

LHC-friendly freeze-in dark matter via the Higgs portal

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It is known that single-field freeze-in dark matter barely leaves footprints in dark matter direct detection and collider experiments. This situation can be altered in the two-field context. In this work we propose a two-field freeze-in dark matter model through the Higgs portal. The observed dark matter relic abundance is obtained by a decay of scalar mediator thermalized in the early Universe. While there is a lack of direct dark matter signals, the scalar mediator is in the reach of High Luminosity LHC either through the vector boson fusion or mono-Z channels. Within the allowed scalar mass window of 10–50 GeV, we use improved cuts to derive both 2σ exclusion and 5σ discovery limits, depending on the value of Higgs portal coupling. If verified, this scalar mediator signal allows us to infer the freeze-in dark matter.

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I. INTRODUCTION

Thermal dark matter (DM) has been severely challenged by a lack of conclusive signals in direct, collider, and indirect detection experiments. For example, the direct detection limits have been close to the neutrino floor within a certain DM mass range; see, e.g., Refs. [1,2] for recent reviews. Inspired by those experimental results, there is a growing interest in nonthermal DM produced in a few ways other than freeze-out mechanism, such as freeze-in mechanism [3].¹ Single-field examples of freeze-in DM (FIDM) include axionlike scalar [4,5], sterile neutrino fermion [6–9], and dark photon [10], some of which are well motivated by new physics beyond the Standard Model (SM). A common feature of these FIDM models is the absence at DM direct and collider detection experiments, as a result of feeble DM coupling to the SM sector. Within the DM mass range larger than the MeV scale, only cosmic [11,12] or astrophysical [13–16] ray data have touched on this type of DM parameter space so far.

The situation changes in two-field FIDM. A two-field FIDM, which contains DM and its force mediator, can be constructed in two different ways, as shown in Fig. 1,

depending on how to assign the feeble interaction. Explicitly,

- (i) In the first case (a), one chooses a feeble interaction between the force mediator and SM sector but a strong coupling between the DM and force mediator. A representative example is dark photon-mediated DM, which has been used to explain EDGES 21-cm anomaly [17,18] and small-scale problem [19], respectively.
- (ii) In the second case (b), one adopts a strong interaction between the force mediator and SM sector but a feeble coupling between the DM and force mediator. In this situation, collider-friendly FIDM models [20–23] can be constructed.

In this work, we investigate a new two-field FIDM of the second type with the following Lagrangian:

$$\mathcal{L}_D = \bar{\psi}(i\gamma^\nu\partial_\nu - m_\psi)\psi + \frac{1}{2}\partial_\mu s\partial^\mu s - \frac{\mu_s^2}{2}s^2 - \kappa s^2|H|^2 - \lambda\bar{\psi}\eta\psi s, \quad (1)$$

where fermion ψ and real singlet scalar s are the DM and force mediator, respectively, and H is the SM Higgs doublet. We impose a $Z_2 \times Z'_2$ dark parity on the dark sector, under which

$$Z_2: \psi \rightarrow -\psi; \quad Z'_2: s \rightarrow -s, \quad (2)$$

and the SM sector is even. In Eq. (1), we assume the breaking of $Z_2 \times Z'_2 \rightarrow Z_2$ by the renormalizable Yukawa

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¹An alternative mechanism not so precisely understood is through gravitational portal (during) after inflation, as the inflaton sector is yet conclusive.

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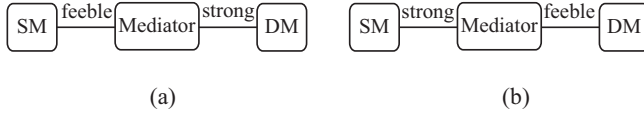


FIG. 1. An illustration of the two-field FIDM model classifications.

interaction in the dark sector among others,² which suggests a small λ by following the argument of naturalness [24]; and the scalar mass squared reads as $m_s^2 = \mu_s^2 + \kappa v^2 \geq \kappa v^2$ with the electroweak scale $v = 246$ GeV. Unlike in the Higgs-portal scalar DM [25–27] based on freeze-out mechanism, where only a narrow scalar mass window near a half of the Higgs mass is not excluded [28,29], the FIDM due to a small λ in Eq. (1) allows us to reopen the scalar mass window above GeV scale. Within this mass window, the scalar mediator which only decays to a pair of DM particles contributes to observable missing energy [30] at the LHC for κ being not far away from unity. This makes our model differ from previous two-field models [20–23], where the mediator therein decays to either SM lepton or hadron final states.

The rest of this paper is organized as follows: in Sec. II, we discuss DM phenomenology including DM production, scattering of SM protons, and annihilation into SM particles. Although there is a lack of DM signals in traditional DM detection experiments, the scalar mediator can leave footprints in the LHC, depending on the value of Higgs portal coupling κ . Section III is devoted to investigating the High Luminosity LHC (HL-LHC) reaches of the scalar mediator in allowed scalar mediator mass range of 10–60 GeV through the Higgs portal-induced vector boson fusion (VBF) and mono-Z channels, where both 2σ exclusion and 5σ discovery limits are presented. The Appendix shows details of how to improve cuts than those of Ref. [30]. Finally, we conclude in Sec. IV.

II. DARK MATTER PHENOMENOLOGY

In this section, we begin with the DM phenomenology in DM parameter space composed of $\{m_\psi, m_s, \lambda\}$ in Eq. (1), with a negligible dependence on κ for κ larger than $\sim 10^{-6}$ required by the freeze-out mechanism.

