A tale of two dark neighbors: WIMP n’ axion

Suman Chatterjee, Anirban Das, Tousik Samui, and Manibrata Sen

a Tata Institute of Fundamental Research, Homi Bhabha Road, Mumbai, 400 005, India
b Regional Centre for Accelerator-based Particle Physics, Harish-Chandra Research Institute, HBNI, Chhatnag Road, Jhunsi, Allahabad – 211 019, India
c Department of Physics, University of California Berkeley, Berkeley, California 94720, USA
d Department of Physics and Astronomy, Northwestern University, Evanston, IL 60208, USA

E-mail: suman.chatterjee@tifr.res.in, anirbandas@theory.tifr.res.in, tousiksamui@hri.res.in, manibrata@berkeley.edu

ABSTRACT: We study the experimental constraints on a model of a two-component dark matter, consisting of the QCD axion, and a scalar particle, both contributing to the dark matter relic abundance of the universe. The global Peccei-Quinn symmetry of the theory can be spontaneously broken down to a residual $Z_2$-symmetry, thereby identifying this scalar as a stable weakly interacting massive particle, i.e., a dark matter candidate, in addition to the axion. We perform a comprehensive study of the model using the latest data from dark matter direct and indirect detection experiments, as well as new physics searches at the Large Hadron Collider. We find that although the model is mostly constrained by the dark matter detection experiments, it is still viable around a small region of the parameter space where the scalar dark matter is half as heavy as the Standard Model Higgs. In this allowed region, the bounds from these experiments are evaded due to a cancellation mechanism in the dark matter-Higgs coupling. The collider search results, however, are shown to impose weak bounds on the model.
1 Introduction

The evidence of Cold Dark Matter (CDM) is overwhelming from the cosmological data, even though its detection and identification continues to be one of the most interesting and challenging problems today [1]. Many particle dark matter (DM) models have been proposed over the last few decades, one of the oldest of them being the Weakly Interacting Massive Particle (WIMP) model [2–5] (for reviews, see [6–8]). In the WIMP scenario, the dark matter relic abundance is obtained through the annihilation of dark matter particles in the early universe with weak scale cross sections, and electroweak scale masses [2, 9–11]. The fact that one expects new physics at the electroweak scale from naturalness arguments makes the WIMP scenario a de facto solution to the dark matter problem [12].

The absence of CP-violation in the strong sector of the Standard Model (SM) is another long-standing puzzle in the particle physics community [13]. The null results of the neutron electric dipole moment measurement experiments so far restrict the value of the coefficient $\theta_{\text{QCD}}$ of the parity-violating $\mathbf{E} \cdot \mathbf{B}$ operator to be less than $10^{-10}$ [14]. In the present form of the SM, this is a fine-tuning problem since there is no symmetry that protects such a small number from large higher order corrections [15]. Therefore, a natural explanation of the smallness of strong CP violation is sought, and an elegant solution to this puzzle is given by the introduction of a global $U(1)$ Peccei-Quinn (PQ) symmetry [16–20]. This symmetry is spontaneously broken at a scale much larger than the electroweak scale by a scalar field, with the axion as the corresponding massless Nambu-Goldstone boson of this $U(1)_{\text{PQ}}$ symmetry. The coefficient $\theta_{\text{QCD}}$ is dynamic in this model and its small value is naturally attained in this way and is inversely proportional to the PQ scale. After the QCD condensation at a temperature of about $T \simeq 200$ MeV, the axion field gains a small
mass inversely proportional to the $U(1)_{PQ}$-breaking scale. In the early universe, axion can be produced non-relativistically through a coherent oscillation of the axion field due to the misalignment of the PQ vacuum. This is known as the misalignment mechanism of axion production [21, 22]. The axion is not completely stable, however, it has very feeble couplings with SM particles, thereby ensuring a lifetime longer than the age of the universe [23]. This makes the axion a very good CDM candidate, although the same feeble couplings make direct detection of these axions challenging [24].

In this work, we study a two-component DM model consisting of a WIMP and the axion as the DM candidates. As a simple realization of this, one can consider the KSVZ (Kim-Shifman-Vainshtein-Zakharov) model [19, 20] of axion with an additional scalar field charged under the $U(1)_{PQ}$ [25]. This additional scalar gets its stability from the residual $Z_2$-symmetry of the broken $U(1)_{PQ}$, and hence becomes a WIMP-like DM candidate [26]. Breaking of the $U(1)_{PQ}$ and the electroweak symmetry leads to a mixing between the Higgs and the radial part of the PQ scalar, which leads to interesting phenomenological consequences. The advantage of this model is that although the axions have very weak interactions with the SM, the coupling between this dark scalar and the SM Higgs doublet provides a portal to test this model in different DM detection experiments, both direct and indirect. The model can also give different signatures at collider experiments. For example, the KSVZ model predicts new colored, electroweak singlet quarks, which can be produced at colliders. Mixing with a scalar affects the properties of the Higgs boson, which can be directly used to constrain the mixing parameters. Furthermore, the dark scalar can also contribute to momentum imbalance in a collision event.

Hence, in the light of recent experiments, we explore the constraints on the WIMP-axion DM model, both from DM search experiments, as well as collider searches. Using the recent limits on DM direct detection from XENON1T×1yr experiment data [27], we find that the phenomenologically interesting mass range of $m_{DM} \gtrsim 100$ GeV is ruled out in such models. However, the stringent bounds from XENON1T×1yr data can be evaded in a small region of the parameter space where the scalar dark matter is half as heavy as the Higgs. This is a direct outcome of the mixing of the Higgs with the scalar, which leads to a cancellation mechanism in the Higgs portal coupling, thereby reducing the DM-nucleon scattering cross-section. As a result, while minimal scalar DM models are mostly ruled out by direct detection bounds [28], such WIMP-axion models can still survive with a reduced parameter space. Collider signals, on the other hand, are highly plagued by the backgrounds from the production of standard model particles, and hence the signals are not significant enough to be observed above the background [29–31].

The paper is organized as follows. Section 2 discusses the model, and the different parameters involved. Section 3 talks about the different experimental bounds, and how they constrain the parameters of the model. In section 4, we summarize the main results, and finally in section 5, we conclude.
|     | $\zeta$ | $\chi$ | $Q_L$ | $Q_R$ |
|-----|---------|--------|-------|-------|
| Spin | 0       | 0      | 1/2   | 1/2   |
| $SU(3)_C$ | 1   | 1       | 3     | 3     |
| $SU(2)_L$ | 1   | 1       | 1     | 1     |
| $U(1)_Y$   | 0   | 0       | $-1/3$| $-1/3$|
| $U(1)_{PQ}$ | 2   | 1       | 1     | $-1$  |

Table 1. New particles in the model and their charges. PQ charges of all the SM particles are zero.

2 The Model

We consider the KSVZ model of the axion, where electroweak singlet quarks $Q_L, Q_R$ and a complex scalar $\zeta$, both transforming under a global $U(1)_{PQ}$ symmetry, are added to the SM [19, 20]. These quarks are vector-like, hence do not introduce any chiral anomaly [32, 33]. We augment this model with a complex scalar $\chi = (\chi_1 + i\chi_2)/\sqrt{2}$ which is a SM singlet but charged under the $U(1)_{PQ}$ symmetry [25]. The axion $a$ is the Nambu-Goldstone mode of the scalar field $\zeta$, which can couple to the vector-like quarks, as well as $\chi$. As in the original KSVZ model, the axion can act as a CDM candidate [24]. The charges and quantum numbers of the new particles are listed in table 1.

