Properties of a Composite Higgs Particle in a Dynamically Broken Formulation of the Standard Model

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We suggest that dynamical breakdown of electroweak symmetry in an SU(3)×SU(2)×U(1) model of a single fermion family without fundamental scalars can reproduce the entire scalar sector of the standard model through fermion bound states and loop-induced couplings. The physical Higgs boson is essentially a $t\bar{t}$ bound state. The top-quark and Higgs-boson masses are approximately 300 and 400 GeV, respectively. The strengths of the induced couplings among the scalars do not indicate a strongly interacting scalar sector.

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All experimental indications are that the standard model (SM) provides an adequate description of strong and electroweak interactions at all energy scales explored to date. Not every aspect of the SM has been tested, however. Perhaps the most crucial untested feature is the mechanism responsible for the spontaneous breakdown of electroweak symmetry. A key step toward verification of the SM would be the discovery of the physical Higgs boson (H), the one directly observable particle of the four fundamental scalars introduced into the theory to explain symmetry breaking.

Should something resembling the H be discovered, however, it will have to be studied closely to determine whether it is indeed a fundamental particle. A composite H could well exist if the symmetry breaking is driven by a dynamical mechanism, rather than by fundamental scalars as in the SM. In models employing dynamical symmetry breaking (DSB),[1] the massless Goldstone states that become the longitudinal components of the $W^{\pm}$ and Z bosons are generally assumed to be fermion-antifermion composites; conceivably the same dynamics that produces these composites could generate one or more composite H's as well.

In this Letter we discuss the properties of a composite Higgs boson that appears to exist in a particular electroweak model incorporating DSB. We argue that the couplings of the composite H to all other particles, when evaluated to a first approximation, are identical to the corresponding tree-level couplings of the fundamental H in the SM. Thus a composite H of this type would be very difficult to distinguish from a fundamental H.

The model we consider[2] is based on the SU(3)×SU(2)×U(1) Lagrangian of the SM, minus all terms involving scalar fields. Note that omitting all reference to scalars means omitting all Yukawa couplings and scalar potential parameters. Thus, this model has fewer free parameters than the SM and, potentially, greater predictive power.

We restrict our discussion to a one-family model in order to avoid complications that arise in the multifamily case,[2,3] but that are not relevant to the present discussion. As will become clear below, in order to make contact with experiment we should identify this single family with the third family (t,b,τ,ν_τ).

Unlike more conventional DSB models, particularly those based on the technicolor idea,[4] our model contains no QCD-like strong interaction at the TeV scale to which electroweak symmetry breaking is ascribed. We nevertheless assume that the symmetry-breaking self-energies are of the electroweak scale. We further assume that the behavior of these self-energies over a range of momenta extending far above the electroweak scale is governed by the SM interactions alone. We refer to this scenario as ultraviolet dynamical symmetry breaking (UVDSB).

This idea is not new.[5-8] The basic premise is that nonperturbative effects give rise to symmetry-violating self-energies (dynamical masses) for the fermions. Ward identities for the broken currents imply that the related proper vertex functions contain zero-momentum poles, which are attributed to three massless bound states (Goldstone bosons) of fermions. The $W^{\pm}$ and Z absorb these states and become massive in a dynamical version of the Higgs mechanism. The gauge masses M_W and M_Z are computed to leading order in the gauge couplings by evaluating fermion-loop integrals involving the fermion-to-Goldstone-boson couplings. The Ward identities relate the latter to the fermion self-energies, with the result that M_W and M_Z can be expressed in terms of fermion masses and gauge couplings.

The symmetry breaking is driven by the fermions and is characterized by the fermion self-energy functions. The latter satisfy a Dyson equation whose solution behaves at large momenta as a linear combination of the "regular" solution, which drops off quickly with increasing momentum, and the "irregular" solution, which drops much more slowly. The regular solution is known to be dominant in the case of a QCD-like Dyson equation containing a sharp ultraviolet cutoff.[9] However, the irregular component dominates if suitable new interactions are assumed to be relevant at the scale of the cutoff[10] or if, alternatively, the ultraviolet cutoff is eliminated altogether. In a UVDSB theory one assumes that the irregular solution is dominant over a range of momenta extending far above the electroweak scale. The gauge-boson mass integrals then receive significant contributions from a very large range of loop momenta, which tends to counteract the smallness of the explicit gauge coupling factors associated with these integrals. M_W and M_Z turn out to be of similar magnitude to the masses of the heaviest fermions to which they couple.

