The Dilaton-like Higgs boson with scalar singlet dark matter

Robyn Campbell, Stephen Godfrey, and Alejandro de la Puente
Ottawa-Carleton Institute for Physics, Carleton University, 1125 Colonel By Drive, Ottawa, Ontario K1S 5B6, Canada
(Dated: July 7, 2016)

We study a model with a Higgs-like dilaton and a Standard Model gauge-singlet scalar dark matter candidate. We begin by updating the status of identifying the observed 125 GeV Higgs-like boson with the pseudo Nambu-Goldstone boson that arises from the spontaneous breaking of scale invariance using recent Higgs boson signal strength measurements by the ATLAS and CMS collaborations. We then constrain the extended model with recent constraints on the Higgs invisible width, the observed dark matter relic abundance and the latest dark matter direct detection limits. We found that the magnitude of the dilaton-$\gamma\gamma$ and dilaton-glue-glue coupling is constrained to be close to the standard model values. The mass of the dark matter candidate is constrained to be greater than half the dilaton mass by relic abundance limits and Higgs invisible width limits. Dark matter direct detection limits allow only small mass regions which will be further constrained by upcoming DEAP measurements.

I. INTRODUCTION

Little is known about what lies behind the mechanism of Electroweak Symmetry Breaking (EWSB). Within the Standard Model EWSB is realized with an elementary scalar field and a negative mass term in the scalar potential. The negative mass term induces an instability that causes the Higgs field to condense leading to spontaneous symmetry breaking of the SU(3)$_c$×SU(2)$_W$×U(1)$_Y$ gauge symmetry down to SU(3)$_c$×U(1)$_{em}$. The scale at which the symmetry breaks and the spectrum depends strongly on the size of the Higgs mass parameter. However, this parameter is strongly sensitive to high energy scales; if the Standard Model is the correct theory up to the Planck Scale, the Higgs mass parameter would receive very large quantum corrections. To keep the Higgs mass parameter close to the electroweak scale, one will have to fine-tune the bare Higgs mass parameter order by order in perturbation theory for the physical Higgs mass to be 125 GeV [1,3]. This naturalness problem has led many to believe that the absence of very large quantum corrections can only be the result of additional symmetries protecting the Higgs mass parameter. One example is supersymmetry, where quadratically divergent contributions to the Higgs mass from top quark loops are canceled by loops of scalar particles with the same gauge quantum numbers as the top quark (See [4] and references therein). However, supersymmetry cannot be an exact symmetry and thus a natural resolution of the naturalness problem would require a supersymmetry breaking scale that is not too large [5].

Another possibility pursued in a number of papers is that the properties of the Standard Model (SM) Higgs boson are mainly fixed by the approximate conformal invariance in the limit when the Higgs potential is turned off. In this case the Higgs vacuum expectation value (vev) will spontaneously break the approximate conformal symmetry as well as the EWS. In this scenario the Higgs is identified with the dilaton which is the Goldstone boson associated with the spontaneously broken conformal symmetry and the Higgs couplings are mainly dictated by this conformal invariance with four parameters for this sector of the theory, the symmetry breaking scale $f$, the couplings to photons and gluons, $c_{\gamma\gamma}$ and $c_{gg}$, and the dilaton self coupling $\lambda$.

A second important problem in particle physics is the nature of dark matter (DM). Any complete model should accommodate DM. Thus, it is possible that these two problems, EWSB and DM, are intertwined and the incorporation of a solution for one will impact the understanding of the other. We thus extend the model of a Higgs-like dilaton with the simplest DM candidate, a scalar singlet which introduces two additional parameters to our model, the DM mass $m_S$ and the DM self coupling $\lambda_S$, although this last parameter remains unconstrained by current experimental measurements.

In the following section we review the motivation for the Higgs-like dilaton and write down the Lagrangian for our model. In Section III we study the constraints that the Higgs properties puts on the parameters of the model using a Markov Chain Monte Carlo fit using the most recent ATLAS CMS combined results for the $\sqrt{s} = 7$ and 8 TeV data sets [3]. In Section IV we study constraints on the dark matter sector of our model using the measured relic abundance and direct detection cross section limits. In section V we further constrain the model using the recent limits on the Higgs invisible width and comment on the dilaton self coupling. Finally in Section VI we summarize our main results and draw conclusions.

* RobynCampbell@cmail.carleton.ca
godfrey@physics.carleton.ca
apuente@physics.carleton.ca
II. A MODEL OF CONFORMAL SYMMETRY BREAKING WITH A SCALAR SINGLET DARK MATTER CANDIDATE

A. Naturalness of the Standard Model

We begin with the observation that at the classical level the Standard Model is scale invariant except for the Higgs mass parameter and a soft mechanism for breaking scale invariance would generate the Higgs mass parameter naturally at the electroweak scale in analogy with supersymmetry. The scale invariance is broken by quantum corrections, that is, by the running of coupling constants. To see this, we can write a general Lagrangian \[ \mathcal{L} = \sum_i g_i(\mu)\mathcal{O}_i(x) \] where \( g_i \) is some coupling constant defined at energy scale \( \mu \) and \( \mathcal{O}_i \) is an operator of dimension \( d_i \). Under scale transformations \( x' = e^{-\lambda}x \), we obtain the following transformations:

\[ \mathcal{O}_i \to e^{\beta \lambda_i} \mathcal{O}_i(\alpha x), \]

\[ \mu \to e^{-\delta} \mu. \]

The variation of the Lagrangian under this transformation is given by

\[ \delta \mathcal{L} = \sum_i g_i(\mu)(d_i + x'\partial_x)\mathcal{O}(x) + \sum_i \beta _i(\mu) \frac{\partial}{\partial g_i} \mathcal{L} \]

where \( \beta_i \) are the \( \beta \) functions of the underlying theory. The above implies that if the dimension of the operator \( d_i = 4 \) and the running is identically zero the theory is scale invariant. From this one obtains the divergence of the scale current \( S^\mu = T^\mu_\nu x^\nu \) where

\[ \partial_x S^\nu = T^\mu_\nu = \sum_i g_i(\mu)(d_i + x'\partial_x)\mathcal{O}(x) + \sum_i \beta _i(\mu) \frac{\partial}{\partial g_i} \mathcal{L}. \]

Since the beta functions vanish at lowest order in perturbation theory, they cannot be responsible for the quadratic divergences in the Higgs mass parameter. This is basically the statement that the quadratic divergences are unrelated to the running of coupling constants and represent a separate explicit source of scale symmetry breaking. However, it has been noted that the appearance of quadratic divergences in a scale free theory is an artifact of the method used to regularize the loop calculations. The explicit appearance of the Higgs mass parameter in the scalar potential of the SM leads to the following trace of the energy momentum tensor

\[ T^\mu_{\mu,\text{tree}} = 2\mu^2 H^\dagger H \]

\[ T^\mu_{\mu,\text{one-loop}} = 2\Delta \mu^2 H^\dagger H + \sum_i \beta_i(\mu)\mathcal{O}(x), \]

where \( H \) are the SM Higgs fields and \( \mu^2 \) the bare Higgs mass parameter with \( \Delta \mu^2 \sim \mu^2 \). It is important to emphasize that the fine-tuning problem may reappear in models where the SM is the low energy description of a more complex theory at high energies and the sensitivity of the Higgs mass parameter to these new scales will depend on the scale invariance properties of these theories.

