Phenomenology of a light scalar: the dilaton

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ABSTRACT: We make use of the language of non-linear realizations to analyze electro-weak symmetry breaking scenarios in which a light dilaton emerges from the breaking of a nearly conformal strong dynamics, and compare the phenomenology of the dilaton to that of the well motivated light composite Higgs scenario. We argue that – in addition to departures in the decay/production rates into massless gauge bosons mediated by the conformal anomaly – characterizing features of the light dilaton scenario (as well as other scenarios admitting a light CP-even scalar not directly related to the breaking of the electro-weak symmetry) are off-shell events at high invariant mass involving two longitudinally polarized vector bosons and a dilaton, and tree-level flavor violating processes. Accommodating both electro-weak precision measurements and flavor constraints appears especially challenging in the ambiguous scenario in which the Higgs and the dilaton fields strongly mix. We show that warped higgsless models of electro-weak symmetry breaking are explicit and tractable realizations of this limiting case.

The relation between the naive radion profile often adopted in the study of holographic realizations of the light dilaton scenario and the actual dynamical dilaton field is clarified in the Appendix.

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1. Introduction

The detection of the Higgs boson is one of the most ambitious aims of the CERN large hadron collider (LHC). Precision measurements strongly suggest the presence of this excitation in the hundred GeV range, but its very existence seems to be in contrast with our understanding of naturalness. Any alternative electro-weak symmetry breaking (EWSB) sector has to explain this tension.

If the forthcoming experiments detect a light CP-even scalar but no accompanying states – the latter being too heavy or too broad to be directly observed – it would turn out to be crucial to understand whether the light scalar is the actual Higgs, namely the physical component of an electroweak doublet. In this letter we would like to address this issue using a phenomenological approach.

A parametric separation between the mass of the light scalar and those of the heavy states is most naturally justified by a strongly coupled EWSB theory in which the scalar emerges as an approximate Goldstone boson of some broken global symmetries. We focus on two classes of models which are consistent with our assumptions. The first class is a deformation of the standard model Higgs sector, in which the CP-even scalar emerges as the fourth excitation of a light composite Higgs doublet. The physics of this strongly coupled light Higgs (SILH) has been studied in detail in [1].

The second class of models include scenarios in which the strong EWSB sector is a nearly conformal field theory (CFT) that admits a light dilaton in the spectrum [2]. In this class the CP-even scalar emerges as an approximate Goldstone boson of the breaking of the conformal symmetry down to the Poincaré group, and it is not directly related to the EWSB sector. Appropriate field theories are expected to reproduce the invoked mechanism
in a quantum mechanical context. However, a generic strong dynamics is not a realistic candidate. An heuristic motivation is that in the latter case the conformal breaking is governed by the running of the gauge coupling; in the IR the coupling is strong and there is no small parameter suppressing the dilaton mass compared to the dynamically generated scale \[3\]. Lattice simulations confirm these conclusions.

This picture changes if the small coupling is taken to be \(1/N\), where \(N\) is the number of fundamental constituents. Evidences in favor of this mainly come from the gauge/gravity correspondence (see \[4\] for a recent example). The Randall-Sundrum model \[5\], as well as several other models for physics beyond the standard model belonging to the same universality class, are well known realizations of the light dilaton scenario. Supersymmetry then offers a simple way to relate the presence of a light \(\eta'\) and the dilaton\(^1\). Indeed, in a supersymmetric theory the scale current and the axial current are part of the same multiplet. If there exists a limit (the large \(N\) limit in a nonsupersymmetric theory!) in which the anomalous axial symmetry is exact, then in a non-chiral vacuum an approximate Goldstone boson (GB), the so called \(\eta'\), would emerge along with a CP-even pseudo-GB partner, i.e. the dilaton. All of these considerations suggest that theories with many fundamental degrees of freedom may represent the natural framework for the realization of the light dilaton scenario. This would have implications in the study of the phenomenological signatures of this class of models.

It is very important to appreciate that, in the case the operators responsible for EWSB and the spontaneous breaking of the CFT do coincide, the Higgs and the dilaton are the very same field. This is what happens at the classical level in the standard model \[6\], and what is expected to occur in well calibrated walking technicolor scenarios \[7\]. In this limit the SILH and the light dilaton effective theories overlap. This intrinsic ambiguity suggests that a phenomenological discrimination between the two scenarios is a highly non-trivial task, in general. Nevertheless, we believe that the identification of features that most naturally characterize one theory as opposed to the other, and in particular a study of how natural is the above limit and what the phenomenological implications are, has physical relevance. This is one of the aims of the present study.

### 2. Phenomenology of a light scalar

A light and chargeless CP-even scalar can arise as the physical excitation of a Higgs doublet, but in principle it may also arise as a generic pseudo Goldstone boson of an appropriately engineered strong dynamics. In the latter case the scalar field would have no direct relations with the Higgs sector. In this paper we address the question whether these two classes of scenarios can be disentangled under the hypothesis that the only new physics detected in the forthcoming experiments is the light CP-even scalar. To be definite we choose a representative of each class: from the first class we consider the strongly interacting light Higgs scenario (SILH), while from the second we consider the strongly interacting light dilaton (SILD) scenario.

\(^{1}\)I am indebted to Sergio Cecotti for an illuminating discussion of this point.
Following the notation introduced in [1], we decide to describe our models by using three free parameters: \( f \), \( g_\rho \), and \( m \). The first parameter is the energy scale \( f \) at which the strong dynamics spontaneously breaks its approximate global symmetry, and hence leading to the appearance of a Goldstone doublet (SILH) or singlet (SILD). The broken symmetry can be either a global one, as in the case of the SILH scenario, or a space-time symmetry, as in the case of the dilaton scenario. The scale \( f \) is basically a measure of the strength of the Goldstone boson interactions. The second parameter, \( g_\rho \), is the typical coupling of the massive resonances, which in particular is expected to enter into the definition of the masses of the heavy composites as a proportionality factor, \( m_\rho \propto g_\rho \). The statement that the new physics is strong compared to the standard model can be expressed via the relation \( g_{SM} \ll g_\rho \), where \( g_{SM} \) is a typical standard model coupling and \( g_\rho = 4\pi \) for a maximally strong dynamics. The third parameter is the mass of the pseudo Goldstone boson, \( m \), and crucially depends on the explicit source of symmetry breaking. The rules of non-linear realizations are expected to hold up to corrections of order \( m^2/f^2 \ll 1 \).

The physics of both the SILH and SILD can be captured by the effective lagrangian of a generic chargeless CP-even scalar \( S \):

\[
L_{\text{eff}} = \frac{1}{2} m_V^2 V_\mu \left( 1 + 2a_1 \frac{S}{v} + a_2 \frac{S^2}{v^2} \right)
+ \bar{\psi}_i \psi_j \left[ m_i \delta^{ij} \left( 1 + b \frac{S}{v} \right) + b^{ij} \frac{S}{v} \right]
+ c \frac{g_{SM}^2}{16\pi^2} F_{\mu\nu}^2 \frac{S}{v} + \ldots
\]  

We allowed the presence of a higher dimensional operator \( S F_{\mu\nu}^2 \), where \( F_{\mu\nu} \) is the photon or gluon field strength, because no such a coupling is present at leading order, and it may play a phenomenological role if the scale of new physics is not too high [3].

The standard model Higgs boson is just a particular case of the theory (2.1) with \( a_1 = a_2 = b = 1 \) and \( b^{ij} = c = 0 \). For completeness we mention that the integration of the top induces a coupling \( c = O(1) \).

