Distinguishing the Higgs Boson from the Dilaton at the Large Hadron Collider

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It is likely that the LHC will observe a color- and charge-neutral scalar whose decays are consistent with those of the standard model (SM) Higgs boson. The Higgs interpretation of such a discovery is not the only possibility. For example, electroweak symmetry breaking could be triggered by a spontaneously broken, nearly conformal sector. The spectrum of states at the electroweak scale would then contain a narrow scalar resonance, the pseudo-Goldstone boson of conformal symmetry breaking, with Higgs-boson-like properties. If the conformal sector is strongly coupled, this pseudodilaton may be the only new state accessible at high energy colliders. We discuss the prospects for distinguishing this mode from a minimal Higgs boson at the LHC and ILC. The main discriminants between the two scenarios are (i) cubic self-interactions and (ii) a potential enhancement of couplings to massless SM gauge bosons.

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The LHC experiment at CERN will begin operations at the end of this year. In the early stages of running, one of its main priorities will be to either rule out or confirm the minimal standard model mechanism of electroweak symmetry breaking (EWSB), generated by the expectation value of an $SU(2)$ doublet Higgs scalar field.

While there are some constraints on the minimal Higgs picture, either from direct collider searches (e.g., the LEP bound $m_H > 114$ GeV), or from precision electroweak physics, the Higgs sector of the standard model (SM) is largely uncharted territory. In anticipation of the LHC, it is important for particle theorists to thoroughly map out the space of theories of EWSB and identify collider signals that might be used to distinguish between different scenarios.

In this Letter, we study the LHC signatures of a class of theories in which EWSB at the scale $v \simeq 246$ GeV is triggered by the spontaneous breaking of scale symmetry at an energy scale $f \geq v$ [1]. Scenarios of this kind typically involve new strongly coupled, nearly conformal dynamics at a scale $\Lambda_{\rm CFT} \sim 4\pi f$ which then feeds into an EW sector composed of new heavy states with masses of order $\Lambda_{\rm EW} \sim 4\pi v$. Examples of such theories can be realized as 4D strongly coupled gauge theories, as in the original models of walking technicolor [2], or via anti–de Sitter/conformal field theory (CFT) correspondence as Randall-Sundrum 5D warped geometries [3] with EWSB through either Higgs VEVs vacuum expectation value or boundary conditions [4]. It is important to note, however, that EWSB through a nearly conformal sector need not involve strong dynamics. A simple example is the minimal Higgs model itself: in the absence of a Higgs sector the SM becomes approximately scale invariant down to the QCD scale. Thus, in addition to breaking EW gauge symmetry at a scale v, the Higgs potential simultaneously breaks conformal invariance both spontaneously at the scale $f = v$, and explictly, due to the presence of a Higgs mass term in the potential.

While the scenarios described above generically predict new heavy states with masses of order $\Lambda_{\rm EW} \sim 1$ TeV associated with EWSB (whose presence is necessary, e.g., to unitarize high energy scattering of longitudinal EW gauge bosons), the detailed spectroscopy is in general highly model dependent. In addition, such states may be too broad to be individually resolved at the LHC. Nevertheless, these theories always contain an electroweak singlet scalar field $\chi(x)$, the *dilaton* mode, in their spectrum. This mode is simply the pseudo-Goldstone boson of spontaneously broken (approximate) scale invariance. The dilaton becomes massless in the limit in which conformal symmetry is recovered. Therefore its mass is naturally light, proportional to the scale f times the parameter that controls deviations from exact scale invariance. A light, narrow resonance is therefore a distinguishing feature of new nearly conformal dynamics.

Since the minimal SM Higgs doublet is a weakly coupled example of EWSB triggered by nearly scale invariant physics, a light Higgs boson itself can be identified with the dilaton through the relation $\chi(x) = \sqrt{H^\dagger H}(x)$. From this point of view, the couplings of the Higgs boson to the SM are simply a consequence of underlying approximate scale invariance (i.e., the soft Higgs theorems of [5]). This observation, however, also indicates that differentiating new conformal physics from the minimal SM at colliders may be difficult.