The recent discovery of a Higgs-like resonance has disfavored technicolor/Higgsless models [3,10], where the electroweak symmetry is broken by strong dynamics. However, there is a well motivated scenario where models of strong dynamics break the electroweak symmetry and yield a light resonance. This observation is due to the fact that the SM is scale invariant in the limit where the Higgs mass parameter goes to zero; and the minimum of the scalar potential has a flat direction, where the vev will spontaneously break the approximate conformal invariance and the electroweak symmetry. In this scenario the Higgs is identified with the massless dilaton with a conformal breaking scale of \( f = 246 \text{ GeV} \). Therefore, if the strongly interacting ultraviolet (UV) theory was also conformal, the condensate breaking the electroweak symmetry would also spontaneously break the conformal symmetry and the dilaton would have properties similar to the Higgs boson [11,15]. It is also possible to study the phenomenology a light dilaton in the presence of a fundamental scalar and this has led to recent work focusing on Higgs-dilatonic mixing and how this scenario can be probed at the LHC [10,15]; however, we do not consider that here.

B. Non-linearly Realized Scale Invariance

If we assume that scale invariance is broken spontaneously by the vev of a dimensionful operator \( \mathcal{O} = f^n \), where \( n \) is the classical dimension of the operator \( \mathcal{O} \), the spontaneous breaking of the scale invariance will imply the existence of a Goldstone boson, the dilaton. The dilaton field transforms non-linearly as [6]

\[ \sigma(x) \to \sigma(\alpha x) + \alpha f, \]

Given Eq. [1], we may incorporate non-linearly realized scale invariance by adding a field \( \chi = e^{\sigma/|f|} \) serving as a conformal compensator with the transformation law

\[ \chi(x) \to e^{\alpha} \chi(\alpha x), \]

and make the replacement

\[ g_i(\mu) \to g_i \left( \mu \frac{\chi}{f} \right) \left( \frac{\chi}{f} \right)^{4-d_i}. \]

In the above equation \( f \) is the order parameter for scale symmetry breaking.

This procedure does not guarantee that once quantum corrections are taken into account the symmetry stays
The work in [13, 14] show that if the operator $O$ given by

$$V(\chi) = \frac{\kappa}{4!} \chi^4.$$  \hspace{1cm} (8)

This term yields a preferred value of $f$, even in the absence of explicit sources of conformal symmetry breaking, that is the flat direction is lost and one must tune $\kappa \to 0$ such that $\langle \chi \rangle = f$ remains undetermined. This issue is relevant since it introduces a fine tuning into the framework of spontaneously broken scale invariance. One way to address this issue is to introduce explicit breaking sources of scale invariance, that is an operator $O(x)$ of scaling dimension $\Delta_O$. In this case the theory is now given by

$$\mathcal{L} = \mathcal{L}_{CFT} + \lambda_O O(x).$$  \hspace{1cm} (9)

The work in [13, 14] show that if the operator $O(x)$ responsible for the breaking of the conformal symmetry is near marginal at the breaking scale, the dilaton can naturally be light and below the scale of the strong dynamics, albeit with some mild fine tuning. In this class of models, the mass of the dilaton is given by $m_\sigma \sim \sqrt{4 - \Delta_O}$ and requires that $\lambda_O \ll 1$, a condition that is not expected to be satisfied in the scenarios of interest for EWSB. However, as stated in [15], it is crucial to have a weakly coupled flat direction available in the theory which is hard to imagine unless the theory is supersymmetric or in the Goldberger-Wise stabilized Randall Sundrum model [19].

### C. The Dilaton Like Higgs with a Scalar Singlet

Models of Technicolor are those for which the electroweak symmetry is broken dynamically by a strongly interacting sector, and where there is no light Higgs. The strongly interacting sector is assumed to be conformal in the far UV. However, the conformal invariance is broken by an operator $O$ which grows large close to the TeV scale triggering EWSB. Within this class of theories if the exact conformal symmetry is broken spontaneously, the low energy effective theory contains a massless scalar, the dilaton associated with the breaking of the conformal symmetry [21, 24].

This class of theories are expected to explicitly violate the conformal invariance since it is necessary to incorporate operators in the conformal sector that become strong in order to induce EWSB, leading to a very massive dilaton. However, in the previous section we mentioned that if these operators are near marginal at the breaking scale, the dilaton can be naturally light and lie below the scale of strong dynamics.

In this work, we are interested in theories where the SM gauge fields are not part of the conformal field theory (CFT), but gauge interactions do constitute a small source of explicit conformal symmetry breaking, since the CFT must be charged under the electroweak group and might also be charged under color. Furthermore, within this framework, the SM fermions may be elementary or composites of the CFT. A well motivated scenario is one where only the right-handed top and the Goldstone bosons needed for EWSB are composites. In addition, we incorporate a scalar dark matter particle, $S$, singlet under the SM gauge group and a composite of the CFT with no direct couplings to the SM. The stability of the dark matter particle is guaranteed by a $Z_2$ parity under which $S$ is odd while all other particles are even. Our framework is similar to the one studied in [25, 26], except that we do not incorporate an elementary Higgs boson candidate and we examine the possibility that the dilaton is the recently discovered spin 0 particle with a mass of 125 GeV.

In our model, to find the couplings of the dilaton to the SM fields and the scalar dark matter candidate we follow the procedure of non-linearly realized scale invariance. We start with the potential

$$V(\bar{\chi}, S) \approx \frac{1}{2} m^2 \bar{\chi}^2 - \frac{\lambda}{3!} \frac{m^2}{f^2} \bar{\chi}^3 + \bar{\chi} \sum_i \left( 1 + \gamma_i \right) m_{\psi_i} \bar{\psi}_i \psi_i + \left( \frac{2\bar{\chi}}{f} + \frac{\bar{\chi}^2}{f^2} \right) \left[ m_{\bar{W}}^2 W^{+\mu} W^- + \frac{1}{2} m_Z^2 Z^\mu Z_\mu \right]$$

$$+ \frac{\alpha_{EM}}{8\pi f} c_{\gamma} \sigma F_{\mu\nu} F^{\mu\nu} + \frac{\alpha_S}{8\pi f} c_{\sigma} \sigma G_{\mu\nu} G^{\mu\nu} + \frac{1}{2} \tilde{m}_S^2 S^2 + \lambda_S S^4$$  \hspace{1cm} (10)

