S by parametrizing (similar to QCD) the resonance and the continuum contributions to R_V and R_A by delta and theta functions, respectively. The leading contributions come from the vector meson resonances (vector meson dominance), which leads to

$$R_V(s) = 12\pi^2 F_{\rho T}^2 \ \delta(s - m_{\rho T}^2) \ , \quad R_A(s) = 12\pi^2 F_{a_1 T}^2 \ \delta(s - m_{a_1 T}^2) \ , \tag{14}$$

where $m_{\rho T}$ and $m_{a_1 T}$ are the Techni-rho and Techni-a_1 masses. With the help of the Weinberg sum rules [9] it is possible to express $F_{\rho T}^2$ and $F_{a_1 T}^2$ as

$$F_{\rho T}^2 = \frac{m_{a_1 T}^2 F_\pi^2}{m_{a_1 T}^2 - m_{\rho T}^2} \ , \quad F_{a_1 T}^2 = \frac{m_{\rho T}^2 F_\pi^2}{m_{a_1 T}^2 - m_{\rho T}^2} \ , \tag{15}$$

where $F_\pi = 250$ GeV. From this one obtains from R_V and R_A the following contribution to S:

$$S = 4\pi \left(1 + \frac{m_{\rho T}^2}{m_{a_1 T}^2} \right) \frac{F_\pi^2}{m_{\rho T}^2} \ . \tag{16}$$

Using large N_c rescaling relations between Technicolor and QCD one finds [10]

$$\frac{m_{\rho T}^2}{m_{a_1 T}^2} = \frac{m_\rho^2}{m_{a_1}^2} \qquad \frac{F_\pi^2}{m_{\rho T}^2} = N_D \frac{N_{TC}}{3} \frac{f_\pi^2}{m_\rho^2} \ . \tag{17}$$

With $f_\pi = 93$ MeV, $m_\rho = 770$ MeV and $m_{a_1} = 1260$ MeV this results in $S \simeq 0.25 N_D N_{TC}/3$, where N_{TC} is the number of Techni-colors and N_D is the number of weak doublets. The continuum contributions to R_V and R_A can be parametrized by theta-functions and lead to logarithmic mass dependencies, while the resonances lead to $1/m_{\rho T}^2$ contributions. These $1/m_{\rho T}^2$ contributions are usually dominant unless logarithmic contributions proportional to N_{TC}^2 become for $N_{TC} > 8$ equally or even more important.

If one evaluates R_V and R_A using a more detailed analysis of QCD data and large N_c rescaling one finds [7]

$$S = 0.3 N_D \, \frac{N_{TC}}{3} \ , \tag{18}$$

which must be confronted with the value of S extracted from global fits to existing data. For such fits it is necessary to specify the top quark mass and the Higgs mass of the Standard Model, and typically the values $m_t = 175$ GeV and $m_H = 300$ GeV are assumed [11]. For comparison with Technicolor models the reference Higgs mass is sometimes shifted to 1 TeV, which leads to

$$S = -0.26 \pm 0.16 \ . \tag{19}$$

The negative sign expresses the fact that the preferred Higgs mass of the Standard Model is much smaller than 1 TeV, and today the best fit is obtained [12] for $m_H \approx 121$ GeV. Comparing equation (19) with the prediction of naive Technicolor, equation (18), one finds the smallest disagreement at the 3 σ level for $N_D = 1$ and $N_{TC} = 3$. More Technicolor doublets and/or a larger Technicolor group lead to even larger deviations from the data.

The above expression for S, equation (12), applies in principle for any theory with composite or fundamental vectors and/or axial-vector states which couple to $J^{\mu 3}$. For example, one could study Z' models in this way and find a $1/m_{Z'}^2$ dependence from the mixing of the Z' with Z. This leads to unacceptable contributions to S if the extra vector state becomes light. However, these contributions are not as severe as in Technicolor since the Z' couplings are usually chosen to be perturbative (i.e. small) and the color factor N_{TC} is absent.

