Phenomenology of Scalar Singlet Dark Matter in BSM Models

by

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Abstract

Approximately 85% of the mass of the universe is classified as dark matter, a type of matter outside the framework of the Standard Model. This thesis studies the phenomenology of three models that incorporate dark matter. Firstly, a singlet scalar dark matter candidate is added to two previously established models: the Higgs-like dilaton and the Georgi-Machacek models. The parameters of these models are constrained using experimental results such as dark matter direct detection measurements, Higgs signal strengths, Higgs invisible width and the measured relic abundance of dark matter in the universe. The notion of self-interacting dark matter, derived from measurements of galaxy cluster Abell 3827, is also studied for SSDM extensions of both the Standard Model and the Two Higgs Doublet model. The parameters of these models are in each case tightly constrained by experiment.

for Maria

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Preface

Included in this thesis is a review of the author’s work done as an M.Sc. student at Carleton University. All work contained within occurred only at Carleton University. Results of this work needed for analysis are included, but details identified as outside the scope of this thesis can be found in References [1–3].

Chapter 4 is entirely the work of the author, except for the Markov Chain Monte Carlo performed by Alejandro de la Puente. The fitting procedure was confirmed by comparison to a simple chi-squared fit to the ATLAS/CMS data.

The plots in Chapter 5 are accredited to Alexandre Poulin, as well are the details of the relic density scan.

All of the work in Chapter 6 was confirmed by the author’s independent calculations, but Figure 6.4 and the ideas of thermal freeze-in and implementing the Two Higgs doublet model are accredited to the collaborators of Reference [1].

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Chapter 1

Introduction

During the press release announcing the detection of gravitational waves, LIGO Scientific Collaboration (LSC) spokesperson Gabriela Gonzalez called the discovery “the beginning of a new era” in gravitational wave astronomy [4]. This is not the only area of fundamental physics that has experienced major discoveries over the past few years. The discovery of the Higgs Boson in 2012 confirmed the 50-year-old prediction of its existence and resulted in a Nobel prize [5]. In late 2015, the Large Hadron Collider (LHC) began operation at 13 TeV, and has already produced a signal exciting theorists, namely the hints of a 750 GeV diphoton excess [6].

Particle identification and interaction are neatly intertwined into a framework called the Standard Model (SM). Current experiment, fuelled by inconsistencies and gaps in this framework, searches for physics beyond the Standard Model (BSM). For example, there exists a question of scale in the SM: why do the masses have the values they do and why are there no intermediate natural scales between the electroweak and Planck scales? One of the most intriguing BSM probes lies in the fields of dark matter and dark energy - humbling detections indicating a lack of knowledge of 95% of the mass of the universe.

Theorists all over the world seek to remedy the ailments of the Standard Model. Some decide to start over and build a new model from scratch, such as string theory. Many take the more conservative approach and build on the existing foundation by introducing new particles and fields, where examples include supersymmetry, extending the Higgs sector, and adding extra dimensions. Of the two approaches this thesis will study the latter, considering select SM extensions that are expanded to include a singlet scalar dark matter (SSDM) candidate.

This thesis is structured as follows:

Chapter 2 is an overview of the Standard Model: its particles, forces, and mechanisms. It also introduces motivation for BSM physics, namely the hierarchy problem and dark matter. Chapter 3 examines the experimental measurements that will be used to constrain the proposed models.

The next three chapters examine three different models that incorporate dark matter candidates.

Chapter 4 considers a model with a Higgs-like dilaton and singlet scalar dark matter. Chapter 5 couples the well-established Georgi-Machacek model to an additional singlet serving as a dark matter candidate. In Chapter 6 self-interacting dark matter is examined, which stems from a dark matter halo offset in galaxy cluster Abell 3827. Results are summarized in Chapter 7.

Chapter 2

Beyond the Standard Model

2.1 Introduction to the Standard Model

The Standard Model (SM) of particle physics provides a framework for interactions of elementary particles. There are four known fundamental interactions: electromagnetic, strong, weak, and gravitational. Properties of the four forces, including relative strength and spin, are listed in Table 2.1.

Interaction	Mediator	Spin	Relative Strength
Gravitational	Graviton	2	10^{-42}
Weak	$W^\pm, Z$	1	10^{-13}
Electromagnetic	Photon	1	10^{-2}
Strong	Gluon	1	10

Table 2.1: Known elementary interactions and their mediators [7].

The Standard Model unifies the strong, weak, and electromagnetic interactions. Individual gravitons have not yet been discovered and there is yet to be a consistent theory of quantum gravity.

The known elementary particles are listed in Figure 2.1. The Dirac equation also predicts for each particle an antiparticle, or one of the same mass and spin but opposite charge. There are three leptons and three lepton neutrinos, of which only the electron is stable. It is believed that charge is preserved in all interactions, as is lepton number apart from neutrino oscillations.

u	c	t	g	H
d	s	b	γ	
e	μ	τ	Z	
ν_e	ν_μ	ν_τ	W	

Figure 2.1: Elementary particles of the Standard Model

Quarks are found only in bound states called hadrons. Hadrons can be classified further as either mesons (a quark and an antiquark) or baryons (three quarks). There is also evidence of four-quark systems, and recently bound states of five quarks, aptly named pentaquarks, may have been discovered [8]. All mesons are unstable, and the only stable baryon is the proton.

There are also four gauge bosons: gluons, which carry the strong interaction, photons, which carry the electromagnetic interaction, and the W and Z bosons, which together carry the weak interaction. Finally, the SM includes the Higgs boson, which serves to give the other particles mass. Details of these mechanisms can be found in the following sections.

2.1.1 Quantum Electrodynamics

Quantum electrodynamics (QED) is an extremely successful dynamical theory, mediating the interactions between photons and matter. Every interaction can be reduced to an elementary vertex of the form $\gamma \rightarrow qq$, where q is any charged particle. The interaction governed by QED of a photon with an electron is shown in Figure 2.2. The QED Lagrangian is

$$\mathcal{L}_{QED} = \bar{\psi}(i\not{D} - m)\psi - \frac{1}{4}(F_{\mu\nu})^2 \quad (2.1)$$

Figure 2.2: An elementary vertex for QED

where ψ is a field of spin-1/2 particles of mass m , and D_μ is the gauge covariant derivative

$$D_\mu = \partial_\mu + ieA_\mu(x). \quad (2.2)$$

Here, A_μ is the electromagnetic vector potential, such that the electromagnetic field tensor $F_{\mu\nu}$ is defined as [9]

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu. \quad (2.3)$$

The Lagrangian is invariant under the gauge transformations

$$\begin{aligned} A_\mu &\rightarrow A_\mu + \partial_\mu \lambda \\ \psi &\rightarrow e^{-ie\lambda} \psi. \end{aligned} \quad (2.4)$$

The electric charge e is typically written in terms of the fine structure constant, α :

$$\alpha = \frac{e^2}{4\pi\hbar c} \approx \frac{1}{137}. \quad (2.5)$$

QED is valid for all electromagnetic phenomena observed in nature, such as Compton scattering, Bhabha scattering, and electron-positron annihilation.

2.1.2 Quantum Chromodynamics

Quantum chromodynamics (QCD) describes the interactions of coloured particles: quarks and gluons. Quarks can either be red (r), green (g), or blue (b), and anti-quarks can be the corresponding anti-colours. Interactions proceed via the exchange of gluons, and is dependent on the strength of the strong coupling constant α_s .

The QCD portion of the Standard Model is as follows:

$$\mathcal{L}_{QCD} = -\frac{1}{4}F_{\mu\nu}^a F^{\mu\nu a} + i \sum_q \bar{\phi}_q^i \gamma^\mu (D_\mu)_{ij} \phi_q^j - \sum_q m_q \bar{\phi}_q^i \phi_{qi} \quad (2.6)$$

where

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a - g_s f_{abc} A_\mu^b A_\nu^c \quad (2.7)$$

and

$$(D_\mu)_{ij} = \delta_{ij} \partial_\mu + i g_s \sum_a \frac{\lambda_{ij}^a}{2} A_\mu^a. \quad (2.8)$$

Here, ϕ_q^i are the quark fields with colour i and flavour q , A_μ^a are the gluon fields, g_s is the strong coupling strength (ie $g_s = \sqrt{4\pi\alpha_S}$), f_{abc} are the SU(3) structure constants, and λ_{ij}^a are the Gell-Mann matrices [9].

2.1.3 The Electroweak Theory and the Higgs Mechanism

This section follows the combined notation of [10] and [11]. The Higgs Mechanism is introduced as the means to give fermions and gauge bosons mass. Consider the Glashow-Weinberg-Salam theory of electroweak interactions, containing three SU(2)_L gauge bosons W_μ^i ($i = 1, 2, 3$), and one U(1)_Y gauge boson B_μ . The Lagrangian is

$$\mathcal{L} = (D^\mu \Phi)^\dagger (D_\mu \Phi) - V(\Phi) - \frac{1}{4} W_{\mu\nu}^i W^{\mu\nu i} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \quad (2.9)$$

where D_μ is the covariant derivative associated to SU(2)_L × U(1)_Y

$$D_\mu = \partial_\mu + i g' \frac{Y}{2} B_\mu + i g \tau^i \cdot \vec{W}_\mu^i. \quad (2.10)$$

and Φ is the complex scalar doublet

$$\Phi = \begin{pmatrix} \Phi_+ \\ \Phi_0 \end{pmatrix}.$$

The simplest potential that is renormalizable and SU(2)_L invariant is

$$V = -\mu^2 |\Phi^\dagger \Phi| + \lambda (|\Phi^\dagger \Phi|)^2. \quad (2.11)$$

A ground state at the minimum of V is chosen, such that

$$|\Phi_0| = \frac{\mu}{\sqrt{2\lambda}}, \quad |\Phi_+| = 0. \quad (2.12)$$

The unitary gauge is now used to avoid explicit Goldstone bosons, and these degrees of freedom are transformed into longitudinal components of massive gauge bosons. The symmetry is broken spontaneously by introducing the vacuum expectation value (vev):

$$\langle \Phi \rangle = \sqrt{\frac{\mu^2}{2\lambda}} \equiv \frac{v}{\sqrt{2}} \quad (2.13)$$

It is then convenient to rewrite Φ as

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h \end{pmatrix}. \quad (2.14)$$

The Lagrangian in Equation 2.9 is invariant under the SU(2) transformation $\Phi \rightarrow -\Phi$, however the Lagrangian is not invariant under the transformation $h \rightarrow -h$. Therefore, the electroweak symmetry is spontaneously broken. After substitution of 2.14 into the Lagrangian, which is then expanded up to second order terms, four physical fields emerge: two charged fields, $W^\pm$, and two neutral fields, Z and γ :

$$W_\mu^\pm = \frac{1}{\sqrt{2}}(W_\mu^1 \mp iW_\mu^2) \quad (2.15)$$

$$Z^\mu = \frac{-g'B_\mu + gW_\mu^3}{\sqrt{g^2 + g'^2}} \quad (2.16)$$

$$A^\mu = \frac{gB_\mu + g'W_\mu^3}{\sqrt{g^2 + g'^2}}. \quad (2.17)$$

The Lagrangian can be rewritten in terms of these new fields

$$\mathcal{L} = (D^\mu \phi)^\dagger (D_\mu \phi) - \frac{1}{8}v^2(g^2(W^+)^2 + g^2(W^-)^2 + (g^2 + g'^2)Z_\mu^2 + 0 \cdot A_\mu^2) \quad (2.18)$$

from which the masses of the gauge bosons can be extracted. They are as follows:

$$m_W^2 = \frac{1}{4}g^2v^2 \quad (2.19)$$

$$m_Z^2 = \frac{1}{4}(g^2 + g'^2)v^2 \quad (2.20)$$

$$m_A = 0. \quad (2.21)$$

This also defines the weak mixing angle θ_W by way of the coupling constants and the electromagnetic strength e :

$$\begin{aligned} e &= g \sin \theta_W \\ e &= g' \cos \theta_W. \end{aligned} \quad (2.22)$$