where the dilaton $(\sigma)$ is parametrized with the non-linear realization as $\chi = f e^{\sigma/f}$, and we expand the dilaton about its VEV as $\bar{\chi} = \chi - f$, and $\tilde{m}_S^2 = m_S^2 e^{2\sigma/f}$. We expand the exponentials to leading orders and obtain the following parametrization for the potential:
where \( m_f \) denotes the scale where the conformal symmetry is spontaneously broken, \( v \), the EWSB scale and \( \epsilon = v/f \). The parameter \( \lambda \) parametrizes explicit conformal symmetry violating effects. In addition, the mass of the dilaton and its self-interactions were obtained assuming that either the operator explicitly breaking the conformal symmetry in eqn. \( \lambda \) is near marginal or \( \lambda \ll 1 \). As outlined in [11], for a single operator with dimension \( \leq 4 \), this leads to a bound on the dilaton cubic coupling \( \lambda < 5 \) which can be relaxed with a more elaborate conformal symmetry breaking sector. We assume that \( f \geq v \) since we want the CFT to induce EWSB. The parameters \( c_{\gamma \gamma} \) and \( c_gg \) can be written as

\[
c_{\gamma \gamma} = b_{\text{IR}}^{\text{EM}} - b_{\text{UV}}^{\text{EM}}, \quad c_gg = b_g^{\text{IR}} - b_g^{\text{UV}}. \tag{12}
\]

The coefficients denoted by \( b_{\text{IR}} \) parametrize the breaking of the conformal symmetry at the quantum level below the breaking scale \( f \). The running in the UV, \( b_{\text{UV}} \), does constitute an explicit source of conformal symmetry breaking and since the gluon and photon are elementary, there is no constraint on \( b_{\text{UV}} \). The structure of the \( b_{\text{UV}} \)'s depends on the details of the CFT and receives contributions from both elementary and CFT states. In practice we have subsumed the SM loop factors [27] into \( c_gg \) and \( c_{\gamma \gamma} \) [13] [29].

In the fermion sector, the leading source of conformal symmetry breaking are the fermion masses and the coupling of the dilaton to fermions depend on whether the latter are elementary fields or composites of the CFT. This choice has been parametrized with the anomalous dimension, \( \gamma_i \) in Eq. (12), which measures the explicit breaking of conformal symmetry that arises from the mixing of elementary and composite fields. Within our framework, we allow for the possibility that the right-handed top quark is a composite of the CFT and let \( \gamma_T \rightarrow 0 \). In the case of elementary fermions, in theories of conformal technicolor, fermion masses arise from couplings of fermions to a scalar operator with the quantum numbers of the Higgs [13] [28]. In our work we will use the parametrization introduced in [13], where the coupling of the dilaton to the top quark is given by

\[
m_t \frac{1}{f} (1 + \delta) \sigma Q_3 v, \tag{13}
\]

where \( \delta = \Delta_H - 1 \), and \( \Delta_H \) is the scaling dimension of the scalar operator in the CFT.

\[V(\sigma, S) \approx \frac{1}{2} m_\sigma^2 \sigma^2 + \epsilon \frac{m_\sigma^2}{v} \left( \frac{1}{2} - \frac{1}{6} \right) \sigma^3 + \frac{\sigma}{f} \sum \left( 1 + \gamma_i \right) m_{\psi_i} \bar{\psi}_i \psi_i + \left( \frac{2\sigma}{f} + \frac{\sigma^2}{f^2} \right) \left[ m_\psi^2 W^{\mu\nu} W_\mu \gamma^\nu + \frac{1}{2} m_\psi^2 Z^\mu Z^\nu \right] \]

\[+ \frac{\alpha_{\text{EM}}}{8\pi f} c_{\gamma \gamma} F_{\mu\nu} F^{\mu\nu} + \frac{\alpha_S}{8\pi f} c_gg \sigma G_{\mu\nu} G^{\mu\nu} + \frac{1}{2} m_S^2 \left( 1 + 2c_\sigma \frac{\sigma}{v} + 2c_{\gamma \gamma} \frac{\sigma^2}{v^2} \right) S^2 + \lambda_S S^4, \tag{11}
\]

\[\text{III. FIT TO HIGGS SIGNAL STRENGTHS USING MARKOV CHAIN MONTE CARLO}\]

In this section we present the results of a general fit to the most up-to-date Higgs data using the 10 Higgs production and decay signal strengths from the most recent ATLAS CMS combined fit [3]. The signal strengths are defined as:

\[
\mu_i = \frac{\sigma_i}{(\sigma_i)_{\text{SM}}} \quad \text{and} \quad \mu^f = \frac{BR^f}{(BR^f)_{\text{SM}}}, \tag{14}
\]

for \( \sigma_i \) with \( i = ggF, VBF, WH, ZH, ttH \) and \( BR_i \) with \( f = ZZ, WW, \gamma \gamma, \tau \tau, bb \). For the \( BR_i \)’s one needs to also evaluate the modified total width where we rescale the SM Higgs partial widths by \( \epsilon^2 \) with the additional rescaling of the \( c_{gg} \) and \( c_{\gamma \gamma} \) factors (see for example Ref. [29]):

\[
\Gamma_{\text{tot}} = \epsilon^2 \left( \sum f \Gamma_{f}^{\text{SM}} + \Gamma_{WW}^{\text{SM}} + \Gamma_{ZZ}^{\text{SM}} + \Gamma_{\gamma \gamma}^{\text{SM}} \right) + \left( \frac{c_{gg}}{c_{gg}^{\text{SM}}} \right)^2 \Gamma_{gg}^{\text{SM}} + \left( \frac{c_{\gamma \gamma}}{c_{\gamma \gamma}^{\text{SM}}} \right)^2 \Gamma_{\gamma \gamma}^{\text{SM}} + \Gamma_{\text{inv}} \tag{15}
\]

where \( \Gamma_i^{\text{SM}} \) are the SM Higgs partial widths and \( \Gamma_{\text{inv}} \) is the dilaton’s invisible width when decays to \( SS \) are kinematically allowed. For the convenience of the reader we quote the numerical values for the signal strengths in Table 1. It should be noted that there are assumptions built into this set of signal strengths in that the production signal strengths assume SM BR’s and the final state signal strengths assume SM production cross sections. In addition to the signal strengths we include an additional observable in the fit which is the bound derived in [30] on the total width of the Higgs boson of \( 6.1^{+3.7}_{-2.9} \) MeV. This result is in agreement with a CMS study that places an upper bound of 22 MeV at 95% CL [31].