Another severe problem of many Technicolor scenarios stems from additional extra Techni-fermion doublets. The problem shows up in the precision variable T defined above and emerges in a more general context for new, left-handed fermionic doublets, i.e., it is not only a genuine Technicolor problem. Since extra fermions are not observed such a $\mathrm{SU}(2)_L$ doublet is required to be massive with masses well above one hundred GeV. When the massive propagators for the fermions in the doublet are written as

$$S_j = \frac{i}{\not p - \Sigma_j(p^2)} \ , \tag{20}$$

where $j = U, D$ for a new doublet and $j = t, b$ for the top and bottom quark, then the contribution to T is given by [13]

$$T = \frac{-N_c}{16\pi^2 \alpha_{em} v^2} \int_0^\infty dk^2 \ \frac{k^4 (\Sigma_U^2 - \Sigma_D^2)^2}{(k^2 - \Sigma_U^2)^2 (k^2 - \Sigma_D^2)^2} \ . \tag{21}$$

Here N_c is the number of colors or Techni-colors and $v \simeq 246$ GeV. If we insert $\Sigma_t = m_t$ and $\Sigma_b = 0$ into equation (21) one finds the well known, leading Standard Model top mass contribution to the T-parameter,

$$T_{SM} = \frac{N_c}{16\pi^2 \alpha_{em}} \frac{m_t^2}{v^2} = \frac{N_c}{16\pi \sin^2\theta_W \cos^2\theta_W} \frac{m_t^2}{M_Z^2} \ . \tag{22}$$

Equation (21) is also valid for new heavy doublets. In Technicolor, for example, ordinary quark and lepton masses (like the top mass) must be generated by so-called Extended Technicolor [14] interactions between quarks and Techni-quarks. The coupled system of gap equations leads, in a first approximation [15], to the relation

$$\Sigma_U - \Sigma_D = \Sigma_t \ , \tag{23}$$

where Σ_U and Σ_D are the mass functions of the Technicolor doublet. Assuming that equation (23) is valid and approximating Σ_i as $\Sigma_i = m_i \, \Theta(\Lambda^2 - p^2)$, equation (21) yields corrections to the Standard Model value of T:

$$\begin{aligned} T_{total} &= T_{SM} + T \\ &= \frac{N_c m_t^2}{16\pi^2 \alpha_{em} v^2} \left(1 + \frac{4}{9} N_{TC} + \frac{m_t^2}{\Lambda^2} + \frac{4}{3} N_{TC} \frac{m_U^2}{\Lambda^2} + \mathcal{O}(m^4/\Lambda^4) \right) \end{aligned} \tag{24}$$

For $\Lambda \to \infty$ this expression becomes the result usually quoted in the literature [15]. For finite Λ we can read off the m_t^2/Λ^2 corrections to equation (24), and in addition a term proportional to m_U^2/Λ^2. For sufficiently large Λ these $1/\Lambda^2$ terms are small and can be omitted. In this limit the value of T is given by $4/9 \cdot N_{TC} \cdot T_{SM}$, which is excluded by phenomenology due to the good agreement of the experimental value of T with its Standard Model value. One might argue that equation (23) is very model dependent and may be generalized by a more complex relation, but it is not easy to explain the top-bottom mass splitting without sizable corrections to T.

As mentioned earlier, the problems of extra left-handed doublets in Technicolor have a more general scope. In fact any theory with extra left-handed doublets and a similar mechanism to explain the top-bottom mass difference faces the same problems. If the mechanism which explains the top-bottom mass splitting induces a $U - D$ splitting proportional to $K \cdot m_t$, where K is typically expected to be of order unity, then this will lead to $T = 4/9 \cdot K^2 N_c \, T_{SM}$. Unless K is very tiny, this leads again to custodial SU(2) violation (measured by T) too large to be reconciled with precision measurements [d].

This discussion shows that two ingredients are disfavoured when constructing models of dynamical symmetry breaking with a Standard Model limit: Low lying vector states lead to undesired contributions to the S-parameter, and additional fermionic $SU(2)_L$-doublets may easily cause excessive violation of custodial SU(2) symmetry.