As an example, consider the fermion masses. The Yukawa Lagrangian is as follows:

$$\mathcal{L}_{\text{Yuk}} = \Gamma_{mn}^u \bar{q}_{m,L} \tilde{\phi} u_{n,R} + \Gamma_{mn}^d \bar{q}_{m,L} \phi d_{n,R} + \Gamma_{mn}^e \bar{l}_{m,L} \phi e_{n,R} + \text{h.c.} \quad (2.23)$$

where the matrices Γ_{mn} contain the Yukawa couplings. Since the mass terms must be without hypercharge, there must be two representations of Higgs fields, ϕ and $\tilde{\phi}$, in order to give mass to the down type and up type fermions separately. In short, the Yukawa Lagrangian can be rewritten for the first generation:

$$\mathcal{L}_{\text{Yuk}} = f_e \bar{l}_L \phi e_R + f_u \bar{q}_L \tilde{\phi} u_R + f_d \bar{q}_L \phi d_R + \text{h.c.} \quad (2.24)$$

which is, as before, in the unitary gauge. Substituting the fields into the Lagrangian gives the convenient form

$$\mathcal{L}_{\text{Yuk}} = \frac{f_e v}{\sqrt{2}} (\bar{e}_L e_R + \bar{e}_R e_L) + \frac{f_u v}{\sqrt{2}} (\bar{u}_L u_R + \bar{u}_R u_L) + \frac{f_d v}{\sqrt{2}} (\bar{d}_L d_R + \bar{d}_R d_L). \quad (2.25)$$

This gives an expression for the mass terms:

$$m_i = -\frac{f_i v}{\sqrt{2}}, \quad i = e, u, d. \quad (2.26)$$

2.2 Motivation for BSM Physics

2.2.1 The Hierarchy Problem

The Higgs mechanism is the driving force behind electroweak symmetry breaking (EWSB), but theorists hesitate to accept it as a complete theory. It does not explain the physical value of the vev ($v=246$ GeV), or why fermions have the masses they do. The most serious problem, however, is that first order loop corrections to the Higgs mass are quadratically divergent [12]. If we define $m_H^2 = m_{H_0}^2 + \delta m_H^2$, the loop correction at the one-loop level reads

$$\delta m_H^2 = \frac{G_F \Lambda^2}{4\sqrt{2}\pi^2} (6m_W^2 + 3m_Z^2 + m_H^2 - 12m_t^2) + \dots \quad (2.27)$$

where Λ^2 is the ultraviolet (UV) scale boundary of the SM. The only natural scale found in the SM is the Planck scale, or $\Lambda = M_{\text{Planck}} \approx 10^{19}$ GeV. In order to produce results comparable to the order of the Higgs mass, fine-tuned counterterms must be added to cancel the divergences.

Alternatively, we can assume that the correction is on the order of the observed 125 GeV Higgs boson, meaning the scale Λ is less than the Planck scale. This implies that Λ is on the order of 1 TeV, so one would expect to see new physics at about this scale. The little hierarchy problem arises due to the fact that much of the new physics one expects to see is ruled out by experimental searches [13].

2.2.2 Dark Matter

There is overwhelming evidence that the matter observed in the Standard Model makes up only approximately 4% of the mass of the universe (see Figure 2.3). The remainder is dark energy ($\sim 73\%$) and dark matter ($\sim 23\%$) [14]. Dark energy is a repulsive force aiding in the acceleration of the expansion of the universe. A number of solutions have been proposed to this mystery, such as alterations to the cosmological constant and additional scalar fields. This thesis does not concern itself with dark energy, but instead focuses on dark matter.

The initial evidence for dark matter is found in a paper written by Fritz Zwicky in 1937 [15]. The argument states that estimates of nebulae masses based on their luminosities and rotation curves are inconsistent. Measurements of the Coma cluster

Figure 2.3: The composition of the universe, according to percentage by mass.

show that its galaxies cannot be gravitationally bound to the cluster by the force of the visible matter alone. Given its radical conclusion, Zwicky’s paper only received about 10 citations in the first forty years after its publication [16].

The notion of dark matter was revived with Vera Rubin in the 1970’s. She found that the rotational speed of a galaxy remains constant past the visible edge of the galaxy, where classical physics expects it to decrease. This indicates that most of the mass in a galaxy is contained in a spherical “dark halo” around the edge of the galaxy [17]. There exists additional evidence for the existence of dark matter, for example, gravitational lensing, the Cosmic Microwave background, and the large-scale structure of the universe [14].

Numerous DM particle candidates have been proposed. Most are Standard Model additions, coming into effect at the electroweak scale. Some hinge off other models, such as supersymmetry or string theory. However, as of yet, a dark matter particle has not been observed by any experiment- the best one can do is interpret what its absence reveals about its nature.

It is known that a dark matter candidate must be weakly interacting but still able to exert a gravitational pull. It does not reflect, emit or absorb light. The particle has a known energy density and galactic velocity, and it freezes out its relic abundance. There are very few limits on the mass of a dark matter particle, but a phenomena known as the “WIMP miracle” dictates that the self annihilation cross section required to obtain the correct abundance of dark matter today suggests a particle in the GeV mass range, which can be probed by colliders [14].

Chapter 3

Experimental Constraints on Models

To explore possibilities beyond the Standard Model, new models are proposed to remedy one or more of the previously posed issues with the current framework. However, it is preferred that these extensions be testable with existing experimental data, or else have the potential to guide future experiments. This chapter introduces the phenomenological tools used in this thesis, used to constrain parameters of the models in the subsequent sections.

3.1 Higgs Signal Strengths

The signal strengths are defined as:

$$\mu_i = \frac{\sigma_i}{(\sigma_i)_{SM}} \quad \text{and} \quad \mu^f = \frac{BR^f}{(BR^f)_{SM}} \quad (3.1)$$

where the σ_i are Higgs boson production cross sections for initial states i and the BR^f are decay branching ratios for final states f . The SM subscripts indicate the SM predictions for the cross sections and branching ratios, such that the expected signal strength values are $\mu_i = \mu^f = 1$.

The ATLAS/CMS collaborations provide numerical results for the values of μ_i and μ^f [18]. They present a combined fit to both the $\sqrt{s} = 7$ and $\sqrt{s} = 8$ data, assuming that the signal strength is the same between the two energies.

The main production modes studied are gluon fusion (ggF), vector boson fusion (VBF), and associated production with vector bosons (VH) or a pair of top quarks (ttH). The decay modes considered are decays to Z bosons, W bosons, photons, τ leptons, bottom quarks, and muons. In general, we assume the measured particle

mass of 125.09 GeV has the properties of the SM Higgs. A comparison can be done between the experimental results and the values predicted by a given model, often invoking a χ^2 fit or other fitting procedure. Since the signal strength results are applied slightly differently between models in this thesis, procedural details will be included in their respective chapters.

3.2 Higgs Invisible Width

The process $h \rightarrow SS$ (Figure 3.1) where S is the dark matter candidate provides a constraint on the dark matter in the lower mass range $m_S < m_h/2$. This process is invisible to the detector apart from large missing energy. The invisible width of a

Figure 3.1: Feynman vertex for invisible Higgs decay

Higgs boson, or the probability of a Higgs decaying invisibly to two DM candidates, is in many cases a simple calculation given the Higgs-DM coupling. It can be written in the generic form

$$\Gamma(h \rightarrow SS) = \frac{g_{hSS}^2}{32m_h^2\pi} \sqrt{m_h^2 - 4m_S^2} \quad (3.2)$$

where ig_{hSS} is the Feynman coupling for the hSS vertex. Experimentally, we are interested in the branching ratio, which also depends on the total width of the Higgs boson to standard model particles $\Gamma(h \rightarrow SM)$:

$$BR(h \rightarrow invis.) = \frac{\Gamma(h \rightarrow SS)}{\Gamma(h \rightarrow SM) + \Gamma(h \rightarrow SS)}. \quad (3.3)$$

The ATLAS experiment bounds the invisible Higgs branching ratio from above.

This constraint comes from direct searches for Higgs decays in dilepton or dijet final states with missing transverse momentum at $\sqrt{s}=8$ TeV. There are three mechanisms considered, combined statistically to give the final result:

- the Higgs boson is produced by itself by VBF, and decays invisibly. The signature of this process is two jets with a large separation in pseudorapidity, and large missing transverse momentum.
- the Higgs boson is produced with a Z boson, which decays to leptons while the Higgs decays invisibly. The signature of this process is opposite signed, same flavoured leptons and large missing transverse momentum.
- the Higgs boson is produced with a vector boson (Z or W), which decays to jets while the Higgs decays invisibly. The signature of this process is two jets where the invariant mass of the jets is consistent with the boson mass, and large transverse momentum

The model independent observed upper limit at the 95% CL is $BR(h \rightarrow invis.) < 0.28$ [19].

3.3 Dark Matter Relic Abundance

In the early universe, there was an abundance of high mass, high energy particles created and annihilated in thermal equilibrium. However, due to the expansion and cooling of the universe, the equilibrium was lost and the number density became constant. This is known as thermal “freeze out” and occurs when the Hubble expansion rate is equal to the reaction rate, which can be found using the Boltzmann equation for the comoving particle density [20]

$$\frac{dn_S}{dt} + 3Hn_S = -\langle\sigma v\rangle(n_S^2 - (n_S^{eq})^2), \quad (3.4)$$

where H is the Hubble parameter, σ is the DM annihilation cross section, v is the relative velocity, n_S is the number density of S , and n_S^{eq} is the equilibrium number density of S . The bracket notation $\langle\sigma v\rangle$ indicates the thermal average of the enclosed cross section.

The freezeout temperature, T_{FO} , can be found by analytically solving the transcendental equation

$$\frac{m_S}{T_{FO}} \approx \ln \left(\frac{0.038 g_S m_S M_{pl}}{g^{\frac{1}{2}} \left(\frac{m_S}{T_{FO}} \right)^{\frac{1}{2}}} \langle \sigma v \rangle \right) \Big|_{T=T_{FO}} \quad (3.5)$$

where the thermalized cross section is given by

$$\langle \sigma v \rangle = \frac{1}{8m_S^4 T K_2^2 \left(\frac{m_S}{T} \right)} \int_{4m_S^2}^{\infty} ds \sigma(s) (s - 4m_S^2) s^{\frac{1}{2}} K_1 \left(\frac{s^{\frac{1}{2}}}{T} \right). \quad (3.6)$$

Here, $K_1(z)$ and $K_2(z)$ are modified Bessel functions of the second kind, and M_{pl} is the Planck mass.

The value of $x = m_S/T_{FO}$ is typically in the range $\{4, 40\}$. The thermalized cross section can then be evaluated at the freeze out temperature and inserted into the expression for relic abundance, which gives

$$\Omega h^2 = 1.65 \times 10^{-10} \left(\frac{1}{\langle \sigma v \rangle} \text{Log} \left(\frac{0.038 g_S m_S M_{pl}}{\sqrt{g}} \langle \sigma v \rangle \right) \right) \Big|_{T=T_{FO}} \quad (3.7)$$

where g_S is the number of degrees of freedom for the dark matter candidate (equals 1 for SSDM) and $g = 86$ is the number of relativistic degrees of freedom at the freeze-out temperature. Here, Ω is the critical energy density of the universe and h is the Hubble parameter.

The Planck Collaboration measured the value to be $\Omega h^2 = 0.1199 \pm 0.0027$ by mapping anisotropies of the cosmic microwave background (CMB) [21]. The predicted value of the relic density for a given model must be less or equal to the experimental value, as a deficiency can potentially be accounted for by adding additional DM candidates.

3.4 Dark Matter Direct Detection

If a dark matter candidate comes into close enough proximity to a nucleon, a momentum transfer may occur by exchange of a Higgs boson (see Figure 3.2). The nuclear recoil, however, is small: a WIMP with a mass of 10-1000 GeV/c² generates a recoil in the range of 1-100 keV [22].

Figure 3.2: The Feynman diagram for DM direct detection through DM/nucleon scattering.

The general form of the direct detection cross section is

$$\sigma = \frac{f_N^2 g_{hSS}^2 m_N^4}{4\pi(m_N + m_S)^2 m_h^4 v^2} \quad (3.8)$$

where f_N is the Feynman coupling for the Higgs-nucleon vertex. The experimental values used for comparison are from the XENON and LUX experiments [23, 24]. Both are underground liquid xenon detectors, generating and recording scintillation light from reactions. A time lapse occurs between the initial reaction and a secondary reaction generated by free electrons, making it possible to localize events in three coordinates. The detectors are placed deep underground to minimize background from cosmic rays [25, 26]. Projected limits are also included from the DEAP 3600 experiment [27].