A. Relic abundance

With κ being large enough, the scalar mediator keeps thermal equilibrium with the SM thermal bath. In contrast, for a tiny λ , the DM is produced in the early Universe through the freeze-in mechanism with the Boltzmann equation

²We neglect the Z'_2 -preserving s^4 term that does not affect either DM or collider phenomenological analysis. In addition, we ignore the Z'_2 -breaking $s|H|^2$ term, leading to tiny mixing effect between s and Higgs, whose magnitude can be of the same order as λ .

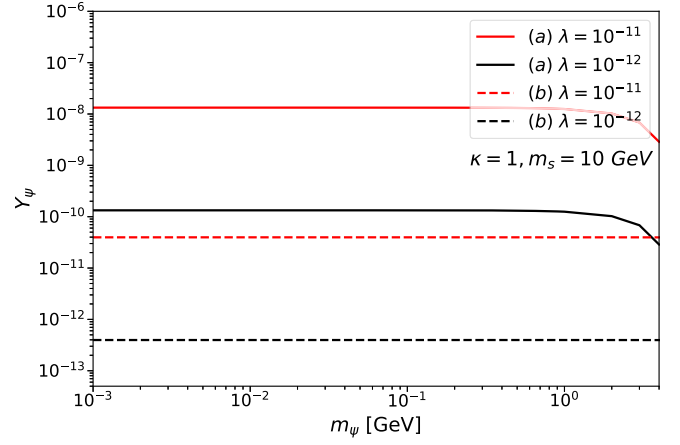


FIG. 2. An illustration of the individual contribution to Y_ψ as function of m_ψ for a set of fixed values $\kappa = 1$, $\lambda = \{10^{-12}, 10^{-11}\}$, and $2m_\psi < m_s = 10$ GeV, where (a) and (b) are shown as solid and dashed lines, respectively.

$$\dot{n}_\psi + 3Hn_\psi = \mathcal{C}, \quad (3)$$

where n_ψ is the DM number density, H is the Hubble rate due to cosmic expansion, and $\mathcal{C}$ is the “collision” term mainly arising from (a) decay $s \rightarrow \psi\bar{\psi}$ in the DM mass range of $m_s > 2m_\psi$ considered here and (b) two-body annihilation $hs \rightarrow \psi\bar{\psi}$ via the κ coupling.

To see which process dominates the $\mathcal{C}$ term in Eq. (3), we present in Fig. 2 the individual contribution to DM yield $Y_\psi \equiv n_\psi/S$ as a function of m_ψ with S the entropy of the SM thermal bath, for a set of fixed values $\kappa = 1$, $\lambda = \{10^{-12}, 10^{-11}\}$, and $2m_\psi < m_s = 10$ GeV. In this plot, (a) and (b) are shown as solid and dashed lines, respectively.

- (i) The decay induced contribution to Y_ψ is consistent with an analytic approximation [3],

$$Y_\psi \approx \frac{135}{4\pi^3 \cdot 1.66 g_{s,*} \sqrt{g_{\rho,*}}} \frac{M_P \Gamma_s}{m_s^2}, \quad (4)$$

where $g_{s,*}$ and $g_{\rho,*}$ are the number of degrees of freedom in entropy and energy density, respectively; M_P is the nonreduced Planck mass; and the scalar mediator decay width is

$$\Gamma_s(s \rightarrow \psi\bar{\psi}) = \frac{\lambda^2}{8\pi} m_s \left(1 - \frac{4m_\psi^2}{m_s^2}\right)^{3/2}. \quad (5)$$

- (ii) The annihilation-induced contribution to Y_ψ is proportional to $\sim \lambda^2 \kappa^2$. As seen in Fig. 2, this contribution is smaller than the decay-induced one by more than 2 orders, which is further reduced by choosing κ smaller than unity.

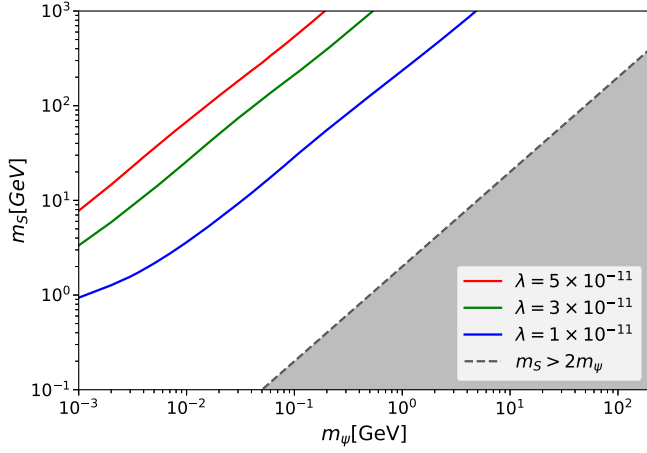


FIG. 3. Contours of the observed DM relic abundance projected to the plane of $m_s - m_\psi$ for various values of λ , where the shaded region is excluded by $2m_\psi < m_s$ and the top left region with $\lambda \sim 10^{-11}$ and $m_s \geq 1$ GeV evades the BBN constraint on τ_s .

Therefore, in the parameter regions with $10^{-6} < \kappa \leq 1$ and $2m_\psi < m_s < 1$ TeV considered in this paper, the DM freeze-in production in the early Universe is dominated by the scalar decay process.

