Indirect detection imprint of a \( CP \) violating dark sector

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We introduce a simple scenario involving fermionic dark matter (\( \chi \)) and singlet scalar mediators that may account for the Galactic center GeV \( \gamma \)-ray excess while satisfying present direct detection constraints. \( CP \) violation in the scalar potential leads to a mixing between the Standard Model Higgs boson and the scalar singlet, resulting in three scalars, \( h_{1,2,3} \), of indefinite \( CP \)-transformation properties. This mixing enables \( s \)-wave \( \chi \chi \) annihilation into discalar states, followed by decays into four-fermion final states. The observed \( \gamma \)-ray spectrum can be fitted while respecting present direct detection bounds and Higgs boson properties for \( m_\chi \approx 60 \sim 80 \) GeV, and \( m_{h_i} \sim m_\chi \). Searches for the Higgs exotic decay channel \( h_i \rightarrow h_3h_3 \) at the 14 TeV LHC should be able to further probe the parameter region favored by the \( \gamma \)-ray excess.

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I. INTRODUCTION

Although the presence of dark matter (DM) has been firmly established by numerous observational data via its gravitational effects, the particle nature of DM remains a mystery. It is imperative to search for DM in every feasible way: direct detection, indirect detection, collider searches, etc. Both direct detection and collider searches have observed null results, which put constraints on particle DM properties. On the other hand, indirect detection offers some hints of the particle nature of the DM. The annihilation or decays of particle DM in the galaxies are expected to produce observable fluxes of cosmic rays, such as antiprotons, positrons, gamma rays, and neutrinos. Of particular interest are gamma rays from the Galactic center and the dwarf spheroidal galaxies, which are able to be detected by the Fermi Large Area Telescope (Fermi-LAT) \[1\].

Recent analyses of the Fermi-LAT data by several groups have identified an excess of gamma rays with several GeV energy and a nearly spherically symmetric distribution in the center of the Milky Way, known as the Galactic center excess (GCE) \[2–6\]. Although astrophysical explanations such as millisecond pulsars have been proposed, the DM annihilation explanation of the GCE has generated much recent attention and is being widely explored. The reason is that the morphology of the GCE is consistent with what is expected from DM annihilation while complying with the observed thermal density.

A common and simple class of DM scenarios that can explain the GCE is the two-body DM annihilation \( \chi\chi \rightarrow f\bar{f} \), where \( f \) represents a Standard Model (SM) fermion. The spectrum of the GCE has been fit well by the DM annihilation into \( bb \) with \( m_\chi \) around \( 31 \sim 40 \) GeV and the thermal averaged cross section \( \langle \sigma v \rangle_{bb} \sim O(1 \sim 3) \times 10^{-26} \) cm\(^3\)/s \[4,7\]. Dark matter annihilation into \( \tau\tau \) provides a acceptable fit to the spectrum with lighter \( m_\chi \) around 10 GeV and smaller annihilation cross section. In these models, the DM \( \chi \) interacts with the SM fermion through a mediator \( \phi \), which could be either scalar, fermion, or vector boson. There are two types of scenarios that explain the GCE through 2 \( \rightarrow 2 \) annihilation: \( s \)-channel models with a neutral mediator \( \phi \) \[7\], and \( t \)-channel models with a charged mediator \( \phi \) \[8,9\]. Usually in both models a heavy mediator \( \phi \) is needed to avoid the LHC constraints, and there should be a mechanism to suppress the spin-independent DM-nucleus scattering cross section.

An interesting case is the fermionic DM model with a light neutral pseudoscalar mediator \[7,10–15\], in which the Lagrangian can be written as

\[ \mathcal{L} \supset g_f \bar{f} i \gamma_5 \phi + g_f \bar{f} \gamma_5 \phi, \]

where \( \chi \) is DM, \( f \) represents SM fermions, and \( \phi \) is a pseudoscalar. The direct detection cross section in this case is purely spin dependent, and thus the direct detection rate is significantly reduced. Thus, pseudoscalar mediated DM models have received much attention in the context of explaining the GCE, with the annihilation channels \( \chi\chi \rightarrow f\bar{f} \) \[7,10,11\]. However, there are some tensions with this model. First, the requirement of the correct relic density \( \Omega_c \) prefers a moderate value of \( g_f \), which is constrained by collider searches \[16,17\]. Second, in many UV complete scenarios, \( \phi \) is degenerate with a \( CP \)-even Higgs boson, which is highly constrained by the LHC heavy Higgs searches. It is not easy to obtain a light pseudoscalar and a heavy real scalar with large mass splitting. More importantly, the latest results on the dwarf spheroidal galaxies \[18\]...
put strong bounds on the two-body fermion final states, including $b\bar{b}$ and $\tau\bar{\tau}$ final states, creating tension between the GCE signal parameter region and the allowed dwarf spheroidal region.

The above issues might be addressed when $\phi$ is lighter than $\chi$. In this case $\chi$ can annihilate into a pair of $\phi$’s which subsequently decay into SM particles: $\chi\chi \rightarrow \phi\phi \rightarrow f\bar{f}f\bar{f}$. In this way, $\Omega_\chi$ does not depend on interactions between $\phi$ and $f$, which can be quite weak. Furthermore, the DM annihilation products are four-fermion final states, which is still compatible with current constraints from dwarf spheroidal galaxies. This is the so-called hidden sector dark matter scenario [19,20]. In the hidden sector scenario, there are large couplings among the hidden sector particles but small couplings of the hidden particles to the SM particles. This scenario can evade the tension between tight constraints from direct detection and LHC searches and a large GCE signature.

However, there are caveats in this particular hidden sector scenario: the $\sigma(\chi\chi \rightarrow \phi\phi)$ is typically a $p$ wave. This gives rise to a negligible indirect detection signature. Let us understand this from the parity transformation property of the initial and final states, and angular momentum conservation. We know that if the annihilation amplitude has zero orbital angular momentum, the annihilation cross section should be $s$-wave annihilation. Under a parity ($P$) transformation, the fermion-antifermion initial state transforms as $(-1)^{L+1}$, where $L$ is the total orbital angular momentum. Depending on the final states, we have

(i) In Fig. 1(a), two identical pseudoscalars in the final state. Since the two boson final state is symmetric under interchange, the $P$ transformation is simply $P = 1$. Therefore, although total angular momentum conservation gives rise to $L = 0$, 1, from the parity we determine that the total angular momentum is $L = 1$, which implies the annihilation cross section is $p$-wave suppressed.