In order to apply direct detection searches, a number of assumptions must be made:

- the halo of DM particles is smooth, and can be described by Maxwell velocity distributions
- the galactic disk rotates, therefore there exists an annual modulation due to the Earth’s orbital motion to compensate for the so-called “WIMP wind”
- the nucleus is a hard sphere, described by Helm form factor
- the scattering is elastic

It is not the recoil energy that is measured directly, but the recoil energy distribution, creating a DM density profile [\[28\]](#).

Chapter 4

Higgs-like Dilaton with Scalar Singlet DM

4.1 Background and Motivation

Classically, the Standard Model is scale invariant apart from the Higgs mass, which can be generated naturally at the electroweak scale by a soft breaking of the scale invariance. However, breaking non-linear scale invariance spontaneously implies the existence of a Goldstone boson, the dilaton. If the dimensionful operators are near marginal at the breaking scale, the dilaton can be naturally light.

The model studied in this chapter considers a Higgs-like dilaton, that is, a dilaton with properties of the 125 GeV Higgs boson [29]. It is further extended by a singlet scalar dark matter candidate. We take the simplest addition, a real scalar field, which only interacts with standard model particles through the dilaton portal.

In this 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. The potential is [2]

$$\begin{aligned} V(\bar{\chi}, S) \approx & \frac{1}{2}m_\sigma^2\bar{\chi}^2 - \frac{\lambda}{3!}\frac{m_\sigma^2}{f}\bar{\chi}^3 + \frac{\bar{\chi}}{f}\sum_i(1+\gamma_{\psi_i})m_{\psi_i}\bar{\psi}_i\psi_i + \left(\frac{2\bar{\chi}}{f} + \frac{\bar{\chi}^2}{f^2}\right) \\ & \times \left[m_W^2 W^{+\mu}W_\mu^- + \frac{1}{2}m_Z^2 Z^\mu Z_\mu \right] + \frac{\alpha_{\text{EM}}}{8\pi f}c_{\gamma\gamma}\sigma F_{\mu\nu}F^{\mu\nu} + \frac{\alpha_S}{8\pi f}c_{gg}\sigma G_{\mu\nu}G^{\mu\nu} \\ & + \frac{1}{2}\bar{m}_S^2 S^2 + \lambda_S S^4 \end{aligned} \quad (4.1)$$

where the dilaton (σ) is parametrized with the non-linear realization as $\chi = fe^{\sigma/f}$, and dilaton is expanded about its vev as $\bar{\chi} = \chi - f$, while $\bar{m}_S^2 = m_S^2 e^{2\sigma/f}$. The

exponentials are expanded to leading orders to obtain the following parametrization for the potential:

$$\begin{aligned}
 V(\sigma, S) \approx & \frac{1}{2}m_\sigma^2\sigma^2 + \epsilon\frac{m_\sigma^2}{v}\left(\frac{1}{2} - \frac{1}{6}\lambda\right)\sigma^3 + \frac{\sigma}{f}\sum_i(1 + \gamma_{\psi_i})m_{\psi_i}\bar{\psi}_i\psi_i + \left(\frac{2\sigma}{f} + \frac{\sigma^2}{f^2}\right) \\
 & \left[m_W^2W^{+\mu}W_\mu^- + \frac{1}{2}m_Z^2Z^\mu Z_\mu\right] + \frac{\alpha_{\text{EM}}}{8\pi f}c_{\gamma\gamma}\sigma F_{\mu\nu}F^{\mu\nu} + \frac{\alpha_S}{8\pi f}c_{gg}\sigma G_{\mu\nu}G^{\mu\nu} \\
 & + \frac{1}{2}m_S^2\left(1 + 2\epsilon\frac{\sigma}{v} + 2\epsilon^2\frac{\sigma^2}{v^2}\right)S^2 + \lambda_S S^4,
 \end{aligned} \tag{4.2}$$

where f denotes the scale where the conformal symmetry is spontaneously broken, v , the EWSB scale and $\epsilon = v/f$. The parameter λ is related to the trilinear dilaton coupling. A Z_2 symmetry is imposed under which S is odd while all other fields are even, ensuring that the dark matter candidate is stable. The parameter λ_S is a free parameter unable to be constrained by experiment at this time.

To summarize, this theory depends on six independent parameters: m_S , λ , λ_S , ϵ , c_{gg} and $c_{\gamma\gamma}$ where

$$c_{gg} = (b_{IR}^3 - b_{UV}^3) \tag{4.3}$$

$$c_{\gamma\gamma} = (b_{IR}^{EM} - b_{UV}^{EM}). \tag{4.4}$$

Here, b_{UV} and b_{IR} represent the ultraviolet and infrared contributions to the running of the EM and strong gauge couplings respectively at a scale above conformal symmetry breaking. At high energies, the underlying theory is not considered and the SM loop contribution to c_{gg} and $c_{\gamma\gamma}$ are absorbed in the definition.

4.2 Phenomenology

4.2.1 Higgs Signal Strengths

The latest data at 8 TeV from the ATLAS and CMS collaborations shown in Tables 4.1 and 4.2 is used to constrain the dilaton with measured Higgs couplings.

	ggF	VBF	WH	ZH	ttH
experimental	1.03	1.18	0.89	0.79	2.3
error	0.15	0.24	0.39	0.37	0.65
theory	$\left(\frac{\epsilon c_{gg}}{c_{gg}^{SM}}\right)^2$	ϵ^2	ϵ^2	ϵ^2	ϵ^2

Table 4.1: Signal strengths for Higgs boson production channels according to ATLAS/CMS collab. [30], including predictions from the model in terms of ϵ and

$c_{\gamma\gamma}$

	$\gamma\gamma$	ZZ	WW	$\tau\tau$	bb
experimental	1.14	1.29	1.09	1.11	0.70
error	0.19	0.25	0.17	0.23	0.28
theory	$\left(\frac{\epsilon c_{\gamma\gamma}}{c_{\gamma\gamma}^{SM}}\right)^2$	ϵ^2	ϵ^2	ϵ^2	ϵ^2

Table 4.2: Signal strengths for Higgs boson decay channels according to ATLAS/CMS collab. [30], including predictions from the model in terms of ϵ and

c_{gg} .

The total width of the dilaton is

$$\Gamma_{tot} = \epsilon^2 \left(\sum_f \Gamma_f^{SM} + \Gamma_{WW}^{SM} + \Gamma_{ZZ}^{SM} + \Gamma_{\gamma Z}^{SM} + \left(\frac{c_{gg}}{c_{gg}^{SM}} \right)^2 \Gamma_{gg}^{SM} + \left(\frac{c_{\gamma\gamma}}{c_{\gamma\gamma}^{SM}} \right)^2 \Gamma_{\gamma\gamma}^{SM} \right) + \Gamma_{inv} \quad (4.5)$$

where Γ_{inv} is the invisible width of the process $\sigma \rightarrow SS$. The fit is done by a Markov Chain Monte Carlo (MCMC) algorithm, using a Metropolis-Hastings algorithm with simulated annealing [2]. There are 7 degrees of freedom derived from the 10 Higgs observables and the three parameters appearing in the dilaton couplings: ϵ , c_{gg} , and $c_{\gamma\gamma}$. The target distribution is a Gaussian likelihood function:

$$\mathcal{L} = \prod_i \exp \left[-\frac{(x_i - \bar{x}_i)^2}{2\sigma_i^2} \right] \quad (4.6)$$

where x_i are the theoretical predictions, $\bar{x}_i$ are the experimental observables, and σ_i are the errors on these observables.

There also exists a fit from electroweak precision data (EWPD) that places a limit

on the value of ϵ , by looking specifically at non-standard Higgs couplings. This places a lower bound on ϵ of about 0.98 [31].

First, we test the effectiveness of the fit by comparing with fits of a Higgs-like dilaton without the dark matter candidate. We note that our fit gives similar large-epsilon behaviour as that of previous work [29].

Figure 4.1: Parameter space of c_{gg} and $c_{\gamma\gamma}$, where the allowed region is $\chi^2 < 5.9$.

The dark matter candidate is then included in the fitting procedure. The fit is further extended by including another observable related to the dark matter mass: the total width of the Higgs boson. The bound derived in Reference [32] is $6.7^{+7.7}_{-2.9}$ MeV. This agrees with a CMS study that places an upper bound of 22 MeV at 95% CL [33]. Figure 4.1 shows the results of the MCMC. The 95% confidence interval for seven degrees of freedom corresponds to a χ^2 of 5.9. The parameter space at 95% confidence is $0.5 \leq |c_{gg}| \leq 1$, while $4 \leq |c_{\gamma\gamma}| \leq 8$. The fourfold symmetry comes from the quadratic dependence on c_{gg} and $c_{\gamma\gamma}$. These compare well to the Standard Model values of $c_{gg}^{SM} = 0.67$ and $c_{\gamma\gamma}^{SM} = -6.5$. It is notable that some regions of parameter space, where m_S is less than half the dilaton mass, can result in a large contribution to invisible decays, discussed in the next section.

4.2.2 Higgs Invisible Width

The decay of the dilaton to two dark matter candidates is kinematically accessible when $m_S < m_\sigma/2$. In this model, the invisible decay width of the dilaton is given by

$$\Gamma(\sigma \rightarrow SS) = \frac{m_S^4 \epsilon^2}{8\pi m_\sigma v^2} \sqrt{1 - 4 \frac{m_S^2}{m_\sigma^2}}, \quad (4.7)$$

where the branching ratio is analogous to Equation 3.3:

$$BR(\sigma \rightarrow inv.) = \frac{\Gamma(\sigma \rightarrow SS)}{\Gamma_{tot}}. \quad (4.8)$$

with Γ_{tot} defined in equation 4.5.

In Figure 4.2 the branching ratio is plotted against the dark matter mass, including the experimental upper bound from ATLAS/CMS of 0.28 [19].

Figure 4.2: The branching ratio for the dilaton decay to invisible particles, with $c_{gg} = 0.67$ and $c_{\gamma\gamma} = -6.5$. The horizontal line is the ATLAS upper bound [19].

The Standard Model values were used in Figure 4.2 for c_{gg} and $c_{\gamma\gamma}$. Given the bounds from signal strengths found in the previous section, the constraint from the invisible width does not greatly depend on c_{gg} , and even less so on $c_{\gamma\gamma}$ given the mass of the DM. This is illustrated for the former case in Figure 4.3, where contours of

the branching ratio is shown for varying m_S and c_{gg} . Selecting a value of m_S gives a number for the branching ratio almost completely independent of the value of c_{gg} .

Figure 4.3: A contour plot of the m_S and c_{gg} . The contours are the values for the branching ratio, with an experimental cutoff of 0.28. The value of c_{gg} is clearly not sensitive to the value of m_S – no new information is gained concerning these parameters.

4.2.3 Dark Matter Relic Abundance

In order to calculate the relic abundance, all reactions of the form $SS \rightarrow xx$, where x is a standard model particle or the dilaton, must be taken into account for the thermalized cross section.

Using the Feynman rules in Appendix A, each annihilation cross section can be calculated individually. We start with the processes with SM particles in the final

Figure 4.4: The Feynman diagrams for the processes $SS \rightarrow xx$, where x is a SM particle.

state, illustrated by Figure 4.4. For the process $SS \rightarrow ZZ$,

$$\sigma_{ZZ}(s) = \frac{1}{8\pi} \frac{\sqrt{s - 4m_Z^2}}{\sqrt{s - 4m_S^2}} \left(\frac{m_S \epsilon}{v} \right)^4 \frac{s}{(s - m_\sigma^2)^2 + \Gamma_\sigma^2 m_\sigma^2} \left(1 + \frac{12m_Z^4}{s^2} - \frac{4m_Z^2}{s} \right) \quad (4.9)$$

and similarly for W bosons in the final state,

$$\sigma_{WW}(s) = \frac{1}{4\pi} \frac{\sqrt{s - 4m_W^2}}{\sqrt{s - 4m_S^2}} \left(\frac{m_S \epsilon}{v} \right)^4 \frac{s}{(s - m_\sigma^2)^2 + \Gamma_\sigma^2 m_\sigma^2} \left(1 + \frac{12m_W^4}{s^2} - \frac{4m_W^2}{s} \right). \quad (4.10)$$

The annihilation cross section into photons is

$$\sigma_{\gamma\gamma}(s) = \frac{(c_{\gamma\gamma} \alpha_{EM})^2}{8\pi^3} \frac{s}{(s - m_\sigma^2)^2 - \Gamma_\sigma^2 m_\sigma^2} \left(\frac{m_S \epsilon}{v} \right)^4 \left(1 - \frac{4m_S^2}{s} \right)^{-\frac{1}{2}}. \quad (4.11)$$

The annihilation cross section into gluons is

$$\sigma_{gg}(s) = \frac{(c_{gg}\alpha_S)^2}{\pi^3} \frac{s}{(s - m_\sigma^2)^2 - \Gamma_\sigma^2 m_\sigma^2} \left(\frac{m_S \epsilon}{v}\right)^4 \left(1 - \frac{4m_S^2}{s}\right)^{-\frac{1}{2}} \quad (4.12)$$

The annihilation cross section to fermions is

$$\sigma_{f\bar{f}}(s) = \frac{2m_f^2 N_C}{\pi s} \left(\frac{m_S \epsilon}{v}\right)^4 \frac{1}{(s - m_\sigma^2)^2 + \Gamma_\sigma^2 m_\sigma^2} \frac{(s - 4m_f^2)^{\frac{3}{2}}}{\sqrt{s - 4m_S^2}} \quad (4.13)$$

where $N_C = 3$ for quarks, else $N_C = 1$.