To scan the parameter space we use a Markov Chain Monte Carlo (MCMC) and implement the Metropolis-Hastings algorithm with simulated annealing [32]. The target distribution is a Gaussian likelihood proposal given by

\[
\mathcal{L} = \prod_i \exp \left[ -\frac{(x_i - \bar{x}_i)^2}{2\sigma_i^2} \right] \tag{16}
\]

where the \( \bar{x}_i \) are the experimental observables, and \( \sigma_i \) their corresponding uncertainties. The parameters allowed to vary in the fit are those characterizing the coupling of the dilaton to the massless gauge bosons given
FIG. 1. Results of the MCMC fit with simulated annealing. The areas enclosed by the black contour lines correspond to the regions of parameter space for $\epsilon \geq 0.98$ allowed at 95% confidence level corresponding to $X < 5.9$.

below and introduced in the previous section, the ratio $\epsilon$ of the EWB scale $v = 246$ GeV to the scale of spontaneous breaking of scale invariance, $f$, and the mass of the dark matter candidate $m_S$:

$$\epsilon^2 = \frac{(v/f)^2}{c_{gg}} = (b_{IR} - b_{UV})^2,$$

where the expressions for $c_{gg}$ and $c_{\gamma\gamma}$ are to leading order and to which we should add the SM loop contributions. In practice we let them float in the fits. In Figure 1 we show the results of the MCMC in the $c_{gg} - c_{\gamma\gamma}$ plane where the fit has been restricted to values of $\epsilon \geq 0.98$ based on the most up-to-date electroweak precision data (EWPD) [33]. Regions allowed at 95% confidence level (CL) are those for which $X < 5.9$ where $X$ is the sum of the terms in the exponential of eqn. 16. In our fit all Yukawa anomalous dimensions have been set to zero, and any information on the IR running of both $\beta$ functions has been diluted within the factor $c_{gg}$ and $c_{\gamma\gamma}$. We find that our results are consistent with results obtained in a similar analysis but with the 7 TeV data set [13, 15, 34].

IV. DARK MATTER

The introduction of the scalar singlet dark matter candidate introduces two additional parameters, the mass of the DM candidate $m_S$ and DM self coupling $\lambda_S$ where the dilaton DM couplings are proportional to $m_S^2$. In addition the dilaton trilinear coupling comes into play through DM self-annihilation to dilaton final states which contributes to the DM relic abundance. $\lambda_S$ does not enter any of the expressions for the observables we are studying so it remains unconstrained. In the following subsections we explore the further constraints that the DM relic abundance and direct detection limits put on the parameters

| Process   | Value   |
|-----------|---------|
| $\mu_{ggF}$ | $1.03^{+0.10}_{-0.14}$ |
| $\mu_{VBF}$ | $1.18^{+0.25}_{-0.23}$ |
| $\mu_{WH}$ | $0.89^{+0.40}_{-0.38}$ |
| $\mu_{ZH}$ | $0.79^{+0.38}_{-0.36}$ |
| $\mu_{ttH}$ | $2.3^{+0.6}_{-0.5}$ |
| $\mu_{\tau\tau}$ | $1.14^{+0.19}_{-0.18}$ |
| $\mu_{ZZ}$ | $1.29^{+0.26}_{-0.23}$ |
| $\mu_{WW}$ | $1.09^{+0.18}_{-0.16}$ |
| $\mu_{bb}$ | $1.11^{+0.24}_{-0.22}$ |
| $\mu_{\nu\nu}$ | $0.70^{+0.29}_{-0.27}$ |

TABLE I. Measured signal strengths and their total uncertainties for different Higgs boson production processes and decay channels. The results are for the combined ATLAS CMS fits for the combined $\sqrt{s} = 7$ and 8 TeV data from ref. [2].
of the model.

A. Relic Abundance

The scalar singlet dark matter candidate $S$ is stabilized by the existence of an unbroken $Z_2$ symmetry and the interactions introduced in eqn. (11) yield a mechanism for its relic abundance. The $SS$ annihilations into SM states proceed through s-channel dilaton exchange. The Feynman diagrams for these annihilation processes are shown in Figure 2. The result is given by

$$
\frac{\langle \sigma v \rangle}{g_*^{1/2} x_{FO}} \approx 3 \times 10^{-6},
$$

where $g_*$ is the number of relativistic degrees of freedom at the freeze out temperature, $g_{S}$ is the number of degrees of freedom of the dark matter candidate which is equal to one for a real scalar singlet and $M_{pl}$ is the Planck mass. The present day relic abundance is then given by

$$
\Omega_{S} h^2 \approx 1.65 \times 10^{-10} \left( \frac{\text{GeV}^2}{\langle \sigma v \rangle} \right) \log \left( \frac{0.038 g_{S} m_{S} M_{pl} \langle \sigma v \rangle}{g_*^{1/2} x_{FO}} \right).
$$

The observed relic abundance can be achieved with a value of $\langle \sigma v \rangle \approx 3 \times 10^{-26} \text{cm}^3/\text{s}$.

The thermally average cross section at temperature $T$ can be calculated from the annihilation cross section $\langle \sigma v \rangle$ after summing over all the kinematically allowed diagrams given in Figures 2 and 3. The result is given by

$$
\langle \sigma v \rangle = \int_{4 m_{S}^2}^{\infty} ds \left( s - 4 m_{S}^2 \right) s^{1/2} K_1(s^{1/2}/T) \sigma (s),
$$

where $K_1(z)$ and $K_2(z)$ are the modified Bessel functions of the second kind.

Our dark matter analysis compliments the analysis by the authors in Ref. [26] in regions of parameter space where $v \approx f$. However in our scenario, annihilations into fundamental scalars are absent and the value of the dilaton mass has been fixed to 125 GeV and plays the role of the Higgs boson. The parameter space of interest is that where the Higgs boson. The parameter space of interest is that where $c_{gg} = -1/3$ and $c_{\gamma \gamma} = -11/9$ are introduced. Both scenarios differ from the SM predictions of $c_{gg}^{SM} \approx 0.67$ and $c_{\gamma \gamma}^{SM} \approx -6.5$.

The parameter space for $c_{gg}$ and $c_{\gamma \gamma}$ allowed at 95% CL is consistent with the SM expectation but allows for val-
values in the ranges $0.5 \lesssim |c_{gg}| \lesssim 1$ and $4 \lesssim |c_{\gamma\gamma}| < 8$. This is particularly interesting since our model is very similar to singlet scalar dark matter extensions of the SM \cite{57,58} albeit with two main differences: The mass of the dark matter candidate and its coupling to the Higgs-like dilaton are related, and the trilinear dilaton coupling, $\lambda$, is essentially a free parameter. However, the conformal algebra and unitarity imply $\lambda \sim 2$. The SM trilinear coupling corresponds to $\lambda = 3$, as long as $\lambda_\sigma \ll f^2$. This is important since annihilations into a pair of dilatons can be enhanced for large values of $\lambda$ and thus suppress the relic abundance leading to smaller direct detection cross sections as we will see in the next section. The study by the authors in \cite{11} only incorporates a single source of explicit conformal symmetry breaking through an operator of dimension $\Delta_\sigma \leq 4$. The bound on $\lambda$ from such an operator is $\lambda \lesssim 5$. In our study we will allow for $\lambda$ to be as large as $4\pi$. This has two implications: The annihilation cross section can be greatly enhanced as well as dilaton pair production at colliders. We will comment on this and also the status of an upcoming trilinear Higgs coupling measurement at the LHC in the following section.