(ii) In Fig. 1(b), two identical scalars in the final state. From the similar argument above, we obtain $L = 1$ and thus the annihilation cross section is $p$ wave.

(iii) In Fig. 1(c), two different scalar bosons in the final state. In this case, there is no such exchange symmetry. Thus the orbital angular momentum could be zero. The annihilation cross section should have $s$ and $p$ waves.

From the above arguments, we note that if the final states have an odd number of pseudoscalars, the annihilation is $s$ wave. Therefore, there are several ways to realize the $s$-wave annihilation cross section in the hidden sector DM.

(i) One way is that three pseudoscalars are produced in the annihilation process, and thus the cross section $\chi\chi \rightarrow \phi\phi\phi$ is $s$ wave [12]. Although phase space suppression exists, if the interactions between the DM and the pseudoscalar are much larger than those between the DM and the SM fermions, the annihilation channel $\chi\chi \rightarrow \phi\phi\phi$ is still larger than the $s$ channel $\chi\chi \rightarrow f\bar{f}f$. The gamma-ray signature comes from the six-fermion final states. However, using this channel it is challenging to obtain both the GCE signature and the correct relic density [12].

(ii) Another possible way [14] is that if there are two light pseudoscalars, the annihilation process $\chi\chi \rightarrow \phi_1\phi_2$ is $s$ wave. This still needs model building efforts to split the masses of the light pseudoscalars from the heavy real ones.

We propose a third alternative. Instead of a pseudoscalar, a complex scalar singlet $S = (s + ia)/\sqrt{2}$ is introduced in the hidden sector. Due to $CP$ violation in the scalar potential, the $CP$-even and $CP$-odd field components will mix with each other and with the Standard Model Higgs boson. Thus, the resulting mass eigenstates, $h_{1,2,3}$, couple to both of the $\chi\chi$ and $\chi\pi_5\pi_5$ bilinears. Assuming $m_{h_1} \ll m_{h_2}(m_{h_3})$, which can be realized via the $CP$ violating terms in the Higgs potential, and $m_{h_1} < m_\chi$, the process $\chi\chi \rightarrow h_1h_3$ is thus kinematically allowed and the cross section can be $s$ wave. The reason is that the amplitude $\chi\chi \rightarrow h_1h_3$ contains the parity odd bilinear $\chi\pi_5\pi_5$. We show that this scenario can readily accommodate the GCE with a thermal relic cross section, and still satisfy the other constraints. Here are the main results:

(i) The annihilation rate $\chi\chi \rightarrow h_1h_3$ depends on the $CP$ violating phases in the scalar potential. The larger the $CP$ violating phase, the larger the annihilation rate in the indirect detection.

(ii) Direct detection depends on both the $CP$ violation strength and the couplings of the scalars to the SM quarks. Since the hidden sector has small couplings to the SM particles, direct detection constraints may be avoided even though there are large $CP$ violating phases.

(iii) To fit the GCE spectrum, the DM annihilation cross section favors the thermal relic rate. This could be realized via sufficient $CP$ violation. We will show that the cascade annihilation $\chi\chi \rightarrow \phi\phi \rightarrow f\bar{f}f\bar{f}$ could explain GCE while still being consistent with dwarf spheroidal constraints.

(iv) If $m_{h_1} < m_{h_1}/2$, the Higgs boson $h_1$ will have an exotic decay channel $h_1 \rightarrow h_3h_3$. This gives us an

![FIG. 1. The Feynman diagrams on DM annihilation to (a) two light pseudoscalar $AA$’s, (b) two light real scalar $SS$’s, and (c) one light real scalar $S$ and another light pseudoscalar $A$.](image)
additional probe on the CP violating phases. Depending on the self-coupling of the complex scalar, the CP violating phases could be probed at the run 2 LHC with high luminosity.

Our discussion of this scenario is as follows. We begin with the description of the CP violating complex scalar singlet model. In Sec. III, we discuss the DM relic density and direction in this model. In Sec. IV, we present constraints on the model from oblique parameters and a Higgs measurement. In Sec. V, we discuss the GCE arising from the cascade annihilation. In Sec. VI, we study signatures of the model at the LHC. We give concluding remarks in Sec. VII.

II. THE COMPLEX CP VIOLATING SCALAR SINGLET MODEL

As discussed in the Introduction, the hidden sector includes a Dirac fermion DM χ and a complex scalar singlet S = (s + iα)/√2. The interaction between the DM and the scalar singlet can be written as

$$\mathcal{L}_{\text{DM}} = \tilde{\chi}^\mu \partial_\mu \chi - m_0 \tilde{S} S - y_{\tilde{\chi} S} L S + H.\text{c.}$$ (2)

In general, the complex scalar S also interacts with the SM Higgs boson. This complex singlet scalar singlet extended SM is referred to as the complex scalar singlet model (cSM) [21,22]. The tree-level scalar potential can be written as

$$V_{c\text{SM}} = -\mu_h^2 H^\dagger H + \lambda_h (H^\dagger H)^2 - \mu_\chi^2 S^\dagger S + \lambda_\chi (S^\dagger S)^2 + \lambda_{\chi h} S^\dagger H^\dagger H + \lambda_C S^4 + \text{H.c.}. $$ (3)

where H is the SM Higgs doublet. Here although there is a Z_2 symmetry in the tree-level scalar potential, this Z_2 symmetry is broken by the Yukawa term in Eq. (2), and thus there is no domain wall problem. The mass term μ_h^2 and the couplings λ_h, λ_C can be treated as the spurious [23], which might trigger explicit or spontaneous CP violations. We assume there is only explicit CP violation for simplicity. To parametrize the CP violating phases, we define the rephrasing invariants as follows:

$$\delta_1 = \text{Arg}(\lambda_B \mu_\chi^2), \quad \delta_2 = \text{Arg}(\lambda_C \mu_\chi^2), $$ (4)

whose expressions in terms of the physical parameters of the model will be given at the end of this section. The SM Higgs doublet and S are written in component form as follows: $H = (G^+, v_h + \hat{h} + iG^0)^T/\sqrt{2}$, and $S \equiv (v_s + \hat{s} + i\tilde{a})/\sqrt{2}$, where $v$ and $v_s$ are the vacuum expectation values of $H$ and $S$, respectively, determined by the tadpole conditions,

$$\frac{\partial V_{c\text{SM}}}{\partial \phi_i} \bigg|_{\phi=(\hat{h}, \hat{s}, \tilde{a})=0} = 0. $$ (5)

The scalar mass matrix in the basis $(\hat{h}, \hat{s}, \tilde{a})$ is

$$-2\text{Im}(\lambda_B) v v_s, \quad -4\text{Im}(\lambda_C) v_s^2, \quad 4\text{Re}(\mu_\chi^2) - 8\text{Re}(\lambda_C) v_s^2 - 2\text{Re}(\lambda_B) v^2$$ (6)

where $h_1$ is identified as the SM Higgs boson, and $h_3$ is the light mediator to the DM. Here $h_3$ is assumed to be very heavy to avoid possible direct detection constraints. The mixing angles $\theta_{13}$ and $\theta_{23}$ parametrize the CP violating phases $\delta_{1,2}$.