Figure 4.5: The Feynman diagrams for the process $SS \rightarrow \sigma\sigma$.

Finally, the annihilation cross section to dilatons must include four different Feynman diagrams, shown in Figure 4.5. The result is given by

$$\begin{aligned}
\sigma_{\sigma\sigma}(s) = & \frac{\sqrt{s-4m_\sigma^2}}{4\pi s\sqrt{s-4m_S^2}} \left(\frac{m_S\epsilon}{v}\right)^4 \left(2 - \frac{m_S^4}{m_\sigma^2 m_S^2 - \frac{1}{4}(m_\sigma^4 - sm_S^2)} + \right. \\
& \frac{3.38m_\sigma^4 - 2m_\sigma^2(m_\sigma^2 - s)(3 + \lambda_p)}{-m_S^4 + (m_\sigma^2 - s)^2 + m_\sigma^2\Gamma_\sigma^2} + \frac{m_S^4 \arctan\left(\frac{\sqrt{4m_\sigma^2 - s}\sqrt{s-4m_S^2}}{2m_\sigma^2 - s}\right)}{2\sqrt{1 - \frac{4m_\sigma^2}{s}}\sqrt{1 - \frac{4m_S^2}{s}}(m_\sigma^2 - s/2)s} + \\
& \frac{m_\sigma^2 m_S^2 (m_\sigma^2 - s)(3 + \lambda_p)}{s(-m_S^4 + (m_\sigma^2 - s)^2 + m_\sigma^2\Gamma_\sigma^2)} \log\left(\frac{2m_\sigma^2 + s\left(-1 + \sqrt{1 - \frac{4m_\sigma^2}{s}}\sqrt{1 - \frac{4m_S^2}{s}}\right)}{2m_\sigma^2 - s}\right) - \\
& \left. \frac{8m_S^2 \log\left(\frac{2m_\sigma^2 + s\left(-1 + \sqrt{1 - \frac{4m_\sigma^2}{s}}\sqrt{1 - \frac{4m_S^2}{s}}\right)}{2m_\sigma^2 - s}\right)}{\sqrt{1 - \frac{4m_\sigma^2}{s}}\sqrt{1 - \frac{4m_S^2}{s}}s} \right) \quad (4.14)
\end{aligned}$$

The sum of the six processes is the total cross section, $\sigma(s)$. This can be inserted into Equation 3.6 to find the thermalized cross section, which is then used in Equation 3.5 to calculate the freeze-out temperature. Finally, evaluating the thermalized cross section at the freeze-out temperature and inserting into Equation 3.7 gives the relic abundance. This is compared with the observed relic abundance, 0.1199 ± 0.0027 .

Figure 4.6 plots the present-day relic abundance against the DM mass, including the aforementioned experimental upper bound. The values for the relic abundance are checked using micrOMEGAs [34]. Included are the curves for $\lambda = 3$, corresponding to SM Higgs self coupling, and $\lambda = 4\pi$, where strong interactions dominate. The model gives the observed DM relic abundance when mass $m_S \approx 50$, and rapidly becomes suppressed at masses close to the resonant region, $m_\sigma/2$.

The singlet scalar dark matter candidate with a Higgs-like dilaton clearly cannot by itself reproduce the observed dark matter relic abundance. It can, however, give insight into the nature of the Higgs self-coupling. When dark matter annihilations into dilatons are kinematically allowed, the self coupling can dictate how dominant the decay is and therefore determine how large the DM abundance in that mass region

Figure 4.6: The DM relic abundance as a function of the DM mass. The $\lambda = 3$ case corresponds to the black line, and the $\lambda = 4\pi$ case corresponds to the lower grey line. The horizontal grey line is the experimental value [21].

can be.

4.2.4 Dark Matter Direct Detection

We introduce the effective nucleon dilaton coupling f_N , which can be written as a sum of the contribution from each quark coupling:

$$f_N = \sum_q \langle N | \bar{q}q | N \rangle \frac{m_q}{m_N} = \sum_q f_N^q \quad (4.15)$$

The contributions from u, d, and s are related to the light quark matrix elements, and values are taken from Reference [35] as follows:

$$\begin{aligned} f_N^u &= 0.024 \\ f_N^d &= 0.033 \\ f_N^s &= 0.042. \end{aligned} \quad (4.16)$$

These values along with the heavy quark contributions give the Standard Model value $f_N^{SM} = 0.30$, which agrees with Reference [35]. Since this model subsumes the SM heavy quark contribution into c_{gg} , there must be a correction to f_N by the gluon contribution

$$f_N = \epsilon(f^u + f^d + f^s + f^G) = \epsilon \left(0.099 + 0.201 \frac{c_{gg}}{c_{gg}^{SM}} \right). \quad (4.17)$$

This expression is used to calculate the direct detection cross section:

$$\sigma(NS \rightarrow NS) = \left(\frac{m_N m_S}{m_\sigma v} \right)^4 \frac{\epsilon^2 f_N^2}{\pi(m_N + m_S)^2}, \quad (4.18)$$

where m_N is the nucleon mass (938.95 MeV).

Figure 4.7: The DM direct detection cross section as a function of the DM mass. The cases for $\lambda = 3$ and $\lambda = 4\pi$ are shown in black and dark grey, respectively. The experimental limits are, starting from bottom: DEAP projection (black), LUX (grey), and XENON (light grey).

The direct detection cross section is shown in Figure 4.7 as a function of m_S . Included are the experimental limits from XENON [23], LUX [24], and the projected limits from DEAP 3600 [27], all of which were extracted using Reference [36]. Note that although negative values of c_{gg} are allowed from the signal strength fit as shown in Figure 4.1, the points included were obtained using $c_{gg} > 0$. It is interesting to

note that some of the points allowed by LUX and XENON would be ruled out by the DEAP expected limit.

4.3 Summary

A model is introduced that treats the SM physical Higgs boson as a pseudo-Goldstone boson of spontaneous conformal symmetry breaking, the dilaton, and introduces a singlet scalar dark matter candidate. The relevant model parameters are fitted to the measured Higgs signal strengths, which restricts c_{gg} and $c_{\gamma\gamma}$ to a parameter space consistent with the SM values. The relic abundance and direct detection constraints limit the possible DM masses to a small range, but is increased to the range 125-300 GeV when the self coupling is changed to be that of a strongly interacting composite region. 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.

Chapter 5

Georgi-Machacek Model with Scalar DM

5.1 Background and Motivation

A natural step for BSM model building is to start with the existing SM framework and add extra doublets or triplets. The Georgi-Machacek (GM) Model is a well-established benchmark model that features an extended Higgs sector and preserves custodial SU(2) symmetry. This Higgs sector consists of a doublet, real triplet, and complex triplet [37]. This chapter studies the GM model extended by a singlet scalar dark matter candidate (SSDM), S. It is assumed for simplicity that S accounts for all of the observed dark matter relic abundance via the standard thermal freeze-out mechanism. This work is currently in preparation for publication [3].

The Georgi-Machacek Model

The scalar sector of the GM model¹ consists of the usual complex Higgs doublet (ϕ^+, ϕ^0) , a real triplet (ξ^+, ξ^0, ξ^-) , and a complex triplet $(\chi^{++}, \chi^+, \chi^0)$. The doublet functions as in the SM Higgs mechanism. This doublet can be written in the form of a bidoublet Φ and the triplets are combined to form a bitriplet X :

$$\Phi = \begin{pmatrix} \phi^{0*} & \phi^+ \\ -\phi^{+*} & \phi^0 \end{pmatrix}, \quad X = \begin{pmatrix} \chi^{0*} & \xi^+ & \chi^{++} \\ -\chi^{+*} & \xi^0 & \chi^+ \\ \chi^{++*} & -\xi^{+*} & \chi^0 \end{pmatrix}. \quad (5.1)$$

The vevs are defined by $\langle \Phi \rangle = \frac{v_\phi}{\sqrt{2}} I_{2 \times 2}$ and $\langle X \rangle = v_\chi I_{3 \times 3}$, where I is the unit

¹For a more detailed introduction to the Georgi-Machacek model, see Reference [38].

matrix and the Fermi constant G_F constrains

$$v_\phi^2 + 8v_\chi^2 \equiv v^2 = \frac{1}{\sqrt{2}G_F} \approx (246 \text{ GeV})^2. \quad (5.2)$$

The most general gauge-invariant scalar potential involving these fields and the real singlet S , while conserving the global $SU(2)_L \times SU(2)_R$ and the Z_2 symmetry $S \rightarrow -S$, is given by

$$\begin{aligned} V(\Phi, X) = & \frac{\mu_2^2}{2} \text{Tr}(\Phi^\dagger \Phi) + \frac{\mu_3^2}{2} \text{Tr}(X^\dagger X) + \lambda_1 [\text{Tr}(\Phi^\dagger \Phi)]^2 + \lambda_2 \text{Tr}(\Phi^\dagger \Phi) \text{Tr}(X^\dagger X) \\ & + \lambda_3 \text{Tr}(X^\dagger X X^\dagger X) + \lambda_4 [\text{Tr}(X^\dagger X)]^2 - \lambda_5 \text{Tr}(\Phi^\dagger \tau^a \Phi \tau^b) \text{Tr}(X^\dagger t^a X t^b) \\ & - M_1 \text{Tr}(\Phi^\dagger \tau^a \Phi \tau^b) (UXU^\dagger)_{ab} - M_2 \text{Tr}(X^\dagger t^a X t^b) (UXU^\dagger)_{ab} \\ & + \frac{\mu_S^2}{2} S^2 + \lambda_a S^2 \text{Tr}(\Phi^\dagger \Phi) + \lambda_b S^2 \text{Tr}(X^\dagger X) + \lambda_S S^4. \end{aligned} \quad (5.3)$$

The last line of the potential in Equation 5.3 contains the terms for the singlet scalar DM, while the rest is from the GM model. The physical fields that emerge can be organized by their transformation properties under the custodial $SU(2)$ symmetry into a fiveplet, a triplet, and three singlets, one of which is S . The fiveplet and triplet states are given by (plus the hermitian conjugate where appropriate)

$$\begin{aligned} H_5^{++} &= \chi^{++}, & H_5^+ &= \frac{(\chi^+ - \xi^+)}{\sqrt{2}}, & H_5^0 &= \sqrt{\frac{2}{3}} \xi^0 - \sqrt{\frac{1}{3}} \chi^{0,r}, \\ H_3^+ &= -s_H \phi^+ + c_H \frac{(\chi^+ + \xi^+)}{\sqrt{2}}, & H_3^0 &= -s_H \phi^{0,i} + c_H \chi^{0,i}, \end{aligned} \quad (5.4)$$

where the vevs are parameterized by

$$c_H \equiv \cos \theta_H = \frac{v_\phi}{v}, \quad s_H \equiv \sin \theta_H = \frac{2\sqrt{2} v_\chi}{v}, \quad (5.5)$$

and the neutral fields are decomposed into real and imaginary parts according to

$$\phi^0 \rightarrow \frac{v_\phi}{\sqrt{2}} + \frac{\phi^{0,r} + i\phi^{0,i}}{\sqrt{2}}, \quad \chi^0 \rightarrow v_\chi + \frac{\chi^{0,r} + i\chi^{0,i}}{\sqrt{2}}, \quad \xi^0 \rightarrow v_\chi + \xi^0. \quad (5.6)$$

The masses within each custodial multiplet can be written in terms of the vevs as

$$\begin{aligned} m_5^2 &= \frac{M_1}{4v_\chi} v_\phi^2 + 12M_2 v_\chi + \frac{3}{2} \lambda_5 v_\phi^2 + 8\lambda_3 v_\chi^2, \\ m_3^2 &= \frac{M_1}{4v_\chi} (v_\phi^2 + 8v_\chi^2) + \frac{\lambda_5}{2} (v_\phi^2 + 8v_\chi^2) = \left(\frac{M_1}{4v_\chi} + \frac{\lambda_5}{2} \right) v^2. \end{aligned} \quad (5.7)$$

The gauge singlet S remains a mass eigenstate, with physical mass-squared given by

$$m_S^2 = \mu_S^2 + 2\lambda_a v_\phi^2 + 6\lambda_b v_\chi^2, \quad (5.8)$$

which is required to be positive to avoid breaking the Z_2 symmetry.