In terms of publicly available code `micrOMEGAs6.0` [31], we show in Fig. 3 contours of the observed relic density $\Omega_\psi h^2 = 0.12 \pm 0.001$ [32] projected to the plane of $m_s - m_\psi$ for various values of $10^{11} \lambda = \{1, 3, 5\}$, where the shaded region is excluded by $2m_\psi < m_s$. The numerical results of Fig. 3 are consistent with the analytic approximation

$$\Omega_\psi h^2 = \frac{m_\psi Y_\psi S_0}{\rho_{\text{crit},0}} \approx 0.1 \left(\frac{\lambda}{5 \times 10^{-13}} \right)^2 \left(\frac{m_\psi}{m_s} \right), \quad (6)$$

derived from Eq. (4), where S_0 is the present value of S and $\rho_{\text{crit},0}$ is the critical energy density. Take an explicit value of $\lambda = 5 \times 10^{-11}$, for example, in which case $m_s/m_\psi \sim 10^4$, as shown by the red curve is verified by Eq. (6). On the other hand, Eq. (6) suggests that a curve with respect to a larger λ should have a larger slope, which is also justified by comparing the curves in this plot.

The contours of $\Omega_\psi h^2$ in Fig. 3 barely depend on κ ; see also Eq. (6). This is true as long as κ is larger than the threshold value $\sim 10^{-6}$ required to keep the scalar mediator in thermal equilibrium with the SM thermal bath. In this situation, κ only affects the DM relic density through altering the values of $g_{s,*}$ and $g_{\rho,*}$ in Eq. (4) by $\sim 1\%$, which can be safely neglected.

Before closing the analysis on the DM relic density, we remind readers that the scalar mediator lifetime is

$$\tau_s \sim 10 \text{ s} \left(\frac{\lambda}{10^{-12}} \right)^{-2} \left(\frac{m_s}{1 \text{ GeV}} \right)^{-1}, \quad (7)$$

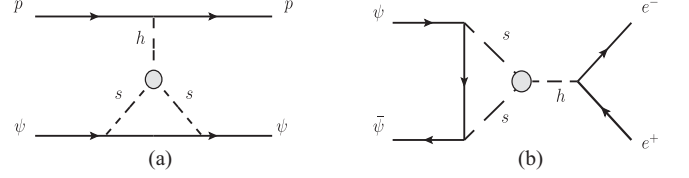


FIG. 4. Feynman diagrams contributing to (a) spin-independent DM direct detection and (b) DM indirect detection, respectively.

using Eq. (5). Equation (7) shows that the lifetime is always shorter than $\sim 10^4$ s, when the big bang nucleosynthesis (BBN) takes place, in the top left region of Fig. 3 with $\lambda \sim 10^{-11}$ and $m_s \geq 1$ GeV, which evades the BBN constraint [33,34].

B. Direct and indirect detection

In our model, the DM scatters off protons or electrons by one-loop Feynman diagram. Consider that the electron Yukawa coupling is small, we focus on the spin-independent DM-proton cross section σ_{SI} as shown in diagram (a) of Fig. 4. It can be estimated by a nonrelativistic effective operator analysis in the low-energy region. Integrating out the scalar mediator in Eq. (1), one finds

$$\mathcal{L}_{\text{eff}} \sim \kappa \lambda^2 \frac{m_\psi^3}{m_s^4} \psi \bar{\psi} |H|^2 + \dots, \quad (8)$$

which gives

$$\sigma_{\text{SI}} \sim 10^{-86} \text{ cm}^2 \cdot \kappa^2 \cdot \left(\frac{\lambda}{10^{-11}} \right)^4, \quad (9)$$

for the reference masses $m_s \sim m_\psi \sim 100$ GeV. Normalized to $\kappa = 1$, the value of σ_{SI} suppressed by λ^4 is roughly ~ 42 orders of magnitude smaller than the latest XENON1T [35], PandaX-4T [36] and LZ [37] bounds. Note that the effective operator analysis in Eq. (8) is valid for a rough estimate on the magnitudes of σ_{SI} . A similar conclusion stands for sub-GeV DM compared to the SENSEI [38] and DAMIC [39] bounds on DM-electron scattering cross section.

As mentioned in the introduction, so far the measurements of astrophysical or cosmic rays place the most stringent constraints on DM annihilation cross sections within the DM mass range of $m_\psi > 1$ MeV. Take the $e\bar{e}$ final states as shown in diagram (b) of Fig. 4, for example. The DM annihilation cross section reads as

$$\sigma_{\text{ann}}(\psi\bar{\psi} \rightarrow e\bar{e}) v_{\text{rel}} \sim 10^{-68} \text{ cm}^3 \text{ s}^{-1} \cdot \kappa^2 \cdot \left(\frac{\lambda}{10^{-11}} \right)^4, \quad (10)$$

for the reference mass $m_\psi \sim 100$ GeV by using Eq. (8), where v_{rel} is the relative velocity of two incoming DM particles. Due to the λ^4 suppression, the value of

$\sigma_{\text{ann}}(\psi\bar{\psi} \rightarrow e\bar{e})v_{\text{rel}}$ in Eq. (10) is roughly ~ 38 orders of magnitude smaller than current cosmic ray bounds; see, e.g., Ref. [11].