The relevant part of the Lagrangian, governing the interactions of $Q_{L,R}, \zeta,$ and $\chi$ with the SM, is given by

\[
\mathcal{L} \supset -\lambda_H \left( |H|^2 - \frac{v_H^2}{2} \right)^2 - \lambda_\zeta \left( |\zeta|^2 - \frac{F_a^2}{2} \right)^2 - \lambda_{H\zeta} \left( |H|^2 - \frac{v_H^2}{2} \right) \left( |\zeta|^2 - \frac{F_a^2}{2} \right) - \lambda_\chi |\chi|^4 \\
-\mu_\chi |\chi|^2 - \lambda_{H\chi} |H|^2 |\chi|^2 - \lambda_{\zeta\chi} |\zeta|^2 |\chi|^2 + \left[ \epsilon_\chi \zeta^* \chi^2 + f_d \chi Q_L d_R + f_Q \zeta Q_L Q_R + \text{h.c.} \right].
\]

Here $H$ is the SM Higgs doublet and $d_R$ represents right-handed down-type quarks in the SM. After electroweak symmetry breaking via the Higgs vacuum expectation value (vev) $v_H$, one has $|H| = (h_0 + v_H)/\sqrt{2}$ where $h_0$ is the Higgs boson. Similarly, using the nonlinear representation, one can write $\zeta = e^{ia/F_a} (F_a + \sigma_0)/\sqrt{2}$, where $F_a$ is the $U(1)_{PQ}$ symmetry breaking scale as well as the axion decay constant, and $\sigma_0$ is the radial excitation of the $\zeta$ field. Constraints from supernova cooling data disfavor values of $F_a$ smaller than $10^{10}$ GeV [34].

After the breaking of both the symmetries, viz. electroweak and PQ symmetries, the interaction term between $H$ and $\zeta$ fields leads to mixing between $h_0$ and $\sigma_0$ with the mass matrix

\[
M^2 = \begin{pmatrix}
2v_H^2\lambda_H & F_a v_H \lambda_\zeta H \\
F_a v_H \lambda_\zeta H & 2F_a^2\lambda_\zeta
\end{pmatrix}.
\]
As a result of the mixing, the scalars in the mass basis are related to those in the flavor basis as
\[
\begin{pmatrix}
h \\
\sigma
\end{pmatrix} = \begin{pmatrix}
\cos \theta & -\sin \theta \\
\sin \theta & \cos \theta
\end{pmatrix}
\begin{pmatrix}
h_0 \\
\sigma_0
\end{pmatrix},
\tag{2.3}
\]
where the mixing angle, in the limit $F_a \gg v_H$, is given by
\[
\sin \theta \simeq \frac{v_H \lambda \zeta}{F_a}.
\tag{2.4}
\]

One obtains the masses of the physical states as
\[
m_h \simeq v_H \sqrt{2 \lambda_H \left(1 - \frac{\lambda_H^2}{4 \lambda_H \lambda \zeta}\right) + \mathcal{O}\left(\frac{v_H}{F_a}\right)},
\tag{2.5}
\]
\[
m_\sigma \simeq F_a \sqrt{2 \lambda \zeta} + \mathcal{O}\left(\frac{v_H}{F_a}\right).
\tag{2.6}
\]

Note that the mass $m_h$ of the mixed state $h$ is no longer $\sqrt{2 \lambda_H v_H^2}$, as predicted by the SM. Since $h$ is the physical state, we fix $m_h$ at 125 GeV and the Higgs vev $v_H$ at 246 GeV to match with the experimentally measured masses of the observed scalar [35, 36] and $W$, $Z$ bosons respectively [37]. The value of $\lambda_H$ is no longer the SM value, $\lambda_H^{\text{SM}} \simeq 0.13$, but is dependent on other parameters in this model and can be calculated using Eq. (2.5).

In fact, if we take $\lambda_H = \lambda_H^{\text{SM}} = \frac{m_h^2}{2 v_H^2}$, from Eq. (2.5) it is evident that $\lambda \zeta$ has to be zero, i.e., the SM Higgs does not mix with $\zeta$, as considered in [25]. Note that there is no underlying symmetry in the theory that allows us to set $\lambda \zeta$ to zero in the Lagrangian. More importantly, although the mixing is very small, the relation of the masses of the physical state with other model parameters plays a major role in imposing constraints on the model. Therefore, we do not neglect the mixing of $h_0$ with $\sigma_0$ in this study.

The masses of $\chi_1$ and $\chi_2$ are given by
\[
m_{\chi_1,2}^2 = \frac{1}{2} \left(2 \mu_\chi^2 + v_H^2 \lambda_H + F_a^2 \lambda \zeta + 2 \sqrt{2} F_a \epsilon_\chi\right),
\tag{2.7}
\]

Without loss of generality, we can take $\epsilon_\chi > 0$ such that $m_{\chi_1} < m_{\chi_2}$, hence $\chi_1$ can be the DM candidate, and we, henceforth, denote the mass of $\chi_1$ as just $m_\chi$. Note that after the PQ-symmetry breaking, the Lagrangian in Eq. (2.1) has a residual $\mathbb{Z}_2$-symmetry which stabilizes $\chi_1$. Note that in Eq. (2.7), $\mu_\chi^2$ is defined to be negative and hence cancels out the large contribution coming from $F_a$. This type of fine-tuning is a general feature of these axion models [25]. Since the fine-tuning is required mainly in the dark sector, we do not explore it further and defer the details to a later work. Furthermore, one can also motivate a tiny value of $\epsilon_\chi$ from naturalness arguments. As $\epsilon_\chi \to 0$, one obtains an extra $U(1)$ symmetry in the theory, apart from the $U(1)_{\text{PQ}}$. This can allow $\epsilon_\chi$ to be naturally small.

The mass of the axion is obtained through non-perturbative QCD effects and is inversely proportional to $F_a$,
\[
m_a \simeq 0.6 \text{ meV} \times \left(\frac{10^{10} \text{ GeV}}{F_a}\right).
\tag{2.8}
\]
The couplings of the axion to SM particles are also suppressed by inverse power of $F_a$, so the decay lifetime of the axion is very large. In fact, if we take the value of $F_a > 10^{10}$ GeV, as allowed by the supernova cooling data [34], its lifetime becomes larger than the age of the universe. Thus, the axion also acts as a viable candidate for CDM in this model. Therefore, both $\chi_1$ and the axion will contribute to the total DM relic density in the universe.

Finally, the vector-like quarks obtain their mass $m_Q = f_Q F_a / \sqrt{2}$, as $\zeta$ develops a vev. If this mass is $\sim O$(TeV), they can be produced at the LHC. This is expected to give direct constraints on this model, however, in order to have a mass of $\sim O$(TeV), the coupling $f_Q$ needs to be extremely tiny $\sim O(10^{-6})$.

The new interactions introduce two portals connecting the SM and the dark sector through the Higgs (via the $h\chi_1\chi_1$) and the down-type quark (via the $\chi_1\bar{Q}_{L,R}$). Of the two, the $h\chi_1\chi_1$ interaction is the more important one and will play a key role in our analysis. The $h\chi_1\chi_1$ coupling is given by

$$g_{h\chi_1\chi_1} = i \left( F_a \lambda_\zeta \sin \theta - v_H \lambda_{\chi H} \cos \theta - \sqrt{2} \epsilon \sin \theta \right).$$  \hspace{1cm} (2.9)

Though $\sin \theta$ is small, the first term cannot be ignored due to the large scale $F_a$. Using the approximation for $\sin \theta$ in Eq. (2.4), we obtain

$$g_{h\chi_1\chi_1} \simeq i v_H \left( \frac{\lambda_{\chi H} \lambda_\zeta}{2 \lambda_\zeta} - \lambda_{\chi H} \right).$$  \hspace{1cm} (2.10)

Note that in the presence of nonzero $\lambda_{\zeta H}$, the $h\chi_1\chi_1$ coupling vanishes at

$$\lambda_{\chi H} = \frac{\lambda_{\chi H} \lambda_\zeta}{2 \lambda_\zeta},$$  \hspace{1cm} (2.11)

as opposed to at $\lambda_{\chi H} = 0$ as in [25]. This shift will play a crucial role in the analyses to follow.