The DM mass is obtained as

$$m_\chi = m_0 + \frac{1}{\sqrt{2}} y_\chi v_s, $$ (9)

and its interaction with the S takes the form $y_{\tilde{\chi} S} (\hat{s} + i\tilde{a}) \chi$. Interactions in the scalar mass eigenbasis can be parametrized as

$$\mathcal{L} \supset \bar{\chi}(\lambda_{\chi i} + \lambda_{pi} y_s^4) \chi + \tilde{f} (g_{si} + g_{pi} i y_s^4) f| h_i, $$ (10)

where $f$ represents SM fermions and the coupling strengths are

$$\lambda_{\chi i} = -hy_\chi U_{2i}/\sqrt{2}, \quad \lambda_{pi} = -iy_\chi U_{3i}/\sqrt{2}. $$ (11)
has a very small coupling to SM particles. This implies a small \( \mu \) between them. Finally, we need a large mass splitting between \( \chi \) and \( \bar{\chi} \).

The Feynman rules for the scalar interactions are given in Table I. Among these, the most relevant couplings are

\[ \bar{\chi} \chi h_3: (i) y_\chi(U_{23} + U_{31} i \gamma_5)/\sqrt{2}; \]

\[ \bar{f} f h_3: (-i) U_{12} m_f/v; \]

\[ \bar{\chi} h_1: (i) y_\chi (U_{21} + i U_{31} \gamma_5)/\sqrt{2}; \]

\[ \bar{f} f h_1: (-i) U_{13} m_f/v. \]

As the hidden scalar mediator, the light mediator \( h_3 \) only has a very small coupling to SM particles. This implies a small \( s_{13} \) is favored. To have a sizable coupling \( \bar{\chi} \chi h_3 \), \( s_{23} \) should be moderately large, which induces a large mixing between \( s \) and \( \bar{s} \). Furthermore, to avoid constraints from the high mass Higgs searches [24], the mixing angle \( s_{12} \) should be small. Finally, we need a large mass splitting between \( h_2 \) and \( h_3 \), which can be realized through the CP violating terms in the potential.

Before proceeding to study constraints on the parameter space of this model, we count scalar sector physical parameters from scalar interactions, which are \( m_{h_1}, m_{h_2}, m_{h_3}, v, v_s, \theta_{ij}, \lambda_s \), and \( \lambda_{sih} \). The mass squared parameters \( \mu^2_\chi \) and \( \mu^2_\bar{\chi} \) can be determined by the tadpole conditions, while other parameters can be reconstructed as

\[ \text{Re}(\lambda_B) = -\frac{\lambda_{1h}}{2} + \frac{1}{2 v v_s} (m^2_{h_1} U_{11} U_{21} + m^2_{h_2} U_{11} U_{22} + m^2_{h_3} U_{11} U_{23}). \]

\[ \text{Im}(\lambda_B) = \frac{1}{2 v v_s} (m^2_{h_1} U_{11} U_{31} + m^2_{h_2} U_{11} U_{32} + m^2_{h_3} U_{11} U_{33}), \]

\[ \text{Re}(\lambda_C) = -\frac{\lambda_2}{2} + \frac{1}{4 v_s^2} (m^2_{h_1} U_{21} U_{31} + m^2_{h_2} U_{22} U_{32} + m^2_{h_3} U_{23} U_{33}). \]

\[ \text{Im}(\mu^2_\chi) = -\frac{1}{2} \text{Im}(\lambda_B) v^2 - \text{Im}(\lambda_C) v^2, \]

\[ \text{Re}(\mu^2_\chi) = -\frac{1}{2} \text{Re}(\lambda_B) v^2 - 2 \text{Re}(\lambda_C) v^2_\chi \]

\[ -\frac{1}{4} (m^2_{h_1} U^2_{31} + m^2_{h_2} U^2_{32} + m^2_{h_3} U^2_{33}), \]

\[ \lambda_h = \frac{1}{2} v^2 (m^2_{h_1} U^2_{11} + m^2_{h_2} U^2_{12} + m^2_{h_3} U^2_{13}). \]

The rephasing invariants can be expressed as

\[ \delta_1 = \text{arctan} \left( \frac{\text{Im}(\mu_B)}{\text{Re}(\mu_B)} \right) - \text{arctan} \left( \frac{\text{Im}(\mu_A)}{\text{Re}(\mu_A)} \right), \]

\[ \delta_2 = \text{arctan} \left( \frac{\text{Im}(\mu_C)}{\text{Re}(\mu_C)} \right) - 2 \text{arctan} \left( \frac{\text{Im}(\lambda_B)}{\text{Re}(\lambda_B)} \right). \]

### III. RELIC DENSITY AND DIRECT DETECTION

In the standard weakly interacting massive particle [25] scenario, \( \chi \) thermally freezes out, leaving a significant relic abundance. In this model, \( \chi \) and \( h_3 \) are assumed to be in the mass range of \( 10 \sim 100 \text{ GeV} \). \( h_2 \) is assumed to be sufficiently heavy that its contribution to the relic density is negligible. In this parameter region, the annihilation processes are \( \chi \bar{\chi} \rightarrow f \bar{f}, \chi \bar{\chi} \rightarrow WW/ZZ, \) and \( \chi \bar{\chi} \rightarrow h_3 h_3 \). We will calculate the thermal relic cross sections in these channels.