**The SILH:** In the SILH model the full Higgs doublet emerges as a Goldstone field, the massive composites being naturally at the scale \( m_\rho = g_\rho f \). The explicit violation of the global symmetry induces the generation of a potential for the doublet and eventually implies EWSB. In the broken electro-weak phase the Higgs boson is the only physical and light relic of the strong dynamics. A generic coupling of the SILH to the standard model vectors and fermions differs by an \( O(v^2/f^2) \) correction with respect to the standard model Higgs, and approaches the latter as the heavy states decouple, \( v \ll f \). Although this regime appears unnatural, it turns out that a mild suppression \( v < f \) is sufficient for fitting the electro-weak precision data. In this model the authors of [1] found that the physics of the light scalar can be effectively described by the lagrangian (2.1) with:

\[
a_1 = 1 - c_H \frac{v^2}{2f^2}; \quad a_2 = 1 - 2c_H \frac{v^2}{f^2}; \quad b = 1 - (c_y + c_H/2) \frac{v^2}{f^2}; \quad b^{ij} = 0; \quad c = 0,
\]  

(2.2)
where \( c_y, H \) are \( O(1) \) parameters defined in [1]. In the rest of the paper we will be mainly concerned with the dilaton physics. We refer the reader to the paper [1] for more details on the SILH scenario.

**The SILD:** In the SILD the only Goldstone mode emerging at a scale \( f \) is the dilaton itself, while EWSB is triggered by a strong dynamics at a somewhat lower scale \( v = 250 \) GeV. For practical purposes the standard model symmetry can be taken to be non-linearly realized below the characteristic mass scale of the strong Higgs sector, \( m_p = g_p v \). Notice that the light scalar \( S \) in this case is not the physical excitation of a Higgs doublet. The couplings of the dilaton to the standard model fields can be derived by applying the theory of phenomenological lagrangians. For completeness we now briefly review the main aspects of the procedure, for a detailed derivation see for instance [9].

The conformal symmetry is an extension of the Poincaré algebra which includes scale invariance as well as the so called special conformal transformations. By definition, a CFT is necessarily scale invariant, but the converse need not be true. Let us first focus on the implications of scale invariance on the effective field theory. Given a local operator \( O(x) \) of scaling dimension \( \Delta \), the transformation under dilatations \( x^\mu \rightarrow e^\lambda x^\mu \), where \( \lambda \) is an arbitrary constant, is given by \( O(x) \rightarrow e^{\lambda \Delta} O(e^\lambda x) \). The theory is invariant under scaling if the lagrangian has dimension \( \Delta = 4 \). For \( \Delta \neq 4 \) the scale symmetry can be locally restored by introducing an appropriate GB field \( \sigma \), the dilaton. This field transforms in-homogeneously under a dilatation \( \sigma \rightarrow \sigma + f\lambda \) and thus signals the spontaneous breakdown of the \( \sigma = 0 \) vacuum. In complete analogy with the non-linear realization of internal symmetries, it is convenient to introduce a field transforming linearly under scale transformations:

\[
\chi(x) \equiv f e^{\sigma(x)/f} \rightarrow e^{\lambda} \chi(e^\lambda x) \tag{2.3}
\]

For practical purposes we can identify \( \chi = \chi - f = \sigma + O(\sigma^2) \) as our dynamical field rather than \( \sigma \) itself. Because \( \sigma \) and \( \chi \) generate the same one-particle states, the S-matrices of the two fields do coincide. A dimension-4 operator can finally be obtained by replacing the dimension \( \Delta \) operator \( O \) with \( O(\chi/f)^{4-\Delta} \).

At the classical level this is basically what the standard model Higgs does [1], [2]. To appreciate this, let us derive the leading couplings of the dilaton to the spin-1, spin-1/2, and the Goldstone \( SU(2) \) matrix \( U \) of the standard model \( SU(2) \times U(1) \rightarrow U(1) \) symmetry breaking pattern by assuming that the standard model fields have a classical scaling:

\[
\mathcal{L} = \frac{\nu^2}{4} Tr|D_\mu U|^2 \chi^2 f + m_\bar{\psi}_i U \bar{\psi}_i \chi f + \ldots \tag{2.4}
\]

\[
= \frac{1}{2} m_\bar{\psi}_i \bar{\psi}_i \left( 1 + \frac{\chi}{f} \right)^2 + m_\bar{\psi}_i \bar{\psi}_i \left( 1 + \frac{\chi}{f} \right) + \ldots
\]

The model (2.4) corresponds to the case

\[
a_1 = a_2 = \frac{v}{f}; \quad b = \frac{v}{f}; \quad b_{ij} = 0 \tag{2.5}
\]

and proves that, apart from a \( v/f \) rescaling, the couplings of the dilaton are formally those of a fundamental Higgs.
The $\Delta > 4$ operators are either irrelevant for our purposes or severely constrained by electro-weak precision data\(^2\). An exception appears to be the coupling to the unbroken gauge bosons. The latter arises as an effect of the scale-dependence of the renormalized coupling $g(\mu)$, i.e. the so called scale anomaly. The prescription for the non-linear realization of the scale symmetry illustrated above may be equivalently restated by saying that the renormalization scale $\mu$ must be compensated by an appropriate power of the dilaton field $\mu \to \mu \chi/f$. Hence, if we assume that the gauge symmetry of the standard model is part of the CFT, the coupling

$$
- \frac{1}{4g_{SM}^2} F_{\mu\nu}^2 = + \frac{\beta_{SM}}{2g_{SM}} F_{\mu\nu}^2 \frac{\chi}{f} + \ldots \quad (2.6)
$$

should be added to (2.4). (Notice that on the right hand side of the equation the vectors have been canonically normalized). This is equivalent to the introduction of a coefficient:

$$
c = \frac{8\pi^2 \beta_{SM} v}{g_{SM}^4} \quad (2.7)
$$

in (2.4). The phenomenology of (2.3) plus (2.7) has been studied by the authors of [2]. Because the coefficient (2.7) can potentially be larger than the top loop contribution, the anomalous term may well represent the most accessible signal of the SILD, especially in the ambiguous regime $v \sim f$ [2].

If the standard model gauge symmetry represents an explicit breaking of the scale invariance, then the dilaton couplings to the unbroken gauge group is no more predicted by the trace anomaly, and $c$ is no more given by (2.7). The latter coupling is more generally mediated by mixing between composites and standard model vectors, and by the integration of charged heavy composites. The resulting operator has been included in (2.4) with a model dependent $c = O(v/f)$ factor. Such a contribution can be in principle sizable if the scale invariant theory has a large number of degrees of freedom at the scale $g_{\rho} f$, and may compete with the top loop irrespective of whether the gauge symmetry is or is not embedded into the broken CFT.

In the next few sections we will also argue that the lagrangian (2.4) is not accurate if the Higgs sector has non-negligible anomalous dimensions, as it is expected from a strongly coupled sector, or if the CFT flavor structure is non-trivial, as it is expected in a generic model addressing the standard model fermions hierarchy. In particular we will see that the relations $a_i^2 = a_2^2$ and $b^{ij} = 0$ are not characteristic features of the light dilaton scenario.