In Figure 4 we show the present day DM abundance as a function of the dark matter mass corresponding to regions of parameter space consistent with the Higgs signal strength fit of the previous section with the value of $\lambda = 3$ denoted by the red line, and $\lambda = 4\pi$ by the blue line (the darker lower line when the two lines don’t overlap). The relic abundance calculation was carried out using MicroMegas \cite{58} with model files generated with FeynRules \cite{59}. These results are in perfect agreement with two independent calculations using Mathematica \cite{60} and a separate computer program using numerical integration. The range $10 < m_S \lesssim 50$ GeV correspond to $\sigma$-$S$ couplings consistent with a coupling in the Higgs-$S$ system, $\lambda_{\nu}H^+HS^2$ (See for example \cite{57}), between 0.001 and 0.05. Within this region of parameter space, the dominant annihilation channels are to $b\bar{b}$ and $\tau^+\tau^-$ pairs and the cross sections are suppressed by factors of $(m_{\nu}\epsilon)/v$. The observed DM abundance is achieved for a value of $m_S \sim 50$ GeV and rapidly becomes greatly suppressed in the resonant region where $m_S = m_\sigma/2$. Beyond this point and below $m_S = m_\sigma$ the dominant annihilation channels are into $W^+W^-$ and $ZZ$ pairs with an enhancement in the abundance below the mass of the $W$ gauge boson. This region of parameter space corresponds to values of $\lambda_\sigma$ in the range $0.05 - 0.3$ and annihilations are enhanced. We also observe that even when annihilations into a pair of dilatons are kinematically allowed, the value of $\lambda$ is SM-like and there is no further enhancement of the annihilation cross section. This situation is very different for $\lambda = 4\pi$ where a very noticeable drop in the abundance can be seen for $m_S$ in the vicinity of 125 GeV. Beyond $m_S \sim m_\sigma$, annihilations into top quarks dominates the thermalized cross section and we see a further reduction in the relic abundance, albeit constant for values of $m_S \gtrsim 300$ GeV. In this region of parameter space $\lambda_\sigma$ can lie between values of 1 and $4\pi$, where the latter corresponds to dark matter masses above $\sim 900$ GeV. Even though these values of the dark matter mass lie below the unitarity bound of $m_S \sim \sqrt{8\pi f}$ \cite{20}, this region of parameter space needs to be rescaled by a Sommerfeld enhancement factor.

A Higgs-like dilaton augmented by a gauge singlet scalar dark matter candidate cannot by itself reproduce the observed abundance in most of the parameter space but it can give us an additional probe into the nature of the Higgs self-coupling. For masses $m_S > 125$ GeV the self coupling can influence how important the DM anni-
hilations into a pair of Higgs-like dilatons are and thus determine how large the DM abundance in that mass region can be.

A final comment is in order. As we saw in the previous section, the region allowed by the MCMC fit is symmetric in four quadrants. The SM expectation is that $c_{gg}$ is positive and $c_{\gamma\gamma}$ negative. However $c_{gg}$ and $c_{\gamma\gamma}$ enter quadratically into the annihilation cross section calculation making it insensitive to the sign of these parameters. We will see in the next subsection that constraints from DM direct detection measurements can further constrain $c_{gg}$ and tell us its sign and hence rule out regions of the $c_{gg}$ parameter space.

B. Direct Detection

In our model the messenger between the dark sector and the SM fields is the Higgs-like dilaton so that only interactions which are spin independent between $S$ and the SM fields are important. The heavy quark contribution resulting in the following expression for $f_N$:

$$f_N = \epsilon(f_u + f_d + f_s + f_G) = \epsilon \left( 0.099 + 0.201 \frac{c_{gg}}{c_{SM}} \right).$$

Using eqns. 22 and 25, the parameters allowed by the fits to the Higgs signal strengths consistent with the DM relic abundance, we obtain the direct detection cross sections shown in Fig. 5 along with the LUX 03, XENON100 04 and projected DEAP 65 direct detection limits. The points shown in Fig. 5 were obtained using $c_{gg} > 0$, consistent with the sign of the SM value, $c_{SM}^{gg}$. If we relax this constrain and allow negative values of $c_{gg}$ from the MCMC fit the resulting value of $f_N$ will be small enough to yield cross sections that evade direct detection constraints in most of the region of parameter space for which $m_S > 50$ GeV, for both small and large $\lambda$.

V. OTHER HIGGS PROPERTIES

A. Invisible Dilaton Width

In the previous sections we analyzed the constraints on the dilaton, which continues to be a viable candidate for the 125 GeV scalar resonance observed at the LHC. We have also analyzed whether within the allowed parameter space, the dilaton can couple to a scalar dark matter particle that does not over close the universe and does not violate current dark matter direct detection constraints. Within our framework, the dilaton can decay to the DM candidate, $S$, and thus contribute to its invisible width. If the dilaton is to mimic the SM Higgs boson, its branching ratio to invisible final states must lie below the most recent 95% CL upper bound from the ATLAS collaboration of $BR(H \rightarrow inv.) < 0.28$ 06. In this model, the invisible decay width of the dilaton is given by

$$\Gamma(\sigma \rightarrow SS) = \frac{m_S^4 \epsilon^2}{8\pi m_s v^2} \sqrt{1 - 4 \frac{m_S^2}{m^2}},$$

with the branching ratio given by

$$BR(\sigma \rightarrow inv.) = \frac{\Gamma(\sigma \rightarrow SS)}{\Gamma(\sigma \rightarrow SM) + \Gamma(\sigma \rightarrow SS)}.$$  

where $\Gamma(\sigma \rightarrow SM)$ can be read from eqn. 15.

In our model, a scalar dark matter particle with a mass $m_S \approx 50$ GeV can saturate the observed dark matter relic abundance. However, this mass corresponds to a $S - \sigma$ coupling that leads to a direct detection cross section above the experimental bounds. In our framework masses below $m_s/2$ (black points in Figure 3) lead to an invisible branching fraction well above the experimental bound given above. This is consistent with the results obtained for a singlet extension of the SM where the Higgs-S couplings, $\lambda_S \gtrsim 0.015$ are ruled out by this observable.
Masses in the range $10 - m_\sigma/2$ correspond to values of $\lambda_p$ in the Higgs-$S$ system in the range $\sim 0.01 - 0.07$.