Below, we use symmetry arguments to work out the pseudodilaton couplings to the SM fields which arise from an underlying conformal sector. These results lead to a definite pattern of collider signals that can be contrasted with those of a minimal Higgs scenario. Briefly, we find the following: Scale invariance implies that the tree-level pseudodilaton couplings are obtained from those of the Higgs boson by replacing the electroweak scale v with the scale f of conformal symmetry breaking. At the loop level, the dilaton has potentially enhanced couplings to massless SM gauge bosons relative to those of a light

Higgs boson. However, the precise values are model dependent and do not constitute a clean method of distinguishing between models. Regardless, if f is close to v and no other light states appear, it would be nearly impossible to distinguish the new strong dynamics from the minimal Higgs model. Finally, the dilaton self-couplings depend on the dimension of operators that explicitly break the conformal symmetry, and this provides an opportunity to differentiate the dilaton from the Higgs boson at the ILC.

Our setup (an expanded version of our results can be found in [6]) is as follows: we adopt a framework with SM gauge bosons and fermions as well as an electroweak neutral scalar, the dilaton, whose mass is protected by approximate scale invariance. All other states responsible either for conformal or electroweak breaking are taken to be roughly heavier than a scale $\Lambda_{\rm EW} \sim 4\pi v$. The interactions among the light fields are described by a low-energy effective Lagrangian with nonlinearly realized $SU(2)_L \times U(1)_Y$ and (approximate) conformal invariance.

A simple way of incorporating the nonlinearly realized scale invariance is to add a field $\chi(x)$ that serves as a conformal compensator. If one writes the original Lagrangian in a basis of anomalous dimension eigenoperators, $\mathcal{L} = \sum_i g_i(\mu)\mathcal{O}_i(x)$, then under scale transformations $x^\mu \to e^\lambda x^\mu$, $\mathcal{O}_i(x) \to e^{\lambda d_i}\mathcal{O}_i(e^\lambda x)$. Assigning the scale transformation law $\chi(x) \to e^\lambda \chi(e^\lambda x)$, we simply need to make the replacement $g_i(\mu) \to g_i(\mu \frac{\chi}{f})(\frac{\chi}{f})^{4-d_i}$, in $\mathcal{L}$. Here $f = \langle\chi\rangle$ is the order parameter for scale symmetry breaking, determined by the dynamics of the underlying strong sector. The Goldstone boson associated with conformal symmetry breaking is parameterized as $\chi(x) = f e^{\sigma(x)/f}$, which transforms nonlinearly under scale transformations, λ:$\sigma(x)/f \to \sigma(e^\lambda x)/f + \lambda$. However, a more convenient parameterization for the fluctuations about the VEV is vacuum expectation value $\bar{\chi}(x) = \chi(x) - f$.

In our scenario, $\mathcal{L}$ above is the electroweak chiral Lagrangian [7]. This theory provides a convenient model-independent description of a strongly interacting EWSB sector, including not only the tree-level couplings of massive SM gauge bosons and fermions but also all precision electroweak observables, which are encoded in the coefficients of higher dimensional operators. In the unitary gauge, the relevant couplings to the dilaton are obtained by making the replacement $v \to v\chi/f$ in the electroweak chiral theory. This gives,

$$\mathcal{L}_{\chi,\rm SM} = \left(\frac{2\bar{\chi}}{f} + \frac{\bar{\chi}^2}{f^2}\right)\left[m_W^2 W_\mu^+ W^{-\mu} + \frac{1}{2}m_Z^2 Z_\mu Z^\mu\right] + \frac{\bar{\chi}}{f}\sum_\psi m_\psi \bar{\psi}\psi, \qquad (1)$$

which is identical in form to the couplings of a minimal Higgs boson.

In the limit of exact scale invariance χ is derivatively self-coupled. Ignoring for the time being terms that explicitly break the symmetry, self-interactions of the dilaton are of the form, e.g., $c_4(\partial_\mu\chi\partial^\mu\chi)^2/(4\pi\chi)^4$, where the constant $c_4 \sim \mathcal{O}(1)$ depends on the details of the underlying CFT. The inverse powers of χ are necessary to ensure that $\mathcal{L}_\chi$ transforms correctly under scalings.