We conclude this section by mentioning that a phenomenologically acceptable realization of both the SILH and SILD scenarios requires the existence of an explicit breaking

\(^2\)In this regard we should mention that a light dilaton is expected to alleviate the tension between the scale $m_\rho$ of new physics and the electro-weak precision measurements by screening the large corrections of the strong dynamics to the precision parameters. The effect is estimated to be small – see the quantitative discussion in [11] – so that the higher dimensional operators can be neglected in our discussion. Yet, it would be interesting to understand the relation between this screening effect and the slight improvement on the electro-weak S-parameter fit that a walking dynamics, as opposed to a generic strongly coupled theory, is expected to cause.
of the global symmetries of the strong dynamics, i.e. the generation of a mass $m$ for the light scalar. The source of explicit breaking is typically the standard model in the SILH scenario \([1]\), while, as already stressed in the introduction, it may be the standard model as well as the strong dynamics itself in the SILD model. A deformation of the CFT generally implies corrections to the dilaton vacuum expectation value with respect to the symmetric solution. The true vacuum can thus differ from $f$, causing a shift in the dilaton couplings. The relation $f < v$ may be in principle allowed in this case, but the validity of our phenomenological approach would come into question. We will comment more on the magnitude of the ratio $v/f$, and its physical meaning, in a following section. For the moment we emphasize that the phenomenology of the SILD approaches that of a SILH in the regime $v \sim f$, whereas that of an ordinary Higgsless theory for $v/f \ll 1$.

**Implications of the special conformal transformations:** The additional constraints coming from the special conformal transformations affect the derivative couplings only. Given a local operator $O$ of weight $\Delta$ and transforming under the Lorentz group as an irreducible representation with infinitesimal generators $S_{\mu\nu}$, we can define a conformally invariant derivative as \([9]\)

$$D_\mu O = [\partial_\mu + (iS_{\mu\nu} - \Delta \eta_{\mu\nu}) \partial_\nu \log \chi] O.$$  

(2.8)

An important point to be noticed is that the non-linear realization of the full conformal group does not require the introduction of additional Goldstone bosons: the broken special conformal transformations do not generate new physical excitations from the vacuum. Contenting ourselves with a leading order expansion in external momenta, the derivative interaction of the dilaton field required by (2.8) has physical impact only if $O$ is a scalar field with non-vanishing vacuum expectation value (see section 5). Being mainly interested in the dilaton couplings to spin 1 and $1/2$, we conclude that the leading low energy effective field theory is not sensitive to whether the UV completion possesses a full conformal invariance or just a scaling symmetry.

**The radion:** The model (2.5) plus (2.7) has a dual interpretation in terms of the RS1 scenario \([3]\), \([11]\) with a heavy Higgs and the standard model fields placed on the IR brane. Yet, as soon as the SM fields move away from the IR brane the relations (2.3) and (2.7) no longer hold.

The presence of the IR brane in the bulk AdS background can be interpreted as a spontaneous breaking of the conformal symmetry of the strongly coupled 4D dual theory, and the gauge/gravity interpretation of the dilaton is given in terms of the radion field \([3]\), \([11]\). The dynamical description of the radion field is encoded in the perturbed line element \([12]\)

$$ds^2 = e^{-2ky+Qe^{2ky}} \eta_{\mu\nu} dx^\mu dx^\nu - \left(1 - Qe^{2ky}\right)^2 dy^2$$  

(2.9)

$$= e^{-2kw} \eta_{\mu\nu} dx^\mu dx^\nu - dw^2 + \ldots$$

where $Q(x)$ represents the 4D field appearing in the effective field theory. The solution $Q = 0$ is the conventional Randall-Sundrum background.
In the second line of (2.9) we made a change of variable to the new coordinate \( w(y,x) \), defined by \( 2kw = 2ky - Qe^{2ky} \). In terms of it the metric can be expressed in a compact form where the dynamical field \( w \) determines the spacetime-dependent proper length of the extra dimension with classical vacuum \( y \). At the quantum level the extra dimension can thus be thought of as a physical degree of freedom, the dilaton of the dual theory. Notice in fact that the model preserves a dilatation symmetry \( x^\mu \rightarrow e^{\lambda} x^\mu \) at the perturbed level if the field \( w \) transforms inhomogeneously \( kw \rightarrow \lambda + kw \), in complete analogy with \( \sigma \) in (2.3).

The second line of (2.9) has the same form as the naive ansatz proposed originally by Randall and Sundrum [5], up to \( O(\partial_\mu w) \) terms having subleading impact on the 4D physics. This clarifies why the physics of the actual dynamical mode \( Q \) agrees (up to derivative couplings and hierarchically suppressed corrections) with the results found for the naive ansatz (see for instance [3],[13],[14],[15], and references therein), although the latter is not a dynamical mode. See the Appendix for a discussion of this point.

3. Light scalars and EWSB

Suppose the LHC detects a light CP-even scalar and no other states: can we tell whether this particle is really the physical excitation of a Higgs doublet?

Thanks to the equivalence theorem, at external momenta \( p^2 \gg m_\nu^2 \) the three Goldstone boson fields \( \pi^a \) belonging to the unitary matrix \( U = e^{i\pi^a \sigma^a/v} \), where \( \sigma^a \) are the Pauli matrices, can be identified with the longitudinally polarized vector bosons \( V_L^a \). Thus, we can directly probe the strong dynamics by focusing on the physics of the scalars \( \pi^a \) and \( S \), and rewriting the first line of (2.1) as

\[
\frac{v^2}{4} Tr \left( \partial_\mu U^\dagger \partial^\mu U \right) \left( 1 + 2a_1 \frac{S}{v} + a_2 \frac{S^2}{v^2} \right) .
\]  

(3.1)

The lagrangian (3.1) should be supplemented with a potential for the scalar \( S \); however, being mainly interested in the high energy behavior of the amplitudes, this term can be discarded.

Typical values of the parameters \( a_1 \) and \( a_2 \) for the SILH are given up to \( O(v^2/f^2) \) in (2.2), and for the SILD in (2.5), under the assumption of negligible anomalous dimensions. In both cases we see that the elastic scattering of longitudinally polarized vectors violates perturbative unitarity at the scale of compositeness of the Higgs sector:

\[
A(\pi^a \pi^a \rightarrow \pi^b \pi^b) = \frac{s}{v^2} (1 - a_1^2),
\]  

(3.2)

with \( s \) the Mandelstam center of mass energy and \( a \neq b \). These processes can be effectively used to discriminate between the two scenarios only if we are able to probe the heavy resonances, at scales of order \( g_\rho v \) and \( g_\rho f \) for the SILD and SILH respectively. Even if the heavy states are not detected, this channel will provide crucial informations regarding the strength of the Higgs sector, see [18] for a recent study.

\[^3\]This section has some overlaps with the analysis presented in [16] and [17].
Another interesting channel is the elastic scattering $2S \rightarrow 2\pi$, for which

$$A(SS \rightarrow \pi^a\pi^a) = \frac{8}{v^2}(a_1^2 - a_2^2).$$  \hspace{1cm} (3.3)

If we neglect the anomalous dimensions of the Higgs sector of the CFT, $a_1^2 = a_2^2$, and the latter amplitude vanishes for $S$ the dilaton, as in the standard model Higgs scenario; the observation of high energy $2S \rightarrow 2\pi$ events in this limit would be a characterizing feature of the SILH. More generally, if we allow for non-zero anomalous dimensions of the CFT Higgs sector (we will see in section 5 that this implies $a_1^2 \neq a_2^2$), then (3.3) represents an indisputable evidence in favor of the compositeness of the EWSB sector.

If we were able to isolate the strong dynamics from the explicit symmetry breaking, though, the Goldstone boson sectors of the SILH and SILD scenarios would look radically different. On the one hand, the SILH scenario, as the fundamental Higgs of the standard model, possesses an $O(4)$ symmetry rotating the 4 scalars $\pi^a, S$. At scales $f^2 \gg p^2 \gg m^2$ the Goldstones $\pi$ effectively recombine with the SILH in the full order parameter, i.e. the Higgs doublet, and the $O(4)$ symmetry manifests itself in relations between the $S, \pi$ scattering amplitudes

$$A(SS \rightarrow 2\pi^a) = A(2\pi^a \rightarrow 2\pi^b) \left(1 + O\left(\frac{v^2}{f^2}\right)\right),$$  \hspace{1cm} (3.4)

(compare (3.3) with (3.2) by making use of (2.2)), and in the suppression of $O(4)$ violating events as $S \rightarrow 2\pi$. On the other hand, no such symmetry is present in the SILD model or any model in which the light scalar $S$ is not part of the weak doublet. In the latter case, relations of the type (3.4) are not observed, whereas events like $\chi \rightarrow 2\pi$ are allowed by the strong dynamics and are therefore enhanced at high energies.