III. LHC PHENOMENOLOGY

While it is unlikely to probe the freeze-in DM at the DM detection facilities, the scalar mediator s which only decays to a pair of DM particles can be probed by the LHC, given the κ coupling to Higgs being large enough. In this section, we firstly discuss the current LHC limit on Higgs invisible decay, then turn to HL-LHC reaches of s with c.m. energy fixed to be 14 TeV.

A. Indirect detection

For $2m_s < m_h$ with m_h the observed Higgs mass, the scalar mediator contributes to the invisible decay of Higgs with decay width

$$\Gamma_h^{\text{inv}}(h \rightarrow ss) = \frac{\kappa^2 v^2}{8\pi m_h} \sqrt{1 - \frac{4m_s^2}{m_h^2}}, \quad (11)$$

which is upper bounded as $\Gamma_h^{\text{inv}} \leq 0.26\Gamma_h$ at 95% CL [40] with the Higgs decay width $\Gamma_h \approx 4.15$ MeV.

B. Direct detection

Given the κ coupling between s and H , the scalar mediator can be produced through both VBF and the mono-Z channels, similar to the Higgs-portal scalar DM [30]. As mentioned in the introduction, unlike in Ref. [30], where only $m_s \sim m_h/2$ is not excluded, our model reopens a larger mass window of $m_s \geq 1$ GeV.

1. Productions

We use FeynRules [41] to generate model files prepared for MadGraph5 [42] which includes Pythia6 [43] for parton showering and hadronization, and the package Delphes3 [44] for fast detector simulation.³

i. VBF. As induced by mixing effect [46], the VBF channel corresponds to $q\bar{q} \rightarrow h^* + jj \rightarrow ss + jj$ as shown in the plot (a) of Fig. 5, where q and q' refer to SM quarks, and s contributes to a large missing transverse momentum. We present in the left plot of Fig. 6 the values of signal cross section $\sigma_{s,\text{VBF}}$ as function of m_s at the 14-TeV LHC, for various values of $\kappa = \{10^{-3}, 10^{-2}, 10^{-1}, 1\}$. In this plot, for a fixed value of κ the magnitudes of $\sigma_{s,\text{VBF}}$ vary over 6 orders from $m_s = 10$ GeV to 100 GeV, which significantly decrease at $m_s \sim m_h/2$. With respect to this VBF channel, the SM background arises from $W(\ell\nu) + \text{jets}$, $Z(\nu\bar{\nu}) + \text{jets}$, with subdominant $t\bar{t} + \text{jets}$ and QCD multijets neglected. For the final states with two jets, the cross section

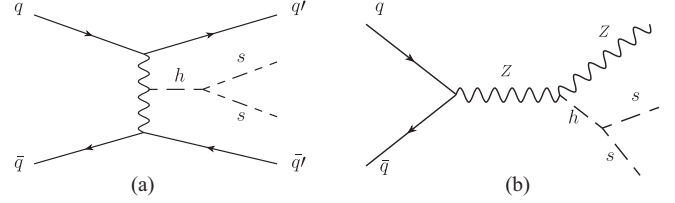


FIG. 5. Leading-order Feynman diagram with respect to the (a) VBF and (b) mono-Z channels at the LHC, respectively.

$\sigma_{b,\text{VBF}} \approx 6.4(1.2) \times 10^4$ pb, implying that a large portion of the parameter regions with $m_s \sim 10\text{--}100$ GeV and $\kappa \sim 10^{-3} - 1$ is in the reach of the HL-LHC.

ii. Mono-Z. The mono-Z channel refers to $q\bar{q} \rightarrow h^* + Z \rightarrow ss + Z(\ell\bar{\ell})$ as shown in the plot (b) of Fig. 5. The right plot of Fig. 6 shows the values of signal cross section $\sigma_{s,\text{MZ}}$ as function of m_s at the 14-TeV LHC for various values of $\kappa = \{10^{-3}, 10^{-2}, 10^{-1}, 1\}$. In this plot, the magnitudes of $\sigma_{s,\text{MZ}}$ have a dramatic decline at $m_s \sim m_h/2$, similar to the left plot of Fig. 6, despite being smaller. Regarding this mono-Z channel, the SM background arises from $Z(\nu\bar{\nu})Z(\ell\bar{\ell})$ with an explicit value of cross section $\sigma_{b,\text{MZ}} \approx 30$ pb, where $\ell = e, \mu$. Compared with the VBF channel, the mono-Z channel is subleading but relatively cleaner. Likewise, a large portion of the parameter regions with $m_s \sim 10\text{--}100$ GeV and $\kappa \sim 10^{-3} - 1$ can be also reached by the HL-LHC.

Compared to the VBF process, while gluon fusion process produces ~ 5 times larger cross section of monojet signal in the mass range of $m_s \sim 10\text{--}50$ GeV, the efficiency of this signal selection is however much smaller [47], making the ability of exclusion weaker; see, e.g., Ref. [48]. As a result, we do not consider the gluon fusion process in the following analysis.

2. Cuts

i. VBF. Reference [30] has applied the earlier CMS cuts in Ref. [49] to a numerical simulation of the scalar DM via the VBF channel. Those cuts give a ratio of $\epsilon_s/\sqrt{\epsilon_b} \sim 5$ for 1 million events, where ϵ_s and ϵ_b refer to signal and background efficiency, respectively. In this work, we make use of recently reported CMS [50] and ATLAS [51] cuts to update the simulation of the VBF channel. Explicitly,

(i) First, we impose the following trigger and selection cuts:

$$\Delta R_{jj} > 0, \quad \eta_j < 5, \quad P_T^j > 25 \text{ GeV}, \quad (12)$$

where $\Delta R_{jj} = \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2}$, with $\Delta\eta$ being the pseudorapidity difference between the two leading jets, and P_T^j and η_j are the transverse momentum and pseudorapidity of jet $j = 1, 2$, respectively.