In addition, the theory may possess explicit sources of scale symmetry breaking. For example, suppose that conformal invariance is broken by the addition of an operator $\mathcal{O}(x)$ with scaling dimension $\Delta_\mathcal{O} \neq 4$ to the Lagrangian, $\mathcal{L}_{\rm CFT} \to \mathcal{L}_{\rm CFT} + \lambda_\mathcal{O}\mathcal{O}(x)$. It is straightforward to include this pattern of symmetry breaking by the introduction of a spurion field into the low-energy effective theory. This spurion constrains the nonderivative interactions of $\chi(x)$ to be of the form

$$V(\chi) = \chi^4 \sum_{n=0}^{\infty} c_n(\Delta_\mathcal{O})\left(\frac{\chi}{f}\right)^{n(\Delta_\mathcal{O}-4)}, \qquad (2)$$

where the coefficients $c_n \sim \lambda_\mathcal{O}^n$ depend on the dynamics of the underlying CFT. By assumption, this dynamics must be such that $V(\chi)$ is minimized at $\langle\chi\rangle = f$ with $m_\chi^2 = d^2V(\langle\chi\rangle)/d\chi^2 > 0$. In general, the coefficients c_n depend on the scaling dimension $\Delta_\mathcal{O}$. We assume that the c_n are nonsingular in the limit $\Delta_\mathcal{O} \to 4$.

It is not possible to make detailed predictions without knowledge of the coefficients c_n in $V(\chi)$ unless there exists a small expansion parameter. Here we are interested in the case where the explicit conformal breaking term above is small. This can be either because the operator $\mathcal{O}$ is nearly marginal ($|\Delta_\mathcal{O} - 4| \ll 1$), as is the case in walking technicolor theories or RS models stabilized by the scenario of [8], or because the coefficient $\lambda_\mathcal{O}$ is chosen to be small in units of f, as in the case of the minimal Higgs model. If this happens, it is possible to obtain definite expressions for the dilaton self-couplings once the parameters m and f are fixed. We find that the potential is $V(\bar{\chi}) = \frac{1}{2}m^2\bar{\chi}^2 + \frac{\lambda}{3!} \times \frac{m^2}{f}\bar{\chi}^3 + \cdots$, where $m^2 \ll f^2$ is proportional to the small symmetry breaking parameter: $m^2/f^2 \propto \lambda_\mathcal{O}$ for $\lambda_\mathcal{O} \ll 1$ (in units of f) and $\Delta_\mathcal{O}$ arbitrary, or $m^2/f^2 \sim |\Delta_\mathcal{O} - 4|$ for $|\Delta_\mathcal{O} - 4| \ll 1$ and $\lambda_\mathcal{O}$ of arbitrary size. The cubic coupling is given by $\lambda = (\Delta_\mathcal{O} + 1) + \mathcal{O}(\lambda_\mathcal{O})$ for $\lambda_\mathcal{O} \ll 1$, or $\lambda = 5 + \mathcal{O}(|\Delta_\mathcal{O} - 4|)$ when $|\Delta_\mathcal{O} - 4| \ll 1$, and is in principle a probe of the scaling dimension of the operator responsible for scale symmetry breaking. We do not expect scale symmetry breaking to occur if $\mathcal{O}(x)$ is an IR irrelevant perturbation. This implies the bound $\lambda \leq 5$ that is saturated near marginality. In addition, the conformal algebra together with unitarity implies $\lambda \geq 2$. Moreover, for $\Delta_\mathcal{O} = 2$ and $\lambda_\mathcal{O} \ll f^2$ this result reproduces the usual Higgs trilinear coupling. Note that when $|\Delta_\mathcal{O} - 4| \ll 1$ the entire potential for χ, up to corrections of order $(\Delta_\mathcal{O} - 4)^2$, is calculable: $V(\chi) = (m^2\chi^4/16f^2)[4\ln\chi/f - 1] + \mathcal{O}(\Delta_\mathcal{O} - 4)^2$. In addition, note that the assumption of scale invariance ensures that our predictions for the lightness of χ and for the cubic coupling are stable against radiative corrections due to SM loops.

In the SM, Higgs couplings to the top quark and the massive gauge bosons induce the couplings $H\gamma\gamma$ and Hgg. The same mechanism induces the dilaton couplings $\chi\gamma\gamma$

and χgg. Because these processes are generated by loop effects, these couplings are also sensitive to contributions from heavy particles present in any extension of the SM. Since these couplings are crucial for collider phenomenology we derive them here and show that the dilaton coupling to gluons and photons can be significantly enhanced under very mild assumptions about high scale physics.