Using Eq. (2.5), $\lambda_\zeta$ can be written in terms of $m_h$, $\lambda_{\zeta H}$, and $\lambda_H$. This gives a family of solutions, satisfying Eq. (2.11). In figure 1, we show four contours of $\lambda_H$ in the $\lambda_{\zeta H} - \lambda_{\chi H}$ plane for a given value of $\lambda_{\chi H} = 0.1$. Any point on these hyperbolae satisfies Eq. (2.11), leading to vanishing $h\chi_1\chi_1$ coupling. The benchmark point chosen for further analysis, $\lambda_{\zeta H} = 0.1, \lambda_{\chi H} = 0.14$ and $\lambda_H = 0.2$, is shown as a black circle on the plot. One can, in principle, probe other values of $\lambda_H$ in this parameter space, and we do not show them here for clarity. However, one should not take $\lambda_H < \lambda_{\chi H}^{SM} \simeq 0.13$ since it leads to negative values of $\lambda_\zeta$, thereby making the potential for $\zeta$ unstable.

Finally, note from Eq. (2.6) that the mass of $\sigma$ is proportional to the $U(1)_{PQ}$-breaking scale $F_a$. So if $\lambda_\zeta \sim O(1)$, $\sigma$ becomes very heavy and decouples from the low energy theory. Therefore, for all practical purposes, $\sigma$ does not play any significant role in present experiments. However, it is possible to have the mass of $\sigma$ at around TeV, but only within a highly fine-tuned region of the parameter space.

One may wonder as to how much fine-tuning might be necessary for this scenario. Without going into details, we provide a back-of-the-envelope estimate here. From Eq. (2.2), if $\lambda_\zeta \sim 10^{-14}$, then both the scalars $h$ and $\sigma$ can have a mass $\sim O(100)$ GeV.
However, in order to keep the physical masses real, i.e., both the eigenvalues of the mass matrix positive, the off-diagonal terms have to be of the same order as the diagonal terms. This requires $\lambda_{\zeta H}$ to be further fine-tuned to values $\sim 10^{-7}$. However, such small values of $\lambda_\zeta$ and $\lambda_{\zeta H}$ will raise the value of $g_{h\chi_1\chi_1}$ (see Eqs. (2.9) and (2.10)) to values $\gg 1$, which makes the whole problem highly non-perturbative. Then, one would again need to choose $\lambda_{\zeta \chi}$ unnaturally small to solve this issue.$^1$

Since the above scenario is fine-tuned, we do not pursue it here. Rather, we consider natural values of all couplings $\lesssim \mathcal{O}(1)$. As a result, in this work, the heavy scalar $\sigma$ decouples early on and does not enter our analysis.

3 Experimental Probes of Dark Matter

Naturally, this model will have vast implications for dark matter search experiments. In addition, the LHC search for heavy vector-like particles, as well as missing energy searches, will also test this model. Using \texttt{FeynRules} [38, 39] to implement the model, we constrain it with the latest results from these experiments. Broadly, three avenues are explored:

1. DM direct and indirect detection experiments set upper limits on the observed cross section of its interaction with ordinary matter. This will be used to constrain the relevant parameters.

$^1$Note that the results given in Eqs. (2.5–2.10) were obtained in the limit $F_a \gg v_H$. This approximation breaks down when the $\lambda$s are set to such small values. Hence one has to start from the mass matrix in Eq. (2.2) and proceed without any approximation to arrive at this conclusion.
2. Mixing between \( h_0 \) and \( \sigma_0 \) changes the couplings of the observed 125 GeV scalar from that of the SM Higgs. This leads to changes in the properties of the observed scalar measured in the collider experiments from that of SM Higgs. This will also constrain the parameters of the model.

3. Since the masses of the DM and the vector-like quarks are lighter or near TeV range, they can potentially be produced at the LHC. Non-observation of such particles will limit the model parameter space.

The rest of the section discusses these types of experimental constraints in details.

### 3.1 Dark Matter Relic Abundance

After the \( U(1)_{PQ} \)-symmetry breaking, the axion \( a \), being a Nambu-Goldstone, enjoys a continuous shift symmetry. This symmetry is broken explicitly as a result of the chiral symmetry breaking in the QCD sector, and a temperature-dependent potential for the axion is generated from non-perturbative QCD effects \[40\]. But the axion field does not start rolling in the potential and remains frozen at its initial value until its mass becomes larger than the Hubble expansion rate \( H(t) = \dot{R}/R \) where \( R(t) \) is the scale factor of the universe. After the epoch when \( m_a(t) \approx H(t) \), the field starts oscillating coherently and the axion particles are produced with non-relativistic speed. They contribute towards the CDM abundance today and their density is approximately given by \[24, 41\],

\[
\Omega_a h^2 \approx 0.18 \theta_a^2 \left( \frac{F_a}{10^{12}\text{GeV}} \right)^{1.19}.
\]  

(3.1)

Here \( \theta_a \) is the initial misalignment angle of the axion field relative to the minimum of the axion potential. For simplicity, we shall assume \( \theta_a \sim 1 \) in the rest of the paper \[42\]. In order that the axions do not overproduce DM in the universe, the PQ breaking scale \( F_a \) has to be less than \( 10^{12} \text{GeV} \). In this work, we will focus on \( 10^{10} \text{GeV} \leq F_a \leq 10^{12} \text{GeV} \).

As already noted, \( \chi_1 \) gains stability from the residual \( Z_2 \)-symmetry and is a DM candidate. In the early universe, \( \chi_{1,2} \) are in chemical equilibrium with the thermal bath of the SM particles. As the temperature of the universe decreases below \( \sim m_\chi/20 \), their rate of interaction drops below the expansion rate and \( \chi_{1,2} \) cease being in equilibrium with the SM particles. The heavier component \( \chi_2 \), however, does not remain stable since it decays to \( \chi_1 \), which then forms the relic abundance \( \Omega_\chi h^2 \). The relic abundance is formed after the freeze-out of \( \chi_1 \chi_1 \) annihilations. The annihilation can be mediated by \( h \) as well as \( \sigma \). However, the \( h \)-mediated process dominates, since \( m_\sigma \gg m_h \). The relic abundance, being governed by \( \chi_1 \chi_1 \rightarrow \text{SM SM} \), depends directly on \( m_\chi \).

We show the dependence of the \( \chi_1 \) relic density as a function of its mass \( m_\chi \) in figure 2. We used \textsc{micrOMEGAs5.0} \[43\] to numerically compute \( \Omega_\chi h^2 \). The behavior for very small and large \( m_\chi \) can be understood as follows. For very small values of \( m_\chi (\approx \text{few GeV}) \), \( \chi_1 \) can annihilate only into light quarks and the cross section is suppressed by the small Yukawa couplings resulting in overabundance of \( \chi_1 \). For \( m_\chi \gg m_t \), the annihilation cross section is \( 1/m_\chi^2 \) suppressed. Since the relic abundance is inversely proportional to the
annihilation cross section, we expect the region around $m_\chi \approx 100$ GeV to give the correct ballpark value of the desired relic abundance.

The sharp dip at $m_\chi \simeq m_h/2 \simeq 62.5$ GeV is due to the $s$-channel resonance from the $h$ propagator. As $m_\chi$ increases further from 62.5 GeV, the cross section falls leading to sharp increase in the relic. When the $\chi_1$ is heavier than $h$, the new annihilation channel $\chi_1\chi_1 \rightarrow hh$ opens up and dominates over all other channels. As a result, the relic abundance decreases, leading to the second dip. As $\chi_1$ becomes more massive, the relic increases again because of the decrease in annihilation cross section with the characteristic $1/m_\chi^2$ suppression. Note that we do not consider $m_\chi > M_Q$, since the colored $Q_L,R$ can become the lightest dark sector particle.