In our one-family model, for example, the symmetry-violating fermion self-energies can be written for large $-p^2$ as[2]

$$\Sigma_f(p) \sim m_f \left(\frac{-p^2}{m_f^2} \right)^{-\epsilon_f} , \tag{1}$$

where f takes on the values t, b, and τ, and the ϵ_f are functions of the $SU(3)\times SU(2)\times U(1)$ gauge couplings g_3, g_2, and g_1:

$$\begin{aligned} \epsilon_t &= (16\pi^2)^{-1}(4g_3^2 + \tfrac{1}{3} g_1^2) , \\ \epsilon_b &= (16\pi^2)^{-1}(4g_3^2 - \tfrac{1}{6} g_1^2) , \\ \epsilon_\tau &= (16\pi^2)^{-1}(\tfrac{3}{2} g_1^2) . \end{aligned} \tag{2}$$

The form for Σ_f in (1) derives from a linearized version of the Dyson equation (in Landau gauge). As a result, the effective fermion masses m_f are undetermined in this approximation and must be put in by hand. Another approximation made in (1) is the neglect of renormalization effects (running couplings). A consistent treatment of these is possible only when the theory is asymptotically free,[6,8,11] which is not the case for the simple model we consider here.[12]

As discussed above, knowledge of the fermion self-energy functions enables us to compute the gauge-boson masses. Using (1) we obtain[2]

$$M_Z^2 \cos^2\theta_W = M_W^2 = \frac{g_2^2}{64\pi^2} \sum_f \frac{m_f^2}{\epsilon_f} \tag{3}$$

(a sum over colors is implied for quarks) from the fermion-loop contribution; the gauge-boson-loop contribution vanishes.[6,8] Equation (3) requires $m_t \approx 300$ GeV in order to give M_W and M_Z their observed values. Note that it appears to be impossible in this model to have a light top and a fourth family that dominates the sum in (3). This follows because experiment seems to rule out a fourth family with a light neutrino,[13] and neutrino mass cannot be accommodated in this model.[2,14]

Whether such a large t-b mass splitting is compatible with experimental bounds on the deviation of $\rho \equiv M_W^2/M_Z^2\cos^2\theta_W$ from unity is not known. Equation (3) gives $\rho = 1$ at lowest order (independent of t-b mass splitting[7]), but there are higher-order corrections to this relation. Part of these corrections come from contributions to the integrals for M_W and M_Z that are nonleading, in the sense that they are not enhanced by factors of ϵ_f^{-1} as in Eq. (3). A rough estimate of these terms gives $\delta\rho \equiv \rho - 1 \approx \epsilon_t \approx 3\%$. It is not at present clear whether a more complete calculation, including all $O(g^2)$ corrections relevant within the framework of our approach, will lead to a result within the experimental limit of $\approx 1\%$.

Until now we have seen nothing to indicate that there is a physical Higgs boson in this model. As for composite states, we know only that there are three unphysical Goldstone bosons ϕ^a $(a=1,2,3)$, which give masses to the $W^\pm$ and Z. Using Ward identities we can show the couplings of the ϕ^a to fermions to be [see Fig. 1(a)]

$$G^a(p+q,p) = -i\frac{g_2}{2M_W}[\tau^a\Sigma(p)P_R - \Sigma(p+q)\tau^a P_L] , \tag{4}$$

where $P_{R,L} = \frac{1}{2}(1 \pm \gamma_5)$ and Σ is a matrix of self-energies Σ_f in weak-isodoublet space.

Unfortunately, the Ward identities tell us nothing about massive states such as a composite H. We might anticipate, however, that a new particle resembling an H should occur in this model. The reason is that the theory as described so far would behave essentially like a theory of massive fermions and massive vector bosons (but no Higgs bosons) over an energy range extending far beyond the electroweak scale, i.e., a theory that violates perturbative unitarity. The Higgs boson plays an important role in restoring perturbative unitarity in the SM —perhaps a composite Higgs boson could perform an analogous function in our DSB model.