Let us begin by recalling the situation for the SM loop-induced Higgs couplings to gluons. The logic is identical for the couplings to two photons. For heavy particles, defined as $m_i \gg m_h$, the Higgs boson couples to gluons through an operator

$$\mathcal{L}_{hGG} = \frac{\alpha_s}{8\pi} \sum_i b_0^i \frac{h}{v} (G_{\mu\nu}^a)^2, \tag{3}$$

where we have expanded $\ln(H^\dagger H/v^2) = 1/2 + h/v + \cdots$ and $G_{\mu\nu}^a$ is the canonically normalized gluon field strength. The sum runs over the heavy fields only and b_0^i is the contribution of each heavy particle to the one-loop QCD beta function, normalized as $\beta_i(g) = b_0^i g^3/16\pi^2$. As expected, this result is independent of the heavy masses. This well known result is modified by the loop contributions of other heavy particles which are not part of the SM. In addition to inducing the term in Eq. (3), heavy states can also generate new dimension-six operators of the form, e.g., $\frac{\alpha_s}{4\pi} c_{hg} H^\dagger H (G_{\mu\nu}^a)^2$, which depending on the size of the coefficient c_{hg} can significantly modify the properties of a light Higgs boson [9].

The dilaton couplings to massless gauge bosons can be simply obtained by making the replacement $\frac{2m_i^2}{v^2} H^\dagger H \to \frac{m_i^2}{f^2}\chi^2$, in the derivation of the Higgs couplings. Again, one can split the sum over all colored particles into sums over light and heavy states, where the dividing scale is given by the dilaton mass. Note that if one assumes that QCD is fully embedded in the conformal sector, one can make UV insensitive predictions, since by conformal invariance $\sum_{\text{light}} b_0 + \sum_{\text{heavy}} b_0 = 0$. Thus the effective coupling at one-loop is

$$\mathcal{L}_{\chi gg} = -\frac{\alpha_s}{8\pi} b_0^{\text{light}} \frac{\bar{\chi}}{f} (G_{\mu\nu}^a)^2, \tag{4}$$

where $b_0^{\text{light}} = -11 + \frac{2}{3} n_{\text{light}}$. The number of light fermions, n_{light}, is either $n_{\text{light}} = 5$ if the dilaton is lighter than the top quark, or $n_{\text{light}} = 6$ otherwise. Equation (4) indicates about a tenfold increase of the coupling strength compared to that of the SM Higgs boson, which could have profound consequences at the LHC. Unlike the Higgs case, corrections to this result from higher dimension operators are negligible. For example, one might consider operators such as $g_s^2 \frac{c_{\chi g}}{(4\pi\chi)^2} D_\alpha G_{\mu\nu}^a D^\alpha G^{\mu\nu a}$. However, such operators are suppressed by powers of $m^2/f^2 \ll 1$ relative to the terms coming from the conformal anomaly.

To summarize the discussion so far, the couplings of the dilaton at energies below the scale $4\pi f$ are given by

$$\mathcal{L}_\chi = \frac{1}{2}\partial_\mu \bar{\chi} \partial^\mu \bar{\chi} - \frac{1}{2} m^2 \bar{\chi}^2 + \frac{\lambda}{3!}\frac{m^2}{f}\bar{\chi}^3 + \mathcal{L}_{\text{SM},\chi} + \frac{\alpha_{EM}}{8\pi f} c_{\text{EM}} \bar{\chi} (F_{\mu\nu})^2 + \frac{\alpha_s}{8\pi f} c_G \bar{\chi} (G_{\mu\nu}^a)^2, \tag{5}$$

with $\mathcal{L}_{\text{SM},\chi}$ as in Eq. (1) and the coefficients c_{EM}, c_G were discussed above. For example, if electromagnetic and strong interactions are embedded in the conformal sector at high scales, $c_{\text{EM}} = -17/9$ if $m_W < m < m_t$ and $c_{\text{EM}} = -11/3$ if $m > m_t$, while $c_G = 11 - 2n_{\text{light}}/3$, where n_{light} is the number of quarks lighter than the dilaton.