In our analysis, we take the Planck (TT, TE, EE, lowP) measurement of the CDM energy density $\Omega_c h^2 = 0.12 \pm 0.0012$ represented by the horizontal line labeled as $\Omega_c h^2$ in figure 2 [1]. The over-abundance region, shown as a gray shade, is disallowed. However, the under-abundance region is allowed since the axion abundance $\Omega_a h^2$ can account for the rest of the relic. Therefore, the observed relic abundance $\Omega_c h^2$

$$\Omega_c h^2 = \Omega_\chi h^2 + \Omega_a h^2 \quad (3.2)$$

We note that $\Omega_\chi$ is virtually independent of $F_a$ due to $v_H/F_a$ suppression in the couplings and mixing angle. Hence, $F_a$ is fixed by Eq. (3.2) via the $\Omega_a h^2$ term.
3.2 Direct Detection of Dark Matter Particles

The DM direct detection (DD) experiments look for scattering between the DM particle and nuclei of the detector material. Any interaction between the DM and the SM quarks/gluons in a given model leads to a possible signal in the direct detection experiments. Non-observation of such a scattering signal in such experiments constrains the parameters of the model. In the present case, the dominant channel of interaction arises again through the $h\chi_1\chi_1$ coupling, since $h$ mediates the DM and SM quark scatterings. A typical behavior of the scattering cross section as a function of $\lambda_{\chi H}$ is shown on the left panel of figure 3.

The cross section $\sigma_{\chi N}$ is constant for very small $\lambda_{\chi H}$ because the coupling becomes independent of $\lambda_{\chi H}$. For very large $\lambda_{\chi H}$, the cross section increases as $\sim \lambda_{\chi H}^2$, as expected. In between, a dip occurs because of the cancellation of two terms appearing in the vertex factor of $h\chi_1\chi_1$ coupling (see Eq. (2.10)). Presently, the most stringent bound on this cross section is given by the XENON1T×1yr experiment [27]. It is most sensitive to the DM mass in the range $10\text{ GeV} - 1\text{ TeV}$ and the strongest upper bound quoted is $\sigma_{\chi N} \simeq 10^{-46}\text{ cm}^2$.

We will show later that due to the stringent constraint, the only experimentally allowed region of DM mass turns out to be around $m_\chi \simeq 62.5\text{ GeV}$.

The right panel of figure 3 shows the $\chi_1$-nucleon scattering cross section $\sigma_{\chi N}$ as a function of $m_\chi$ for two different values of $\lambda_{\chi H}$. Note that in this model, $\chi_1$ forms only a fraction $f_\chi(\equiv \Omega_\chi/\Omega_c)$ of the present dark matter abundance. Therefore the XENON1T bound is to be accordingly divided by $f_\chi$ before applying to this model.

All the above bounds apply for $\chi_1$ as the DM candidate. However, direct detection experiments for axion need to follow a different search strategy because of its ultra-low...
mass. There have been a few experimental efforts to look for axionic dark matter. For example, the ADMX experiment [44] uses RF cavity to look for its interaction with the electromagnetic field. In the KSVZ model, this interaction strength is given by [19, 20]

$$g_{a\gamma} = -1.92 \frac{\alpha}{2\pi F_a},$$

(3.3)

where $\alpha$ is the fine structure constant. Presently, ADMX rules out a narrow region of the parameter space above $g_{a\gamma} \simeq 10^{-15}$ GeV$^{-1}$ ($F_a \simeq 10^{12}$ GeV) around $m_a \simeq 2 \mu$eV. For higher mass axion, the bound is even weaker. Another proposed experiment is CASPER-Electric which will probe $F_a \gtrsim 10^{12}$ GeV for lighter axions [45]. Moreover, we should remember that these bounds assume that 100% CDM abundance is given by axion which is not be true in our model. These bounds are weaker than the upper limit on $F_a$ from the dark matter relic abundance, even after adjusting for correct factor to cancel out the assumption, hence does not need a special attention.

### 3.3 Dark Matter Annihilation Signal

Various astrophysical observations hint that the present day universe consists of galaxies sitting inside halo-like structures formed by gravitational clustering of DM particles [47]. At the center of these halos, the DM density is high enough to scatter with each other and annihilate into SM particles. These final state particles would further decay and give rise to gamma-ray signals from various astrophysical objects, such as dwarf galaxies, the Milky Way center etc. We focus on bounds arising from gamma-ray signals due to such annihilations of DM particles.

---

**Figure 4.** The annihilation rate of $\chi_1\chi_1$ into $b\bar{b}$ in this model as a function of the mass of $\chi_1$ for two values of $\lambda_{\chi H}$. The sharp peak is due to the s-channel resonance from the SM Higgs. Most stringent upper bound on this cross section is provided by the dwarf galaxy observation of the Fermi-LAT satellite data which is shown as the gray shaded region [46].
We pay more attention to the DM mass around $m_\chi \simeq m_h/2 = 62.5$ GeV which is still allowed by the direct detection experiment data. The total annihilation is dominated by the $b\bar{b}$-channel ($\sim 90\%$), which is shown in figure 4. Note that here also the annihilation cross section is enhanced due to the s-channel resonance from the SM Higgs propagator. Hence the largest annihilation signal is predicted at this mass. The dependence on $\lambda_{\chi H}$ comes through the $g_{h\chi_1\chi_1}$ coupling. The Fermi-LAT constraint becomes ineffective for the value of $\lambda_{\chi H}$ for which this coupling vanishes (the red curve in figure 4) as is discussed in section 2.

There have been many experiments which have looked for DM annihilation signals from various astrophysical objects [46, 48–50]. At present, the most stringent upper bounds on the thermally averaged DM annihilation cross section $\langle \sigma v \rangle$ is given by the DES-Fermi-LAT joint gamma-ray search from the satellite galaxies of the Milky Way [46]. It is derived from 6 years observation of 45 such objects by the LAT. They have relatively less amount of visible baryonic matter and the DM population is expected to dominate their matter density. In figure 4, we show this upper bound on the annihilation cross section as the gray shaded region. This does not rule out most part of our parameter space, except a region of $m_\chi$ around Higgs mass. In passing we also note that the DM mass needed for the resonantly enhanced annihilation signal in the $b\bar{b}$-channel matches the result of the galactic center excess analysis done in ref. [51] within $1\sigma$ C.L. (also see [52]).

### 3.4 New Physics Searches at the LHC

In this subsection, we will focus on various signatures of the model at the LHC. The model has an extended scalar sector: apart from the SM Higgs boson $h_0$, there exists a scalar DM candidate $\chi_1$ and its heavier counterpart $\chi_2$, and another scalar field $\sigma_0$, which is the radial component of $\zeta$. As discussed earlier, $h_0$ and $\sigma_0$ mixes with each other giving rise to physical states $h$ and $\sigma$. The mixing between $\sigma_0$ and $h_0$ changes the properties of $h$ from that of the SM Higgs via its coupling to SM particles as well as to the new states present in this model. Since various properties of the observed scalar particle at the LHC resemble that of the SM Higgs boson, we expect some constraints on the parameter space of the model from the measurement of the properties of the observed 125 GeV scalar.

One of the measurements that provides relevant information about the properties of the observed 125 GeV scalar is its signal strength. If the scalar decays to $X \in \{\ell^\pm, q, g, Z, W\}$ and its conjugate, $\bar{X}$, its signal strength is defined as

$$\mu_{X\bar{X}} = \frac{\sigma_{\exp}(pp \rightarrow h) \times BR_{\exp}(h \rightarrow X\bar{X})}{\sigma_{\SM}(pp \rightarrow h) \times BR_{\SM}(h \rightarrow X\bar{X})},$$

(3.4)

where $\sigma_{\exp}$ stands for the experimentally observed cross section of the process $pp \rightarrow h$, and $BR_{\exp}$ is the experimentally observed branching ratio of the process $h \rightarrow X\bar{X}$. Similarly, $\sigma_{\SM}$ and $BR_{\SM}$ in Eq. (3.4) stands for the corresponding values predicted in the SM. We compare observed $\mu_{X\bar{X}}$ with the theoretically calculated $\mu_{X\bar{X}}$ from the model in different decay channels.