These observations reveal that a discrimination between the two models may be feasible in the limit in which the corrections in (3.4) are sufficiently small. For example, if the SILH had, say, a $v^2/f^2$ of the order of a few percent, while the SILD $v/f \sim 1$ at the ten percent level, then the two scalars would look very much like the fundamental Higgs boson for what concerns the couplings to the standard model fields, but they would present a radically different UV physics for what concerns the couplings to the unphysical Goldstone bosons. First, the SILH would satisfy the crucial relation (3.4), where the amplitudes are enhanced at high energy $A \sim s/f^2$ as opposed to a fundamental Higgs. This would lead to interesting observable consequences, like the sum rule eq.(4.19) in [1]. Second, observations of energy enhanced off-shell $O(4)$ violating processes $S \rightarrow 2V_L$ (in vector boson fusion as well as associated production events) would be a distinctive signature of the dilaton model. These processes may be for example tested as excesses in $t\bar{t} \rightarrow \chi \rightarrow V_LV_L'(\rightarrow 4l, ll\nu\nu)$ events at large invariant mass.

The deviations in the hard process $t\bar{t} \rightarrow S \rightarrow V_LV_L$ are potentially significant, and reach the 100% at moderately high energy scales $\sqrt{s} > 3m$ even for $O(10\%)$ departures from the standard model couplings. Given such a sensitivity in the hard process, the potentialities of these events in the extraction of the relevant Higgs couplings are remarkable. At

\footnote{The $O(v^2/f^2)$ corrections are next to leading order in the expansion adopted in [1], but are already present at the 2-derivative level.}
the LHC, even after a relatively low luminosity \( \sim 30 \text{ fb}^{-1} \), it would be possible to achieve accuracies as high as \( 5 - 10\% \), thus approaching the noise of a \( \sim 2\sigma \) deviation [20].

Processes involving only the light scalar \( S \) are generally more difficult to observe but are also useful to probe the strong dynamics. At low energies the scalar self-couplings are dominated by the potential interactions, and the extrapolations of trilinear \( (g_{SSS}) \) and quartic \( (g_{SSSS}) \) couplings would provide useful informations on the source of explicit breaking of the Goldstone symmetry [1], [2]. Large \( O(1) \) deviations from the standard model Higgs trilinear coupling may in principle be observed at the LHC in a limited mass range and after an integrated luminosity of about \( 300 \text{ fb}^{-1} \) is attained. An international linear collider (ILC) would certainly be more adequate, and may be able to test the triple coupling with a much larger accuracy, of order some \( 10\% \).

At sufficiently large energies these measurements are, however, substantially altered by the strong dynamics. In the absence of an explicit breaking of the global invariance the scalar self-couplings are derivatively induced. The SILH self-couplings are governed in the high energy regime by the \( O(4) \) invariance and, for instance:

\[
\mathcal{A}_{\text{SILH}}(2S \rightarrow 2S) = c_h \frac{s}{f^2},
\]

(3.5)

The SILD lagrangian is, up to \( O(p^4) \),

\[
\frac{(\partial_\mu \chi)^2}{2} + \frac{c_3}{(4\pi)^2} \frac{\partial_\mu \chi \partial_\nu \chi \partial^\mu \partial^\nu \chi}{\chi^3} + \frac{c_4}{(4\pi)^2} \frac{(\partial_\mu \chi \partial_\nu \chi)^2}{\chi^4} + \ldots,
\]

(3.6)

where for brevity we neglected operators that give vanishing contribution on shell. The coefficients \( c_i \) can be estimated using naive dimensional analysis\(^5\), and are expected to be of order unity in a strong dynamics. Both \( 2\chi \rightarrow 2\chi \) and the (“\( O(4) \) violating”) process \( 2\chi \rightarrow \chi \) start at \( O(p^4) \) and are suppressed in the regime of validity of our effective description.

In both the SILH and the SILD, the strong dynamics modifies significantly the scalar vertices at energies of the order of \( f \) or bigger. This observation is potentially relevant for the SILD physics at a linear collider, especially in the regime \( v \sim f \), as the \( O(p^4) \) terms would come into play at the scales \( (500 - 1000 \text{ GeV}) \) at which the ILC operates. As an instance, for \( v/f = 1 + 10\% \) a perturbative description of the \( V_LV_L \) scattering may be delayed up to \( \sim 4 \text{ TeV} \), whereas a strong \( \chi \) self-coupling would enter around 1 TeV.

### 4. Dilaton couplings to standard model fermions

Tree level flavor violations mediated by a composite Higgs are typically suppressed if the Higgs is a SILH [21]. We now show that this conclusion does not apply to many realistic realizations of the SILD scenario.

\(^5\)As opposed to the chiral lagrangian the interactions start at \( O(p^4) \) – although this property is a bit obscured in terms of the variable \( \sigma = f \log \chi/f \). In a cutoff regularization scheme, we can estimate the coefficients by requiring for example that the operators \( c_i \) (\( i = 3, 4 \)) induce small corrections on the kinetic term (this is not a relevant correction on shell, but our conclusions are completely general). This condition requires \( c_i < (4\pi/g_\rho)^4 \), where we cutoff the loop integrals at the physical scale \( g_\rho f \). For \( g_\rho \sim 4\pi \) our estimate \( c_i = O(1) \) agrees with the rules proposed in [4].
The dilaton couplings to the standard model fermions are expected not to mediate flavor changing neutral currents (FCNC) only if the standard model fermions couple to CFT operators with flavor universal scaling dimensions. This is certainly realized if the standard model fermions have all the same scaling representation. If this is the case, the most general Yukawa coupling for the dilaton can be written after EWSB as

\[ F_{ij} \bar{\psi}_i \psi_j \left( \frac{\chi}{f} \right)^\Delta = m_{ij} \bar{\psi}_i \psi_j \left( 1 + \Delta \frac{\chi}{f} + \ldots \right) \tag{4.1} \]

where \( \Delta \) is an arbitrary number, \( i, j \) are flavor indeces, and \( F_{ij} \) is a dimensionless scale invariant function, which we can identify with the Yukawa matrix. If the CFT operators do not have a flavor universal scaling, or if the CFT is broken, in the low energy dynamics it would emerge a non-scale invariant function of the dilaton \( F_{ij} = F_{ij}^{(0)} + F_{ij}^{(1)} \frac{\chi}{f} + \ldots \), with \( F_{ij}^{(a)} \)'s naturally of the same order. Equivalently, the dimension \( \Delta \) in (4.1) would become flavor-dependent and would generally lead to flavor violating events.

