New Constraints on Higgs-Portal Scalar Dark Matter

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Abstract

The simplest Higgs-portal dark matter model, in which a real scalar singlet is added to the standard model, has been comprehensively revisited, by taking into account the constraints from perturbativity, electroweak vacuum stability in the early Universe, dark matter direct detection, and Higgs invisible decay at the LHC. We show that the resonant mass region is totally excluded and the high mass region is reduced to a narrow window \(1.1 \text{ TeV} \leq m_s \leq 2.55 \text{ TeV}\), which is slightly reduced to \(1.1 \text{ TeV} \leq m_s \leq 2.0 \text{ TeV}\) if the perturbativity is further imposed. This high mass region can be fully detected by the Xenon1T experiment.
1 Introduction

The Standard Model (SM) as the effective theory below the electroweak (EW) scale has been established after the discovery of SM Higgs scalar with mass around 125 GeV \cite{1,2} at the Large Hadron Collider (LHC). Unfortunately, there is no viable candidate for a dark matter (DM) in the SM, which implies that an extension beyond the SM is necessary.

Among viable extensions the simplest Higgs-portal dark matter (HDM) \cite{3,4,5,6,7} is the most economic and of special interest from both the DM- and LHC-phenomenology. In this extension, only a real singlet scalar $s$ is added to the SM. With the aid of $Z_2$ parity which is needed for the stability of singlet DM, the number of model parameters is reduced to three. It implies that the minimal model is very predictable. In the subsequent studies on the HDM \cite{8,9,10,11,12,13,14,15,16,17,18,19,20,21}, the parameter space is analyzed in terms of experimental constraints as follows.

- Dark matter relic density \cite{22}, which shows that the low mass region, resonant mass region, and high mass region are all viable.

- Dark matter direct detection such as Xenon100 \cite{23}, Xenon1T \cite{24} and LUX \cite{25} experiments, which are able to effectively detect the high mass region up to $\sim 3$ TeV.

- Limit on the Higgs invisible decay $h \rightarrow ss$, which can effectively detect the low mass region with $m_s$ of order a few GeV.

- Indirect detection such as limit on $\gamma$-ray lines from the Fermi-LAT \cite{27} is rather efficient to constrain the resonant mass region \cite{28,29,30}.

However, these constraints are not sufficient, as some possibly strong theoretic ones have been ignored. The object of this paper is to take into account all known experimental and theoretic constraints. Simultaneously, experiments constraints will be updated, if available. In compared with the previous works, e.g., a comprehensive review in \cite{21}, our main differences are the following ones.

- The SM EW vacuum suffers from large quantum fluctuations in the early Universe as induced by the inflation, which imposes strong constraint \cite{32,33,34} on the model parameters. SM vacuum stability must be taken into account to constrain the parameter space. For an earlier attempt, see \cite{35,36}.
• Assume the HDM as an effective theory between the EW and Plank scale, perturbative analysis should be considered seriously, which leads to strong upper bounds on the quartic coupling constant of $s^2 |H|^2$ and the DM self-coupling constant. This effective theory can be applied, e.g., to S-inflation, where the scalar singlet DM is identified as the inflaton \[37\].

• Very recently, the ATLAS collaboration has reported the latest limit about the Higgs invisible decay width \[26\] in the low and resonant mass regions. This will be updated in our analysis.

The paper is organized as follows. In Sec.2, we introduce notation and conventions in the HDM. In this section the constraint on the model parameters from the requirement of perturbativity is studied. In Sec.3, we discuss the constraint on the model parameters from the SM vacuum stability for 125 GeV Higgs mass. No additional mechanism for stabilizing the SM vacuum stability is employed except the DM scalar.

In Sec.4 we re-examine the DM phenomenology. In particular, we calculate the scalar relic density and the spin-independent nucleon-DM scattering cross section by using MicrOMEGAs \[38\], which are consistent with results in some earlier literature. In Sec.5, we update the LHC phenomenology about scalar DM, by focusing on the latest limit on the Higgs invisible decay.