Due to the mixing, the physical scalar $h$ will have a $\cos \theta$ component in all the couplings with the SM. An additional decay mode of $h$ to $\chi_1\chi_1$ is possible if $m_\chi < m_h/2$. If the
Table 2. Measured values of the signal strengths of the 125 GeV observed scalar. The superscripts represent the production modes and the subscripts indicate the decay modes of the observed scalar $h$. The measurements are done by ATLAS and CMS at the LHC with $\sim 36$ fb$^{-1}$ luminosity at $\sqrt{s} = 13$ TeV.

| Production Mode | ATLAS       | CMS        |
|-----------------|-------------|------------|
| $(ggF)$          | $1.21^{+0.22}_{-0.21}$ [53]  | $1.38^{+0.21}_{-0.24}$ [54]  |
| $(ggF)$          | $1.11^{+0.23}_{-0.22}$ [55]  | $1.20^{+0.22}_{-0.21}$ [56]  |
| $(ggF+VH+VBF+ttH)$ | $0.99^{+0.15}_{-0.14}$ [57]  | $1.18^{+0.17}_{-0.14}$ [58]  |
| $(VH)$           | $1.20^{+0.42}_{-0.36}$ [59]  | $1.06^{+0.31}_{-0.29}$ [60]  |
| $(ggF+VH+VBF)$   | $1.43^{+0.43}_{-0.37}$ [61]  | $1.09^{+0.27}_{-0.26}$ [62]  |

The partial decay width of the new decay modes of $h$ is $\Gamma_{\text{new}}$, the signal strength of $h$ decaying to any SM particle pairs $XX$ can be written as

$$\mu_{XX} = \frac{\cos^2 \theta}{1 + \frac{\Gamma_{\text{new}}}{\cos^2 \theta \Gamma_{\text{SM}}}}$$  \hspace{1cm} (3.5)$$

where $\Gamma_{\text{SM}}$ is the total decay width of SM Higgs boson.

In table 2, we tabulate the recent measurements of signal strength of the observed scalar $h$ by both ATLAS and CMS collaboration at 13 TeV with $\sim 36$ pb$^{-1}$ integrated luminosity in different decay channels of $h$. The superscripts in the $\mu_{XX}$ represent the production mode of the scalar $h$. For our analysis, we constrain the parameter space by imposing the value to be at 95% C.L. of the measured values, i.e., with $\pm 2\sigma$ around the measured central value. Since, in the model, $\mu_{XX}$ is always below unity, it is the lower bound at 95% C.L. which will actually put constraints on the parameters.

In the left panel of figure 5, we show the variation of the signal strength of $h$ in $WW^*$ channel as a function of $\lambda_{1}H$ for two different masses of $\chi_1$. As expected from Eq. (3.5), the variation is a Lorentzian, with a narrow width governed by $\Gamma_{\text{SM}}$ and $m_\chi$. Since the coupling for $h$ to $\chi_1\chi_1$, as given in Eq. (2.10), vanishes at $\lambda_{1}H = 2\lambda_{1}\chi_1 \left(\lambda_H - \lambda_{1}^M\right) / \lambda_H$ ($\approx 0.14$ for the chosen benchmark point), the decay mode for the $h$ vanishes at that point, and hence the $\mu_{XX}$ becomes 1 around that point. The gray (green) shaded region shows the area disallowed at 95% C.L. by the measurements by CMS (ATLAS) as indicated in the plot, and the allowed region is shown in white. Although the measurements for different decay channels of $h$ are listed in table 2 for completeness, we only plotted $\mu_{WW^*}^{(ggF)}$, which gives the strongest bounds from the signal strength measurement.

We also study the bounds from the invisible decay of $h$ which arises from the decay
Figure 5. Bounds arising from (left) the Higgs signal strength in $WW^*$ channel and (right) the invisible decay of the Higgs. The gray (green) shaded regions in both the plots are excluded by CMS (ATLAS) measurement at 95% C.L. The allowed regions are shown in white.

| BR($h \rightarrow \text{inv}$) | ATLAS   | CMS     |
|-------------------------------|---------|---------|
| 0.67 [63]                     | 0.24 [64]|

Table 3. Observed upper limit on the branching ratio of invisible decay of the scalar $h$.

channel $h \rightarrow \chi_1\chi_1$ for $m_\chi < m_h/2$ in this model. The BR of the decay can be written as

$$
\text{BR}(h \rightarrow \chi_1\chi_1) = \frac{1}{1 + \cos^2 \theta \frac{\Gamma_{\text{tot}}^{\text{SM}}}{\Gamma_{\text{new}}}}.
$$

(3.6)

The dependence of BR($h \rightarrow \chi_1\chi_1$) with the parameter $\lambda_{\chi H}$ is plotted in the right panel of figure 5 for two different masses of $\chi_1$. As in the case with the signal strength, the BR($h \rightarrow \chi_1\chi_1$) vanishes at the point where the coupling of $h$ to $\chi_1\chi_1$, given by $g_{h\chi_1\chi_1}$ (see Eq. (2.10)), goes to zero. This feature is evident from the plot in the right panel of figure 5. Away from this point, the BR increases in both sides, tending to unity for high value of $g_{h\chi_1\chi_1}$, which indicates that $\Gamma_{\text{new}}$ is the dominant decay mode, and all other modes are suppressed.

Non-observation of this decay mode of the observed 125 GeV scalar at the LHC, therefore, places upper limit on the invisible decays of $h$. These upper limits are tabulated in table 3. In the right panel of figure 5, the gray (green) shaded region is the area disallowed at 95% C.L. by CMS [64] (ATLAS [63]) measurements on the invisible decays of 125 GeV scalar. It is therefore clear that only a small range of $\lambda_{\chi H}$, for which the BR curves fall within the white region, is allowed by current measurements.

At this point, it is worth mentioning that the trilinear coupling of $h$ is also modified due to the mixing with $\sigma_0$, which will change the di-Higgs production rate. Measurements
for the trilinear coupling of $h$ as well as di-Higgs production have been carried out by both ATLAS\cite{65} and CMS\cite{66} in the di-Higgs channel. However, the upper bounds are well above the SM prediction due to lack of signal in the di-Higgs channel. Hence, much of the parameter space, especially the region of interest, of the model is not constrained by the measurement of trilinear coupling of $h$.

The model also predicts new particles at around GeV-TeV range, which can potentially be observed in a TeV collider. One such particle is the DM candidate, $\chi_1$, which is weakly interacting and does not decay within the detector. If it is produced in the collider, it will not be detected and will contribute to the missing momentum in an event. The other particles, within the observable range of TeV collider, are the vector-like quarks $Q_L$ and $Q_R$. Since these quarks are colored, they can be produced in a hadron collider and subsequently decay to a down-type quark and a $\chi_1$. Presence of $\chi_1$ will again contribute to the missing energy in the detector. Lack of agreement of such signals with those predicted at the TeV colliders will also put bounds on the parameter space of the model in consideration.

Now, we turn to the discussion of direct production of the new particles at the LHC. The new particles, being charged under a PQ symmetry, should be produced in pairs. There are three different pairs of new particles that can be directly produced: $Q\bar{Q}$, $Q\chi_1$, and $\bar{Q}\chi_1$. Hence, these processes will contribute to the following final states: dijet ($2j$)+MET in case of $Q\bar{Q}$ production, and monojet ($j$)+MET final state in case of $Q\chi_1$ and $\bar{Q}\chi_1$ production, where MET stands for missing transverse energy. In the rest of this section, we will discuss the constraints on the parameter space in view of the observation of the above-mentioned final states at collider.