The other mass eigenstates are given by

$$h = \cos \alpha \phi^{0,r} - \sin \alpha H_1^{0r}, \quad H = \sin \alpha \phi^{0,r} + \cos \alpha H_1^{0r}, \quad (5.9)$$

where

$$H_1^{0r} = \sqrt{\frac{1}{3}} \xi^0 + \sqrt{\frac{2}{3}} \chi^{0,r}. \quad (5.10)$$

The mixing angle and masses are given by

$$\begin{aligned} \sin 2\alpha &= \frac{2\mathcal{M}_{12}^2}{m_H^2 - m_h^2}, \quad \cos 2\alpha = \frac{\mathcal{M}_{22}^2 - \mathcal{M}_{11}^2}{m_H^2 - m_h^2}, \\ m_{h,H}^2 &= \frac{1}{2} \left[\mathcal{M}_{11}^2 + \mathcal{M}_{22}^2 \mp \sqrt{(\mathcal{M}_{11}^2 - \mathcal{M}_{22}^2)^2 + 4(\mathcal{M}_{12}^2)^2} \right], \end{aligned} \quad (5.11)$$

where the definition $m_h < m_H$ is chosen, and

$$\begin{aligned} \mathcal{M}_{11}^2 &= 8\lambda_1 v_\phi^2, \\ \mathcal{M}_{12}^2 &= \frac{\sqrt{3}}{2} v_\phi [-M_1 + 4(2\lambda_2 - \lambda_5) v_\chi], \\ \mathcal{M}_{22}^2 &= \frac{M_1 v_\phi^2}{4v_\chi} - 6M_2 v_\chi + 8(\lambda_3 + 3\lambda_4) v_\chi^2. \end{aligned} \quad (5.12)$$

The Feynman rules for this model are listed in Appendix A. To summarize, the GM+SSDM model has a total of 13 free parameters: λ_a , λ_b , λ_1 , λ_2 , λ_3 , λ_4 , λ_5 , λ_S , M_1 , M_2 , μ_2^2 , μ_3^2 , and μ_S^2 . These will be constrained in the following section.

5.2 Phenomenology

This section will look at the constraints on the parameters of the GM+SSDM model. First, theoretical constraints are applied. Then, we include experimental constraints from relic density, direct detection, Higgs invisible width and Higgs signal strengths. A plot in a given section will cumulatively include constraints from all those imposed before the constraints added in the section.

5.2.1 Theoretical Constraints

A number of theoretical constraints can be applied to the singlet scalar DM extension of the GM model. Firstly, perturbative unitarity imposed on the $2 \rightarrow 2$ scattering amplitudes is applied, then the vacuum stability constraints (meaning the potential is bounded from below), then conditions to avoid false minima. For more information, see References [3, 39]. Experimental constraints applied to the generated points but not included in the body of this thesis are the limit on $b \rightarrow s\gamma$ and dwarf spheroidal galaxy constraints [38].

5.2.2 Dark Matter Relic Density

To calculate the relic density, the dark matter annihilation cross sections is considered for all possible processes. The annihilation into non-scalar SM particles (A,B) (Figure 5.1), gives a thermalized cross section of

Figure 5.1: Feynman diagram for $SS \rightarrow A + B$ where A, B are SM particles

Figure 5.2: Feynman diagram for $SS \rightarrow H_3 V^{(*)}$.

$$\begin{aligned} \sigma v_{\text{rel}}(SS \rightarrow A + B) &= \frac{4}{\sqrt{s}} \left(\frac{g_{SSh} g_{hAB}}{s - m_h^2} + \frac{g_{SSH} g_{HAB}}{s - m_H^2} \right)^2 \\ &\times \frac{1}{g_{h_{SM}AB}^2} \Gamma(m_{h_{SM}} = \sqrt{s}, h_{SM} \rightarrow A + B) \end{aligned} \quad (5.13)$$

where Γ is the partial width for the Higgs decaying into SM particles, evaluated when $m_{h_{SM}} = \sqrt{s}$. We also must include annihilation products of one scalar and one SM vector boson, and also the final state $V^* H_3$, where the gauge boson is off-shell (Figure 5.2). The thermalized cross section for the first process is

$$\begin{aligned} \sigma v_{\text{rel}}(SS \rightarrow H_3 V) &= \frac{(s^2 - 2s(m_{H_3}^2 + m_V^2) + (m_{H_3}^2 - m_V^2)^2)^{\frac{3}{2}}}{8\pi s^2 m_W^2} \\ &\times \left(\frac{g_{SSh} g_{hH_3V}}{s - m_h^2} + \frac{g_{SSH} g_{HH_3V}}{s - m_H^2} \right)^2 \end{aligned} \quad (5.14)$$

while that for the second is [40]

$$\sigma v_{\text{rel}}(SS \rightarrow H_3 V^*) = \left(\frac{g_{SSh} g_{hH_3V}}{s - m_h^2} + \frac{g_{SSH} g_{HH_3V}}{s - m_H^2} \right)^2 \delta_V \frac{3m_V^2}{4\pi^3 v^2} G_{ij}, \quad (5.15)$$

where

$$G_{ij} = \frac{1}{4} \left(2(-1 + k_j - k_i) \sqrt{\lambda_{ij}} \left(\frac{\pi}{2} + \arctan \left(\frac{k_j(1 - k_j + k_i) - \lambda_{ij}}{(1 - k_i) \sqrt{\lambda_{ij}}} \right) \right) \right. \\ \left. + (\lambda_{ij} - 2k_i) \log k_i + \frac{1}{3}(1 - k_i) \left(5(1 + k_i) - 4k_j + \frac{2\lambda_{ij}}{k_j} \right) \right) \quad (5.16)$$

$$\delta_W = \frac{3}{2} \quad (5.17)$$

$$\delta_Z = 3 \left(\frac{7}{12} - \frac{10}{9} s_W^2 + \frac{40}{27} s_W^4 \right). \quad (5.18)$$

Here, $k_i = m_{H_3}^2/s$, $k_j = m_V^2/s$, and $\lambda_{ij} = -1 + 2k_i + 2k_j - (k_i - k_j)^2$.

Figure 5.3: Feynman diagrams for $SS \rightarrow h_1 h_2$

Next we consider DM annihilation into any of the following combinations of the Higgs triplet scalar fields:

$$(H_3^0, H_3^0), (H_3^+, H_3^-), (H_5^0, H_5^0), (H_5^+, H_5^-), (H_5^{++}, H_5^{--}). \quad (5.19)$$

In this case we must take into account an additional four-point diagram.

The cross section for identical particles is given by:

$$\sigma_{CM} v_{\text{rel}}(SS \rightarrow h_1 h_2) = \frac{(s - 4m_{h_1}^2)^{\frac{1}{2}}}{32\pi s^{\frac{3}{2}}} \left(g_{SSH_1 h_2} + \frac{g_{SSH} g_{hh_1 h_2}}{s - m_h^2} + \frac{g_{SSH} g_{HH_1 h_2}}{s - m_H^2} \right)^2 \quad (5.20)$$

and for non-identical particles:

$$\sigma_{CM} v_{\text{rel}}(SS \rightarrow h_1 h_2) = \frac{(s - 4m_{h_1}^2)^{\frac{1}{2}}}{16\pi s^{\frac{3}{2}}} \left(g_{SSH_1 h_2} + \frac{g_{SSH} g_{hh_1 h_2}}{s - m_h^2} + \frac{g_{SSH} g_{HH_1 h_2}}{s - m_H^2} \right)^2 \quad (5.21)$$

where h_1 and h_2 are any of the scalar pairs listed in Eqn. 5.19.

Figure 5.4: Feynman diagrams for process $SS \rightarrow h_i h_j$

The cross sections with (h, H) , (h, h) or (H, H) in the final state include an additional t-channel and u-channel diagram as well as the s-channel and the four point interaction in Figure 5.4. For non identical h_i, h_j in the final state, the end result is

$$\frac{d\sigma}{d\Omega} v_{rel} = \frac{\sqrt{s^2 + (m_{h_j}^2 - m_{h_i}^2)^2 - 2s(m_{h_j}^2 + m_{h_i}^2)}}{32\pi^2 s^2} \left[a + \frac{b}{c^2 - \cos^2(\theta)d^2} \right]^2 \quad (5.22)$$

where we define for convenience

$$\begin{aligned} a &= g_{SSh_i h_j} + \frac{g_{SSh_i} g_{h_i h_i h_j}}{s - m_{h_i}^2} + \frac{g_{SSh_j} g_{h_i h_j h_j}}{s - m_{h_j}^2} \\ b &= 4g_{SSh_i} g_{SSh_j} (s - (m_{h_i}^2 + m_{h_j}^2)) \\ c^2 &= (s - m_{h_j}^2 - m_{h_i}^2)^2 \\ d^2 &= \left(1 - \frac{4m_S^2}{s}\right) \left((s - m_{h_j}^2 + m_{h_i}^2)^2 - 4sm_{h_i}^2\right) \end{aligned}$$

and

$$g(s) = \int_{-1}^1 d\cos(\theta) \left[a - \frac{b}{c^2 - \cos^2(\theta)d^2} \right]^2 = 2a^2 + \frac{b^2}{c^2(c^2 - d^2)} + \frac{b(b - 4ac^2)}{c^3d} \tanh^{-1} \left(\frac{d}{c} \right).$$

For non identical h_i, h_j , the cross section is

$$\sigma v_{rel} = \frac{\sqrt{s^2 + (m_{h_j}^2 - m_{h_i}^2)^2 - 2s(m_{h_j}^2 + m_{h_i}^2)}}{16\pi s^2} g(s)$$

and for identical particles, is

$$\sigma v_{rel} = \frac{\sqrt{s - 4m_{h_i}^2}}{32\pi s^{3/2}} g(s)$$

This is applied in a slightly different way than that of Section 4.2.3. The thermalized cross section is calculated using equation 3.3 as used for the dilaton. We use the thermally averaged total annihilation cross section as input for the usual Boltzmann equation:

$$\frac{dn_S}{dt} + 3Hn_S = -\langle\sigma v\rangle \left[n_S^2 - (n_S^{\text{eq}})^2 \right]. \quad (5.23)$$

We solve this equation numerically to obtain the number density n_S today which translates to a value for the relic abundance by multiplying by m_S and dividing by the critical density of the universe today.

The relic density is a constraint on a combination of λ_a and λ_b . To implement this, we first note that with all other parameters fixed, there is an oval like shaped set of (λ_a, λ_b) for which we obtain the correct relic density. In the spirit of the numerical scans, we would like to be able to randomly select one such point. This is done by randomly generating an angle θ_λ in $[-\frac{\pi}{2}, \frac{\pi}{2}]$. Then we set $\lambda_b = \lambda_a \tan \theta_\lambda$, generate a mass m_S , and find the λ_a that gives the correct relic abundance, which also in turn gives μ_S . Searching for the correct value of λ_a can be done easily since σv_{rel} is an increasing function of λ_a and $\sigma v_{\text{rel}} \rightarrow \infty$ when $\lambda_a \rightarrow \infty$ and $\sigma v_{\text{rel}} \rightarrow \infty$ when $\lambda_a \rightarrow 0$.