It is instructive to show how these conclusions arise by looking at specific UV realizations. One can identify two distinct ways to generate the standard model flavor structure. The first is the one implemented in composite Higgs models and (extended) technicolor models, and contains the minimal flavor violation (MFV) class. This is based on the coupling

\[ L_1 = y_{ij,a} \bar{\psi}_i \psi_j O_a, \tag{4.2} \]

where the CFT operators \( O_a \) have generally different scaling dimensions \( \Delta_a \). After EWSB the low energy effective theory is obtained by integrating out the energetic fluctuations of the CFT and reads

\[ L_1 = y_{ij,a} \bar{\psi}_i \psi_j \langle O_a \rangle \left( \frac{\chi}{f} \right)^{\Delta_a} + \frac{y^2}{m^2_\rho} \bar{\psi} \psi \bar{\psi} \psi + \ldots \tag{4.3} \]

\[ = m_{ij} \bar{\psi}_i \psi_j \left( 1 + \Delta_{ij} \frac{\chi}{f} + \ldots \right) + \frac{y^2}{m^2_\rho} \bar{\psi} \psi \bar{\psi} \psi + \ldots \]

The coefficients of the four fermion operator is a symbolic representation for the sum \( y^2/m^2_\rho = y_{ij,a} D_{ab} y_{kl,b} \), where the \( D \) matrix originates from the propagator of the operators \( O_a \). It is important to emphasize that the FCNC couplings are there even if the standard model fermions are part of the CFT, for example if \( \Delta_a = 4 - \Delta_i - \Delta_j \). For a universal dimension \( \Delta_a = \Delta \) we recover (4.1), that is \( \Delta_{ij} = \Delta \).

The second class approaches the flavor problem by employing a seesaw type mechanism. This is the class usually implemented in the Randall-Sundrum scenarios with standard model fermions in the bulk [22], see [23] for the dual interpretation. The bare lagrangian now reads:

\[ L_2 = \lambda^i_L \psi_L O^i_R + \lambda^j_R \psi_R O^j_L \tag{4.4} \]

and leads to an effective field theory of the form

\[ L_2 = \lambda^i_L \chi^j_R \langle D^{ab} \rangle \psi^i_L \psi^j_R \left( \frac{\chi}{f} \right)^{\Delta_2 + \Delta_4 - 4} + \frac{\lambda^4}{m^2_\rho} \bar{\psi} \psi \bar{\psi} \psi + \ldots \tag{4.5} \]
\[
\begin{align*}
&= \lambda_L^{ia} \lambda_R^{jb} \langle D^{ab} c \rangle \bar{\psi}_L^i \psi_R^j \left( 1 + (\Delta_R^a + \Delta_L^b - 4) \frac{\chi}{f} + \ldots \right) + \frac{\lambda^4}{m^2_F} \bar{\psi}_F \psi_F + \ldots \\
&= m_{ij} \bar{\psi}_i \psi_j \left( 1 + \Delta_{ij} \frac{\chi}{f} + \ldots \right) + \frac{\lambda^4}{m^2_F} \bar{\psi}_F \psi_F + \ldots
\end{align*}
\]

The appearance of the power of the dilaton follows from the expansion of the operator \( \int d^4 y \mathcal{O}_R^K(x) \mathcal{O}_L^K(y) \) in the process of CFT integration. Again, if the dimensions \( \Delta_{L,R} \) were not to depend on the family label \( a,b \), the SILD coupling would be flavor diagonal \( \Delta_{ij} = \Delta_R + \Delta_L - 4 \). For a derivation of eq. (4.3) in the Randall-Sundrum model see [24].

If the CFT is flavor anarchic we expect \( \Delta_{ij} \) in (4.3) and (4.3) to be a generic matrix with no hierarchical relations between the components. We will assume this is the case in the following. The leading FCNC effects can be thus described by the following lagrangian:

\[
\bar{\psi}_i \psi_j \left[ m_i \delta_{ij} \left( 1 + b \frac{\chi}{f} \right) + \sqrt{m_i m_j} b_{ij} \frac{\chi}{f} \right] + C_{ijkl} \bar{\psi}_i \psi_j \bar{\psi}_k \psi_l,
\]

where the factor \( \sqrt{m_i m_j} \) in front of the flavor violating term characterizes both classes of models illustrated above.

The flavor mixing term \( b_{ij} \) is proportional to the anomalous dimensions of the CFT operators, and it is expected to be an \( O(1) \) parameter in a strongly coupled dynamics. This parameter is severely constrained by FCNC bounds. Tree level exchanges of the dilaton lead to 4 fermions interactions with coefficients \( C \sim (mb^j/f)^2/m_\chi^2 \). The strongest bound for a generic left-right mixing and CP violating operator comes from \( K\bar{K} \) mixing. The UTFit analysis [23] reports the bound \( Im(C) \lesssim (10^5 \text{ TeV})^{-2} \) (we are neglecting subleading logarithmic running effects), which translates into

\[
|b_{ij}| \frac{\chi}{f} \lesssim 10^{-2} \left( \frac{m_\chi}{100 \text{ GeV}} \right).
\]

If \( v < f \), the bound (4.7) may be saturated for \( b_{ij} = O(1) \) and a light dilaton. In this case the dominant decay channel is the \( \chi \to b\bar{b} \), and flavor violating processes may have relatively large branching ratios \( BR(\chi \to b\bar{s}) \sim 5\% \). In the regime \( v \sim f \), the \( b_{ij} \) should be neglected compared to the flavor diagonal couplings. More generally, the flavor violating coupling becomes phenomenologically relevant at relatively large dilaton masses. In the latter case, however, the dominant decay mode is expected to be into massive vectors, as for a fundamental Higgs, and branching ratios into fermions become much less accessible. One can estimate flavor violating events to be down by an order \( BR(\chi \to t\bar{c}) \sim 10^{-4} \) in this limit. Production channels like \( t \to \chi c \) may still play a role [24].

The 4-fermions contact terms in (4.6) have coefficients of the order \( C_{ijkl} \sim \sqrt{y_i y_j y_k y_l} / m_\rho^2 \), where \( m_i \sim y_i v \) is a standard model fermion mass. The phenomenological bounds on these latter can be obtained from the bounds on the dilaton couplings by noting the correspondence \( m_\rho \sim f m_\chi / (|b_{ij}| v) > 10 \text{ TeV} \). In a generic theory with heavy composites around the scale \( 4\pi v \), this bound cannot be satisfied and one should find a way to push the flavor violating effects at higher scales. This can be done by adjusting the theory in such a way that the only physics at the scale \( g_\rho v \sim 4\pi v \) is the strong Higgs sector, for which FCNC effects are negligible. If our model can account for a splitting between the CFT breaking scale
and the electro-weak vacuum, the 4-fermion coupling \( C_{ijkl} \) of (4.6) would be suppressed by a high scale \( m_\rho = g_\rho f \) and the constraint (4.7) would translate into \( f \sim \text{TeV} \) for a maximally strong dynamics \( (g_\rho \sim 4\pi) \).

To conclude, models of type (4.2) are more adequate to embed MFV, while those of type (4.4) seem more appropriate to explain the fermion hierarchy. From the above qualitative estimates we see that \( v/f \lesssim \text{few} \times 0.1 \) in generic models, while \( v \sim f \) is compatible with FCNC observations only if MFV is at work, in which case \( b_{ij} = 0 \) by definition.

5. Higgs-dilaton mixing

As far as scale invariance is not explicitly broken, a hierarchical separation \( v \ll f \) is unnatural in a CFT. Under this assumption, the dilaton is an exact Nambu-Goldstone boson and parameterizes a flat direction. Its effective lagrangian, obtained after integrating out all the other fields, contains no potential (see eq. (3.6)): the only CFT invariant candidate, \( V = a\chi^4 \), would imply an unbroken vacuum \( \langle \chi \rangle = 0 \), that would contradict our original hypothesis \( \langle \chi \rangle = f \). The integration of light particles at the scale \( v \) must therefore exactly compensate the contribution of the heavy composites at the scale \( f \). Because by dimensional analysis the former give \( a_{\text{light}} \sim (v/f)^\Delta \), the smaller the ratio \( v/f \), the larger the fine-tuning required to accommodate a vanishing potential. This explains why, if the full standard model is embedded into a spontaneously broken CFT, the natural cut-off scale for new physics cannot be far from the weak scale. The same conclusion is found in the context of the Randall-Sundrum scenario if the standard model is placed on the IR brane, where the cut-off suppressing higher dimensional operators is naturally at the weak scale irrespective of the mass of the Kaluza-Klein modes.