³Apart from such sophisticated methods, new machine learning methods such as XGBoost [45] are also viable.

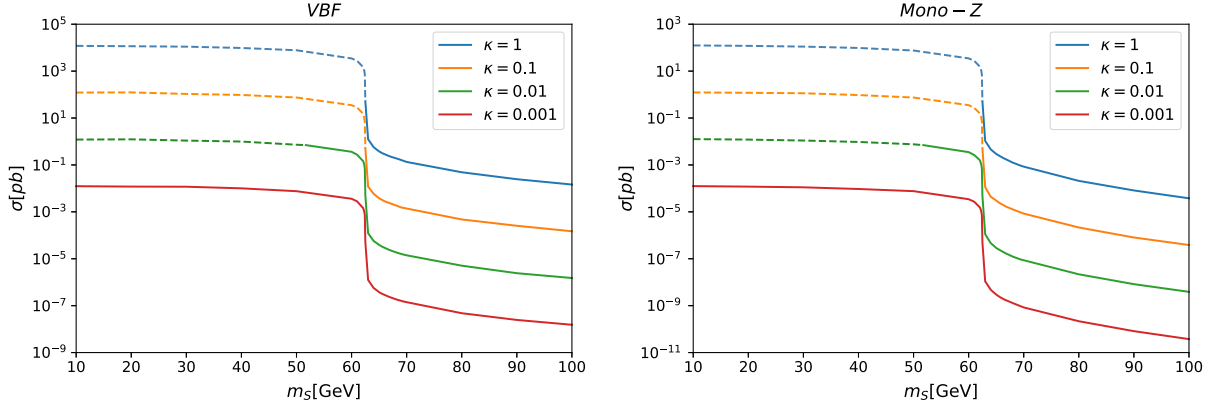


FIG. 6. Values of signal cross section $\sigma_{s,\text{VBF}}$ (left) and $\sigma_{s,\text{MZ}}$ (right) as function of m_s for various values of $\kappa = \{10^{-3}, 10^{-2}, 10^{-1}, 1\}$ at the 14-TeV LHC, which have a significant decline at $m_s \sim m_h/2$. With respect to each κ , the dotted line is excluded by the current bound on Γ_h^{inv} in Eq. (11).

- (ii) Second, we use MadAnalysis [52] to analyze the distribution of both signal and background events as a function of relevant parameters in Fig. 8 of the Appendix, where $P_T^{j_{1(2)}}$ and $\eta_{j_{1(2)}}$ are the transverse momentum and pseudorapidity of the first (second) leading jet, respectively; M_{jj} is the invariant mass of the two leading jets; and E_T^{miss} is the missing transverse energy.
- (iii) Third, we choose the explicit values of those cut parameters one by one by referring the cuts used in Refs. [50,51]. During this process, when a new added cut such as azimuthal angle $\Delta\phi$ between the two jets does not offer an effective increase in the ratio of $\epsilon_s/\sqrt{\epsilon_b}$ but obviously reduces the number of signal events, we neglect that cut parameter for simplicity.

Table I shows the final cuts applied to the VBF channel, in terms of which the value of $\epsilon_s/\sqrt{\epsilon_b}$ is enhanced by ~ 3 –4 times compared to that of Ref. [30].

ii. Mono-Z. Using the earlier CMS cuts in Ref. [53] that give a ratio of $\epsilon_s/\sqrt{\epsilon_b} \sim 15$ for 1 million events, Ref. [30] also analyzed the scalar DM reach via the mono-Z channel at the HL-LHC. Now we repeat this analysis for the scalar mediator by referring to updated cuts in Refs. [54–56]. To do so,

- (i) First, as in Eq. (12), we impose the following trigger and selection cuts:

$$\Delta R_{\ell\ell} > 0.4, \quad \eta_\ell < 2.5, \quad P_T^\ell > 10 \text{ GeV}, \quad (13)$$

where $\Delta R_{\ell\ell}$ is the separation between two leptons and P_T^ℓ and η_ℓ are the transverse momentum and pseudorapidity of lepton $\ell = 1, 2$, respectively.

- (ii) Second, we pick cut parameters specified for the lepton pair $\ell\bar{\ell}$ with $\ell = e, \mu$ and neglect those related to $\ell = \tau$ and jets. We use MadAnalysis [52] to carry out the distribution of both signal and background events as function of the selected cuts in Fig. 9 of the Appendix, where $P_T^{\ell\ell}$ is the dilepton transverse momentum, $M_{\ell\ell}$ the dilepton mass, $\Delta\phi(\vec{P}_T^{\ell\ell}, \vec{P}_T^{\text{miss}})$ the azimuthal separation between dilepton and missing momentum, and $\Delta R_{\ell\ell}$ the separation between two leptons.
- (iii) Third, repeating the same process as in the VBF channel, we choose the explicit values of those cut parameters one by one.

Table I shows the final cuts applied to the mono-Z channel, in term of which one obtains an enhancement on the ratio of $\epsilon_s/\sqrt{\epsilon_b}$ by ~ 2 –3 times compared to Ref. [30].