5.2.3 Dark Matter Direct Detection

In our model, direct detection proceeds via exchange of a virtual h or H , as shown in Figure 5.5.

Figure 5.5: Feynman diagram for direct detection

The spin independent direct detection cross section is

$$\sigma = \left(\frac{f_N g_{SSh} \cos(\alpha)}{m_h^2 \cos(\theta_H)} + \frac{f_N g_{SSH} \sin(\alpha)}{m_H^2 \cos(\theta_H)} \right)^2 \frac{m_N^4}{4\pi(m_N + m_S)^2 v^2} \quad (5.24)$$

where f_N is the nucleon vertex factor,

$$f_N = \sum_q f_q = \sum_q \frac{m_q}{m_N} \langle N | \bar{q}q | N \rangle \quad (5.25)$$

and the sum is over all quark flavours. The value of f_N used is 0.30 ± 0.03 , and the nucleon mass is $m_N = 938.95$ MeV [35]. The upper limits on the DM-Higgs vertices g_{SSh} and g_{SSH} were calculated in a previous paper [1]. Limits are included from the experiments LUX and XENON, as well as projected limits from DEAP 3600 [23, 24, 27, 36]. Results are shown in Figure 5.6 for both the low and high mass regions. Points which fall above the experimental line are excluded, therefore direct detection effectively eliminates much of the low mass region.

Figure 5.6: Direct detection cross section as a function of m_S . The left plot is a zoom of the low- m_S region while the right plot shows the full m_S range scanned. Experimental bounds include LUX (magenta), XENON10 (purple), XENON1T (brown), DEAP3600 (green), and the neutrino floor (yellow).

5.2.4 Dark Matter Invisible Width

When $m_S < m_{h,H}/2$, the decay of the Higgs boson to two dark matter candidates is kinematically accessible. The expression for the width of this process is

$$\Gamma_{SSH} = \frac{g_{SSH}^2}{32m_h^2\pi} \sqrt{m_h^2 - 4m_s^2}. \quad (5.26)$$

To compare with experiment, the ratio

$$\frac{\sigma_{VBF} BR_{inv}}{\sigma_{VBF}^{SM}} = \frac{(\kappa_V)^2 \Gamma(h \rightarrow inv)}{\Gamma_{tot}}, \quad (5.27)$$

is considered, written in terms of the vector boson fusion (VBF) production cross section and the invisible branching ratio, where κ_V is defined in Equation 5.30. Here,

$$\begin{aligned} \Gamma_{tot} = & (\kappa_f^h)^2 [\Gamma(h \rightarrow b\bar{b}) + \Gamma(h \rightarrow c\bar{c}) + \Gamma(h \rightarrow \tau^+\tau^-) + \Gamma(h \rightarrow t\bar{t})] \\ & + (\kappa_W^h)^2 \Gamma(h \rightarrow WW^*) + (\kappa_Z^h)^2 \Gamma(h \rightarrow ZZ^*) + \Gamma(h \rightarrow invs). \end{aligned} \quad (5.28)$$

The ratio in Equation 5.27 is shown in Figure 5.7, plotted against m_S in the kinematically allowed region. The experimental limit [19, 41] is the horizontal solid blue line, and excludes much of the remaining low mass DM parameter space.

Figure 5.7: Constraint from invisible decays of the Higgs for a dark matter mass range between 0 GeV and 125/2 GeV. The blue points are allowed while the red points are excluded by either direct detection constraints, $b \rightarrow s\gamma$ constraints, constraints on the invisible decays of the Higgs, dwarf spheroidal galaxy constraints, or mass constraints on the triplet and fiveplet. Anything above the blue line is excluded.

5.2.5 Higgs Signal Strengths

It is convenient in this model to look at the signal strengths slightly differently. The κ 's are defined by

$$\kappa_f = \frac{g_{hf\bar{f}}}{g_{hf\bar{f}}^{\text{SM}}} \quad (5.29)$$

and

$$\kappa_V = \frac{g_{hVV}}{g_{hVV}^{\text{SM}}}, \quad (5.30)$$

as ratios of the coupling of h to fermions and to vector bosons. The GM model dictates that κ_f is the same for all f , and similarly for κ_V for all V . Rewriting the

signal strengths in terms of these parameters, there are five different possibilities:

$$\mu(AB \rightarrow h \rightarrow CD) = \frac{\Gamma_{Tot}^{SM}(h)}{\Gamma_{Tot}(h)} \begin{cases} \kappa_f^4 & (1) \\ \kappa_V^4 & (2) \\ \kappa_f^2 \kappa_V^2 & (3) \\ \kappa_f^2 \frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma^{SM}(h \rightarrow \gamma\gamma)} & (4) \\ \kappa_V^2 \frac{\Gamma(h \rightarrow \gamma\gamma)}{\Gamma^{SM}(h \rightarrow \gamma\gamma)} & (5) \end{cases} \quad (5.31)$$

where (1) occurs if the production mode is gluon gluon fusion or top associated production and the final state are fermions or two gluons, (2) occurs if the production mode is vector boson fusion or vector boson associated production and the Higgs decays to WW^* or ZZ^* , (3) occurs if you mix cases (1) and (2), (4) occurs if you have the production modes of (1) and a final state of two photons and (5) occurs if you have the production modes of (2) and a final state of two photons.

We look at the ATLAS/CMS results using the mechanisms of the form [30]

$$\frac{\sigma(gg \rightarrow H \rightarrow ZZ)}{\sigma(gg \rightarrow H \rightarrow ZZ)_{SM}} = \kappa_f^2 \kappa_V^2 \frac{\Gamma_h^{SM}}{\Gamma_h}. \quad (5.32)$$

A chi-squared fit with the condition $|\Delta\chi^2| < 4$ results in a parameter space for κ_f and κ_V shown in Figure 5.8

The parameters κ_f and κ_V in general do not deviate far from their standard model values ($\kappa_f = \kappa_V = 1$). The red points excluded have high enough values to contribute to high κ values near the Higgs poles, therefore small couplings are needed to pass the relic abundance constraint.

5.3 Summary

The Georgi-Machacek model is extended by singlet scalar DM. After applying all constraints, singlet scalar DM masses are excluded below about 50 GeV. The tightest constraints on this model are the observed dark matter relic density and direct detection limits. There does exist, above the Higgs pole, a substantial region of parameter space yielding the correct dark matter relic density that can escape the near-future direct detection experiments DEAP and XENON 1T for dark matter

Figure 5.8: Parameter space of κ_f and κ_V . The blue points are allowed by all previously imposed constraints.

masses as low as 125 GeV.

Chapter 6

Self-Interacting Dark Matter

6.1 Background and Motivation

Recently, a study of the galaxy cluster Abell 3827 revealed a spacial offset of the dark matter halo from at least one of its central galaxies [42]. This cluster is unique in that the remnants of four elliptical galaxies are within the core radius of approximately 10 kpc. Using both Hubble Space Telescope (HST) imaging (see Figure 6.1) and integral field spectroscopy of the gravitationally lensed system through the core, the distribution of associated dark matter can be profiled. The dark matter halo in this galaxy cluster was found to be spatially offset by $\Delta = 1.62^{+0.47}_{-0.49}$ kpc.

This offset could possibly be described by dark matter self interaction. In a high-momentum scenario, we cannot assume that the halo effects can be described by a drag force akin to classical physics. Most of the dark matter particles within the halo will not interact, but those that do will have momentum away from the core, separating from the halo. On a small time scale these will become observationally separable from the dark matter halo.

Reference [43] finds that the offset, although sensitive to the definition of the centroids of the galaxy and halo, gives an approximate cross section

$$\frac{\sigma_{DM}}{m} \sim 1.5 \text{ cm}^2 \text{g}^{-1}. \quad (6.1)$$

In the case of contact interactions, this result is in tension with previously cited values from other astronomical systems, which give an upper bound of [44]

$$\frac{\sigma_{DM}}{m} \leq 1 \text{ cm}^2 \text{g}^{-1}. \quad (6.2)$$

Figure 6.1: Abell 3827 HST image from Reference [42], showing the F160W (red), F160W (green) and F336W (blue) optical bands on a logarithmic colour scale. Labels N1 through N4 are the four bright central galaxies, while N5 is a star.

Both of these values will be used to constrain and thus give a range for the estimated self-interaction cross section.

The extension of the SM by a singlet scalar dark matter candidate, S , is considered, which is stable by imposing a Z_2 symmetry. This couples to the rest of the SM through a “Higgs portal” operator of the form $SS\Phi^\dagger\Phi$, where Φ is the $SU(2)_L$ doublet SM Higgs field.

6.2 Minimal Singlet Extension of the Standard Model

The SM plus a real singlet scalar has the scalar potential

$$V = -\mu^2\Phi^\dagger\Phi + \frac{\mu_S^2}{2}SS + \lambda_h(\Phi^\dagger\Phi)^2 + \lambda_p SS\Phi^\dagger\Phi + \lambda_S S^4, \quad (6.3)$$

where

$$\lambda_h = \frac{m_h^2}{2v^2} \simeq 0.129 \quad (6.4)$$

and

$$\mu^2 = \lambda_h v^2 \simeq (88.4 \text{ GeV})^2. \quad (6.5)$$

To ensure the stability of the singlet (i.e. it does not acquire a vev), it is required that

$$m_S^2 = \mu_S^2 + \lambda_p v^2 > 0, \quad (6.6)$$

where m_S is the physical mass of the singlet scalar in the electroweak-breaking vacuum. It is necessary that $\mu_S^2 > 0$ to ensure that the potential does not develop a second minimum. To ensure boundedness from below, the conditions $\lambda_S > 0$ and $\lambda_p > -2\sqrt{\lambda_h \lambda_S}$ must also hold.

The relevant Feynman rules for the singlet vertices can be found in Equation 6.7, and the Feynman diagrams for the processes in Figure 6.2.

$$\begin{aligned} SSSS &: -24i\lambda_S \\ SSh &: -2i\lambda_p v. \end{aligned} \quad (6.7)$$

There will be three DM scenarios considered: that of high mass, low mass, and mass near the Higgs resonance. Note that we will compare to the measured values in particle physics units, or, rewriting Equations 6.1 and 6.2,

$$\sigma_{SM}/m \approx (4.7 - 7.0) \times 10^3 \text{ GeV}^{-3}. \quad (6.8)$$

6.2.1 Heavy Dark Matter, $m_S \gg m_H$

At a high dark matter mass, the momentum of the internal Higgs is suppressed by momentum factors, so the four-point diagram dominates. Therefore the scattering cross section per unit of dark matter mass is given by

$$\frac{\sigma_{DM}}{m_S} \simeq \frac{9\lambda_S^2}{2\pi m_S^3}. \quad (6.9)$$

However, perturbative unitarity sets an upper bound on λ_S of $\frac{2\pi}{3}$, due to identical particles in the final state and requirements of the zeroth partial wave amplitude. This results in a self interaction cross section that is inconsistent with the measured

Figure 6.2: Feynman diagrams for the interaction $SS \rightarrow SS$

value in Equation 6.8. In fact, reproducing the measured dark matter interaction cross section requires low-mass dark matter, or $m_S \leq 0.1\text{GeV}$, thus ruling out the heavy DM scenario.

6.2.2 Dark matter near the Higgs resonance

At a mass near the Higgs resonance ($m_S \simeq m_h/2$), the s-channel exchange diagram dominates as in Figure 6.3. In this case, the cross section per unit mass is

$$\frac{\sigma_{DM}}{m_S} \simeq \frac{g_{SSh}^4}{128\pi m_S^3 (4m_S^2 - m_h^2)^2} \quad (6.10)$$

$$\simeq \frac{v^4 \lambda_p^4}{16\pi m_h^5 (m_S - m_h/2)^2} \quad (6.11)$$

where ig_{SSh} is the hSS Feynman coupling, and $m_S \approx m_h/2$. An additional constraint comes from perturbative unitarity, yielding $|g_{SSh}| < 8\pi v$. This requires severe tuning of m_S , such that $|m_S - m_h/2| \leq 0.11$ GeV.

Figure 6.3: Feynman diagrams for the annihilation of DM into SM particles in the real singlet scalar extension of the SM.