Finally, we present our main conclusions in Sec.6. We find that the resonant mass region is totally excluded and the high mass region is reduced to a narrow window $1.1 \text{ TeV} \leq m_s \leq 2.4 \text{ TeV}$. If the requirement of perturbativity is further imposed, the allowed mass range is slightly reduced to $1.1 \text{ TeV} \leq m_s \leq 2.0 \text{ TeV}$. In either case the mass window can be fully detected by the Xenon 1T experiment.

2 HDM within Perturbative Region

The Lagrangian for the simplest HDM with a $Z_2$ parity, under which $s$ is odd and SM particles are even, is given by,

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \frac{1}{2} (\partial s)^2 + V(s, H)$$

where

$$V(s, H) = \frac{\lambda}{2} \left( |H|^2 - \frac{\nu^2}{2} \right)^2 + \frac{\mu_s^2}{2} s^2 + \frac{\lambda_s}{2} s^4 + \frac{\kappa_s}{2} s^2 |H|^2 .$$
In Eq. (2.2) the first term denotes the Higgs potential with EW scale \( v \simeq 246 \text{ GeV} \), and the \( \lambda_s \)-term and \( \kappa_s \)-term refers to \( s \) self-interaction and Higgs-DM interaction, respectively. Expand the fields as \( H = (v + h)/\sqrt{2} \) and \( s = \langle s \rangle + s = 0 + s \) along the vacuum structure for positive \( \lambda_s \) and \( \kappa_s \), one obtains from Eq. (2.2),

\[
V(s, h) = \frac{1}{2} m_s^2 s^2 + \frac{\lambda_s}{2} s^4 + \frac{\kappa_s v}{2} s^2 h + \frac{\kappa_s}{4} s^2 h^2. \tag{2.3}
\]

where \( m_s = \mu_s^2 + \kappa_s v^2/2 \). Remarkably, there are only three new parameters in the HDM. It implies that this model can be studied concretely.

In some situations perturbativity should be imposed on the model. This is true when it is assumed to the effective theory below some high energy scale such as Plank mass scale. See, e.g. [37] for an application of such idea to the S-inflation.

In Fig. 1 we show the upper bounds on the EW values of \( \kappa_s \) and \( \lambda_s \) by the requirement of perturbativity

\[
0 < \kappa_s(\mu) < \sqrt{4\pi}, \quad 0 < \lambda_s(\mu) < \sqrt{4\pi}, \tag{2.4}
\]

for Renormalization Group (RG) scale \( \mu \) between the EW and Plank mass scale. We have used two-loop RG equations [41],

\[
\begin{align*}
\beta_{\kappa_s} &= \frac{1}{(4\pi)^2} \left[ - \left( \frac{9}{10} g_1^2 + \frac{9}{2} g_2^2 - 6\lambda - 6 y_t^2 - 12 \lambda_s \right) \kappa_s + 4 \kappa_s^2 \right] + \frac{1}{(4\pi)^4} \cdots \\
\beta_{\lambda_s} &= \frac{1}{(4\pi)^2} \left( \kappa_s^2 + 36 \lambda_s^2 \right) + \frac{1}{(4\pi)^4} \cdots \tag{2.5}
\end{align*}
\]

where \( g_i \) refer to the SM gauge coupling, and \( y_t \) denotes the top Yukawa coupling. As expected from Eq. (2.5), the correlation between the upper bounds on \( \kappa_s \) and \( \lambda_s \) is mildly sensitive to \( s \) scalar mass, but strongly to each other. Our scan shows that the critical values are given by \( \kappa_s < 0.6 \) and \( \lambda_s < 0.131 \).

Of course, these bounds can be relaxed when the high energy scale above which the model is not effective is smaller than our reference value.