3. Sensitivities of HL-LHC

Collecting the previous results about the cross sections and efficiencies, we now present the 2σ exclusion (5σ discovery) defined as

$$\frac{S}{\sqrt{B}} \left(\frac{S}{\sqrt{S+B}} \right) = 2(5), \quad (14)$$

where S and B are the number of signal and background events, respectively. In this work, we neglect systematic uncertainties of order a few percent level such as those in

TABLE I. Cuts adopted for numerical simulations of the VBF (left) and mono-Z (right) channels, respectively.

VBF	Mono-Z
$N(\text{jet}) \geq 2$	$N(\ell) = 2, \ell = e, \mu$
$N(\ell) = 0$	$P_T^{\ell\ell} > 100 \text{ GeV}$
$P_T^{j_1} > 60 \text{ GeV}$	$ M_{\ell\ell} - m_Z > 10 \text{ GeV}$
$ \eta_{j_2} > 3$	$E_T^{\text{miss}} > 180 \text{ GeV}$
$E_T^{\text{miss}} > 100 \text{ GeV}$	$ E_T^{\text{miss}} - P_T^{\ell\ell} /P_T^{\ell\ell} < 1$
$M_{jj} > 1100 \text{ GeV}$	$\Delta\phi(\vec{P}_T^{\ell\ell}, \vec{P}_T^{\text{miss}}) > 2.5 \text{ rad}$
$\Delta R_{jj} > 5.5$	$\Delta R_{\ell\ell} < 1.2$

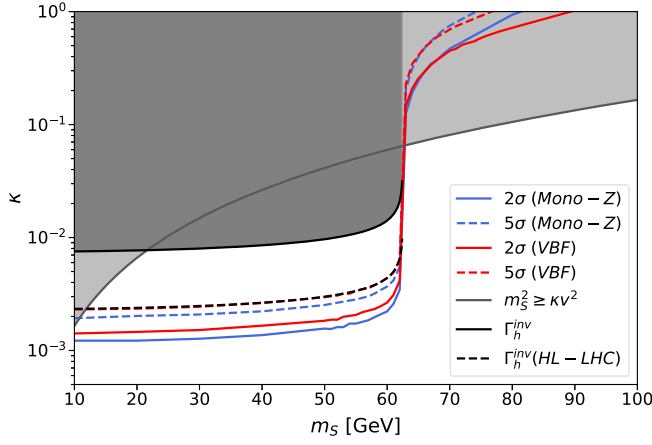


FIG. 7. Sensitivities of the HL-LHC with 2σ exclusion and 5σ discovery with respect to the integrated luminosity $\mathcal{L} = 3 \text{ ab}^{-1}$ for the VBF (red) and mono-Z (blue) channels, respectively, which are projected to the plane of $m_s - \kappa$ evading the BBN constraint on τ_s in Eq. (7). Here, the shaded regions are excluded by the current LHC bound on Γ_h^{inv} (in black) and the scalar mass bound (in gray). The future HL-LHC limit $\Gamma_h^{\text{inv}} \leq 2.5\% \Gamma_h$ [57] (black dashed line) is shown for comparison.

parton showers and the proton PDFs; see Refs. [51,56] for details.

i. VBF. The red curves of Fig. 7 shows the 2σ exclusion and 5σ discovery of HL-LHC projected to the plane of (m_s, κ) for the integrated luminosity $\mathcal{L} = 3 \text{ ab}^{-1}$, where significant changes occur at $m_s \sim m_h/2$, as seen in Fig. 6. In this plot, the black and gray shaded regions are excluded by the bound on Γ_h^{inv} in Eq. (11) and the scalar mediator mass bound $m_s^2 \geq \kappa v^2$, respectively, pointing to an allowed parameter region with $1 \text{ GeV} < m_s < m_h/2$ and $\kappa \leq 7.5 \times 10^{-3}$. In this region, the range of $\kappa \geq 2.4(1.5) \times 10^{-3}$ with respect to $m_s \sim 10\text{--}60 \text{ GeV}$ can be discovered (excluded), implying a stronger ability of exclusion than the future HL-LHC limit on the Higgs invisible decay width $\Gamma_h^{\text{inv}} \leq 2.5\% \Gamma_h$ [57].

ii. Mono-Z. The blue curves of Fig. 7 presents the 2σ exclusion and 5σ discovery of the HL-LHC projected to the plane of (m_s, κ) for $\mathcal{L} = 3 \text{ ab}^{-1}$, where both the significant changes in these curves and the shaded regions are the same as above. Compared to the VBF process and the future HL-LHC limit on Γ_h^{inv} , the mono-Z process provides the strongest signal reaches. Explicitly, the range of $\kappa \geq 2.0(1.25) \times 10^{-3}$ with respect to $m_s \sim 10\text{--}60 \text{ GeV}$ can be discovered (excluded).

Once detected by the HL-LHC either through the VBF or mono-Z channel, the scalar signal within the mass window⁴ of $m_s \sim 10\text{--}50 \text{ GeV}$ favors an existence of the nonthermal DM as studied in Sec. II, whose mass can

⁴This scalar mediator mass window does not overlap with the narrow mass window of thermal scalar DM.

be inferred in terms of Fig. 3 with respect to an explicit value of κ .