### B. The dilaton self coupling

Here we comment on the dilaton self coupling. The observation of what appears to be a SM-like Higgs boson has given us an insight into the mechanism of EWSB. In our work, we have focused on a particular form for the dilaton potential and precise measurement of its parameters is necessary. Within the SM, the cubic and quartic Higgs self-couplings are uniquely defined:

$$\lambda^\text{SM}_{HHH} = \frac{m_H^2}{2\sigma^2}$$

Double and trilinear Higgs production can help probe the cubic and quartic self-interactions governing the Higgs potential. Any deviation from the SM value will be a sign of new physics and force us to rethink and possibly reformulate the recipe for EWSB. In our work, we have incorporated the potential introduced in ref. [11] and given in eq. [11]. With this parametrization, the cubic coupling is given by

$$\lambda_{\sigma\sigma\sigma} = \epsilon \left(\frac{m_\sigma}{v}\right)^2 \left(1 - \frac{1}{6}\lambda\right),$$

with $\lambda$ directly related to the operator given in eq. [9]

$$\lambda \approx \begin{cases} \Delta_\sigma + O(\lambda_\sigma) & \text{when } \lambda_\sigma \ll 1, \\ 5 + O(|\Delta_\sigma - 4|) & \text{when } |\Delta_\sigma - 4| \ll 1, \end{cases}$$

where $\Delta_\sigma$ is the scaling dimension of the symmetry breaking operator. A SM-like coupling corresponds to a scaling dimension of $\Delta_\sigma = 2$ for a case where $\lambda_\sigma/f^2 \ll 1$. More elaborate ways of breaking scale invariance, and in particular with near-marginal operators can lead to values of $\lambda$ closer to its non-perturbative value of $4\pi$; however this may involve a certain degree of fine tuning to generate a dilaton mass below the conformal symmetry breaking scale $[13, 14]$. For $\lambda = 4\pi$, the corresponding value of the cubic coupling is approximately a factor of 1.5 larger than $\lambda^\text{SM}_{HHH}$. Thus, we can get partial insight into the dynamics leading to EWSB and whether the observed scalar is indeed the one predicted by the SM or a dilaton by measuring double Higgs production at the LHC; its prospects which have been actively studied and that show that a combination of channels as well as large data sets are needed to overcome the SM backgrounds [67, 70]. A detailed review that includes beyond the LHC colliders in [71], shows how a 1 TeV lepton collider such as the ILC can achieve a precision of $\sim 30\%$ on the determination of the cubic self-coupling. In addition, a recent CMS study places an upper bound on non-resonant double Higgs production of 200 times the SM rate [72] using center of mass energies of $\sqrt{s} = 13$ TeV and 2.7 fb$^{-1}$ of data. This limit is still much weaker than the limit obtained by ATLAS with the full data set at $\sqrt{s} = 8$ TeV. The observed limit in this case is approximately 70 times the SM rate [73].

### VI. SUMMARY

We have studied the phenomenology of a model that treats the physical Higgs boson of the SM as a dilaton, the pseudo-Goldstone boson of spontaneous conformal symmetry breaking and introduces a scalar singlet dark matter candidate. We fitted the parameters of the model to the measured Higgs boson properties and found that the parameters of the model sensitive to these observables are forced to be consistent with those of the standard model to a high degree modula a sign ambiguity.
for two of them: $c_{gg}$ and $c_{\gamma\gamma}$. Using the allowed region of the parameter space we then study constraints on the model from DM constraints. We find that $m_S$ is ruled out for masses less than $\sim 50$ GeV by the allowed values of DM relic abundance. When direct detection limits by the XENON and LUX experiments are imposed only a small mass region remains assuming a SM-like dilaton self coupling. When one assumes the self coupling is more like a strongly interacting composite theory a second region survives in the mass region of $\sim 125 - 300$ GeV. In addition, the non-linear nature of the coupling between the dilaton and $S$ yield large couplings for masses $m_S < m_{\sigma}/2$ and thus DM candidates below this limit are ruled out by the current upper limit of the Higgs invisible width. The higher mass region will be probed in the near future by the DEAP DM search experiment. We also find that the region of parameter space allowed by direct detection experiments is enhanced when one allows for a negative value of $c_{gg}$. When the Higgs was first discovered, a slight enhancement in the $H \rightarrow \gamma\gamma$ channel led to work that could potentially explain this excess with an effective negative Higgs-gluon coupling \cite{[74]}, due to an overall negative coupling between the Higgs and top quark or new physics running in the loop. In this work we see that the structure of the CFT dictates the sign of $c_{gg}$, but we also observe that whatever this structure is, a fit to the Higgs data leads to an absolute value consistent with the SM. Furthermore, we observe that future precision measurements on the parameters governing the Higgs potential can tell us if we are indeed observing the scalar responsible for the Brout-Englert-Higgs EWSB mechanism or a dilaton. If these measurements deviate from the SM expectation, we can learn more about the CFT for energies $f \gtrsim 246$ GeV and the mechanism responsible for a light dilaton, given that its natural mass will be on the order of $4\pi f$.

One important lesson to be learned from this work is that if we give up the idea of explaining with a single model the complete nature of the dark matter and simply work on a first order theory producing a fraction of the abundance with a dark matter component with spin independent interactions with nuclei, one can learn in a way complimentary to precise measurements of the Higgs scalar potential parameters the nature of the UV theory leading to EWSB.

ACKNOWLEDGMENTS

The authors have benefited from conversations with Heather Logan, Alex Poulin and Rouzbeh Yazdi. This research was supported in part the Natural Sciences and Engineering Research Council of Canada under Grant No. 121209-2009 SAPIN.

Appendix A: SS Self Annihilations Cross Sections

In this appendix we include the expressions for the $SS$ annihilation cross sections used to calculate the DM relic abundance.

\begin{align}
\sigma(SS \rightarrow ZZ) &= \frac{1}{8\pi} \sqrt{s - 4m_S^2} \left( \frac{m_S \epsilon}{v} \right)^4 \frac{s}{s - m_S^2} \left( 1 + \frac{12m_S^4}{s^2} - \frac{4m_S^2}{s} \right) \\
\sigma(SS \rightarrow W^+W^-) &= \frac{1}{4\pi} \sqrt{s - 4m_W^2} \left( \frac{m_S \epsilon}{v} \right)^4 \frac{s}{s - m_S^2} \left( 1 + \frac{12m_W^4}{s^2} - \frac{4m_W^2}{s} \right) \\
\sigma(SS \rightarrow \gamma\gamma) &= \frac{(c_{\gamma\gamma}\alpha_{EM})^2}{64\pi^3} \frac{s}{s - m_S^2} \left( \frac{m_S \epsilon}{v} \right)^4 \left( 1 - \frac{4m_S^2}{s} \right)^{-\frac{1}{2}} \\
\sigma(SS \rightarrow gg) &= \frac{(c_{gg}\alpha_{s})^2}{8\pi^3} \frac{s}{s - m_S^2} \left( \frac{m_S \epsilon}{v} \right)^4 \left( 1 - \frac{4m_S^2}{s} \right)^{-\frac{1}{2}} \\
\sigma(SS \rightarrow f\bar{f}) &= \frac{m_f^2 N_C}{2\pi s} \left( \frac{m_S \epsilon}{v} \right)^4 \frac{1}{s - m_f^2} \sqrt{s - 4m_S^2}
\end{align}

where $N_C = 1$ for leptons and 3 for quarks.