A stronger argument for the existence of a composite H can be obtained by appealing to the effective-action formalism for composite operators.[11] It has been argued in the context of simple DSB models[11,15] that translationally noninvariant fluctuations of certain fermion bilinears can be viewed as fluctuations of a massive scalar field. The composite particle associated with this field couples to fermions in proportion to their self-energies and may be regarded as a physical Higgs boson. An

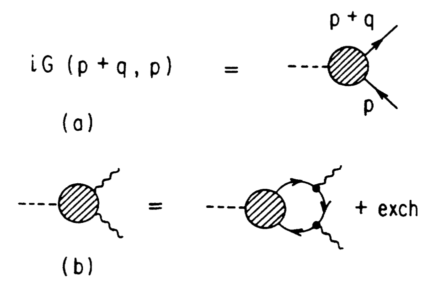

FIG. 1. Couplings of composite scalar bosons to (a) fermions and (b) gauge bosons. Dashed lines represent either Goldstone or Higgs bosons.

analysis is presently under way[16] for more realistic models, including the one considered here, where similar results appear to hold. For the purposes of this Letter, we simply assume there exists a composite H that couples to fermion self-energies, then derive consequences of this assumption.

Specifically, we assume that the H-to-fermion coupling matrix at large $-p^2$ is [cf. (4)]

$$G(p,p) \approx -\frac{g_2}{2M_W}\Sigma(p) . \tag{5}$$

We have not attempted to distinguish between the two fermion momenta p and $p+q$, since, in all situations with which we will be concerned, $\Sigma(p+q) \approx \Sigma(p)$. Note that if $\Sigma(p)$ is replaced by the fermion mass matrix M [i.e., if the weak momentum dependence of $\Sigma(p)$ is removed entirely], then (5) is just the tree-level H-to-fermion coupling in the SM. If the same replacement is made in (4), then together with (5) it becomes equivalent to a term in the Lagrangian of the form $-(g_2/2M_W)\bar{\psi}(H+i\tau^a\phi^a\gamma_5)M\psi$, where ψ is a weak isodoublet. Thus our assumption (5) implies that the H and the ϕ^a couple to fermions essentially as would a complete $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ multiplet of four fundamental scalars. Since the H couples to the isosinglet combination $\bar{\psi}M\psi$ and $m_t \gg m_b, m_\tau$, the H is essentially at $t\bar{t}$ bound state.

All we assume about the composite H is that it couples to fermions as in (5). However, it then follows that there are induced couplings of the H to other particles in the theory through their interactions with fermions. These induced couplings are calculable as loop effects.

For example, the coupling of the H to W^+W^- is given approximately by the one-loop graphs of Fig. 1(b). Using (1) and (5), we find upon integration

$$G^{\mu\nu}_{HWW} = \frac{g_2^3}{64\pi^2 M_W} g^{\mu\nu} \sum_f \frac{m_f^2}{\epsilon_f} . \tag{6}$$

Only terms that generate factors of ϵ_f^{-1} when integrated have been kept in (6); this has excluded all terms proportional to external momenta. The $O(\epsilon_f^{-1})$ terms in (6), as well as in (3), are consequences of the weak momentum dependence of the dynamical masses. In fact, except for constant factors, the integrals leading to (3) and (6) are identical in lowest order, and (6) can be rewritten as $G^{\mu\nu}_{HWW} = g_2 M_W g^{\mu\nu}$, independent of the fermion masses. This is precisely the tree-level HWW coupling in the SM, obtained here as a loop effect.

One-loop HZZ, $HHWW$, and $HHZZ$ couplings can be computed similarly, and in each case the leading $[O(\epsilon_f^{-1})]$ contribution is identical to the corresponding tree-level SM vertex. Remarkably, the only couplings that receive ϵ_f^{-1} enhancements are those that occur at tree level in the SM.