The attempt to construct a viable model of dynamical electro-weak symmetry breaking may be viewed as a bottom-up approach, guided by the current phenomenological situation. As discussed, the data so far favour models with a Standard Model limit, and we will identify now the ingredients of symmetry breaking that *naturally* lead to such a limit. In this context it is important to carefully distinguish between the Higgs mechanism (i.e. the Goldstone bosons "eaten" by gauge bosons, serving as their longitudinal degrees of freedom) and a physical Higgs particle. While the scalar Goldstone bosons are simply a consequence of broken global symmetries, irrespective of the nature of symmetry breaking (perturbatively or tree level), the existence or non-existence of a Higgs particle is a feature of the particular field theory considered.

The Higgs mechanism for a specific gauge symmetry breaking pattern requires only an operator $\hat{O}$ with the following properties:

- $\langle \hat{O} \rangle \neq 0$, i.e. a "condensate"
- Lorentz invariance of the vacuum requires that $\hat{O}$ must be scalar
- $\hat{O}$ must transform non-trivially under the gauge group to be broken, and as a singlet under the desired unbroken subgroups.

Note that $\hat{O}$ does not have to be a fundamental scalar Higgs field. For the case of the Standard Model gauge group $\mathrm{SU}(3)_c \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ the operator $\hat{O}$ should be a doublet of $\mathrm{SU}(2)_L$ with $Y = 1$ for the hypercharge [e]. The condensate $\langle \hat{O} \rangle \neq 0$ then breaks the global $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ symmetry, which implies the existence of Goldstone bosons which can be eaten by W's and Z's. In unitary gauge the operator can be expanded as

$$\hat{O} = \left(\langle \hat{O} \rangle + \delta \hat{O} \right) e^{i \varphi_a T_a} \,, \tag{25}$$

[d] It is possible to introduce left-handed doublets together with right-handed partners such that custodial SU(2) is affected much less [44].

[e] $\hat{O}$ could in principle have different quantum numbers, but phenomenological evidence like the smallness of $\Delta\rho$ lead to strong constraints.

and if a $|D\hat{O}|^2$ term is present in the effective Lagrangian with $\langle\hat{O}\rangle \neq 0$, the Goldstone bosons φ_a are absorbed by the corresponding gauge bosons. The Higgs mechanism requires only the existence of a suitable condensate $\langle\hat{O}\rangle \neq 0$ and operates on the basis of symmetries, it does not depend on either the fundamental or non-fundamental nature of $\hat{O}$ and/or the presence of $\delta\hat{O}$.

The interactions of the Goldstone bosons can thus be understood in terms of the involved symmetries and the corresponding Ward-Identities, and the details of the interaction responsible for symmetry breaking become almost irrelevant. This is well established in QCD, where even a Nambu–Jona-Lasinio description of chiral symmetry breaking (which has almost nothing to do with QCD, but can be arranged to break the chiral symmetries correctly) leads to a remarkable good description of pion interactions.

On the other hand, $\delta\hat{O}$ or other non-scalar excitations of the vacuum are not related to the symmetries of the theory and may or may not include a fundamental or composite physical scalar Higgs particle H. If $\hat{O}$ is fundamental and $\delta\hat{O}$ is omitted completely, one arrives at the non-linear sigma model. This is fine with respect to symmetry breaking, but renormalizability is lost, which is an essential feature of a fundamental scalar theory. Unless new physics is very close, Goldstone boson scattering amplitudes are unbounded and unitarity is violated. One way out is to postulate a fundamental scalar Higgs field H which can be grouped together with the Goldstone bosons in a $SU(2)_L$ doublet Φ, and where H can act as "unitarity partner" for the Goldstone bosons. Unless new physics (i.e. some extra scale) is close renormalizability requires in addition that the Higgs potential $V(H)$ has only a mass term and $\lambda\Phi^4$ interactions.