We can try to reconcile this scenario using the observed dark matter relic density. The value of the corresponding annihilation cross section is

$$\sigma v_{rel} \approx 2.2 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1} \approx 1.9 \times 10^{-9} \text{ GeV}^{-2} \quad (6.12)$$

where v_{rel} is the relative velocity between the two dark matter particles in the centre of momentum frame. The annihilation cross section can be written as

$$\sigma v_{rel} = \frac{2g_{SSh}^2 \Gamma_{tot}^{h_{SM}}}{E_{cm}(E_{cm}^2 - m_h^2)^2} \quad (6.13)$$

where $E_{cm} = 2m_S + \frac{1}{4}m_S v_{rel}^2$ is the centre of mass energy and $\Gamma_{tot}^{h_{SM}}$ is the total width of the SM Higgs (4.1 MeV) [45]. If the singlet scalar is heavier than half the Higgs mass, the thermal motion pushes the annihilation process far from resonance. Integrating the annihilation cross section over the Boltzmann distribution under the assumption $T = m_S/20$ gives

$$\langle \sigma v_{rel} \rangle \approx 0.03 \text{ GeV}^{-2}. \quad (6.14)$$

This is many orders of magnitude greater than the required cross section from Equation 6.12, so that the relic abundance would be too small to account for the

observed dark matter. If the singlet scalar is lighter than half the Higgs mass, the cross section becomes even larger and forces the relic abundance smaller, due to the large hSS coupling and the tuning of the singlet mass.

6.2.3 Low-mass dark matter, $m_S \ll m_h$

At low mass ($m_S \ll m_h$), all four of the Feynman diagrams in Figure 6.2 will contribute to dark matter self annihilation. At low energies the Higgs boson can be integrated out, which yields an effective 4-point coupling with Feynman rule $-24i\lambda_{eff}$, where

$$\lambda_{eff} = \lambda_S - \frac{g_{SSh}^2}{8m_h^2}. \quad (6.15)$$

Again, bounds from perturbative unitarity must be applied, giving the upper bound $|\lambda_{eff}| < 2\pi/3$. The DM scattering cross section as in Equation 6.9 becomes

$$\frac{\sigma_{DM}}{m_S} \approx \frac{9\lambda_{eff}^2}{2\pi m_S^3} \quad (6.16)$$

These results are consistent with Reference [43] if the dark matter mass is very light, ie below about 0.1 GeV, as shown in Figure 6.4.

6.2.4 Phenomenology

Invisible Higgs Decays

In the low mass range $m_S \lesssim 0.1$ GeV, the parameter λ_p is constrained by a lack of observation of invisible Higgs decays. The expression for the invisible partial width is

$$\Gamma(h \rightarrow SS) = \frac{g_{SSh}^2}{32\pi m_h} \sqrt{1 - \frac{4m_S^2}{m_h^2}} \simeq \frac{g_{SSh}^2}{32\pi m_h} \quad (6.17)$$

where the last approximation holds for $m_S \ll m_h$. Again, the strongest constraint on the branching ratio (Equation 3.3) is $\text{BR}(h \rightarrow \text{invis.}) < 0.28$, coming from ATLAS bounds on vector boson fusion with a 125GeV Higgs and a SM production cross section [19]. Therefore the upper bound on g_{SSh} is

Figure 6.4: The allowed region for m_S and $|\lambda_{eff}|$ to obtain the desired range in Equation 6.8. This includes the constraints from perturbative unitarity and thresholds for DM annihilation to electrons, muons, and neutral pions.

$$g_{SSh}^2 < 21 \text{ GeV}^2 \quad \text{or} \quad |\lambda_p| < 9.3 \times 10^{-3}, \quad (6.18)$$

which is strong when compared to the unitary constraint $|\lambda_p| < 4\pi$.

An upper bound on g_{SSh} can also be found from the measured signal strengths of the Higgs boson. In the model SM + scalar singlet, all signal strengths are multiplied by the factor

$$\mu_i = \frac{\Gamma_{tot}^{hsm}}{\Gamma_{tot}^{hsm} + \Gamma(h \rightarrow SS)}. \quad (6.19)$$

ATLAS and CMS measure an overall signal strength of

$$\begin{aligned} \mu &= 1.18^{+0.15}_{-0.14} & (\text{ATLAS}) \\ \mu &= 1.00^{+0.14}_{-0.13} & (\text{CMS}) \end{aligned} \quad (6.20)$$

However, there is an additional constraint on this model in particular, that $\mu \leq 1$, so the $\Delta\chi^2 = 4$ lower bounds are

$$\begin{aligned}\mu &\geq 0.85 && (\text{ATLAS}), \\ \mu &\geq 0.74 && (\text{CMS}).\end{aligned}\tag{6.21}$$

These lead to upper bounds on g_{ssh}^2 :

$$\begin{aligned}g_{SSH}^2 &\leq 9.1 \text{ GeV}^2 \quad \text{or} \quad |\lambda_p| \leq 6.1 \times 10^{-3} && (\text{ATLAS}), \\ g_{SSH}^2 &\leq 18 \text{ GeV}^2 \quad \text{or} \quad |\lambda_p| \leq 8.6 \times 10^{-3} && (\text{CMS}).\end{aligned}\tag{6.22}$$

Relic Density

For small $m_S \sim 0.1$, the annihilation cross section must be [46]

$$\sigma v_{rel} \simeq 4.8 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1} \simeq 4.1 \times 10^{-9} \text{ GeV}^{-2}\tag{6.23}$$

At this low mass, the dark matter annihilation is predominantly s-channel, as in Figure 6.3. This can be expressed in terms of the total width of the SM Higgs if the dark matter were half the Higgs mass:

$$\sigma v_{rel} \simeq \frac{4v^2 \lambda_p^2}{m_h^4} \frac{\Gamma_{tot}^{hSM}(m_h = 2m_S)}{m_S}\tag{6.24}$$

As shown in Figure 6.4, the highest annihilation cross section is found at the high end of the singlet mass range, when annihilations into muons are kinematically accessible, and $\Gamma_{tot}^{hSM}(m_h = 2m_S)$ can be approximated by the width into muons:

$$\Gamma(h^* \rightarrow \mu\mu) = \frac{1}{8\pi} \frac{m_\mu^2}{v^2} (2m_S) \left(1 - \frac{m_\mu^2}{m_S^2}\right)\tag{6.25}$$

which, substituting into Equation 6.24, gives

$$\sigma v_{rel} \simeq (1.5 \times 10^{-11} \text{ GeV}^{-2}) \lambda_p^2 \left(1 - \frac{m_\mu^2}{m_S^2}\right)^{\frac{3}{2}}\tag{6.26}$$

A large enough value of the annihilation cross section cannot be obtained from constraints on perturbative unitarity alone, not to mention constraints on the muon mass and $|\lambda_p|$ from LHC Higgs data. In fact, the constraint on $|\lambda_p|$ from invisible Higgs decays and signal strengths results in the annihilation cross section to be six orders

of magnitude too small to yield the correct relic abundance. Furthermore, the low mass scenario with $SS \rightarrow \mu\mu$ is ruled out by relic abundance constraints.

Thermal Freeze-In

The correct dark matter relic abundance for thermal freezeout cannot be obtained by this model, for the cross section to SM particles cannot be made large enough. However, there is an alternate way to obtain the relic abundance: through thermal freeze-in [47, 48]. Here, the cross section of dark matter annihilating to SM particles is extremely small, so that thermal equilibrium is never reached. Instead, dark matter is produced through annihilations of SM particles, until the universe becomes too cool to produce dark matter, freezing in the current amount of dark matter. The annihilation of dark matter to SM particles become kinematically irrelevant.

This scenario was previously studied in references [47, 49]. In the low mass region $m_S \leq 0.1 \text{ GeV}$, the correct dark matter relic abundance is obtained for $|\lambda_p| \sim (2 - 3) \times 10^{-11}$. This is consistent with the upper bound $|\lambda_p| < (6.1 - 9.3) \times 10^{-3}$ from LHC signal strengths and invisible width branching ratio limits. The only problem is in the phenomenological sense: the dark matter becomes impossible to detect through non-gravitational interactions with SM particles.

6.3 Two Higgs Doublet Model with Singlet DM

The thermal freeze-out scenario is now considered for the Two Higgs doublet model (2HDM), again with singlet scalar dark matter. The 2HDM is characterized by a SM Higgs with mass 125 GeV, and a second CP-even scalar, H , which could be lighter or heavier than h .

Previously, the resonant-scattering produced a cross section far too large for DM annihilating through a Higgs portal. For this new scenario, the DM interacts resonantly through the H , where the thermal DM annihilation cross section can be suppressed through the coupling of H to SM particles.

6.3.1 Type I 2HDM

The scalar potential of the 2HDM with a singlet scalar DM extension is

$$V = \frac{\mu_S^2}{2} S^2 + \lambda_S S^4 + \lambda_{p1} SS\Phi_1^\dagger \Phi_1 + \lambda_{p2} SS\Phi_2^\dagger \Phi_2 + V_{2HDM}. \quad (6.27)$$

Here, Φ_1 and Φ_2 are the two Higgs doublets, and V_{2HDM} is the scalar potential of the 2HDM.

The class of 2HDMs can be divided into “types” according to symmetries and couplings. The soft breaking of the Z_2 symmetry ($\Phi_1 \rightarrow -\Phi_1$) forces the right-handed fermions to only couple to one of the Higgs doublets [50]. This transformation is classified as a “Type I” 2HDM.

The vacuum expectation values are as follows:

$$\langle \Phi_1 \rangle = \begin{pmatrix} 0 \\ v_1/\sqrt{2} \end{pmatrix}, \quad \langle \Phi_2 \rangle = \begin{pmatrix} 0 \\ v_2/\sqrt{2} \end{pmatrix}, \quad (6.28)$$

where normalization gives $v_1^2 + v_2^2 = v^2$, and $\tan\beta = v_2/v_1$ by definition.

In terms of the real neutral components of the two doublets $\Phi_i^0 \equiv \phi_i^{0,r}/\sqrt{2}$, the physical CP-even neutral scalars h and H are defined as

$$\begin{aligned} h &= -\sin\alpha \phi_1^{0,r} + \cos\alpha \phi_2^{0,r}, \\ H &= \cos\alpha \phi_1^{0,r} + \sin\alpha \phi_2^{0,r}, \end{aligned} \quad (6.29)$$

where α is a mixing angle.

The first constraint comes from perturbative unitarity of the scattering amplitudes $SS \rightarrow \Phi_{1(2)}\Phi_{1(2)}$, which constrains the Higgs-DM couplings:

$$|\lambda_{p1}| < 4\pi, \quad |\lambda_{p2}| < 4\pi, \quad (6.30)$$

The Feynman rules for the SSh and SSH vertices are ig_{SSh} and ig_{SSH} , respectively, where

$$\begin{aligned} g_{SSh} &= 2v(\lambda_{p1} \cos\beta \sin\alpha - \lambda_{p2} \sin\beta \cos\alpha) \\ g_{SSH} &= 2v(-\lambda_{p1} \cos\beta \cos\alpha - \lambda_{p2} \sin\beta \sin\alpha). \end{aligned} \quad (6.31)$$

The low-energy effective $SS \rightarrow SS$ coupling for low m_S is given by

$$\lambda_{eff} = \lambda_S - \frac{g_{SSh}^2}{8m_h^2} - \frac{g_{SSH}^2}{8m_H^2} \quad (6.32)$$

The constraint from the previous section $m_S \lesssim 0.1$ GeV is still valid here, so we do not continue the low-mass analysis.

To avoid experimental constraints from h coupling measurements, we tune $\sin(\beta - \alpha) \simeq 1$, implying $\sin \alpha \simeq -\cos \beta$ and $\cos \alpha \simeq \sin \beta$. For $\sin(\beta - \alpha) \simeq 1$, the coupling between H and pairs of vector bosons is heavily suppressed, as it is proportional to $\cos(\beta - \alpha)$. In this Type-I model, fermion Yukawa couplings can be suppressed by making $\cot \beta$ small.

The H coupling to fermions is proportional to the SM Higgs coupling, with scaling factor

$$\kappa_F^H = \frac{\sin \alpha}{\sin \beta} \simeq -\cot \beta, \quad (6.33)$$

where the approximation corresponds to the $\sin(\beta - \alpha) \rightarrow 1$ limit.

If $m_S \approx m_H/2$, dark matter scattering $SS \rightarrow SS$ is dominated by the s-channel H exchange diagram. The resulting cross section per unit mass is

$$\sigma_{DM}/m_S \simeq \frac{g_{SSH}^4}{128\pi m_S^3 (4m_S^2 - m_H^2)^2}. \quad (6.34)$$

Given the constraints on the vertex factors, the singlet mass must be tuned very close to $m_H/2$ to obtain a cross section compatible with the halo displacement from Reference [43].