If the weak scale is somewhat smaller than \( f \), and if the CFT is not badly broken, it makes sense to consider an effective theory for the strong EWSB sector of the SILD scenario and the dilaton itself. We focus on a simplified theory in which the Higgs sector is described by an interpolating Higgs doublet and for simplicity we further assume that the anomalous dimensions can be algebraically summed\(^6\). Our results are completely general, though.

Under our simplifying assumptions the most general action for the strongly coupled Higgs is constructed out of the conformally covariant derivative (2.8):

\[
\left( D_\mu - \Delta \frac{\partial_\mu \chi}{\chi} \right) H, \tag{5.1}
\]

where \( D_\mu = \partial_\mu + iA_\mu \) is the usual gauge covariant derivative, and the potential

\[
V(\chi, H) = \chi^4 \hat{V} \left( \frac{H^\dagger H}{\chi^{2\Delta}} \right). \tag{5.2}
\]

Both kinetic and potential terms induce a mixing between the Higgs after EWSB. For the case of a dimension \( \Delta = 1 \) Higgs, the kinetic mixing coincides with the one induced

\(^6\)This can be made rigorous for appropriate local operators at leading \( 1/N \) order and for a certain class of SUSY theories with R-symmetry.
in warped models by the non-minimal gravitational coupling \( \xi R H^\dagger H \) with \( \xi = 1/6 \) \cite{26}, see \cite{14} for a detailed analysis. A deviation from the conformal factor \( \xi = 1/6 \) accounts for the breaking of special conformal symmetries, and it amounts to including an additional
\[
(1 - 6 \xi) H^\dagger D_\mu H \partial^\mu \chi/\chi + h.c.
\]
operator in our 4D language.

In the conformal limit, the system can be easily diagonalized by defining the scale-invariant field \( \hat{H} = H(f/\chi)^\Delta \). The mass eigenstates are then the physical component of the scale-invariant field, \( \hat{H}^\dagger = (0, v + \hat{h}) \), and the dilaton \( \bar{\chi} = \chi - f \). Notice that, in order to be compatible with the background solution \( \chi = f \), \( |H|^2 = v^2 \), the potential must satisfy \( \hat{V} = \hat{V}' = 0 \) on the vacuum. These conditions automatically ensure that no mass term is generated for the dilaton.

The leading coupling between the dilaton and the physical Higgs \( \hat{h} \) is dictated by scale invariance:
\[
\frac{m^2}{2} \hat{h}^2 \left( 1 + 4 \frac{\bar{\chi}}{f} + \ldots \right),
\]
where \( m^2 = \hat{V}'' v^2 \). The mass of the physical composite Higgs \( \hat{h} \) is of order \( m_\rho \lesssim 4\pi v \) in a strong dynamics and its fluctuation may be neglected at scales \( E < m_\rho \). Because of the mixing \( H - \chi \), all the couplings of the gauge eigenvector \( H \) induce vertices for the SILD. In terms of the mass eigenstates, the electromagnetic singlet component of the doublet \( H \) reads:
\[
h = (\hat{h} + v) \left( 1 + \Delta \frac{\bar{\chi}}{f} + \frac{1}{2} \Delta(\Delta - 1) \frac{\bar{\chi}^2}{f^2} + O(\bar{\chi}^3) \right).
\]
(5.3)

If we integrate out the heavy state \( \hat{h} \), we see that the dilaton couples as a Higgs up to a universal \( \Delta v/f \) scaling and modulo universality violations induced by the anomalous dimension \( \Delta - 1 \) of the strong Higgs. In particular, we see that the relation \( a_1^2 = a_2^2 \) in (2.1) characterizes the SILD scenario in the limit of small anomalous dimensions for the Higgs sector.

In principle, one may expect a large dimension \( \Delta \) to compensate the suppression \( v/f \) so that \( \Delta v/f \sim 1 \); however, the dimension of the Higgs cannot be arbitrary large if we require our model to be self-consistent. For example, assuming that the top Yukawa coupling remains perturbative up to the naive dimensional analysis scale \( m_\rho \sim 4\pi v \) requires \( \Delta < 2 \) \cite{27}.

Once sources of explicit symmetry breaking of the CFT are taken into account, a non-zero dilaton mass \( m \) is generated and the decoupling procedure described above is no more valid. As far as \( m^2 \ll m_\rho^2 \) applies, we expect our results to be accurate.

**The dilaton as the Higgs:** The quantity \( v/f \) is an actual measure of the Higgs contribution to the spontaneous CFT breaking, and parameterizes the amount of mixing \( h - \bar{\chi} \). This is explicit in (5.3) and it can be equivalently derived as a general consequence of the algebra of the generators, see eq. (19) in \cite{28}. For \( v \sim f \) the dilaton-Higgs mixing is maximal and in the extreme scenario \( v = f \) the Higgs itself is the dilaton.

The scenario in which the Higgs and the dilaton coincide naturally account for a parametric separation between the Higgs mass and the naive dimensional analysis scale \( 4\pi v \).
However, both electro-weak precision parameters and flavor measurements are potentially crucial tests for this class. It is therefore of some interest to study the viability of this scenario by considering explicit and computable models in which this feature is realized. A tractable candidate may be found by considering the Randall-Sundrum scenario.

In principle, in order to realize our program we should be able to control the physics responsible for the IR brane generation. This physics is related to operators with large dimensions – \( \Delta = O(N) \) – which are dual to heavy stringy modes that we cannot control on the 5D side \[3\]. Nevertheless, consistently with our leading \( 1/N \) approximation we find that warped higgsless theories \[29\] behave as models in which the Higgs field can be identified with the dilaton. To see this, we first introduce the conformal coordinate \( z = \log y \), in terms of which the AdS geometry can be written as

\[
ds^2 = \frac{1}{z^2} \left( \eta_{\mu\nu} dx^\mu dx^\nu - dz^2 \right),
\]

where the curvature has been normalized to 1 for simplicity. In higgsless models the Higgs sector can be idealized as a 5D field \( \phi \) infinitely peaked on the IR brane, which we assume to be placed at some finite \( z = z_{IR} \). Let us see how this translates in the dual strong dynamics.

Any bulk field \( \phi \) on the gravity side is believed to be dual to some operator \( O \) on the gauge side. The 4D Fourier transform of a scalar field \( \phi \) of 5D mass \( m^2 > 0 \) on the AdS geometry is given by \( \phi(p,z) = c(p)z^2(J_{\nu}(pz) + \beta_\nu Y_{\nu}(pz)) \), where \( \nu^2 = 4 + m^2 \) and \( \beta_\nu \) is determined by the boundary conditions. In the UV region \( z \to 0 \) we have \( \phi \to Az^{\Delta} + \varphi_0 z^{4-\Delta} \), where \( \Delta = 2 + \nu \) represents the scaling dimension of the dual scalar operator \( O \) and corresponds to the positive root of \( m^2 = \Delta(\Delta - 4) \). The AdS/CFT correspondence instructs us to identify the parameters \( A \) and \( \varphi_0 \) as the vacuum expectation value and the source of the operator \( O \), respectively \[30\].