Note what happens if $\delta\hat{O}$ corresponds to a composite operator. In this case the Goldstone bosons are composite as well, and the Higgs mechanism will work as before. However, the remaining spectrum of the theory will typically be rich with an effective interaction Lagrangian which need not be renormalizable[f]. Whatever the spectrum of the theory is, the effective Lagrangian must contain some resonances which act as unitarity partner for the Goldstone bosons. The simplest scenarios would be either a composite Higgs like state (i.e. an effective Standard Model) or suitable vector resonances with the remaining spectrum located close to the continuum, i.e., at some high scale Λ. The details of the spectrum remain a dynamical issue, depending on the interactions of the underlying model. In QCD-like theories, for example, rho-like vector resonances emerge in the effective Lagrangian while a physical Higgs particle is absent. One can, however, choose interactions which lead to a composite Higgs H instead. Note that the existence of a composite scalar H allows a richer interaction potential and other resonances in addition. Thus a composite Higgs does not automatically lead to an effective Standard Model.

We can now discuss how the Standard Model limit can be reached systematically in dynamical symmetry breaking models:

- At first a suitable symmetry breaking pattern must be considered in order to arrive at the correct spectrum of Goldstone bosons. Care should be taken to avoid problems with so-called pseudo Goldstone bosons, since, being light, they lead easily to phenomenological problems.
- The existing data for the precision variable S favours a spectrum where a Higgs-like scalar plays the role of the unitarity partner of the Goldstone bosons, as explained in section 2. Vector-like states with SU(2)-quantum numbers are disfavoured, and thus one is lead to consider scenarios which differ significantly from QCD.
- Custodial SU(2) violation measured by $T \simeq m_t^2$ agrees very well with the Standard Model value. This strongly disfavours scenarios which have sizable extra custodial SU(2) violating effects, such

[f] An example is again chiral symmetry breaking in QCD in the limit where pions are massless. Only the underlying theory – here QCD – should be based on fundamental fields and be renormalizable.

as additional fermionic doublets beyond the Standard Model. Unless special care is taken the mechanism responsible for the top-bottom mass splitting will lead to huge effects via extra doublets. Thus, in order to avoid fine-tunings or special choices, we avoid extra fermionic doublets beyond the Standard Model.

- The absence of flavour changing neutral currents (FCNC's) beyond the Standard Model easily becomes a problem in models where the heavy top mass is explained *after* the electro-weak symmetry is broken. A well known example is the generation of quark masses in extended Technicolor. This might suggest that the heavy top mass is intimately related to electro-weak symmetry breaking, which leads to the idea of top-condensates.

Dynamical symmetry breaking models along these guidelines should be phenomenologically much more viable than, e.g., naive Technicolor since the model has a Standard Model limit. The above choices are not very artificial and the physically most interesting point is probably to understand which sort of dynamics produces the scalar spectrum.

There is a nice prototype model which implements the Standard Model limit in a minimal way: The so-called BHL model [16] of electro-weak symmetry breaking. The idea is here to eliminate the fundamental Higgs field in the Standard Model and to introduce instead a new attractive interaction, which leads to the formation of a $t\overline{t}$-condensate. In terms of the discussion given above, the elementary Higgs field of the Standard Model gets replaced by the bi-fermion composite scalar

$$\hat{O} \simeq \overline{Q}_L t_R \ , \tag{26}$$

with exactly the Standard Model-Higgs quantum numbers. We found previously that scalars are favoured as unitarity partners of the Goldstone bosons due the constraints on the S parameter. In addition, in this approach the generation of the large top mass occurs within the process of gauge symmetry breaking. Therefore m_t is naturally of the order of the electro-weak scale and need not be artificially generated at the cost of possibly large FCNC's.