The annihilation \( \chi \bar{\chi} \rightarrow f \bar{f} \) is through the \( s \)-channel \( h_i \) \((i=1,2,3)\) exchange, shown in Fig. 2(a). The \( s \)-wave part of this \( s \)-channel thermal cross section is

\[ (\sigma v)_{\chi \bar{\chi} \rightarrow s f} = N_c \sum_{i=1,3} \frac{\lambda^2_{pi} m^2_{\chi} (m^2_{h_i} - 4 m^2_{\chi} + m^2_i \Gamma^2_i)}{2 \pi m^4_{\chi} (m^2_{h_i} - m^2_{h_{i-1}})^{3/2}}, \]

where \( N_c \) is the number of colors for \( f \), and \( \Gamma_i \) is the total decay width of \( h_i \). Note that this thermal cross section

![FIG. 2. Two dominant DM annihilation processes. (a) The \( s \)-channel DM annihilation process. (b) The \( t \)-channel DM annihilation process.](https://example.com/f2.png)
is proportional to the coupling strengths $\lambda_{pi}$ and $g_{si}$. We consider that new scalars couple to the DM significantly but have negligible couplings to the SM particles, which implies small couplings $g_{si}$. Thus we expect a small thermal cross section arising from Eq. (24), except in the presence of resonant enhancements from the $s$-channel mediator $h_1$ or $h_3$. When $m_s > m_V$, there exists the $\chi\bar{\chi} \to WW$ and $\chi\bar{\chi} \to ZZ$ channels, in which the situation is similar to that of $\chi\bar{\chi} \to f\bar{f}$. Apart from the resonance enhanced region, to obtain the correct relic density, the $t$-channel annihilation should be dominant over $s$-channel processes. Here we should mention that the couplings $g_{si}$ could not be very small. To keep the hidden sector particles in the thermal bath, the rates for the decay and inverse decay processes must be faster than the Hubble rate during freeze-out. This puts a lower bound on the magnitude of $g_{si}$.

When $m_s > m_h$ and $g_{si}$ is small, the dominant channel will be the $t$-channel process $\chi\bar{\chi} \to h_3h_3$, shown in Fig. 2(b). The relevant thermal cross section is

$$\langle \sigma v \rangle_{\chi\bar{\chi} \to h_3h_3} = \frac{\lambda_{si}^2}{12 \pi} \frac{1}{m_t} \sqrt{m_t^2 - m_h^2} \left( \frac{m_t^2 - m_h^2}{2} \right)^{3/2} \left( \frac{m_t^2}{m_h^2} \right)^{1/2} \langle v^2 \rangle$$

$$+ \left( \lambda_{s3}^4 + \lambda_{s3}^4 \right) \frac{m_x^2 (m_x^2 - m_h^2)}{12 \pi (m_t^2 - m_h^2)^4} \langle v^2 \rangle$$

$$- \lambda_{s3}^4 \lambda_{s3}^4 \frac{m_x^2 (32 - 52 x + 20 x^2 - 3 x^3)}{48 \pi (m_t^2 - m_h^2)(2 m_x^2 - m_h^2)^4} \langle v^2 \rangle,$$

(25)

where $x = m_h^2 / m_t^2$. The cross section only depends on the DM couplings. From Eq. (25), we notice that when there are both scalar and pseudoscalar interactions of $h_3$ with $\chi$, the annihilation cross section could be $s$ wave. If the $h_3$ is purely scalar or pseudoscalar, the annihilation cross section will be only $p$ wave. Thus the $h_3$ needs to be a mixture of the scalar $\hat{s}$ and the pseudoscalar $\hat{a}$, which comes from the $CP$ violating terms in the scalar potential.

We calculate the relic density and direct detection cross section numerically using micrOMEGAs [26], which solves the Boltzmann equations numerically and utilizes CalcHEP [27] to calculate the cross section. Mixing angles $s_{12} = s_{13} = 0.05$ and $s_{23} = 0.2$, and $s_{23} = 1/\sqrt{2}$ are assumed.

![FIG. 3. Contours of the coupling strength $y_\chi$ in the $(m_x, m_h)$ plane, which give rise to the correct relic abundance. Mixing angles $s_{12} = s_{13} = 0.05$ and (left panel) $s_{23} = 0.2$ and (right panel) $s_{23} = 1/\sqrt{2}$ are assumed.](image-url)
contributes to the spin-independent (SI) scattering cross section as shown in Fig. 5. The tightest bounds on the SI cross section come from the LUX data \[28\]. In this model, the SI cross section is written as

\[
\sigma_{\text{SI}} = \sum_{i=1,3} \frac{\mu_{\chi h_i}^2 U_{1i}^2}{m_{h_i}^2 v_n^2} \left( \frac{\lambda_{si}^2}{2} + \frac{\mu_{\chi N}^2 v_n^2}{m_{h_i}^2} \right) v^2 \times [Z f_p + (A-Z)f_n]^2, \tag{26}
\]

where \(\mu_{\chi N}\) is the reduced mass, \(f_{p,n}\) are the form factors of the proton and neutron, and \(v \sim 10^{-3}\) is the velocity of the DM. From Eq. (26), we note that although both the scalar and the pseudoscalar interaction of the \(h_3\) to the DM contribute to the SI cross section, the scalar interaction is dominant and the pseudoscalar interaction exhibits velocity suppression. Fixing the coupling strength \(y_\chi\) by \(\Omega_\chi\), we calculate the SI cross sections for different \((m_\chi, m_{h_3})\)’s. Figure 4 shows the exclusion limit on the \((m_\chi, m_{h_3})\) parameter space, given the central value of the observed \(\Omega_\chi\). From Fig. 4, we see that away from the Higgs resonance region the mediator \(h_3\) cannot be very light. This can be seen from the \(1/m_{h_3}^4\) dependence in the SI cross section. Near the Higgs resonance, the coupling strength \(y_\chi\) is quite small, and thus the \(m_{h_3}\) mass could be light.

**IV. ELECTROWEAK PRECISION AND HIGGS CONSTRAINTS**

Typically the nonobservation of permanent electric dipole moments (EDMs) \[29\] of neutral atoms, molecules, neutron, and electron place severe constraints on the strengths of \(CP\) violation. In our model, these constraints are negligible, because \(CP\) violation in the scalar potential cannot lead to any pseudoscalar-type Yukawa interaction of the SM fermion, which plays a key role in generating nonzero EDMs via the two-loop Barr-Zee diagram \[30\]. Constraints mainly come from the LHC Higgs measurements, electroweak precision measurements, and DM direct detections.