IV. CONCLUSION

In this work, we have revisited the Higgs-portal DM scenario with the DM being nonthermal instead of thermal, as previously considered in the literature. Specifically, we have realized this idea in the two-field context, where one is the DM and the other is the scalar mediator coupled to Higgs, with the feeble interaction between them as a result of the Z'_2 -parity breaking.

We have divided the model phenomenology analysis into two parts. From the perspective of DM phenomenology, the DM obtains the observed relic density mainly through the decay of scalar mediator thermalized with the SM thermal bath, and has negligible spin-independent scattering and annihilation cross sections, as expected in a typical FIDM model. From the perspective of LHC phenomenology, for the Higgs portal coupling being not too far away from unity the scalar mediator can be detected by the LHC. Consider that m_s larger than $\sim \text{GeV}$ scale is favored by the BBN constraint, we have reported the HL-LHC reaches of the scalar mediator within the allowed mass window of $m_s \sim 10\text{--}50 \text{ GeV}$ through both the VBF and mono-Z channels. If verified, this scalar signal allows us to infer the existence of a nonthermal DM such as ψ proposed here, as interpreting it as the thermal DM has been excluded.

Our study, together with the previous works [20–23], serves as an illustration of rich collider phenomenology in two-field FIDM, compared to the single-field situation. Instead of the Higgs portal considered here, one can use a neutrino or vector portal to construct alternative two-field FIDM models, following our classifications in the introduction. Different portal may provide unique signals. We left this point for future work.

DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX: SELECTION OF CUTS

After imposing the cuts in Eq. (12), Fig. 8 shows the distribution of the VBF signal and background events as functions of $p_T^{j(2)}$, $|\eta_{j(2)}|$, E_T^{miss} , M_{jj} , $\Delta\phi_{jj}$, and ΔR_{jj} respectively, which illustrates the cut values chosen for the VBF channel in Table I. For example, imposing $\Delta\phi_{jj}$ or $P_T^{j(2)}$ indeed increases the value of $\epsilon_s/\sqrt{\epsilon_b}$ more or less but also leads to a serious suppression on the number of signal events simultaneously. On the contrary, other cuts such

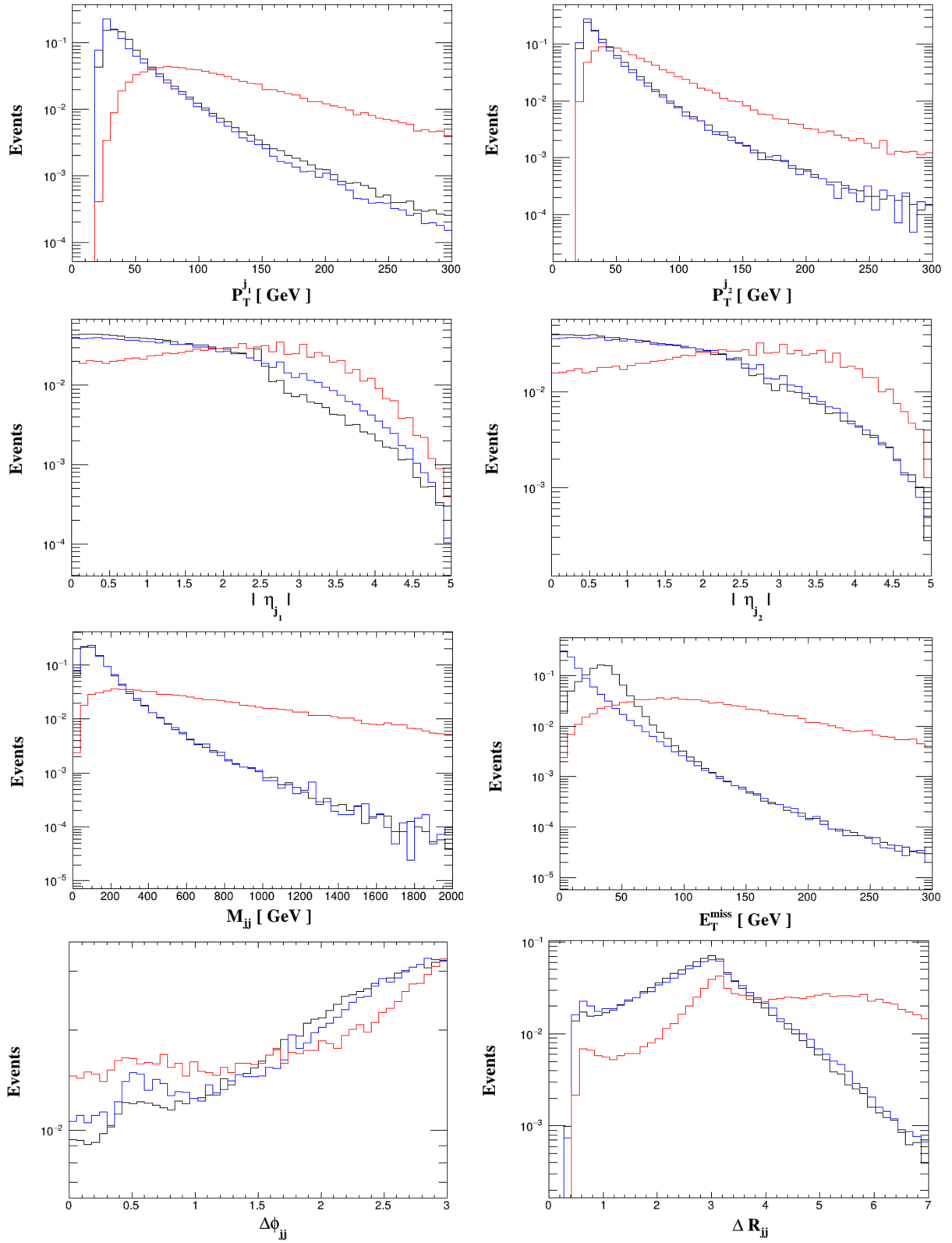


FIG. 8. Distribution of the VBF signal and background events, where the signal, W + jets and Z + jets are shown in red, black, and blue, respectively.