The dark matter annihilation cross section is given for annihilation through the H resonance by

$$\sigma v_{rel} \simeq \frac{2g_{SSH}^2 \Gamma_{tot}^{H \rightarrow SM}}{E_{cm} (E_{cm}^2 - m_H^2)^2} \quad (6.35)$$

which is of the same form as the cross section in Equation 6.34. For the scalar singlet extension of the SM, the problem arose from the dark matter self-interaction cross section causing the annihilation cross section to be too large. Therefore g_{SSH} must be large while the H producing SM particles are also suppressed.

It is assumed that $m_H > m_h$, so that $m_S > m_h/2$ and invisible decay are kinematically forbidden. To suppress the W and Z decay channels, it is assumed that $\sin(\beta - \alpha) \rightarrow 1$, and the fermions can be suppressed taking $\cot \beta \ll 1$.

The decay width of H to SM particles (assuming $m_H < 2m_t$) is

$$\Gamma_{tot}^{H \rightarrow SM} \simeq (\kappa_f^H)^2 \left(\frac{m_H}{125 \text{ GeV}} \right) [1 - BR(h_{SM} \rightarrow VV)] \Gamma_{tot}^{h_{SM}}, \quad (6.36)$$

where

$$BR(h_{SM} \rightarrow VV) \equiv BR(h_{SM} \rightarrow WW) + BR(h_{SM} \rightarrow ZZ) \simeq 0.24 \quad (6.37)$$

and

$$\Gamma_{tot}^{H \rightarrow SM} \propto (\kappa_f^H)^2 \simeq \cot^2 \beta. \quad (6.38)$$

Comparing this expression with the annihilation cross section in Equation 6.34, the parameter combination $(\kappa_f^H)^2 / g_{SSH}^2$ must be extremely suppressed. However, taking into account the approximations, the SSH coupling is

$$g_{SSH} \simeq 2v \sin \beta \cos \beta (-\lambda_{p1} + \lambda_{p2}), \quad (6.39)$$

which is proportional to $\cos \beta$, and is therefore not suppressed by $\cot \beta \ll 1$. Once the upper bounds from perturbative unitarity are imposed ($|g_{SSH}| < 8\pi v$), the situation is no better than the previous SM+singlet scenario.

6.3.2 Type II 2HDM

The above problems could be avoided if the scalar potential contained the term $(SS\Phi_1^\dagger\Phi_2 + h.c.)$, breaking the Z_2 symmetry. This is referred to as the Type II 2HDM, and allows fermions to couple to both Higgs. We add an additional term to the potential in Equation 6.3:

$$V \supset (\lambda_{p3} SS\Phi_1^\dagger\Phi_2 + h.c.), \quad (6.40)$$

where λ_{p3} is real to avoid CP violation. The Feynman rule parameters become

$$\begin{aligned} g_{SSH} &= 2v[\lambda_{p1} \cos \beta \sin \alpha - \lambda_{p2} \sin \beta \cos \alpha - \lambda_{p3}(\cos \beta \cos \alpha - \sin \beta \sin \alpha)] \\ g_{SSH} &= 2v[-\lambda_{p1} \cos \beta \cos \alpha - \lambda_{p2} \sin \beta \sin \alpha - \lambda_{p3}(\cos \beta \sin \alpha + \sin \beta \cos \alpha)] \end{aligned} \quad (6.41)$$

Tuning $\sin(\beta - \alpha) \simeq 1$ to make the h couplings SM-like reduces these to

$$\begin{aligned}
g_{SSh} &\simeq 2v[\lambda_{p1} \cos^2 \beta - \lambda_{p2} \sin^2 \beta - 2\lambda_{p3} \sin \beta \cos \beta] \\
g_{SSH} &\simeq 2v[(-\lambda_{p1} + \lambda_{p2}) \sin \beta \cos \beta - \lambda_{p3}(\sin^2 \beta - \cos^2 \beta)].
\end{aligned} \tag{6.42}$$

Now, a suppression of the $(\kappa_f^H)^2$ factor in Equation 6.35 need not be counteracted by a suppression of g_{SSH}^2 . Choosing $\kappa_f^H \sim 10^{-4}$ gives an annihilation cross section in the right range for thermal freezeout.

6.4 Summary

The apparent observation of dark matter self-interactions in the cluster Abell 3827 are used to constrain models of singlet scalar dark matter. Thermal freeze-out in the simple case of the SM with a singlet scalar DM candidate produces an annihilation cross section six orders of magnitude too small to account for the observed relic abundance. The most natural solution for this deficiency is very light dark matter (below about 0.1 GeV), whose relic abundance is set by freeze-in through extraordinarily weak interactions with the Higgs doublet(s) through a quartic coupling of order 10^{-11} . Two-Higgs doublet models are examined in hopes of a better result. The solution is, however, quite fine tuned: the singlet mass must be tuned to a value very close (of order a part per mille) to half the mass of the second CP-even neutral Higgs boson.

Chapter 7

Summary

The question of dark matter continues to be one to unite experiment and theory in the search for a solution. This thesis considers models of BSM physics that incorporate a singlet scalar DM candidate, and uses theoretical and experimental constraints to constrain model parameters.

In Section 4 the phenomenology of a model is studied that treats the physical Higgs boson of the SM as a dilaton, the pseudo-Goldstone boson of spontaneous conformal symmetry breaking, and additionally introduces a scalar singlet dark matter candidate. Three of the six parameters of the model were fitted to the properties of the 125 GeV Higgs boson properties. Constraints from DM experiments were then implemented. It was found that the magnitudes of c_{gg} and $c_{\gamma\gamma}$ are constrained to be close to the standard model values. It was also found that masses less than ~ 50 GeV are excluded 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 125–300 GeV. The higher mass region will be probed in the near future by the DEAP DM search experiment. The limited region of parameter space allowed by direct detection experiments is enhanced when one allows for a negative value of c_{gg} . Furthermore, future precision measurements on the parameters governing the Higgs potential can indicate whether the scalar observed is the scalar responsible for the Brout-Englert-Higgs mechanism of electroweak symmetry breaking or a dilaton. If these measurements deviate from the SM expectation, they reveal more about the conformal field theory 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 $4\pi f$.

In Section 5 the Georgi-Machacek model with a SSDM candidate is studied. The GM model extends the Standard Model Higgs sector with the addition of isospin-triplet scalar fields in such a way as to preserve the custodial symmetry. The singlet scalar is required to account for the full relic abundance of dark matter in the universe. The combination of the observed dark matter relic density and direct detection constraints exclude singlet scalar masses below about 50 GeV. It is also found that, in the parameter space above the Higgs pole, the correct dark matter relic density can escape the near-future direct detection experiments DEAP and XENON 1T for dark matter masses as low as 125 GeV. This is in contrast to the singlet scalar dark matter extension of the Standard Model, for which these future experiments are expected to exclude dark matter masses above the Higgs pole up to the multi-TeV range.

Finally, in Section 6 the evidence for dark matter self-interactions based on the observation of a spatial offset between the dark matter halo and the stars in a galaxy in the cluster Abell 3827 is considered. Interpreting the offset as due to dark matter self-interactions leads to a cross section measurement of $\text{DM}/m=(1-1.5)\text{cm}^2\text{g}^{-1}$, where m is the mass of the dark matter particle. This observation is used to constrain singlet scalar dark matter extensions of the standard model and two-Higgs-doublet models. We show that the most natural scenario in this class of models is very light dark matter, below about 0.1 GeV, whose relic abundance is set by thermal freeze-in. It is also shown that the dark matter abundance can be established through the usual thermal freeze-out mechanism in the singlet scalar extension of the Yukawa-aligned two-Higgs-doublet model, but that it requires severe fine tuning of the singlet scalar mass.

Future work may include reconsidering the Higgs-like dilaton with other options for dark matter (i.e. composite DM). There are also other existing models that could be put through the set of constraints outlined in this thesis to study the addition of a dark matter candidate.

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Appendix A

Feynman Rules

A.1 Higgs-like Dilaton with Scalar Singlet Dark Matter

The Feynman rules in this section are written in terms of the vertex $-ig_{abc(d)}$.

Feynman rules for couplings to S

The Feynman rules for the coupling of the dilaton to S are

$$\begin{aligned} g_{SS\sigma} &= \frac{-2i\epsilon m_S^2}{v^2} \\ g_{SS\sigma\sigma} &= \frac{-4i\epsilon^2 m_S^2}{v^2}. \end{aligned} \tag{A.1}$$

Feynman rules for couplings to σ

The Feynman rules for the dilaton coupling to SM particles are as follows:

$$\begin{aligned}
g_{\sigma\gamma\gamma} &= \frac{ic_{\gamma\gamma}\epsilon p_1^{\mu_2} p_2^{\mu_1}}{2\pi\alpha_{EM}v} - \frac{ic_{\gamma\gamma}\epsilon g_{\mu_1\mu_2} p_1 \cdot p_2}{2\pi\alpha_{EM}v} \\
g_{\sigma gg} &= \frac{i\alpha_{SC} c_{gg} \epsilon p_1^{\mu_2} p_2^{\mu_1} \delta_{a_1, a_2}}{2\pi v} - \frac{i\alpha_{SC} c_{gg} \epsilon \delta_{a_1, a_2} g_{\mu_1\mu_2} p_1 \cdot p_2}{2\pi v} \\
g_{\sigma WW} &= \frac{2i\epsilon m_W^2 g_{\mu_1\mu_2}}{v} \\
g_{\sigma ZZ} &= \frac{2i\epsilon m_Z^2 g_{\mu_1\mu_2}}{v} \\
g_{\sigma ff} &= \frac{i\epsilon m_f \delta_{m_1, m_2} \delta_{s_1, s_2}}{v}
\end{aligned} \tag{A.2}$$

where f is a fermion. The Lorentz indices μ_i run from 0 to 3, and the gauge indices a run from 1 to 8. The self-coupling of the dilaton is

$$g_{\sigma\sigma\sigma} = -\frac{3i\epsilon m_\sigma^2}{v} + \frac{i\epsilon\lambda m_\sigma^2}{v}. \tag{A.3}$$

A.2 Georgi-Machacek Model with Scalar DM

Triple scalar couplings

The Feynman rules for couplings to S are given by $-ig_{SSH_i}$ with all particles incoming and the couplings defined as follows:

$$\begin{aligned}
g_{SSH} &= -4 \left(\lambda_a s_\alpha v_\phi + \sqrt{3} \lambda_b c_\alpha v_\chi \right) \\
g_{SSH} &= -4 \left(\lambda_a c_\alpha v_\phi + \sqrt{3} \lambda_b s_\alpha v_\chi \right).
\end{aligned} \tag{A.4}$$

Here the shorthand notation $s_\alpha \equiv \sin \alpha$ and $c_\alpha \equiv \cos \alpha$ is used.

Four scalar couplings

The Feynman rules for couplings to S are given by $-ig_{SSs_1s_2}$ with all particles incoming and the couplings defined as follows:

$$\begin{aligned}
g_{SShh} &= -4(\lambda_a c_\alpha^2 + \lambda_b s_\alpha^2) \\
g_{SSHH} &= -4(\lambda_a s_\alpha^2 + \lambda_b c_\alpha^2) \\
g_{SShH} &= -4s_\alpha c_\alpha (\lambda_a - \lambda_b) \\
g_{SSG^0G^0} &= g_{SSG^+G^{+*}} = -4(\lambda_a c_H^2 + \lambda_b s_H^2) \\
g_{SSH_3^0H_3^0} &= g_{SSH_3^+H_3^{+*}} = -4(\lambda_a s_H^2 + \lambda_b c_H^2) \\
g_{SSG^0H_3^0} &= g_{SSG^+H_3^{+*}} = g_{SSH_3^+G^{+*}} = -4s_H c_H (\lambda_b - \lambda_a) \\
g_{SSH_5^0H_5^0} &= g_{SSH_5^+H_5^{+*}} = g_{SSH_5^{++}H_5^{++*}} = -4\lambda_b.
\end{aligned} \tag{A.5}$$

Here the shorthand notation $s_H \equiv \sin \theta_H$ and $c_H \equiv \cos \theta_H$ is used.