As can be seen from Table I, couplings of the SM-like Higgs to all SM particles are rescaled by the factor \(c_{12} c_{13}\), the square of which is equal to the signal rates \(\mu_{*=*}\) associated with Higgs measurements relative to SM Higgs expectations. In this section, we independently perform the universal Higgs fit \[31\] to the Higgs data from both ATLAS \[32\] and CMS \[33,34\], where couplings of \(h_1\) to pairs of \(t, b, \tau, W, Z, \gamma\) equal \(r_t, r_b, r_\tau, r_W, r_Z, r_\gamma\) in units of the SM Higgs couplings. The \(\chi^2\) is a quadratic function of \(\epsilon_i\), where \(\epsilon_i \equiv r_i - 1\), and can be written as

\[
\chi^2 = \sum_{i,j} (\epsilon_i - \mu_i)(\sigma_{ij}^{-1}(\epsilon_j - \mu_j), \tag{27}
\]

where \(\mu_i\) is the mean value of \(\epsilon_i\), \(\sigma_{ij}^{-1} = \sigma_i \rho_{ij} \sigma_j\) with \(\sigma_i\) being the error of \(\epsilon_i\) and \(\rho\) the correlation matrix. The result is

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**FIG. 5.** DM-nucleon interaction that would generate a spin-independent direct detection signal.
shown in Fig. 6. The red solid and blue dotted lines correspond to constraints at the 68% and 95% C.L., respectively.

We consider the electroweak precision constraints, utilizing bounds on the oblique observables [35,36], which are defined in terms of contributions to the vacuum polarizations of gauge bosons. The dependence of $S$ and $T$ parameters on the new scalars can be approximately expressed by the following one-loop terms [37]:

$$
\Delta S = \sum_{k}^{2,3} \frac{U_{1k}^2}{24\pi} \{ \log R_{Xh} + \tilde{G}(M_{Z'}^2, M_{\tilde{Z}}^2) - \tilde{G}(m_h^2, M_{\tilde{Z}}^2) \} 
$$

(28)

$$
\Delta T = \sum_{k}^{2,3} \frac{3U_{1k}^2}{16\pi s_W^2 M_W^2} \left\{ M_Z^2 \left[ \log \frac{R_{Zx}}{1 - R_{Zx}} - \log \frac{R_{Zh}}{1 - R_{Zh}} \right] + M_W^2 \left[ \log \frac{R_{Wx}}{1 - R_{Wx}} - \log \frac{R_{Wh}}{1 - R_{Wh}} \right] \right\},
$$

(29)

(30)

where $c_W = \cos \theta_W$, with $\theta_W$ the weak mixing angle, $R_{\zeta \xi} \equiv M_{\tilde{Z}}^2 / M_{\tilde{Z}}^2$, and

$$
\tilde{G}(M_{Z'}^2, M_{\tilde{Z}}^2) = \frac{79}{3} + 9R_{\zeta \xi} - 2R_{\zeta \xi}^2 + (12 - 4R_{\zeta \xi} + 2R_{\zeta \xi}^2) \tilde{F}_{\zeta \xi} + \left( -10 + 18R_{\zeta \xi} - 6R_{\zeta \xi}^2 + R_{\zeta \xi}^3 + 9 \frac{1 + R_{\zeta \xi}}{1 - R_{\zeta \xi}} \right) \times \log R_{\zeta \xi}.
$$

(31)

The most recent electroweak fit (by setting $m_{H,ref} = 126$ GeV and $m_{H,ref} = 173$ GeV) to the oblique parameters performed by the Gfitter group [38] yields

$$
S = \Delta S^0 \pm \sigma_S = 0.03 \pm 0.10
$$

(33)

$$
T = \Delta T^0 \pm \sigma_T = 0.05 \pm 0.12.
$$

(33)

Constraints can be derived by performing a $\Delta \chi^2$ fit to the data in Eq. (33), where the $\Delta \chi^2$ is given as

$$
\Delta \chi^2 = \sum_{ij}^{2} (\Delta O_{ij} - \Delta O^{0}_{ij}) (\sigma_{ij}^{-1})^{-1} (\Delta O_{ij} - \Delta O^{0}_{ij}),
$$

(34)

in which $O_1 = S$, $O_2 = T$, and $\sigma_{ij} = \sigma_{ij} \rho_{ij} \rho_{ij}$, with $\rho_{11} = \rho_{22} = 1$ and $\rho_{12} = 0.891$ [38].

We show in Fig. 6 oblique parameter constraints on the mixing angles in the $\theta_{12}-\theta_{13}$ plane, by setting $m_{h_1} = 60$ GeV and $m_{h_2} = 500$ GeV. The cyan solid and black dotted lines correspond to exclusion limits at the 95% and 68% C.L., respectively. It is clear that the oblique parameter constraints are much weaker than the ones from Higgs measurements. It is worth mentioning that the constraint of the oblique parameters could be stronger by varying the initial inputs of $m_{h_1, h_2}$.

As was discussed in the last section, there is a lower bound on $g_{\phi,i}$ from the requirement that the hidden sector particles need to be kept in the thermal bath before the freeze-out. This sets a lower limit on the mixing parameter. Assuming $m_{h_1} \sim m_{\phi} \approx 70$ GeV, we obtain the lower limit on the mixing angle $\theta_{13} > \mathcal{O}(10^{-4})$. The lower limit is typically quite small and will not affect the indirect detection signature discussed in the following. Furthermore, DM direct detection experiments also constrain the size of mixing angles. By requiring the DM-nucleus scattering cross section to lie below the exclusion limit set in the LUX experiment, one gets the yellow allowed region in the $\theta_{12}-\theta_{13}$ plane in Fig. 6, where we have assumed $m_{\phi} \approx 70$ GeV, $m_{h_1} \approx 60$ GeV, $m_{h_2} \approx 500$ GeV, and $Y_{\phi} \approx 0.28$ so as to give rise to a correct relic density. The direct detection cross section is approximately proportional to $(1 - 2 \cos 2\theta_{23})$, such that there are two regions allowed in Fig. 6. The reasoning of the $\theta_{13}$ being sensitive to the DM direct detection is that we set a small $m_{h_2}$, which is crucial for explaining the GCE.
V. GALACTIC CENTER GEV GAMMA-RAY EXCESS

Although the direct detection and Higgs measurements put constraints on the model parameters, there still exist large parameter regions that yield indirect detection signatures. Indirect detection experiments search for the products of the DM annihilation or decay. Unlike the DM annihilation during the freeze-out, only the DM annihilation process with an $s$-wave contribution contributes to the indirect detection signature. Off the Higgs resonance, the dominant $s$-wave contribution comes from the $\chi\chi \rightarrow h_3 h_3$. To see the indirect detection signature, the $h_3 ff$ coupling cannot be zero, and so $h_3$ could decay to the SM particles. Thus, CP violation is needed to have indirect detection signature. From previous sections, we learn that the $h_3$ has a very small coupling to the SM fermion. This does not affect the annihilation cross section when the $h_3$ is produced on shell because the branching ratio of $h_3$ does not depend on the $h_3 ff$ coupling. Therefore, the indirect detection signatures mainly come from the cascade decay of the two on-shell $h_3$’s: the four-fermion final states via $\chi\chi \rightarrow h_3 (\rightarrow ff) h_3 (\rightarrow ff)$ as shown in Fig. 7. Among these final states, the dominant decay channel will be four-$b$ final states, because of the relatively large bottom Yukawa coupling. This cascade annihilation and subsequent shower and hadronization produce various measurable signatures, such as gamma ray, etc. In the following, we will study the gamma-ray spectrum of this cascade annihilation.