The interaction responsible for triggering the condensation is assumed to have its origin in new yet unspecified physics above some high-energy scale Λ. At lower energies this sector of new physics manifests itself through non-renormalizable interactions between the usual fermions and gauge bosons, where for energies $E \ll \Lambda$ the lowest dimensional four-fermion operators are most important. Thus, at the scale Λ the Lagrangian can be given by the gauge-kinetic terms for the known chiral fermions and gauge bosons, plus a gauge invariant four-fermion interaction term

$$\mathcal{L} = \mathcal{L}_{kin}(g, f) + \mathcal{L}_{4f}^{BHL} \ . \tag{27}$$

The gauge invariant four-fermion operator which is needed for a condensation of the operator $\hat{O}$ [see (26)] in the $t\overline{t}$-channel is given by [17,18,16,19]

$$\mathcal{L}_{4f}^{BHL} = G(\overline{Q}_{Li} t_R)(\overline{t}_R Q_{Li}) \ . \tag{28}$$

The BHL model and most attempts to modify the BHL model have a Standard Model limit, but they also have a common problem: They tend to produce a top mass which is unacceptably high, even for very high values for the scale of new physics. This is not an accident, but occurs systematically in scenarios of dynamical electro-weak symmetry breaking driven by a top condensate alone. If a top condensate breaks the electro-weak symmetry, then for an asymptotically free theory the dynamically generated top propagator can be written as

$$S_t = \frac{i}{\not{p} - \Sigma_t(p^2)} \ , \tag{29}$$

where

$$\Sigma(p^2) \stackrel{p^2\to\infty}{\longrightarrow} 0 \ . \tag{30}$$

The dynamically generated self energy Σ_t can be related to the Goldstone boson decay constant $F_\pm$, which is induced by the condensate, by virtue of the so-called Pagels–Stokar relations [29]:

$$F_\pm^2 = \frac{N_c}{32\pi^2} \int dk^2 \frac{\Sigma_t^2(k^2)}{(k^2 - \Sigma_t^2(k^2))} \ . \tag{31}$$

These relations are derived primarily on the basis of Ward identities and are therefore valid beyond the approximations used to derive $\Sigma_t(p)$. Note that the integral in (31) is formally log divergent, but finite if the asymptotic behaviour of equation (30) is taken into account. Equation (31) is a very powerful relation between the dynamically generated top mass and the Goldstone boson decay constant, and also the W mass, since after the charged Goldstone boson is absorbed by W one obtains $m_W^2 = g_2^2 F_\pm^2$. It is instructive to observe that the integral on the right hand side of the Pagels–Stokar relation, equation (31), feels the structure of Σ_t only on a logarithmic scale. Thus, for large values of Λ, the integral is dominated by contributions coming from scales both far away from Λ and the electro-weak scale, i.e. from regions where Σ should (up to logarithmic RGE running) essentially be flat. In this case it is a good approximation to express Σ by its 'mean' height and extension:

$$\Sigma_t(p^2) = m_t \Theta(\Lambda^2 - p^2) \ . \tag{32}$$

To be more precise, this should be a valid approximation *except* for small values of Λ, but those values anyway lead to values of m_t that are by far too large. Inserting equation (32) into equation (31) and solving for the top mass, one finds

$$m_t^2 = \frac{32\pi^2 m_W^2}{N_c \ g_2^2 \ \ln(\Lambda^2/m_t^2)} \ , \tag{33}$$

which is exactly the relation obtained in the BHL model in bubble approximation. This is consistent, since the structure of the ansatz (32) corresponds to a Nambu–Jona-Lasinio gap equation. Corrections to this relation come, like in the BHL model, from other, weak gauge contributions and are expected to be moderate and model dependent. Thus, the Pagels–Stokar relation explains why many variants of the BHL model, like two Higgs doublets or the supersymmetric version, produce a similar top mass, a value typically too high to fit the data. On the other hand, inserting $m_t = 175$ GeV in equation (33) one finds a W mass which is too small for any Λ. Therefore, in order to obtain a viable relation between Λ and the top mass, one is led to consider more complex symmetry breaking scenarios, with more condensates and/or more Yukawa couplings.