By switching off the source \( \varphi_0 \) and taking the limit of very large dimension \( \Delta \) (i.e. large 5D mass, \( m^2 \to \infty \)) we find \( \phi(p,z) = c(p)z^2 J_{\nu}(pz) \to Az^{\nu+2} = A(p)z^{\Delta} \). Using the identification introduced in \[30\] we finally write:

\[
\phi = \frac{\langle O \rangle}{2\Delta - 4} z^\Delta.
\]

For a translational invariant vacuum, no propagating mode exists in the \( m \to \infty \) limit and the bulk field reduces to a non-dynamical background. Now, if the operator \( O \) acquires a vacuum expectation value \( \langle O \rangle = cf^\Delta \), then the corresponding 5D profile in the limit \( \Delta \to \infty \) can be approximated by a step function:

\[
\phi \to \begin{cases} 
0 & \text{for } zf < 1, \\
\infty & \text{for } zf > 1.
\end{cases}
\]

This is exactly the physics at the origin of the IR brane in the RS model, provided we identify \( z_{IR} = 1/f \) \[3\].

The argument can be generalized to the case in which a number of operators \( O_i \) contribute to the spontaneous breaking of the CFT by an amount \( \langle O_i \rangle = c_i f^{\Delta_i} \). As far as
their dimension is very large $\Delta_i = O(N)$, the excitations of these operators are not captured by our leading $1/N$ approximations, and the information contained in the coefficients $c_i$ is lost. In this limit the effective theory is equivalent to having a single operator $O_1$ with vacuum $f$. Any field $\phi$ exactly localized on the IR brane can be interpreted as the dual description of $O_1$. In particular, without loss of generality, we can identify the Higgs as our candidate and conclude that the above model is equivalent to a higgsless scenario of EWSB with $v = f$, as anticipated. This model is characterized by standard model fields masses of order $g_{SM}f$ and CFT composites at the scale $g_\rho f$. The well known tension between perturbativity of the strong sector and the fit with precision data – in particular the flavor constraints outlined in the previous section and the electro-weak precision parameter – is manifest. The question whether subleading corrections in the $1/N$ expansion can alleviate this tension while keeping a relatively light dilaton remains open.

6. Discussion and conclusions

If the forthcoming experiments discover a light and chargeless CP-even scalar but no additional heavy states, it would be a priority to understand whether the detected spin-0 particle is the physical excitation of a light Higgs doublet, being it composite or fundamental, or not. In order to answer this question we focused on two scenarios of strong EWSB. The first includes the fundamental Higgs doublet model and it is known as the SILH class [1]. The second scenario describes a broken CFT with an emerging light dilaton [2], the strongly-interacting light dilaton (SILD) class.

The phenomenology of the latter has been extensively studied in the context of the Randall-Sundrum scenario, where it is generally accompanied by a light composite Higgs (for a recent study see [3]). We find it useful to attack the problem from the broader perspective of 4D effective field theories.

Our primary aim was the identification of model-independent signals characterizing the dilaton with respect to the SILH scenario, questions regarding the compatibility of the models with the electro-weak precision parameters (in particular $\hat{S}$) were not considered. We used a phenomenological approach and characterized the models by using three parameters: the scale at which the (approximate) symmetry of the strong dynamics is spontaneously broken ($f$), the strength of the interactions among the resonances ($g_\rho$), and the mass of the light scalar ($m$). An explicit breaking of the global symmetries of the strong sector is necessary to generate the mass $m$. The physically relevant aspect is how much the standard model is involved in the breaking.

Given a competitive low energy phenomenology for the light scalars (in particular, similar couplings to the standard model fermions), the two scenarios are in principle distinguishable at high energies, as the heavy states are expected at parametrically different scales, $g_\rho v \lesssim 4\pi v$ for the SILD and $g_\rho f \lesssim 4\pi f$ for the SILH. In the low energy regime, however, a discrimination between a dilaton and a Higgs field, either fundamental or composite, is possible in specific regions of the parameter space only as the result of combined fits and sufficiently high statistics. A detailed estimate of the actual experimental significance of our conclusions is a model-dependent issue and it has not been addressed here.
At the LHC, the measurements of rates times branching ratios ($\sigma \times BR$) in many channels is expected to provide useful pieces of information regarding the couplings of the light scalar for a sufficiently high integrated luminosity. Under some reasonable assumptions on the new physics model$^7$ and for an intermediate scalar mass $m$, it should be possible to probe relative deviations $\delta$ from the standard model Higgs couplings in the range $|\delta| > 10\%$, see for instance [20] and [32]. With these accuracies, if the SILD has an $f$ within $10\%$ or less of $v$, the LHC will not be able to tell the difference between the two scenarios [8]. A more promising environment is thus provided by a linear collider like the ILC, at which a precision of the percent level in the extraction of $\delta$ from $\sigma \times BR$ may be achieved.

Strong deviations from the SILH physics are likely to be observed in the dilaton couplings to gluons or photons $\chi \rightarrow 2g, 2\gamma$ mediated by the conformal anomaly [26, 2]. If the standard model gauge symmetry is an explicit breaking of the CFT the situation is a bit more subtle because the coefficients entering in such processes are very much model-dependent. Nevertheless, we emphasize that observable departures in the production/decay into massless vectors are expected to be relevant in theories with a large number of fundamental constituents, which is quite a natural framework for obtaining a light dilaton.

A characteristic feature of the SILD – or any model in which the scalar is not part of a weak doublet – would be the observation of energy enhanced $O(4)$-violating processes in $V_LV_L \rightarrow \chi$ events (vector boson fusion as well as associated production) at high virtuality, for instance in the promising channel $gg \rightarrow \chi \rightarrow V_LV'_L \rightarrow 4l, l^\pm \nu \nu$. To reach sufficiently high energies, $p^2 \gg m^2$, a hadron collider seems more adequate, but the precision may not be sufficient due to QCD uncertainties and the large background.

We found that the dilaton generally mediates tree level flavor violating processes in the effective theory. In contrast, a SILH does not if the alignment mechanism described in [21] applies. The flavor violating parameters are directly related to the flavor non-universality of the CFT operators that couple to the standard model fermions, and are present irrespective of whether the standard model is or is not embedded into the CFT. Potentially testable branching ratios for flavor violating processes like $\chi \rightarrow b\bar{s}$ are natural features of a SILD. Other events characterizing the dilaton model and involving standard model fermions were considered in [28].

The CFT necessarily introduces 4-fermions contact interactions in the low energy effective theory, as well, and the current bounds on flavor violation can be translated into a bound on the ratio $v/f$. Except for somewhat ad-hoc scenarios in which the minimal flavor violation paradigm is at work, for which no real explanation of the standard model fermion hierarchy is given, we found that $v/f \lesssim$ few $0.1$ is a realistic estimate. Additional phenomenological bounds on this ratio come from LEP, and apply to a light dilaton $m < 110$ GeV. These are discussed in [2].

In the preferable case $v < f$, a somewhat lighter Higgs sector compared to the scale $m_{\rho} = g_{\rho}f$ is required. We analyzed the implications of this assumption on the Higgs-

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$^7$One of the assumptions that characterize the analysis of [20] and [32] is that the couplings of the light scalar to two $W$'s or $Z$'s are bounded by above by those of the standard model Higgs. These assumptions are satisfied in both the dilaton scenario and the SILH (see [23]), as far as our perturbative description holds.
dilaton mixing and studied the decoupling of the physical excitations in the limit in which the explicit CFT breaking source is negligible. It is intriguing to further consider the possibility of a large Higgs-dilaton mixing, $v \sim f$, especially in light of a possible realization in a walking dynamics context [7]. We showed that the existing warped higgsless models of EWSB [29] are tractable incarnations of this idea.

In summary, if no significant deviations in the magnitude of the standard model couplings to the light scalar (especially in the gauge sector) are observed, no $O(4)$-violating process is detected, as well as no tree level FCNC exchange is measured, then it would be fair to say that the observed scalar is a (composite) Higgs.

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A. Parameterizations of the radion field

In this Appendix we review the derivation of the radion profile in the Randall and Sundrum (RS) model and clarify its relations with the naive ansatz proposed by RS.