Consider a cascade annihilation $\chi\chi \rightarrow \phi(\rightarrow ff) \phi(\rightarrow ff)$, where $\phi$ is an on-shell mediator in general. The gamma-ray spectrum can be obtained from the boost of the gamma-ray spectrum $dN_f(\phi \rightarrow ff)/dE_f$ in the $\phi$ rest frame. This spectrum can be easily obtained from PYTHIA [39]. After boosting into the lab frame, the gamma-ray spectrum is written as [14]

$$dN_f(\chi\chi \rightarrow \phi\phi)/dE_f = \frac{1}{2\beta\gamma} \int_{E_f(1-\beta)}^{E_f(1+\beta)} dE'_f dN_f(\phi \rightarrow ff)/dE'_f,$$

where $\beta = (1 - \gamma^{-2})^{1/2}$, with the boost factor $\gamma = m_{\phi}/m_{\chi}$. Finally, we arrive at the photon flux from the DM annihilation

$$d\Phi(b, \ell)/dE_f = \frac{\langle \sigma v \rangle_{\chi\chi \rightarrow \phi\phi}}{4\pi m_{\phi}^2} \sum_f \text{Br}_{\phi \rightarrow ff} \frac{dN_f(\phi \rightarrow ff)/dE_f}{dE_f} \times \int_{\text{LOS}} dx \rho^2(r(b, \ell, x)).$$

where $r(b, \ell, x) = \sqrt{x^2 + R^2 - 2xR \cos \ell \cos b}$ is the distance from the Galactic center with galactic coordinates $(b, \ell)$, and $\rho(r)$ is the DM density profile, which is commonly taken to be the generalized Navarro-Frenk-White (NFW) shape [40]. Here the $J$ factor is defined as $J = \int_{\text{LOS}} dx \rho^2(r(b, \ell, x))$, where LOS denotes the line of sight integration. $\text{Br}_{\phi \rightarrow ff}$ is the decay branching ratio of the mediator $\phi$ to the final state $f$, which is taken to be similar to the SM Higgs branching ratio.

In Fig. 8, we show the gamma-ray spectrum for a different $m_\chi$ and a different $m_\phi$. The spectrum shown in Fig. 8 has been normalized to corresponding to the $J$ factor $J = 9.09 \times 10^{23}$ GeV$^2$/cm$^5$ with $\gamma = 1.2$ in the NFW profile. We find that the spectrum is very sensitive to $m_\chi$ but not to $m_{h_1}$. The reason is that $m_\chi$ determines the hardness of the spectrum. To obtain the favored parameter space, we define the $\chi^2$ statistic by summing over the bins,

$$\chi^2 = \sum_i \frac{(\Phi^i_{\text{data}} - \Phi^i_{\text{th}}(m_\chi, m_\phi, \langle \sigma v \rangle))^2}{\sigma^2_i},$$

where $\Phi^i_{\text{data}}$ and $\sigma_i$ are the observed flux and the error on the data given in Ref. [4] for the bin $i$, and $\Phi^i_{\text{th}}$ is the theoretical prediction which depends on $(m_\chi, m_\phi, \langle \sigma v \rangle)$. Then we perform a global $\chi^2$ fit. In Fig. 9 (left panel), we show the favored region of the $(m_\chi, m_\phi)$ parameter space at 68% and 95% C.L. For each $(m_\chi, m_{h_1})$, the annihilation cross section is taken to be its best fit value. We also show the direct detection bounds for one of the two benchmark choices: $s_{23} = 1/\sqrt{2}$, $s_{12} = s_{13} = 0.05$. We find that to fit with the GeV gamma-ray spectrum, $m_\phi$ is preferred to be around $60 \sim 85$ GeV, which is still allowed by the tight direct detection bound. The global fit favors the degenerate mass region for $m_{h_1}$ and $m_\phi$ with mass range $60 \sim 85$ GeV. The best fit is on the parameter point $(m_\chi, m_{h_1}) = (72, 70)$ GeV. Similarly, Fig. 9 (right panel) shows the favored region of the $(m_\chi, \langle \sigma v \rangle)$ plane at 68% and 95% C.L. For each $(m_\chi, \langle \sigma v \rangle)$, the $m_{h_1}$ is taken to be its best fit value. The allowed annihilation cross section at
95% C.L. is around $2.4 \sim 3.5$ pb. Therefore, the annihilation $\chi \bar{\chi} \rightarrow h_3h_3$ could simultaneously explain both the GCE and $\Omega_{\chi}$.

The gamma-ray signature at the Galactic center is expected to appear in other galaxies, such as dwarf galaxies. Fermi-LAT experiments investigated the dwarf galaxies but found a null result [18]. This puts bounds on the gamma-ray signatures from the DM annihilation, such as $b\bar{b}$, $\tau\tau$, and other channels. However, the current Fermi-LAT does not put limits on the four-fermion final states with a light mediator. In principle, it is possible to reanalyze the Fermi-LAT data and obtain the limit on the four-fermion final states. We leave this analysis for future study. In the following, we will comment upon the uncertainties on the signature and the current dwarf bounds. As shown in Fig. 2 of Ref. [18], if the Calore et al. data [5] are used, there is still a small parameter region which is allowed by the current dwarf bounds [41]. As was pointed
out in Ref. [41], for a significantly larger integrated \( J \) factor, which corresponds to an extreme high concentration/contraction Milky Way halo model, the signal could escape dwarf galaxy limits.

In this model, the experimental constraints imply that the hidden sector has a nonvanishing coupling to the SM sector. At the same time, the DM could have large coupling to the complex scalar. It might be possible to have self-interacting DM.