Given the similarity to minimal Higgs-boson physics, it is possible to use existing studies of Higgs-boson properties at colliders to understand the physics of a light dilaton as a function of the model parameters m, f, and the couplings λ, c_{EM}, c_G.

At LEP, the main production channel for dilaton production is, as for the Higgs boson, associated production with a virtual Z boson, $e^+e^- \to HZ^*$. The cross section for dilaton production is suppressed by a factor $(v/f)^2$ relative to the corresponding Higgs cross section at the same mass. The LEP collaborations have combined their data to search for the Higgs, including a search for Higgs particles with an anomalous (non-SM) HZZ coupling [10]. This result is immediately applicable to the bounds on the dilaton mass and coupling.

Figure 10 in Ref. [10] summarizes the bound on the dilaton mass and decay constant, where in our case $\xi^2 = (v/f)^2$. Roughly, the dilaton with mass 90 GeV $< m <$ 110 GeV is excluded if $(v/f)^2 > 0.1$ and with mass 12 GeV $< m <$ 90 GeV it is excluded for $(v/f)^2 > 0.01$. These limits predominantly come from the $b\bar{b}$ decay channel, which is kinematically suppressed below 12 GeV. Other available decay channels have been employed for very light masses [11]. Values $m <$ 12 GeV are excluded if $(v/f)^2 > 0.1$; see Fig. 5 in Ref. [11].

The dilaton decay width into quarks and leptons is also suppressed by the factor $(v/f)^2$. However, this discrepancy is not relevant for the LEP search as the branching ratios to fermions remain unchanged. For $(v/f)^2 < 10^{-2}$, LEP is not able to detect the dilaton irrespective of its mass, while for $(v/f)^2 > 10^{-2}$ the suppression of the width is not observable. In this latter case the dilaton decays very promptly and does not have a displaced decay vertex. Therefore its signatures are identical to Higgs-boson signatures.

There are four important production channels for the dilaton at hadron colliders: gluon fusion $gg \to \chi$, associated production with vector bosons $q\bar{q} \to W/Z + \chi$, vector boson fusion $qq \to qq + \chi$, and associated production with the top quark gg, $q\bar{q} \to t\bar{t} + \chi$. The first process, $gg \to \chi$, is likely to be sensitive to new heavy states as we discussed above. For example, assuming that QCD is embedded in the conformal sector, we find a large enhancement of the χgg coupling. This could easily overcome the suppression factor v/f when compared with the $gg \to h$

cross section. The cross sections for the remaining processes scale as $(v/f)^2$ compared to the Higgs-boson production cross sections. Higher order QCD corrections, which are often sizable, do not alter the scaling of the cross sections with $(v/f)^2$ since each of these processes contains just one vertex involving the dilaton. As we already discussed, the dominant branching ratios of the dilaton are the same as for the Higgs boson so most of the Higgs-boson search strategies can be applied directly. The only caveat is in the mode $\chi \to \gamma\gamma$. The width of this decay is likely to be modified by physics beyond SM and scaling the results obtained for the Higgs boson may not be reliable.

Given the simple scaling of the cross section we can estimate the reach of LHC as a function of the dilaton mass and the decay constant f. The statistical significance of the Higgs signal at ATLAS has been presented, for example, in Refs. [12,13]. The significance of the dilaton signal can be obtained from the significance of the Higgs signal by rescaling $(\frac{S}{\sqrt{B}})_\chi = c_G^2 \frac{v^2}{f^2}(\frac{S}{\sqrt{B}})_{\rm Higgs}$, where we assume that the production cross section is dominated by the gluon fusion process. It is easiest to discover a heavy dilaton when the decays to WW and ZZ dominate. Very crudely, with a 100 fb^{-1} of integrated luminosity a discovery is possible when $c_G^2(v/f)^2 > 1/8$ for $m > 160$ GeV. For a lighter dilaton the statistical significance decreases with mass, so if $c_G^2(v/f)^2 \sim 1/10$ one may have to wait to collect about 300 fb^{-1} worth of data for detection. For details see Fig. 3.49 in Ref. [13].