\footnote{The vacuum is given by \( \langle H \rangle = v/\sqrt{2} \) and \( \langle s \rangle = 0 \) for both positive \( \lambda_s \) and \( \kappa_s \), which may be violated when both of them are negative. However, as we will shown in Sec.3, neither negative \( \kappa_s \) and nor negative \( \lambda_s \) are favored by the EW vacuum stability. In this paper we choose both positive \( \lambda_s \) and \( \kappa_s \) for our discussion.}
3 Electroweak Vacuum Stability

It was shown in [31] that the EW vacuum is metastable for the observed Higgs mass $m_h \simeq 125$ GeV, as seen from the RG equation for $\lambda$. Since the sign of beta function $\beta_\lambda$ is always negative in the SM, $\lambda$ becomes negative above some critical RG scale $\mu_*$. $\mu_*$ is approximately equal to the value $h_*$, at which the Higgs potential is maximal. For the central values of top quark mass $m_t = 173.2$ GeV and structure constant of QCD gauge coupling $\alpha = 0.1184$, $h_* \simeq 10^9 - 10^{10}$ GeV at the two-loop level. The situation can be improved to $h_* \simeq 10^{19}$ GeV for lighter $m_t$ within 1$\sigma$ uncertainty [32]. So, in the most of SM parameter space EW vacuum is only metastable.

For EW vacuum being metastable, it suffers from large quantum fluctuation in the early Universe, especially during and after inflation, which leads to strong theoretical constraint on the model parameters.

Let us firstly discuss the survival probability of EW vacuum during inflation. Denote $P(h, t)$ the probability for Higgs field with the value $h$ at time $t$, which can be determined via solving the Fokker-Plank equation [32, 33, 34],

$$
\frac{\partial P}{\partial t} = \frac{\partial}{\partial h} \left[ \frac{H_c^3}{8\pi^2} \frac{\partial P}{\partial h} + \frac{V'(h)}{3H_c} P \right].
$$

(3.1)
Here \( H_c \) refers to the Hubble parameter\(^2\), which is approximate constant during inflation. \( V' \) denotes derivative of Higgs potential over Higgs field. For the case \( h < h_* \) the quantum fluctuation-\( H_c^3 \) term dominates the classical one-the last term in Eq.(3.1). If one ignores this classical term, Eq.(3.1) can be easily solved in terms of separating variants and reasonable boundary conditions. See \([33, 34]\) for details. Then, the probability for staying at the domain \( h > h_* \) other than EW vacuum at the end of inflation is given by,

\[
P(|h| > h_*) = 1 - \int_{-h_*}^{h_*} P(h, \frac{N}{H_c}) dh,
\]

(3.2)

where \( N \approx 55 - 60 \) is the number of e-folds.

Figure 2: The value of \( h_* \) at which the Higgs potential is maximal in the parameter space of \( \kappa_s - m_s \). The green curves correspond to the case of SM. The black curves show the critical value as required by EW vacuum stability. It clearly shows that the value of \( h_* \) can be uplifted into the region surrounded by the two black curves.

If the probability for staying at the metastable EW vacuum is not zero as required, \( P(|h| > h_*) \) should be smaller than \( e^{-3N} \), which gives rise to,

\[
h_* > \sqrt{\frac{3N}{2\pi}} H_c \approx 25H_c \approx 2.5 \times 10^{17}\text{GeV}.
\]

(3.3)

\(^2\)The magnitude of Hubble constant is directly related to the scalar-to-tensor ratio \( r \), with \( H_c \approx 1 \times 10^{16} \times (r/0.01)^{1/4} \) GeV. We take \( H_c = 10^{16} \) GeV as the reference point, which is a good approximation unless \( r \) is extremely small.
Obviously, this condition cannot be satisfied in the most of parameter space of SM. The constraint in Eq. (3.3) can be relaxed in the context of SM by introducing a new interaction $\xi_H R^2 |H|^2$ between the Higgs and Ricci scalar $R$.