\begin{align}
\sigma(SS \rightarrow \sigma\sigma) &= \frac{1}{16\pi s} \sqrt{s - 4m_S^2} \left( \frac{m_S \epsilon}{v} \right)^4 \int_{-1}^{+1} d\cos\theta \left[ \frac{m_\sigma^4(3 - \lambda)^2}{s - m_\sigma^2} \right. \\
&\quad \left. + \frac{4m_\sigma^2m_\sigma^2 (3 - \lambda)(s - m_\sigma^2)}{(s - m_\sigma^2)^2 + \Gamma_\sigma^2m_\sigma^2} \left( \frac{1}{t} + \frac{1}{u} \right) + \frac{4(3 - \lambda)m_\sigma^2(s - m_\sigma^2)}{(s - m_\sigma^2)^2 + \Gamma_\sigma^2m_\sigma^2} \right] \\
&\quad \left. + 4m_\sigma^2 \left( \frac{1}{t} + \frac{1}{u} \right)^2 + 8m_S^2 \left( \frac{1}{t} + \frac{1}{u} \right) + 4 \right]
\end{align}
where $s$, $t$ and $u$ are the appropriate Mandelstam variables.

[1] G. Aad et al. [ATLAS Collaboration], “Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC,” Phys. Lett. B 716, 1 (2012) arXiv:1207.7214 [hep-ex];

[2] S. Chatrchyan et al. [CMS Collaboration], “Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC,” Phys. Lett. B 716, 30 (2012) arXiv:1207.7235 [hep-ex].

[3] G. Aad et al. [ATLAS and CMS Collaborations], “Measurements of the Higgs boson production and decay rates and constraints on its couplings from a combined ATLAS and CMS analysis of the LHC $pp$ collision data at $\sqrt{s} = 7$ and 8 TeV,” arXiv:1606.02266 [hep-ex].

[4] S. P. Martin, “A Supersymmetry primer,” Adv. Ser. Direct. High Energy Phys. 21, 1 (2010) [Adv. Ser. Direct. High Energy Phys. 18, 1 (1998)] hep-ph/9709356.

[5] S. Dimopoulos and G. F. Giudice, “Naturalness constraints in supersymmetric theories with nonuniversal soft terms,” Phys. Lett. B 357, 573 (1995) doi:10.1016/0370-2693(95)00961-J hep-ph/9507282.

[6] S. Coleman, “Dilatations,” in Aspects of Symmetry, Cambridge University Press, 1985.

[7] A. W. Bardeen, “On naturalness in the standard model,” FERMILAB-CONF-95-391-T, C95-08-27.3.

[8] W. A. Bardeen, “Nonlinear realizations of space-time symmetries. Scalar and tensor gravity,” Annals Phys. 62, 98 (1970). doi:10.1016/0003-4916(71)90269-7

[9] J. R. Ellis, “Aspects of conformal symmetry and chirality,” Nucl. Phys. B 22, 478 (1970) Erratum: [Nucl. Phys. B 25, 639 (1971)]. doi:10.1016/0550-3213(70)90439-1

[10] J. R. Ellis, “Phenomenological actions for spontaneously-broken conformal symmetry,” Nucl. Phys. B 26, 536 (1971). doi:10.1016/0550-3213(71)90193-3

[11] W. D. Goldberger, M. B. Wise, “Modulus stabilization with bulk fields,” Phys. Rev. Lett. 83, 4922 (1999) doi:10.1103/PhysRevLett.83.4922 hep-ph/9907447.

[12] A. Salam and J. A. Strathdee, “Nonlinear realizations. 2. Conformal symmetry,” Phys. Rev. 184, 1760 (1969). doi:10.1103/PhysRev.184.1760

[13] C. J. Isham, A. Salam and J. A. Strathdee, “Spontaneous breakdown of conformal symmetry,” Phys. Lett. B 31, 300 (1970). doi:10.1016/0370-2693(70)90177-2

[14] C. J. Isham, A. Salam and J. A. Strathdee, “Nonlinear realizations of space-time symmetries. Scalar and tensor gravity,” Annals Phys. 62, 98 (1970). doi:10.1016/0003-4916(71)90269-7

[15] J. R. Ellis, “Aspects of conformal symmetry and chirality,” Nucl. Phys. B 22, 478 (1970) Erratum: [Nucl. Phys. B 25, 639 (1971)]. doi:10.1016/0550-3213(70)90422-0, 10.1016/0550-3213(71)90439-1

[16] T. Abe, R. Kitano, Y. Konishi, K. y. Oda, J. Sato and S. Sugiyama, “Minimal Dilaton Model,” Phys. Rev. D 86, 115016 (2012) doi:10.1103/PhysRevD.86.115016 arXiv:1209.4544 [hep-ph].

[17] J. Cao, Y. He, P. Wu, M. Zhang and J. Zhu, “Higgs Phenomenology in the Minimal Dilaton Model after Run I of the LHC,” JHEP 1401, 150 (2014) doi:10.1007/JHEP01(2014)150 arXiv:1311.6661 [hep-ph].

[18] D. W. Jung and P. Ko, “Higgs-dilaton(radion) system confronting the LHC Higgs data,” Phys. Lett. B 732, 364 (2014) doi:10.1016/j.physletb.2014.04.005 arXiv:1401.5586 [hep-ph].

[19] W. D. Goldberger and M. B. Wise, “Modulus stabilization with bulk fields,” Phys. Rev. Lett. 83, 4922 (1999) doi:10.1103/PhysRevLett.83.4922 hep-ph/9907447.

[20] S. Weinberg, “Implications of Dynamical Symmetry Breaking in the Weinberg-Salam Theory,” Phys. Rev. D 20, 2619 (1979). doi:10.1103/PhysRevD.20.2619

[21] S. Weinberg, “Implications of Dynamical Symmetry Breaking,” Phys. Rev. D 13, 974 (1976). doi:10.1103/PhysRevD.13.974

[22] S. Weinberg, “Implications of Dynamical Symmetry Breaking: An Addendum,” Phys. Rev. D 19, 1277 (1979). doi:10.1103/PhysRevD.19.1277

[23] W. D. Goldberger, B. Grinstein and W. Skiba, “Distinguishing the Higgs boson from the dilaton at the Large Hadron Collider,” Phys. Rev. Lett. 100, 111802 (2008) doi:10.1103/PhysRevLett.100.111802 arXiv:0708.1463 [hep-ph].

[24] L. Vecchi, “Phenomenology of a light scalar: the dilaton,” Phys. Rev. D 82, 076009 (2010) doi:10.1103/PhysRevD.82.076009 arXiv:1002.1721 [hep-ph].