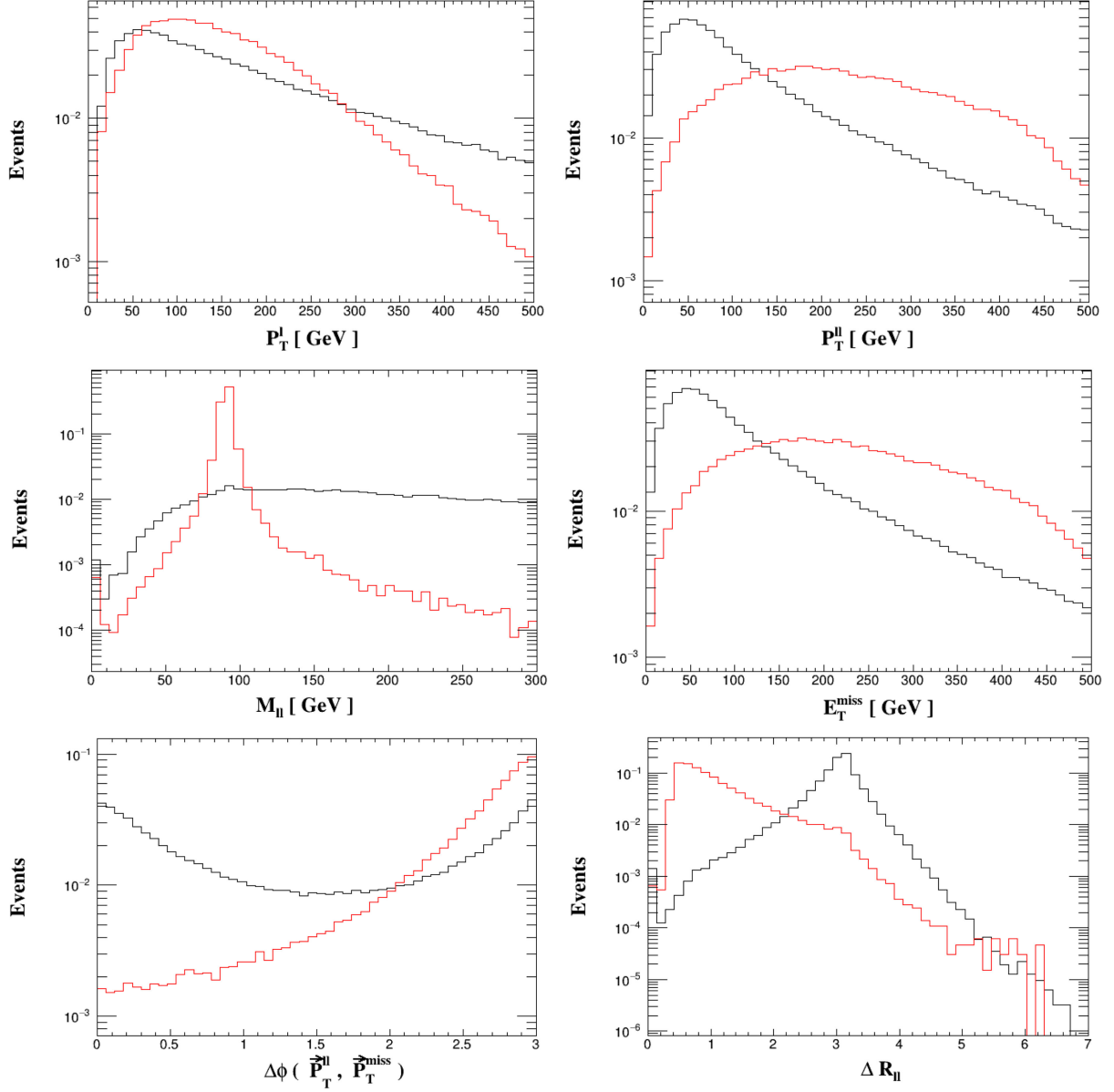


FIG. 9. Distribution of the mono-Z signal and background events, which are shown in red and black, respectively.

$P_T^{j_1} > 60$ GeV can do so with the suppression being affordable.

Likewise, after imposing the cuts in Eq. (13), Fig. 9 shows the distribution of mono-Z signal and background events as functions of P_T^ℓ , $P_T^{\ell\ell}$, $M_{\ell\ell}$, E_T^{miss} , $\Delta\phi(\vec{P}_T^{\ell\ell}, \vec{P}_T^{\text{miss}})$,

and $\Delta R_{\ell\ell}$ respectively, which explains the cut values chosen for the mono-Z channel in Table I. For example, imposing P_T^ℓ is unable to increase the value of $\epsilon_s/\sqrt{\epsilon_b}$ in an efficient way, as explained above, but other cuts such as $P_T^{\ell\ell} > 150$ GeV can do so.

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