With no need of parameter $\xi_H$, this reduction can be alternatively realized in the HDM. It is easy to understand in terms of the new contribution to beta function $\beta_\lambda$, which reads at the two-loop level,

$$\delta \beta_\lambda = \beta_\lambda^{(\text{HDM})} - \beta_\lambda^{(\text{SM})} = \frac{\kappa_s^2}{(4\pi)^2} + \frac{1}{(4\pi)^4}(-5\lambda\kappa_s^2 - 4\kappa_s^3).$$

(3.4)

The value of $h_s$ can be uplifted for small $\kappa_s$, as the sign of the new contribution in Eq. (3.4) can be positive in this region. In contrast, $h_s$ may be reduced in the region of large $\kappa_s$, where the sign of the new contribution in Eq. (3.4) is negative. The green curve in Fig. 2 clearly shows the critical value $\kappa_s^c \sim 0.9$. Below this critical value $h_s$ can be larger than SM value $h_s^{\text{SM}}$. More interesting, as shown by the red region corresponding to the range $\kappa_s \simeq 0.25 - 0.8$, the required value $h_s^c = 2.5 \times 10^{17}$ GeV in Eq. (3.3) can be satisfied.

The effect on the vacuum stability of the singlet scalar potential due to quantum fluctuations is more easily understood, in comparison with the Higgs potential. Since $\langle s \rangle = 0$ is the true minimal of singlet scalar potential, quantum fluctuation with magnitude of order $\delta s \sim H$ may push $s$ towards to an unstable vacuum from the origin. But it rapidly returns to the origin vacuum through classical tuning process. This understanding holds as long as $\lambda_s$ is always positive between the EW and Plank scale.

Once inflation ends, (p)reheating begins immediately. The quantum effect on the EW vacuum during (p) reheating can be studied in terms of the equations of motion for $s$ and Higgs in principle. Unfortunately, a concrete constraint on the model parameters in HDM can not be derived without knowledge of reheating temperature $T_{\text{re}}$, which is an important physical quantity involved \cite{34}. To determine $T_{\text{re}}$ interactions between the inflaton scalar and SM fields should be given explicitly.

4 Dark Matter Phenomenology

In this section we firstly consider the constraints arising from the DM relic density and direct detection limits at the LUX and Xenon100 experiments. Then we discuss the prospect at the Xenon1T experiments.

\textsuperscript{3}For negative $\kappa_s$, $|\kappa_s^c|$ is larger in comparison with positive $\kappa_s$. However, its magnitude is more seriously bounded above by the analysis of perturbativity.
4.1 Relic Density

The Plank and WMAP 9-year data have measured the DM relic density in high precision \( \Omega_{\text{DM}} h^2 = 0.1199 \pm 0.0027 \) \[ \text{(4.1)} \]

With the assumption that the relic density of \( s \) scalar is totally produced by the thermal freeze-out process, we have

\[ \Omega_s h^2 \sim 0.1 \text{ pb/} \langle \sigma_{\text{ann}} v \rangle \]

where \( \sigma_{\text{ann}} \) is the total annihilation cross-section for \( s \) and \( v \) is the relative velocity.

Figure 3: Contours of \( s \) scalar relic density \( \Omega_s h^2 \) as required by the DM projected to the two-parameter plane of \( (m_s, \kappa_s) \).

The \( ss \) annihilation mainly proceeds through Higgs-mediated process in the \( s \) channel. Sub-dominant process include \( s \) exchange in the \( t \) channel and annihilation into \( hh \). So, \( \sigma_{\text{ann}} \) is sensitive to the model parameters \( \kappa_s \) and \( m_s \).

Instead of analytic method we use the code MicrOMEGAs \[ \text{[38]} \] to calculate the \( s \) scalar relic density \( \Omega_s h^2 \), with the Feynman rules generated by the package LanHEP 3.2.0 \[ \text{[39]} \]. In Fig.3 we project the contour of relic density \( \Omega_s h^2 \) as required by Eq.(4.1) to the two-parameter plane of \( (m_s, \kappa_s) \). As expected the required value of \( \kappa_s \) is minimal for \( m_s \) near
$m_h/2$, because near this mass region the annihilation cross section is resonantly enhanced. Plot similar to Fig. 3 is also shown, e.g., in [21].