The RS model is constructed on a 5D space-time ranging from $y_+ \leq y \leq y_-$ and possessing a $Z_2$ reflection symmetry around $y_+$ and $y_-$. Two branes are placed on the fixed points. The action of the model is

$$\mathcal{L} = 2 \int_{y_+}^{y_-} dy \sqrt{-g} \left[ -M^3 R - V \right] + \sqrt{-g_+} \left[ -V_+ \right] + \sqrt{-g_-} \left[ -V_- \right] + \ldots \quad (A.1)$$

The dots stand for possible additional higher order terms and – more importantly – for the Gibbons-Hawking term, necessary to properly define the field theory on a space-time with boundaries. $g_+, -$ denote the determinants of the induced metrics at the points $y = y_+, -$ respectively, whereas the factor of 2 in front of the integral over the extra dimension accounts for the orbifold symmetry.

We focus on background solutions that preserve the 4D Poincare invariance. The most general line element with these properties reads

$$ds^2 = e^{2A} \eta_{\mu\nu} dx^\mu dx^\nu - dy^2, \quad (A.2)$$

where $A = A(y)$. The equations of motion for a line element of the form (A.2) reduce to the following independent equations:

$$A'^2 = -\frac{V}{12M^3} \quad A'' = \frac{V_+}{6M^3} \delta(y - y_+) - \frac{V_-}{6M^3} \delta(y - y_-), \quad (A.3)$$
where a prime means derivative with respect to the fifth coordinate $y$. The RS background is defined by tuning the cosmological constants in such a way that $V = -12k^2M^3$, $V_+ = -V_- = 12kM^3$. Under these hypothesis our background solution is defined by $A = -k|y|$.

The gravitational perturbations around the RS background describe spin-2 and spin-0 fields, the spin-1 components being killed by the orbifold projection. The spin-2 fields form a tower of massive states with a single massless mode following from the preserved local 4D diffeomorphism invariance. The latter describes the fluctuations of the 4D background $\eta_{\mu\nu}$ and is thus properly identified with the 4D graviton. The spin-0 field, called radion, is a classically massless mode describing the fluctuations of the interbrane distance. Its mass gets non-trivial corrections if we go beyond Einstein gravity, namely the interbrane distance is generally stabilized by the introduction of higher dimensional operators. In terms of the dual variables $\tilde{g}$, the dilaton gets a mass once subleading $(1/N)$ corrections are taken into account.

We are interested in studying the scalar degree of freedom. By an appropriate choice of gauge, a general line element can be put in the form

$$ds^2 = W^2\hat{g}_{\mu\nu}dx^\mu dx^\nu - Y^2dy^2,$$

where $W$ and $Y$ are arbitrary scalar functions of the 5D coordinates describing the dilaton fluctuations around the vacuum $\langle W \rangle = e^{-k|y|}$ and $\langle Y \rangle = 1$. The traceless and transverse components of $\hat{g}_{\mu\nu} - \eta_{\mu\nu}$ describe the spin-2 modes. The standard procedure used to identify the dynamical modes consists in solving the linearized equations of motion for the perturbations $W - \langle W \rangle$, $Y - \langle Y \rangle$, and $\hat{g}_{\mu\nu} - \eta_{\mu\nu}$. The variation of the $\mu\nu$ part leads to the nontrivial condition

$$\frac{1}{2} \frac{\partial_\mu \partial_\nu Y^2}{Y^2} - \frac{\partial_\mu \partial_\nu W^2}{W^2} + \hat{g}_{\mu\nu}(\ldots) = 0.$$  (A.5)

The $\hat{g}_{\mu\nu}$ term will not be reproduced here for brevity. It suffices to say that cancellation of this term, as well as of the $\mu 5$ and 55 components of the linearized equations, is ensured by imposing the constraint

$$W' = -kWY,$$  (A.6)

and by requiring that the linearized scalar 4D fields satisfy the massless Klein-Gordon equation. One can show that the two conditions (A.5) and (A.6) are sufficient to derive the radion profile (2.9). The naive ansatz for the radion proposed by Randall and Sundrum, obtained by identifying the fifth coordinate $y$ with the radion, satisfies (A.6) but not (A.5), and thus it does not describe a dynamical mode.

The constraint (A.6) will play the major role on our discussions to follow. To understand its physical meaning we focus on the (non-linear) physics of the radion and the 4D graviton. Without loss of generality we can assume $\hat{g}_{\mu\nu} = \hat{g}_{\mu\nu}(x)$ and write down the 5D lagrangian for these fields. Using the line element (A.4) we find:

$$-\sqrt{-g}R = \sqrt{-\hat{g}} \left[ -6\hat{g}^{\mu\nu} \partial_\mu (WY) \partial_\nu W - W^2Y\hat{R} + 12\frac{W^2}{Y}W'^2 + (\ldots)' + \partial_\mu B^\mu \right],$$  (A.7)
where $\hat{R}$ is the 4D scalar curvature constructed out of $\hat{g}_{\mu\nu}$. The last term in (A.7) is a 4D boundary term, it has no perturbative implications and can be safely discarded in our treatment. The total derivative in $y$ exactly cancels the Hawking-Gibbons term, and thus can be ignored as well. The first and second terms describe at the quadratic level the radion and the graviton kinetic terms, respectively, and at the non-linear order the interactions of the two fields. The third term contains no 4D derivatives of the fields $W, Y$ and would represent a potential energy for the radion. This latter term must vanish once the contribution of the cosmological constants are included, and this gives (A.6).

By solely imposing (A.6) we thus find the form of the 4D effective action for the radion and the graviton:

$$
\mathcal{L}_{eff} = \sqrt{-\hat{g}} \frac{M^3}{k} \left[ 6(\partial W)^2 + W^2 \hat{R} \right]_{y^+} - y^+, \quad (A.8)
$$

with $[f(y)]_{y^-} = f(y^-) - f(y^+)$. The action (A.8) has been used by Rattazzi and Zaffaroni to spell out the dual interpretation of the physics of the IR brane $\mathcal{B}$. These authors studied the effective theory (A.8) in the absence of the UV brane regulator ($y^+ \to -\infty$) and argued that the presence of the IR brane in the AdS background can be interpreted as a spontaneous breaking of the conformal symmetry of the strongly coupled 4D dual theory. In this limit the physical radion decouples from the UV (because of its composite nature) and can be identified with the dilaton $W(y^-)$.

From our derivation, it is clear that the relevant degree of freedom turns out to be $W$, rather than the actual dynamical 4D field $Q$ appearing in (2.9). Any parametrization of the radion field having a small overlap with the UV brane ($\partial W(y^+) \ll \partial W(y^-)$) would be a suitable candidate to effectively describe the radion, up to hierarchically suppressed corrections. Furthermore, any parametrization of the radion would lead to the same non-derivative couplings for the 4D field. In particular, any parametrization would couple with IR strength to the IR brane, and would develop the same mass as the actual radion once explicit CFT breaking sources are added. This follows again from the condition (A.6).

Indeed, by the change of variables $y \to w(x^\mu, y)$, where

$$
w = -\frac{1}{k} \log W \quad (A.9)
$$

the line element (A.2) becomes

$$
ds^2 = e^{-2kw} \hat{g}_{\mu\nu} dx^\mu dx^\nu - dw^2 + O(\partial_\mu w). \quad (A.10)
$$

Hence, up to subleading derivative couplings and hierarchically suppressed corrections, the radion can be parametrized by any field $W$ satisfying (A.6). Under these approximations, the physics of the naive ansatz and that of the dynamical mode $Q$ do coincide. Non-negligible corrections are expected to come into play only for small hierarchies, in which case the interpretation of the radion as a dilaton would fail.

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