[25] B. Bellazzini, C. Csaki, J. Hubisz, J. Serra and J. Terning, “A Higgslike Dilaton,” Eur. Phys. J. C 73, no. 2, 2333 (2013) doi:10.1140/epjc/s10052-013-2333-x arXiv:1209.3299 [hep-ph].

[26] Z. Chacko and R. K. Mishra, “Effective Theory of a Light Dilaton,” Phys. Rev. D 87, no. 11, 115006 (2013) doi:10.1103/PhysRevD.87.115006 arXiv:1209.3022 [hep-ph].

[27] Z. Chacko, R. Franceschini and R. K. Mishra, “Resonance at 125 GeV: Higgs or Dilaton/Radion?,” JHEP 1304, 015 (2013) doi:10.1007/JHEP04(2013)015 arXiv:1209.3259 [hep-ph].

[28] V. Khachatryan et al. [CMS Collaboration], “Constraints on the Higgs boson width from off-shell pro-
duction and decay to Z-boson pairs," Phys. Lett. B 736, 64 (2014) doi:10.1016/j.physletb.2014.06.077 [arXiv:1405.3455 [hep-ex]].

[32] E. A. Baltz, M. Battaglia, M. E. Peskin and T. Wizansky, “Determination of dark matter properties at high-energy colliders,” Phys. Rev. D 74, 103521 (2006) doi:10.1103/PhysRevD.74.103521 [hep-ph/0602187].

[33] M. Ciuchini, E. Franco, S. Mishima and L. Silvestrini, “Electroweak Precision Observables, New Physics and the Nature of a 126 GeV Higgs Boson,” JHEP 1308, 106 (2013) doi:10.1007/JHEP08(2013)106 [arXiv:1306.4644 [hep-ph]].

[34] J. McDonald, “A higgs-like dilaton: viability and implications,” EPJ Web Conf. 60, 17005 (2013) doi:10.1051/epjconf/20136017005 [arXiv:1312.0259 [hep-ph]].

[35] E. W. Kolb and M. S. Turner, “The Early Universe,” Front. Phys. 69, 1 (1990).

[36] P. A. R. Ade et al. [Planck Collaboration], “Planck 2013 results. XVI. Cosmological parameters,” Astron. Astrophys. 571, A16 (2014) doi:10.1051/0004-6361/201321591 [arXiv:1303.5076 [astro-ph.CO]].

[37] S. Baek, P. Ko and W. I. Park, “Invisible Higgs Decay Width vs. Dark Matter Direct Detection Cross Section for Singlet Scalar Dark Matter: LUX, invisible Higgs decays and gamma-ray lines,” JHEP 1503, 045 (2015) [arXiv:1412.1105 [hep-ph]].

[38] A. Djouadi, O. Lebedev, Y. Mambrini and J. Smirnov, “Scalar Dark Matter: Direct vs. Indirect Detection,” arXiv:1509.04282 [hep-ph].

[39] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, “micrOMEGAs: A program for calculating dark matter observables,” Comput. Phys. Commun. 185, 960 (2014) doi:10.1016/j.cpc.2013.10.016 [arXiv:1305.0237 [hep-ph]].

[40] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, “Update on scalar singlet dark matter,” Phys. Rev. D 88, 055025 (2013) [arXiv:1306.4710 [hep-ph]].

[41] J. A. Oller, “Vacuum Stability, Perturbativity, and Scalar Dark Matter: monochromatic gamma rays and metastable vacua,” Phys. Rev. D 90, 055014 (2014) [arXiv:1405.3530 [hep-ph]].

[42] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, “Update on scalar singlet dark matter,” Phys. Rev. D 88, 055025 (2013) [arXiv:1306.4710 [hep-ph]].

[43] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, “Update on scalar singlet dark matter,” Phys. Rev. D 88, 055025 (2013) [arXiv:1306.4710 [hep-ph]].

[44] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, “Update on scalar singlet dark matter,” Phys. Rev. D 88, 055025 (2013) [arXiv:1306.4710 [hep-ph]].

[45] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, “Update on scalar singlet dark matter,” Phys. Rev. D 88, 055025 (2013) [arXiv:1306.4710 [hep-ph]].
[astro-ph.CO].

[65] P.-A. Amaudruz et al. [DEAP Collaboration], “DEAP-3600 Dark Matter Search,” Nucl. Part. Phys. Proc. 273-275, 340 doi:10.1016/j.nuclphysbps.2015.09.048 [arXiv:1410.7673 [physics.ins-det]].

[66] G. Aad et al. [ATLAS Collaboration], “Constraints on new phenomena via Higgs boson couplings and invisible decays with the ATLAS detector,” JHEP 1511, 206 (2015) doi:10.1007/JHEP11(2015)206 [arXiv:1509.00672 [hep-ex]].

[67] M. J. Dolan, C. Englert and M. Spannowsky, “Higgs self-coupling measurements at the LHC,” JHEP 1210, 112 (2012) doi:10.1007/JHEP10(2012)112 [arXiv:1206.5001 [hep-ph]].

[68] M. J. Dolan, C. Englert and M. Spannowsky, “New Physics in LHC Higgs boson pair production,” Phys. Rev. D 87, no. 5, 055002 (2013) doi:10.1103/PhysRevD.87.055002 [arXiv:1210.8166 [hep-ph]].

[69] A. J. Barr, M. J. Dolan, C. Englert and M. Spannowsky, “Di-Higgs final states augMT2ed – selecting hh events at the high luminosity LHC,” Phys. Lett. B 728, 308 (2014) doi:10.1016/j.physletb.2013.12.011 [arXiv:1309.6318 [hep-ph]].

[70] M. J. Dolan, C. Englert, N. Greiner and M. Spannowsky, “Further on up the road: hhjj production at the LHC,” Phys. Rev. Lett. 112, 101802 (2014) doi:10.1103/PhysRevLett.112.101802 [arXiv:1310.1084 [hep-ph]].

[71] G. Panico, “Prospects for double Higgs production,” Frascati Phys. Ser. 61, 102 (2016) [arXiv:1512.04316 [hep-ph]].

[72] CMS Collaboration [CMS Collaboration], “Search for non-resonant Higgs boson pair production in the bbτ+τ− final state,” CMS-PAS-HIG-16-012.

[73] G. Aad et al. [ATLAS Collaboration], “Searches for Higgs boson pair production in the hh \rightarrow bbττ, γγWW∗, γγbb, bbbb channels with the ATLAS detector,” Phys. Rev. D 92, 092004 (2015) doi:10.1103/PhysRevD.92.092004 [arXiv:1509.04670 [hep-ex]].

[74] M. Reece, “Vacuum Instabilities with a Wrong-Sign Higgs-Gluon-Gluon Amplitude,” New J. Phys. 15, 043003 (2013) doi:10.1088/1367-2630/15/4/043003 [arXiv:1208.1765 [hep-ph]].