Fig. 3 indicates that a wide mass range for s scalar is viable. However, it also shows that $\kappa_s$ bigger than $0.1$ is required when $m_s$ is far away from the resonant mass region. For example, large $\kappa_s \geq 0.5$ is needed for $m_s$ above 3 TeV, which is not favored neither by the EW vacuum stability nor the perturbative analysis. In the next subsection, we discuss the direct detection in the allowed parameter space.

![Figure 4: Spin-independent nucleon-DM scattering cross section as function of $m_s$. Various experimental limits are also shown for comparison.](image)

### 4.2 Direct Detection

The spin-independent nucleon-DM scattering cross section is given by,

$$\sigma_{SI} = \frac{\kappa_s^2 f_N^2 \mu^2 m_N^2}{4\pi m_h^4 m_s^2},$$

(4.2)

where $m_N$ is the nucleon mass, $\mu = m_s m_N/(m_s + m_N)$ is the DM-nucleon reduced mass, and $f_N \sim 0.3$ [21] is the hadron matrix element.

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4 Although these plots match very well, we would like to mention that contributions to $\sigma_{ann}$ due to three- and four-body final states or QCD corrections are not included in MicOMEGAs, with which the shape around $W$ boson mass is smoother [41]. We thank the referee for reminding us this point. Since DM mass region near 80 GeV has been excluded by the LUX experiment, these corrections are expected to not affect our main results in Sec.6.
The current best limit on $\sigma_{\text{SI}}$ comes from the Xenon100 \cite{23} and LUX \cite{25} experiments. In Fig.4 we plot the spin-independent nucleon-DM scattering cross section as function of $m_s$, with various experimental limits for comparison. The red curve corresponds to the DM relic density, which indicates that there are three viable regions,

\begin{align*}
\text{low mass region} & : \quad 1 \text{ GeV} \leq m_s \leq 6 \text{ GeV}, \\
\text{resonant mass region} & : \quad 56 \text{ GeV} \leq m_s \leq 66 \text{ GeV}, \\
\text{high mass region} & : \quad m_s \geq 185 \text{ GeV}.
\end{align*}

(4.3)

By combing the constraint from SM vacuum stability, one finds that

- The \textit{resonant mass region} is totally excluded, which is more powerful than the indirect constraint from the Fermi-LAT \cite{27}.

- The \textit{high mass region} is further reduced to more narrow mass window, which is the subject of Xenon 1T \cite{24} experiment for $m_s$ below 4 TeV.

- The \textit{low mass region} can be consistent with the constraint from SM vacuum stability. However, we will show in Sec.V that the \textit{low mass region} is excluded by the Higgs invisible decay $h \rightarrow ss$ at the LHC experiment (see also some literature in Refs.[8-18]).

5 LHC Phenomenology

Since the scalar singlet only couples to the Higgs directly, the signal searches for this scalar at the LHC mainly focus on the Higgs invisible decay $h \rightarrow ss$ in the \textit{low mass region} and \textit{resonant mass region} with $m_s < m_h/2$. Note that for scalar singlet DM there is no mixing effect between $s$ and $h$, as clearly shown in Eq.(2.1) and Eq.(2.2).