Since the $Q$s are colored, the cross section for the production of $Q\bar{Q}$ will be similar

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure6.png}
\caption{Variation of total production cross section for $Q\bar{Q}$ (in red) and for $Q\chi_1$ and $\bar{Q}\chi_1$ (in blue) in dijet+MET and in monojet+MET channels respectively at the LHC at $\sqrt{s} = 13$ TeV as a function of $M_Q$.}
\end{figure}
to that of the SM quarks and will be suppressed for higher masses. Figure 6 shows the variation of total production cross section for $Q\bar{Q}$ (in red) and for $Q\chi_1$ and $\bar{Q}\chi_1$ (in blue) in $2j+\text{MET}$ and in $j+\text{MET}$ channels respectively at the LHC at 13 TeV. The production cross section of $Q\bar{Q}$ in $2j+\text{MET}$ channel have negligible dependence on $f_{d,s,b}$ since the dominant parton-level process for the production is $gg \rightarrow Q\bar{Q}$, which is independent of $f_{d,s,b}$. Hence, the two red curves, solid for $f_{d,s,b} = 0.1$ and dashed for $f_{d,s,b} = 1$ coincide with each other. However, the cross section for $Q\chi_1$ and $\bar{Q}\chi_1$ in $j+\text{MET}$ channels scales as $f_{d,s,b}^2$ since the parton level process involved in the production is $gq, g\bar{q} \rightarrow Q\chi_1, \bar{Q}\chi_1$, whose amplitude is proportional to $f_{d,s,b}$. Note that the only possible decay mode of $Q$ is to a down-type quark and a $\chi_1$.

To estimate the signature of our model in collider experiments, events have been gen-
erated at partonic level using \texttt{MadGraph5} \cite{MadGraph5} with NNPDF2.3LO parton distribution function \cite{NNPDF2.3LO} using the UFO files generated by \texttt{FeynRules} \cite{FeynRules,FeynRules2} at center-of-mass energy of 13 TeV; partons in the final state have been showered and hadronized using the parton shower in \texttt{PYTHIA 8.210} \cite{PYTHIA} with 4C tune \cite{4C}. Stable particles have been clustered into anti-kT \cite{anti-kT} jets of size 0.4 (used by both ATLAS and CMS) using \texttt{FastJet} \cite{FastJet} software package; only the jets with $P_T$ more than 30 GeV have been considered for further analysis.

In figure 7, we present some important and representative differential distribution of some observables as are considered by experimental collaborations to search for signals. The top-left panel in the figure shows distribution of $P_T$ of the leading jet while the panel in top-right shows the distribution for $P_T$ of the second jet. In the bottom-left panel, we show the distribution of missing transverse energy ($\not{E}_T$). The bottom-right panel shows the distribution of $H_T = \sum_{j \in \text{jets}} |\vec{p}_T|$, which is the scalar sum of $P_T$ of all the jets. The major sources of the SM backgrounds for jets+MET are from the production of $Z$ decaying to $\nu\bar{\nu}$ and $W$ decaying to $\tau\nu$ in events with jets. Also QCD events are potential sources to contribute to the same final state. The distribution for these three backgrounds are plotted in four panels of figure 7. SM background samples have been generated with at leading order (LO) using \texttt{MadGraph5} \cite{MadGraph5} with NNPDF2.3LO parton distribution function \cite{NNPDF2.3LO} at center-of-mass energy of 13 TeV and \texttt{PYTHIA 8.210} \cite{PYTHIA}, with the same 4C tune \cite{4C} as used for generation of the signal sample, has been used for the simulation of fragmentation, parton shower, hadronisation and underlying event. The distribution for QCD, $W+$jets, and $Z+$jets backgrounds are plotted with gray, purple, and green respectively with the same color convention in all the four panels. From the figure, it is quite clear that the bumps for signals will not be significant enough to be observed above the expected fluctuation of the background.

Following the distribution in the experimental references \cite{ExperimentalReferences}, we carried out our analysis with the same distribution. As discussed earlier, the direct production of new particles will contribute to $2j+$MET and $j+$MET signals. There are few dedicated search in these channels to search for dark matter signals \cite{ExperimentalReferences}. Few other models, especially SUSY in R-parity conserving scenario, also lead to these kinds of signals. These searches have also been done by both CMS \cite{CMS_SUSY, CMS_SUSY2} and ATLAS \cite{ATLAS_SUSY, ATLAS_SUSY2}. Though the results are given in terms of SUSY parameters or effective theory parameters, one can recast the result for a given model and check for its consistency. But these searches do not yield any further constraint in the parameter space in the model. A dedicated search for this model may give a stronger constraint, but the analysis of such search is beyond the scope of this work.

4 Results

Our main results are summarized in figure 8. The relevant bounds coming from the different experiments are imposed on the region satisfying the DM relic density in the $\lambda_{\chi H} - m_\chi$ plane. The gray shaded region is ruled out by the relic constraints. We allow for both $\chi_1$ as well as the axion to contribute to the DM relic density. Hence the white region, corresponding to the $2\sigma$ bound $\Omega_c h^2 < 0.12$, represents the allowed parameter space,
Figure 8. Allowed regions in the parameter space for the two-component axion-WIMP DM model. The gray shaded region shows the area ruled out by DM relic abundance constraint corresponding to the 2σ bound $\Omega_c h^2 < 0.12$ [1]. The black hatched lines show the regions of parameter space ruled out by the DM direct detection bounds from XENON1T×1 yr experiment [27]. The hatched region within the red curve is ruled out by the DM annihilation data from DES-Fermi-LAT experiment [46]. The blue shaded region show the bounds imposed due to the invisible decay modes of the Higgs, which is roughly 25% of its branching ratio [53, 54]. The bound coming from the signal strength of the Higgs is shown in orange [63, 64]. The white, unshaded region represents the allowed parameter space in this model.

satisfying the relic density. As explained before, near $m_\chi \approx m_h/2$, the DM annihilation cross section is enhanced from the Higgs resonance, thereby decreasing the relic density of DM. This explains why the allowed region from relic is centered around $m_\chi = m_h/2$. Furthermore, there is a particular set of parameters for which $h\chi_1\chi_1$ coupling vanishes, leading to a rise in the relic density. This accounts for the peak-like structure in figure 8, which occurs at $\lambda_{\chi H} \sim 0.14$ for our choice of parameters.

The black hatched lines show the regions of parameter space ruled out by the direct detection bounds from XENON1T×1 yr experiment. The hatched region within the red curve is ruled out by DES-Fermi-LAT joint gamma-ray search data from the Milky Way satellite galaxies. As is clearly seen, most of the allowed regions are ruled out, leaving behind a tiny window around in the $m_\chi - \lambda_{\chi H}$ plane. Clearly, this window is centered around $m_\chi \approx m_h/2$ and the value of $\lambda_{\chi H}$ for which the $h\chi_1\chi_1$ coupling vanishes.

The blue shaded region shows the bounds imposed due to the invisible decay modes of the Higgs, which is roughly 25% of its branching ratio. More stringent bounds are imposed from the signal strength of the Higgs, which is shown in orange. These also help to rule
out extra regions of the parameter space for larger as well as smaller values of $\lambda_{\chi H}$. We have also checked that the LHC bounds from production of $Q\bar{Q}$ are relatively weak, hence they do not impose any extra constraint on the model.

Thus, from the above figure, one concludes that only a small fraction of the model can still be accommodated from existing experimental bounds. This region, however, enjoys the advantage of an accidental cancellation of the couplings near $m_h/2$, thereby making it extremely difficult to rule out experimentally. This tiny window provides a breathing space for the model to survive.

5 Summary and Discussions

In this paper, we have performed a comprehensive study of a two-component dark matter model, consisting of the QCD axion, and an electromagnetic charge neutral scalar particle, both contributing to the relic density. The theory is symmetric under a global Peccei-Quinn symmetry, which can be spontaneously broken down to a residual $\mathbb{Z}_2$ symmetry. For concreteness, we have considered a specific model: the KSVZ model of the axion, augmented with an additional complex scalar. After spontaneous breaking of the PQ symmetry, the residual $\mathbb{Z}_2$ symmetry allows the lightest component of the complex scalar to be a DM candidate, apart from the axion. We have tested the model in the light of recent data from DM direct and indirect search experiments. Furthermore, we have also studied the different collider signatures of this model.