We can similarly determine induced H self-couplings. Doing so can be viewed as fixing the parameters of an effective scalar potential appropriate to the broken-symmetry phase. The $3H$ and $4H$ self-couplings are approximated by graphs consisting of a fermion loop with three and four of the couplings (5) attached. Once again, there is in each case a $O(\epsilon_f^{-1})$ part that does not involve external momenta [there are no $O(\epsilon_f^{-1})$ parts for couplings of more than four H's]. Neglecting the remaining terms gives

$$G_{4H} = \frac{g_2}{2M_W} G_{3H} = \frac{3g_2^4}{128\pi^2 M_W^4} \sum_f \frac{m_f^4}{\epsilon_f} . \tag{7}$$

The coupling G_{4H} is to be identified with the tree-level quartic coupling λ of the SM (normalized so that the bare $4H$ vertex has value $i\lambda$). The crucial difference is that λ is a free parameter in the SM, whereas G_{4H} is related to masses and gauge couplings as in (7). Taking $m_t \approx 300$ GeV, as discussed above, gives $G_{4H} \approx 10$, which is about 20% of the critical value $\lambda_c = 16\pi$ above which perturbative unitarity is violated in the SM.[17] Note that G_{3H} and G_{4H} are related through the same proportionality constant $(2M_W/g_2)$ as the corresponding tree-level couplings in the SM.

Couplings of the H to the composite Goldstone bosons ϕ^a can be computed in direct analogy with the H self-couplings, and once again the $O(\epsilon_f^{-1})$ results are identical to the corresponding tree-level couplings in the SM. The point we wish to stress is that the entire scalar sector of the SM is reproduced here through fermion bound states and loop-induced couplings. Thus, lowest-order scattering amplitudes for processes involving composite Higgs or Goldstone particles are identical to their tree-level counterparts in the SM. An important consequence is that amplitudes for processes like $f\bar{f} \to W_L^+ W_L^-$ remain below their unitarity bounds at high energies in our approximation. Amplitudes for processes like $W_L^+ W_L^- \to W_L^+ W_L^-$, which are proportional to the square of the Higgs-boson mass M_H, likewise satisfy unitarity if $M_H \lesssim 1$ TeV.[17]

A direct evaluation of the Higgs-boson mass M_H does not appear to be possible in the present framework. In the SM, M_H is related in a simple way (at tree level) to M_W and the quartic scalar coupling λ, or, equivalently, to the $3H$ and $4H$ couplings. No such relationship is evident here. However, in light of the demonstrated close correspondence between our model and the SM, it does not seem unreasonable to suppose that the mass of the composite Higgs boson is related to the H self-couplings in the same way as in the SM. This gives

$$M_H^2 = \frac{G_{3H}^2}{3G_{4H}} = \frac{g_2^2}{32\pi^2 M_W^2} \sum_f \frac{m_f^4}{\epsilon_f} , \tag{8}$$

or $M_H \approx \sqrt{2} m_t \approx 400$ GeV. A detailed analysis of the relevant effective action is expected to yield a more reliable estimate of M_H.[16]

While the model we have described encompasses only the heaviest family of quarks and leptons, it is not unreasonable to consider its phenomenological implications, since in any similar multifamily model all of the dynamically generated quantities derived above would continue to be dominated by contributions from the heaviest fermions. As shown above, the lowest-order couplings of the composite H to other particles and to itself are identical to the corresponding couplings in the SM, so to this order there are no phenomenological differences between this model and the SM. Differences might arise at higher orders, but it remains to be shown that the UVDSB program can be extended consistently beyond the lowest-order approximation.

Nevertheless, the present model does make certain predictions that can be tested in future experiments, namely, the relationships among masses and couplings derived above. For example, the values obtained for G_{4H} and M_H indicate that the scalar sector is not strongly interacting, and would not be consistent with large excess W^+W^- production at the Superconducting Super Collider.[18] Another prediction is that the fermion and gauge-boson masses are related as in Eq. (3). Thus, this model could be ruled out if a relatively light ($\ll 300$ GeV) top quark were discovered or if very heavy ($\gg 300$ GeV) weak-isodoublet fermions were found to exist. The discovery of a light Higgs boson ($\ll 400$ GeV) would also provide evidence against this model, although the prediction (8) for M_H is somewhat less reliable than the others.

A fundamental difficulty with the above predictions, and with the UVDSB program as a whole, is that the evaluation of the involved loop integrals requires knowledge of the form of the fermion self-energy functions Σ_f at extremely short distance scales. New physics at very high scales could significantly affect the ultraviolet-dominated loop integrals.[19] This would change the expressions for M_W, M_Z, and the induced scalar-sector couplings in terms of fermion masses. However, our central result, that the UVDSB model reproduces the SM's scalar sector, is independent of fermion masses and may well survive the introduction of new physics.

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