Very recently, the ATLAS Collaboration has reported the latest data about the upper bound on the Higgs invisible decay width \cite{26},

$$\Gamma_{\text{inv}} \leq 0.29 \Gamma_{\text{SM}},$$

(5.1)

\footnote{If $s$ scalar is not DM, one may discuss its phenomenology by introducing mixing effects. For a recent discussion, see \cite{44}.}
Figure 5: Contours of invisible decay width $\Gamma(h \rightarrow ss)$ in the two-parameter plane of $m_s - \kappa_s$. The contour in black denotes the latest upper bound in Eq. (5.1). Region above this line is excluded.

which is obviously stronger than what was used in some earlier literature. An update is thus meaningful. In Eq. (5.1) $\Gamma_{SM} \approx 4.15$ MeV for 125 GeV Higgs mass, and the invisible decay width $\Gamma_{inv}$ is given by,

$$\Gamma(h \rightarrow ss) = \frac{\kappa_s^2 v^2}{32\pi m_h} \sqrt{1 - \frac{4m_s^2}{m_h^2}}. \quad (5.2)$$

We show the decay width $\Gamma(h \rightarrow ss)$ in Fig. 5. The contour in black denotes the latest upper bound in Eq. (5.1), and region above it is excluded. In particular, this bound excludes the whole low mass region and the resonant mass region with $m_s \leq 62.5$ GeV in Eq. (4.3). In contrast, the high mass region is less constrained at the LHC in compared with the low mass region.

6 Conclusions and Discussions

In this paper we have reconsidered the simplest HDM. By combining the constraints arising from DM relic density, direct detection limit on spin-independent nucleon-DM scattering cross section from the LUX, and limit on the Higgs invisible decay from the
Figure 6: Mass regions consistent with various constraints. In the light of the SM vacuum stability, only mass region $1.1 \text{ TeV} \leq m_s \leq 2.5 \text{ TeV}$ is viable, and it slightly reduces to $1.1 \text{ TeV} \leq m_s \leq 2.0 \text{ TeV}$ by the additional requirement of perturbativity. For the explanation about $h_*$ and $\kappa_{\text{max}}$, see previous discussions.

LHC, we have obtained two viable mass regions (see Fig.6):

resonant mass region : $62.5 \text{ GeV} \leq m_s \leq 66 \text{ GeV}$,

high mass region : $m_s \geq 185 \text{ GeV}$.  \hspace{1cm} (6.1)

We also show that if the constraint from the SM vacuum stability is imposed, we arrive at the main conclusion- the resonant mass region is totally excluded and the high mass region in Eq.(6.1) is reduced to a narrow window,

high mass region : $1.1 \text{ TeV} \leq m_s \leq 2.55 \text{ TeV}$, \hspace{1cm} (6.2)

as shown in Fig.6. If one further imposes the constraint due to perturbativity, the mass range in Eq.(6.2) is slightly reduced to,

high mass region : $1.1 \text{ TeV} \leq m_s \leq 2.0 \text{ TeV}$.  \hspace{1cm} (6.3)

Fortunately, as shown in Fig.6 the mass region in Eq.(6.3) can be totally detected by the Xenon 1T experiment in the near further. Similarly to Xenon 1T experiment, this high mass region can be also directly detected at colliders in principle. As a result of
small production cross section of order less than 1 fb for the s scalar at the 14-TeV LHC \[19\], a 5σ discovery acquires luminosity at least of order \( \mathcal{O}(10) \) ab\(^{-1}\).

Independently, effects on large scale structure induced by DM self-interaction is a new window for the detection. In compared with the abelian vector- or fermion-like DM, s scalar with quartic self-interaction is very distinctive. DM self-interaction has been used to interpret \[47\] the observations on four galaxies in the core of galaxy cluster Abell 3827 \[46\]. If the new observations indeed arise from the DM self-interaction other than astrophysical artifact, the high mass region is highly constrained, see, e.g., for recent discussion \[48\].

In variety of contexts the high mass region in Eq. \[6.2\] may be modified. For example, it can be relaxed if one employs additional mechanism to stabilize the SM vacuum by introducing more new parameters, such as the non-minimal interaction between the Higgs and Ricci scalar with non-negligible coupling constant. In contrast, this high mass region may be alternatively strengthened if one assumes \[21\] that the DM relic abundance is partially other than totally saturated by the scalar singlet. In this case scalar s may play a positive role in the electroweak baryogenesis \[40\].

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