Given the success of renormalizable gauge theories, it appears attractive to relate such extra condensates to more complex symmetry breaking patterns of extended gauge sectors. An appealing class of models where such ideas can be exemplified are dynamically broken left-right-symmetric theories. In such models, usually a condensate in the leptonic sector breaks left-right symmetry as a first step, and a second (or more) condensate(s) cause electro-weak symmetry breaking. Dynamical symmetry breaking can indeed be implemented in left-right-symmetric theories [32,33] along the guidelines presented above to give a model with a Standard Model limit. In particular, a modified relation between m_t and Λ which is phenomenologically acceptable - even for low values of Λ is found.

In summary we would like to emphasize again, that neither the success of the MSSM nor the problems of naive Technicolor stem generically from a supersymmetric or dynamical mechanism of electro-weak symmetry breaking. The phenomenological viability or failure of those models is to a large extent simply the presence or absence of the Standard Model limit. Obviously such a Standard Model limit implies that no extra light particles exist which should have already been detected. There are further direct restrictions for extra couplings and masses which stem, for example, from rare decays and FCNC limits. Additional indirect limitations arise from radiative corrections. The agreement between the top mass value derived from radiative corrections with its experimental value leaves hardly any room for further custodial SU(2) violating effects in the electro-weak precision variable T. This most likely requires that the electro-weak symmetry breaking operator is a doublet under $\mathrm{SU}(2)_L$ with hypercharge $Y = 1$. Due to non-decoupling effects this also disfavours models where extra left-handed fermionic doublets contribute to T via loops. Furthermore the "size" of the symmetry breaking sector is limited by the electro-weak precision variable S, which roughly counts (with suitable weights) the number of degrees of freedom. The smallness of S probably points already towards the existence of a scalar Higgs particle as unitarity partner of the Goldstone bosons, since typically the contributions of vector like states (which mix and lead to $1/M^2$ contributions) are considerably larger than those of scalars, which only lead to logarithmic contributions. But even if there is a hint for a scalar Higgs particle, this does not yet determine whether this scalar is composite or fundamental. Therefore, the current data contains no evidence about the nature of the solution to the hierarchy problem and the related new physics beyond the Standard Model at TeV scales.

Models which are sucessful can be understood in terms of the Standard Model limit which allows them essentially to hide behind the Standard Model. We pointed out that it is possible to build models of dynamical symmetry breaking which are systematically viable due to a Standard Model limit. We presented model ingredients which can lead to such scenarios. The most interesting aspect is that dynamics quite different from QCD is required in order to produce a scalar resonance instead of a rho-like vector resonance. An example for a model which has a Standard Model limit is the so-called BHL model of electro-weak symmetry breaking. The BHL model is however unacceptable, since it cannot accommodate the correct top mass. With the help of the Pagels–Stokar relation we argued that this has systematic reasons and that the experimental top mass value is too small for a scenario with just a single top condensate. We argued therefore for a sequential breaking of a extended gauge group with more complex relations between the vacuum expectation values and masses. Finally we pointed to a left-right-symmetric model which realizes the required features given above and which is phenomenologically fully viable. This model has effective (composite) Higgs states, instead of Techni-rho like vector resonances and it includes a Standard Model limit. It is therefore not surprising that the model turns out to be phenomenologically acceptable.

Future precision measurements should ultimately find deviations between the Standard Model and the data (e.g. in quantities like R_b, etc.). In models with Standard Model limit such small deviations can then be explained more easily by departing from the limit, i.e., by lowering some of the additional states, which are expected in most scenarios. Such deviations will also rule out many particular models, since the ability to hide details beyond the Standard Model is lost. In any case, the underlying physics will most likely lead to restrictions in the mass of whatever assumes the role of the Higgs boson in the Standard Model. In supersymmetric models one expects a light Higgs boson due to the relations between the quartic Higgs self-couplings and the gauge couplings. In dynamical symmetry breaking the Higgs is typically at or above the value of the electro-weak vacuum expactation value. The value of the Higgs mass might thus act as a discriminator and give important hints on physics beyond the Standard Model even before any deviations from the Standard Model limit are found.

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