Although the observational and experimental constraints are found to be very restrictive, a synergy of the enhancement of DM annihilation from the Higgs resonance, and the vanishing of the coupling between the Higgs and the dark matter leave room for future experimental investigation of this model. A large portion of the parameter space predicts overabundance of $\chi_1$ in the universe and hence is not viable. In the remaining underabundant region of $\chi_1$, the axion can form the dominant part of the CDM. The viability of the axion being the CDM is being tested in several ongoing experiments. The latest dark matter direct and indirect detection experiments results further constrain this model. Moreover, these results are expected to improve the bounds by few orders of magnitudes over the next few years which will subject this model to even tighter constraints. Although the bounds from the measurements of the properties of the Higgs at collider experiments are relatively weak, they still help to rule out an additional part of the parameter space. Future measurements of vector-like quarks at high luminosity and high energy operating modes of the LHC can shed further light on the viability of this model.

Nevertheless, it is possible to add new particles to this simplistic model, e.g., an additional scalar, to enrich its phenomenology and evade some of the experimental bounds. This leaves room for future scopes of model-building and investigation of observable signatures in high energy experiments. In this work, we have calculated the prediction from our model with some natural choice for the couplings to compare with experimental data; but other values of the couplings can be explored to test the validity of the model on the basis of available experimental results. In conclusion, the two-component dark matter
model, consisting of the WIMP and the axion, continues to survive, in spite of being tightly constrained.

Acknowledgments

We thank the Workshop on High Energy Physics Phenomenology 2017 for providing us with an environment for lively and fruitful discussions. We thank Sabyasachi Chakraborty for relevant discussions in the initial stages of this project, and Basudeb Dasgupta for useful inputs to the project and comments on the manuscript. We also thank the grid computing facility in DHEP, TIFR, which has been used for part of the Monte Carlo generation work. MS acknowledges support from the National Science Foundation, Grant PHY-1630782, and to the Heising-Simons Foundation, Grant 2017-228. TS acknowledges financial support from the Department of Atomic Energy, Government of India, for the Regional Centre for Accelerator-based Particle Physics (RECAPP), Harish-Chandra Research Institute.

References

[1] Planck Collaboration, P. A. R. Ade et al., Planck 2015 results. XIII. Cosmological parameters, Astron. Astrophys. 594 (2016) A13, [arXiv:1502.01589].

[2] G. Jungman, M. Kamionkowski, and K. Griest, Supersymmetric dark matter, Phys. Rept. 267 (1996) 195–373, [hep-ph/9506380].

[3] H. Pagels and J. R. Primack, Supersymmetry, Cosmology and New TeV Physics, Phys. Rev. Lett. 48 (1982) 223.

[4] E. W. Kolb and R. Slansky, Dimensional Reduction in the Early Universe: Where Have the Massive Particles Gone?, Phys. Lett. 135B (1984) 378.

[5] J. R. Ellis, J. S. Hagelin, D. V. Nanopoulos, K. A. Olive, and M. Srednicki, Supersymmetric Relics from the Big Bang, Nucl. Phys. B238 (1984) 453–476.

[6] G. Bertone, D. Hooper, and J. Silk, Particle dark matter: Evidence, candidates and constraints, Phys. Rept. 405 (2005) 279–390, [hep-ph/0404175].

[7] Bergström, Lars, Nonbaryonic dark matter: Observational evidence and detection methods, Rept. Prog. Phys. 63 (2000) 793, [hep-ph/0002126].

[8] G. Bertone and D. Hooper, A History of Dark Matter, Submitted to: Rev. Mod. Phys. (2016) [arXiv:1605.04909].

[9] P. Gondolo and G. Gelmini, Cosmic abundances of stable particles: Improved analysis, Nucl. Phys. B360 (1991) 145–179.

[10] K. Griest and D. Seckel, Three exceptions in the calculation of relic abundances, Phys. Rev. D43 (1991) 3191–3203.

[11] G. Steigman, B. Dasgupta, and J. F. Beacom, Precise Relic WIMP Abundance and its Impact on Searches for Dark Matter Annihilation, Phys. Rev. D86 (2012) 023506, [arXiv:1204.3622].

[12] G. Steigman and M. S. Turner, Cosmological Constraints on the Properties of Weakly Interacting Massive Particles, Nucl. Phys. B253 (1985) 375–386.
[13] R. D. Peccei, *The Strong CP problem and axions*, Lect. Notes Phys. 741 (2008) 3–17, [hep-ph/0607268].

[14] R. J. Crewther, P. Di Vecchia, G. Veneziano, and E. Witten, *Chiral Estimate of the Electric Dipole Moment of the Neutron in Quantum Chromodynamics*, Phys. Lett. 88B (1979) 123. [Erratum: Phys. Lett.91B,487(1980)].

[15] J. E. Kim and G. Carosi, *Axions and the Strong CP Problem*, Rev. Mod. Phys. 82 (2010) 557–602, [arXiv:0807.3125].

[16] R. D. Peccei and H. R. Quinn, *CP Conservation in the Presence of Instantons*, Phys. Rev. Lett. 38 (1977) 1440–1443.

[17] F. Wilczek, *Problem of Strong P and T Invariance in the Presence of Instantons*, Phys. Rev. Lett. 40 (1978) 279–282.

[18] S. Weinberg, *A New Light Boson?*, Phys. Rev. Lett. 40 (1978) 223–226.

[19] J. E. Kim, *Weak Interaction Singlet and Strong CP Invariance*, Phys. Rev. Lett. 43 (1979) 103.

[20] M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, *Can Confinement Ensure Natural CP Invariance of Strong Interactions?*, Nucl. Phys. B166 (1980) 493–506.

[21] M. Kuster, G. Raffelt, and B. Beltran, *Axions: Theory, cosmology, and experimental searches. Proceedings, 1st Joint ILIAS-CERN-CAST axion training, Geneva, Switzerland, November 30-December 2, 2005, Lect. Notes Phys. 741 (2008) pp.1–258.

[22] D. J. E. Marsh, *Axion Cosmology*, Phys. Rept. 643 (2016) 1–79, [arXiv:1510.07633].

[23] J. E. Kim, *Axion as a CDM component*, in *Lepton and photon interactions at high energies. Proceedings, 23rd International Symposium, LP2007, Daegu, South Korea, August 13-18, 2007*, pp. 408–420, 2007. [arXiv:0711.1708].

[24] L. F. Abbott and P. Sikivie, *A Cosmological Bound on the Invisible Axion*, Phys. Lett. 120B (1983) 133–136.

[25] B. Dasgupta, E. Ma, and K. Tsumura, *Weakly interacting massive particle dark matter and radiative neutrino mass from Peccei-Quinn symmetry*, Phys. Rev. D89 (2014), no. 4 041702, [arXiv:1308.4138].

[26] L. M. Krauss and F. Wilczek, *Discrete gauge symmetry in continuum theories*, Phys. Rev. Lett. 62 (Mar, 1989) 1221–1223.

[27] XENON Collaboration, E. Aprile et al., *Dark Matter Search Results from a One Ton Year Exposure of XENON1T*, arXiv:1805.12562.

[28] GAMBIT Collaboration, P. Athron et al., *Status of the scalar singlet dark matter model*, Eur. Phys. J. C77 (2017), no. 8 568, [arXiv:1705.07931].

[29] ATLAS Collaboration, M. Aaboud et al., *Search for new phenomena in final states with an energetic jet and large missing transverse momentum in pp collisions at √s = 13 TeV using the ATLAS detector*, Phys. Rev. D94 (2016), no. 3 032005, [arXiv:1604.07773].

[30] CMS Collaboration, A. M. Sirunyan et al., *Search for dark matter produced with an energetic jet or a hadronically decaying W or Z boson at √s = 13 TeV*, JHEP 07 (2017) 014, [arXiv:1703.01651].
[31] CMS Collaboration, A. M. Sirunyan et al., Search for new physics in final states with an energetic jet or a hadronically decaying W or Z boson and transverse momentum imbalance at $\sqrt{s} = 13$ TeV, Phys. Rev. D97 (2018), no. 9 092005, [arXiv:1712.02345].

[32] S. L. Adler, Axial vector vertex in spinor electrodynamics, Phys. Rev. 177 (1969) 2426–2438.

[33] J. S. Bell and R. Jackiw, A PCAC puzzle: $\pi^0 \to \gamma\gamma$ in the $\sigma$ model, Nuovo Cim. A60 (1969) 47–61.