**VI. HIDDEN SCALAR SEARCHES AT THE LHC**

Given the favored parameter space to fit the GCE signature, we would like to know whether the LHC data are able to probe this parameter region. Collider searches provide us with a complementary way to explore the GCE favored parameter space. However, in this model, due to the small coupling between the SM sector and the hidden sector, it is difficult to utilize the typical DM search channels, such as the monojet, mono-\( X \) plus missing energy, and other pair production of the DM final states. Fortunately, one could still investigate the Higgs exotic decays at the LHC.

If \( m_{\tilde{h}} \) is lighter than half of the Higgs boson mass, the Higgs boson \( h \) will decay to \( h_1 \rightarrow h_3 h_3 \) as shown in Fig. 10. Similarly, if \( m_{\chi} \) is lighter than \( m_{\tilde{h}}/2 \), the Higgs invisible decay channel \( h_1 \rightarrow \tilde{\chi} \tilde{\chi} \) opens. Both ATLAS and CMS looked for the Higgs invisible decay, but found a null result, which puts additional constraint on the parameter space.

Let us investigate such Higgs exotic decay rates. Assuming \( m_h > 2m_{\chi} \), the decay rate of \( h_1 \rightarrow h_3 h_3 \) can be written as

\[
\Gamma(h_1 \rightarrow h_3 h_3) \approx \sqrt{m_h^2 - 4m_{\tilde{h}}^2} |C|^2 / 32\pi m_{\tilde{h}}^2. \tag{38}
\]

where an extra factor \( 1/2 \) comes from the identical particles in the final state and the effective coupling takes the form

\[
C = |\lambda_{sh} + 2\text{Re}(\lambda_B)| \left( v(2U_{21}U_{23}U_{13} + U_{11}U_{13}^2) + v_s(U_{21}U_{13}^2 + 2U_{11}U_{13}U_{23}) \right) \\
+ |\lambda_{sh} - 2\text{Re}(\lambda_B)| \left( v(U_{11}U_{13}^2 + 2U_{31}U_{33}U_{13}) \right) \\
- \text{Im}(\lambda_B) \left( 4v(U_{21}U_{33}U_{13} + U_{31}U_{23}U_{13} + U_{11}U_{13}U_{23}) + 2v_s(U_{31}U_{13}^2 + 2U_{11}U_{13}U_{33}) \right) \\
+ 6v\lambda_{sh}U_{11}U_{13}^2 + \lambda_s v_s(U_{21}U_{23}^2 + 2U_{21}U_{33}U_{23} + 4U_{31}U_{33}U_{23}) \\
+ 12\text{Re}(\lambda_C)(U_{21}U_{23}^2 - U_{23}U_{33}^2 - 2U_{23}U_{21}U_{23}) \\
+ 12\text{Im}(\lambda_C)(U_{31}U_{33}^2 - U_{31}U_{23}^2 - 2U_{31}U_{21}U_{23}). \tag{39}
\]

In the limit of small \( s_{12} \) and \( s_{13} \), one has

\[
\Gamma(h_1 \rightarrow h_3 h_3) \approx \frac{v^2c_{12}c_{13}^2}{32\pi m_{h_1}^2} \left( \lambda_{sh} + 2\text{Re}(\lambda_B) \right) \cos 2\theta_{23} - 2\text{Im}(\lambda_B) \sin 2\theta_{23}^2. \tag{40}
\]

For the case \( m_{h_1} > 2m_{\chi} \), the Higgs to invisible decay rate can be written as

\[
\Gamma(h_1 \rightarrow \tilde{\chi}\tilde{\chi}) = \frac{(1 - c_{12}^2c_{13}^2)v^2}{8\pi m_{h_1}^2} (m_{h_1}^2 - 2m_{\chi}^2)^{3/2} \tag{41}
\]

The Higgs invisible decay has been studied at both ATLAS and CMS [42]. The current upper limits on the Higgs invisible decay branching ratio is \( \text{Br}(h_1 \rightarrow \text{invisible}) < 0.23 \) at the 95\% C.L. This limit puts constraints on the model parameters when the DM mass is lighter than half of the Higgs boson mass. Because \( h_1 \) is lighter than \( \chi \), if the invisible decay channel opens, the new exotic decay channel \( h_1 \rightarrow h_3 h_3 \) will also exist. Both the invisible and the new exotic channel could be classified as the undetected channel in the Higgs measurements. Assuming an undetected channel, the Higgs coupling measurements put a limit on the branching ratio of the Higgs boson decaying to invisible or undetected states.
BR($h_1 \rightarrow \text{new states}$). In Ref. [42], under the assumption $\kappa_V \leq 1$ on the Higgs boson total width, BR($h_1 \rightarrow \text{new states}$) < 0.49 at the 95% C.L. is obtained. We plot in Fig. 11 contours of the branching ratio of Higgs to invisible decay branching ratio in the $m_\psi$--$m_{h_3}$ plane by assuming $\theta_{12} = \theta_{13} = 0.05$, $\theta_{23} = 0.3$, $\lambda_{sh} = 0.1$, $\lambda_s = 0.01$, and $y_\chi = 0.4$. The red curve is the upper bound on the Higgs to invisible decay branching ratio. The region surrounded by the cyan curve is excluded by the upper bound on the branching ratio of the Higgs to new states. Parameter space outside the two shaded regions satisfies the current upper limit of Higgs to invisible decays.

From the Fig. 11, we learn that the Higgs invisible decay could not explore the favored GCE parameter region, which favors $m_\chi \sim 60$–80 GeV. On the other hand, the Higgs exotic decay branching ratio might be able to probe the GCE parameter region. In the Fig. 11, we show that when $\lambda_s$ is small, the theoretical constraint BR($h \rightarrow \text{new states}$) < 0.49 is not sensitive to the GCE parameter region. In the following paragraph, we will show how the branching ratio depends on the parameter $\lambda_s$.

Except for the theoretical constraints on the branching ratio of $h_1 \rightarrow h_1 h_3$, both ATLAS and CMS also studied the exotic decay channel $h_1 \rightarrow h_3 h_3$ and set limits on the BR($h_1 \rightarrow h_3 h_3$). The ATLAS Collaboration has searched for a Higgs boson decaying into $h_3 h_3$ in the $\tau \tau \mu \mu$ channel with $\sqrt{s} = 8$ TeV, and the observed limit with the expected $\pm 1\sigma$ band was given in Fig. 6 of Ref. [43]. We show in Fig. 12 (left panel) the scatter plots of the signal strength of this model as the function of the pseudoscalar mass, which are far below the current observed limit given by ATLAS. With a higher center of mass energy and increased luminosity at the LHC run 2, we expect that better sensitivity to the parameter space can be obtained. Future LHC sensitivities have been studied in Refs. [44–49]. Reference [49] focuses on the $b\bar{b}\mu\mu$ final states, and makes use of techniques of the $b$-tagging and the jet substructure with mass drop tagger to suppress the irreducible $b\bar{b}\mu\mu$.