In addition to discovery, the LHC will be able to measure the dilaton couplings to gauge bosons and the top quark by measuring event rates in different channels. Depending on the mass and the production channel one expects a 10% to 30% accuracy for the extraction of the couplings. The measurement of the cubic self coupling seems hopeless at the LHC even if $f \approx v$ and it only gets harder for $f > v$ [13].

A linear collider with $\sqrt{s} = 500$–1000 GeV would provide an ideal environment for the study of the dilaton and for distinguishing the dilaton from the Higgs boson. A number of precision measurements can be performed. See Ref. [13] for a comprehensive review.

First, the couplings to the gauge bosons and several branching ratios can be measured at a one to few percent level. In addition to determining f it would be a test of whether or not different coupling are scaled by the universal factor v/f.

Second, for $m > 200$ GeV and $f \approx v$, the dilaton would be broad enough for its total width to be observed directly. This would provide yet another check of the universal rescaling of the couplings relative to the Higgs boson. For $f > v$ such a measurement is possible if $m > (f/v)^{2/3}200$ GeV since, in this mass range, the total width increases as mass cubed. (The Higgs-boson mass reach of a linear collider is approximately $0.8\sqrt{s}$.)

It would be fascinating if $f \approx v$ to a degree of accuracy that previously described coupling measurements would not distinguish the dilaton from the Higgs boson. The cubic coupling may then provide the only probe of how conformal symmetry is broken. If the conformal symmetry is broken by a nearly marginal operator we expect a slight enhancement of the cubic coupling by a factor of $5/3$. This is large enough to be probed by the ILC if the dilaton is light enough. The limiting factor for this measurement is the small production cross section of Higgs-boson–dilaton pairs. For $f \approx v$, we can adopt the results of studies on Higgs pair production, which estimate that the cubic coupling can be measured to within, e.g., 20% if $m = 120$ GeV and to within 30% if $m = 140$ GeV [13]. Note that for $f \gg v$ the error in the cubic coupling determination should be multiplied by an additional factor of $(f/v)^2$ if the limitation to the measurement is assumed to be purely statistical and by $(f/v)^4$ under the assumption that it is background dominated. Although neither of these extremes is likely to capture the true experimental situation, it does indicate that for f significantly larger than v the pair production signal is not an effective probe of dilaton physics.

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[1] E. Gildener and S. Weinberg, Phys. Rev. D **13**, 3333 (1976).
[2] B. Holdom, Phys. Lett. B **150**, 301 (1985); T. W. Appelquist, D. Karabali, and L. C. R. Wijewardhana, Phys. Rev. Lett. **57**, 957 (1986).
[3] L. Randall and R. Sundrum, Phys. Rev. Lett. **83**, 3370 (1999).
[4] K. Agashe, A. Delgado, M. J. May, and R. Sundrum, J. High Energy Phys. 08 (2003) 050; K. Agashe, R. Contino, and A. Pomarol, Nucl. Phys. **B719**, 165 (2005); C. Csaki, C. Grojean, L. Pilo, and J. Terning, Phys. Rev. Lett. **92**, 101802 (2004); M. A. Luty and T. Okui, J. High Energy Phys. 09 (2006) 070.
[5] M. A. Shifman, A. I. Vainshtein, M. B. Voloshin, and V. I. Zakharov, Sov. J. Nucl. Phys. **30**, 711 (1979); A. I. Vainshtein, V. I. Zakharov, and M. A. Shifman, Sov. Phys. Usp. **23**, 429 (1980).
[6] W. D. Goldberger, B. Grinstein, and W. Skiba, arXiv:0708.1463.
[7] T. Appelquist and C. W. Bernard, Phys. Rev. D **22**, 200 (1980).
[8] W. D. Goldberger and M. B. Wise, Phys. Rev. Lett. **83**, 4922 (1999).
[9] A. V. Manohar and M. B. Wise, Phys. Lett. B **636**, 107 (2006).
[10] R. Barate *et al.* (LEP Working Group for Higgs boson searches), Phys. Lett. B **565**, 61 (2003).
[11] P. D. Acton *et al.* (OPAL Collaboration), Phys. Lett. B **268**, 122 (1991).
[12] S. Asai *et al.*, Eur. Phys. J. C **32S2**, s19 (2004).
[13] A. Djouadi, arXiv:hep-ph/0503172, and Refs. therein.