[34] G. Raffelt and D. Seckel, Bounds on Exotic Particle Interactions from SN 1987a, Phys. Rev. Lett. 60 (1988) 1793.

[35] ATLAS Collaboration, G. Aad et al., Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC, Phys. Lett. B716 (2012) 1–29, [arXiv:1207.7214].

[36] CMS Collaboration, S. Chatrchyan et al., Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC, Phys. Lett. B716 (2012) 30–61, [arXiv:1207.7235].

[37] Particle Data Group Collaboration, M. Tanabashi et al., Review of Particle Physics, Phys. Rev. D98 (2018), no. 3 030001.

[38] N. D. Christensen and C. Duhr, FeynRules - Feynman rules made easy, Comput. Phys. Commun. 180 (2009) 1614–1641, [arXiv:0806.4194].

[39] A. Alloul, N. D. Christensen, C. Degrande, C. Duhr, and B. Fuks, FeynRules 2.0 - A complete toolbox for tree-level phenomenology, Comput. Phys. Commun. 185 (2014) 2250–2300, [arXiv:1310.1921].

[40] D. J. Gross, R. D. Pisarski, and L. G. Yaffe, QCD and Instantons at Finite Temperature, Rev. Mod. Phys. 53 (1981) 43.

[41] K. J. Bae, J.-H. Huh, and J. E. Kim, Update of axion CDM energy, JCAP 0809 (2008) 005, [arXiv:0806.0497].

[42] P. Sikivie, Of Axions, Domain Walls and the Early Universe, Phys. Rev. Lett. 48 (1982) 1156–1159.

[43] G. Bélanger, F. Boudjema, A. Goudelis, A. Pukhov, and B. Zaldivar, micrOMEGAs5.0 : freeze-in, arXiv:1801.03509.

[44] S. J. Asztalos, G. Carosi, C. Hagmann, D. Kinion, K. van Bibber, M. Hotz, L. J. Rosenberg, G. Rybka, J. Hoskins, J. Hwang, P. Sikivie, D. B. Tanner, R. Bradley, J. Clarke, and ADMX Collaboration, SQUB-Based Microwave Cavity Search for Dark-Matter Axions, Physical Review Letters 104 (Jan., 2010) 041301, [arXiv:0910.5914].

[45] D. Budker, P. W. Graham, M. Ledbetter, S. Rajendran, and A. Sushkov, Proposal for a Cosmic Axion Spin Precession Experiment (CASPEx), Phys. Rev. X4 (2014), no. 2 021030, [arXiv:1306.6089].

[46] DES, Fermi-LAT Collaboration, A. Albert et al., Searching for Dark Matter Annihilation in Recently Discovered Milky Way Satellites with Fermi-LAT, Astrophys. J. 834 (2017), no. 2 110, [arXiv:1611.03184].

[47] D. Clowe, M. Bradac, A. H. Gonzalez, M. Markevitch, S. W. Randall, C. Jones, and D. Zaritsky, A direct empirical proof of the existence of dark matter, Astrophys. J. 648 (2006) L109–L113, [astro-ph/0608407].
H.E.S.S. Collaboration, H. Abdallah et al., Search for dark matter annihilations towards the inner Galactic halo from 10 years of observations with H.E.S.S, Phys. Rev. Lett. 117 (2016), no. 11 111301, [arXiv:1607.08142].

MAGIC Collaboration, M. L. Ahnen et al., Indirect dark matter searches in the dwarf satellite galaxy Ursa Major II with the MAGIC Telescopes, JCAP 1803 (2018), no. 03 009, [arXiv:1712.03095].

IceCube Collaboration, M. G. Aartsen et al., Search for Neutrinos from Dark Matter Self-Annihilations in the center of the Milky Way with 3 years of IceCube/DeepCore, Eur. Phys. J. C77 (2017), no. 9 627, [arXiv:1705.08103].

W. C. Huang, A. Urbano, and W. Xue, Fermi Bubbles under Dark Matter Scrutiny. Part I: Astrophysical Analysis, arXiv:1307.6862.

F. Calore, I. Cholis, and C. Weniger, Background Model Systematics for the Fermi GeV Excess, JCAP 1503 (2015) 038, [arXiv:1409.0042].
[65] ATLAS Collaboration, T. A. collaboration, Combination of searches for Higgs boson pairs in pp collisions at 13 TeV with the ATLAS experiment., ATLAS-CONF-2018-043.

[66] CMS Collaboration, A. M. Sirunyan et al., Search for Higgs boson pair production in the $\gamma\gamma b\bar{b}$ final state in pp collisions at $\sqrt{s} = 13$ TeV, arXiv:1806.00408 CMS-HIG-17-008, CERN-EP-2017-343.

[67] J. Alwall, R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, O.Mattelaer, H. S. Shao, T. Stelzer, P. Torrielli, and M. Zaro, The automated computation of tree-level and next-to-leading order differential cross sections, and their matching to parton shower simulations, JHEP 07 (2014) 079, [arXiv:1405.0301].

[68] NNPDF Collaboration, R. D. Ball et al., Parton distributions for the LHC Run II, JHEP 04 (2015) 040, [arXiv:1410.8849].

[69] T. Sjostrand, S. Ask, J. R. Christiansen, R. Corke, N. Desai, P. Ilten, S. Mrenna, S. Prestel, C. O. Rasmussen, and P. Z. Skands, An Introduction to PYTHIA 8.2, Comput. Phys. Commun. 191 (2015) 159–177, [arXiv:1410.3012].

[70] M. Cacciari, G. P. Salam, and G. Soyez, The Anti-k(t) jet clustering algorithm, JHEP 04 (2008) 063, [arXiv:0802.1189].

[71] M. Cacciari, G. P. Salam, and G. Soyez, FastJet User Manual, Eur. Phys. J. C72 (2012) 1896, [arXiv:1111.6097].

[72] ATLAS Collaboration, M. Aaboud et al., Search for dark matter produced in association with bottom or top quarks in $\sqrt{s} = 13$ TeV pp collisions with the ATLAS detector, Eur. Phys. J. C78 (2018), no. 1 18, [arXiv:1710.11412].

[73] ATLAS Collaboration, M. Aaboud et al., Search for dark matter and other new phenomena in events with an energetic jet and large missing transverse momentum using the ATLAS detector, JHEP 01 (2018) 126, [arXiv:1711.03301].

[74] CMS Collaboration, V. Khachatryan et al., Search for dark matter, extra dimensions, and unparticles in monojet events in protonproton collisions at $\sqrt{s} = 8$ TeV, Eur. Phys. J. C75 (2015), no. 5 235, [arXiv:1408.3583].

[75] CMS Collaboration, A. M. Sirunyan et al., Search for supersymmetry in multijet events with missing transverse momentum in proton-proton collisions at 13 TeV, Phys. Rev. D96 (2017), no. 3 032003, [arXiv:1704.07781].

[76] CMS Collaboration, A. M. Sirunyan et al., Search for new phenomena with the $M_{T2}$ variable in the all-hadronic final state produced in protonproton collisions at $\sqrt{s} = 13$ TeV, Eur. Phys. J. C77 (2017), no. 10 710, [arXiv:1705.04650].

[77] ATLAS Collaboration, M. Aaboud et al., Search for supersymmetry in final states with missing transverse momentum and multiple $b$-jets in proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, JHEP 06 (2018) 107, [arXiv:1711.01901].

[78] ATLAS Collaboration, M. Aaboud et al., Search for squarks and gluinos in final states with jets and missing transverse momentum using 36 fb$^{-1}$ of $\sqrt{s}=13$ TeV pp collision data with the ATLAS detector, arXiv:1712.02332 CERN-EP-2017-136.

[79] ATLAS Collaboration, M. Aaboud et al., Search for a scalar partner of the top quark in the
jets plus missing transverse momentum final state at $\sqrt{s}=13$ TeV with the ATLAS detector, JHEP 12 (2017) 085, [arXiv:1709.04183].