FIG. 11. Contours of constant Higgs to invisible branching ratio in the $m_\psi$--$m_{h_1}$ plane, setting $\theta_{12} = \theta_{13} = 0.05$, $\theta_{23} = 0.3$, $\lambda_{sh} = 0.1$, $\lambda_s = 0.01$, and $y_\chi = 0.4$. The red curve is the current upper bound on the Higgs to invisible decay branching ratio. The region surrounded by the cyan curve is excluded by the upper bound on the branching ratio of the Higgs to new states.

FIG. 12. Scatter plots of the signal strength as the function of $m_{h_3}$. (Left panel) The red curve is the observed limit given by the CMS Collaboration [42]. (Right panel) The red solid curve and blue dashed curves are the theoretical projections at the 14 TeV LHC [49].

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FIG. 13. Given the mixing angles $\theta_{13} = \theta_{12} = 0.05$, $\lambda_{sh} = 0.1$, and the light scalar mass $m_{h_3} = 50$ GeV, the contours of the branching ratio $\text{Br}(h_1 \to h_3 h_3)$ in the $\theta_{23}$ and the $\lambda_s$ plane are shown. The red curve and blue dashed curve are the expected limit at the 14 TeV HL-LHC with 300 and 3000 fb$^{-1}$ data [49]. The green short-dashed curve is the theoretical limit from Ref. [42].

$jj(\mu \mu)$, and $tt$ backgrounds. We utilize the projected sensitivities to $\text{Br}(h_1 \to h_3 h_3)$ at the 14 TeV HL-LHC in Ref. [49] and recast their results to our model. Figure 12 (right panel) shows the scattering plots of the signal strength $3\sigma$ to $5\sigma$ in the region of $2.3$ GeV. In Fig. 13, we show the contours of the branching ratio $\text{Br}(h_1 \to h_3 h_3)$ in the $\theta_{23}$ and the $\lambda_s$ plane, given the mixing angles $\theta_{12} = \theta_{13} = 0.05$, $\lambda_{sh} = 0.1$, and the light scalar mass $m_{h_3} = 50$ GeV. In Fig. 13, we note that as the branching ratio $\text{Br}(h_1 \to h_3 h_3)$ becomes smaller, smaller values of $\theta_{23}$ and $\lambda_s$ are needed. Therefore, as we accumulate more data, we could probe smaller values of the parameters $\theta_{23}$ and $\lambda_s$.

We know that to have an $s$-wave DM annihilation cross section, a moderately large $\theta_{23}$ is preferred to obtain the GCE signature. With 300 and 3000 fb$^{-1}$ data we can reach the branching ratio to be as small as 0.52 and 0.10, which corresponds to $\theta_{23}$ in the region of $0.28 \sim 0.53$, which is the interesting parameter region.

VII. CONCLUSIONS

We have investigated a hidden dark matter scenario, in which the hidden sector includes a fermion DM $\chi$ and a complex scalar $S$. The complex scalar mixes with the SM Higgs boson with suppressed coupling. This solves the possible tension between tight constraints from direct detection and LHC searches and a large indirect detection signature. To obtain a large DM annihilation rate, we propagate that there are CP violations in the scalar potential. The CP violations could mix the real and pseudoscalar parts of the complex scalar, and induce a large mass splitting for the mass eigenstates: a light $h_3$ and a much heavier $h_2$. We focus on an interesting parameter region: the light scalar $h_3$ is lighter than the DM $\chi$. This allows the process $\bar{\chi}\chi \to h_3 h_3$ as the dominant DM annihilation channel with an $s$-wave cross section. This annihilation channel gives rise to significant indirect detection signature and could explain the existing Galactic center gamma-ray excess.

The relevant physical parameters are the DM and light scalar masses, the mixing angles among the Higgs boson and the real and imaginary part of the complex scalar boson: $\theta_{12}$, $\theta_{13}$, and $\theta_{23}$, the Yukawa coupling between the DM and the complex scalar $y_\chi$. To obtain the needed $s$-wave enhancement of DM annihilation, we require the mixing angle between the real and imaginary part of the complex scalar boson $\theta_{23}$ to be large. On the other hand, the direct detection bounds imply small mixing angles $\theta_{12}$ and $\theta_{13}$, and large mass splitting between $h_2$ and $h_3$. We found that the EDM constraint is negligible, but the constraints from the DM direct detection, the electroweak precision, and the Higgs coupling measurements are tight. Both the electroweak precision and the Higgs coupling measurements prefer small mixing angles $\theta_{12}$ and $\theta_{13}$. These constraints force us to consider the hidden DM scenario in this model.

To explain the Galactic center excess, the DM annihilates into the four-fermion (mainly four-$b$) final states via the cascade decay $\chi\chi \to h_3 (\to ff) h_3 (\to ff)$. To fit the observed gamma-ray spectrum, the $m_x$ is preferred to be in the 60 to 80 GeV region. And $m_{h_3} = m_x$ is favored. Moreover, the annihilation cross section is fitted to be in the region to have the correct relic density. In short, this hidden DM model explains the gamma-ray excess. Although the dwarf galaxies might place additional constraints, it is still possible to be compatible with the current bounds if the extreme integrated $J$ factor is adopted. Because the scalar $h_3$ cannot be too light, it is unlikely to have self-interacting DM in this hidden scalar scenario.

We also found that constraints from the Higgs invisible decay and Higgs width imply that the $\chi$ and the light scalar $h_3$ cannot be very light. This constraint is consistent with the observed gamma-ray signature. Although the hidden sector has a small coupling to the SM Higgs boson, if $h_3$ is lighter than half of the Higgs mass, the $h_1 \to h_3 h_3$ is a
golden channel to investigate. The 8 TeV LHC results on the exotic decay $h_1 \rightarrow h_2 h_3$ cannot probe the favored parameter region. However, we show that the future 14 TeV studies could be sensitive to the mixing angle $\theta_{23}$, which controls the DM annihilation rate. Thus we expect that with 300 and 3000 fb$^{-1}$ data the exotic decay $h_1 \rightarrow h_2 h_3$ process will be able to probe the parameter region favored by the Galactic center gamma-ray excess.

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