Dark CP violation through the Z portal

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Despite great agreement with experiment, the Standard Model (SM) of particle physics lacks a viable dark matter (DM) candidate and sufficient amount of CP violation to account for the observed baryon excess in the universe. Nonminimal Higgs frameworks are economic extensions of the SM which could remedy these shortcomings. Within the framework of a three Higgs doublet model, we introduce an extended dark sector which accommodates both DM and CP violation. Such dark sources of CP violation do not contribute to the electric dipole moments and are therefore unconstrained. We present a novel mechanism in which the CP -violating dark particles only interact with the SM through the gauge bosons, primarily the Z boson. Such Z -portal dark CP violation is realized in the regions of the parameter space where Higgs-mediated (co)annihilation processes are subdominant and have negligible contributions to the DM relic density. We show that such Z -portal CP -violating DM can still thermalize and satisfy all experimental and observational bounds and discuss the implications of such phenomena for electroweak baryogenesis.

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I. INTRODUCTION

The Standard Model (SM) of particle physics has been extensively tested and is in great agreement with experiment, with its last missing particle, the Higgs boson, discovered at the LHC in 2012 [1,2]. No significant deviation from the SM has been detected at the LHC so far and the properties of the observed scalar are in agreement with those of the SM-Higgs boson [3,4]. However, SM falls short of explaining several aspects of nature, such as the observed baryon asymmetry in the universe and a viable candidate for dark matter (DM). Many astrophysical observations call for a DM particle which is stable on cosmological timescales, cold, nonbaryonic, neutral, and weakly interacting. A particle with such characteristics does not exist in the SM. Moreover, the amount of CP violation provided by the SM is many orders of magnitude smaller than what is needed to generate the observed baryon excess [5–7].

Therefore, it is widely accepted that one needs to consider beyond SM (BSM) scenarios in pursuit of the

ultimate theory of nature. The simplest BSM scenarios which aim to conquer these shortcomings are nonminimal Higgs frameworks, suggesting that the observed scalar at the LHC is just one member of an extended scalar sector. The scalar potential is the least constrained sector in the SM and could provide new sources of CP violation if extended. Also, nonminimal Higgs sectors with discrete symmetries could naturally accommodate weakly interacting massive particles (WIMPs) [8–11], the most well-studied DM candidates. The stability of the WIMP is ensured by the conservation of the discrete symmetry yielding a relic abundance in agreement with the Planck experiment [12] through the freeze-out mechanism.

Extensive studies have been carried out to an advanced level in one singlet or one doublet scalar extensions of the SM, namely the Higgs portal models and two Higgs doublet models (2HDMs) (see e.g., [13,14] and references therein). These models, however, by construction can only partly provide a solution to the SM shortcomings. Frameworks with a further extended scalar sector, such as three Higgs doublet models (3HDMs), contain viable DM candidates, provide new sources of CP violation and strongly first order phase transition as the origin of the baryon asymmetry, contain inflaton candidates driving the inflation process in the beginning of the universe, and provide a solution to the fermion mass hierarchy problem, all in one framework owing to different symmetries and symmetry breaking patterns realizable in the scalar

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potential, with the symmetry breaking patterns determining the number of active [developing a vacuum expectation value (VEV)] and inert (without a VEV) multiplets.

The well-known Z_2 symmetric Higgs portal model [15] and the inert doublet model (IDM) [16] are in agreement with direct and indirect DM searches. However, they are severely constrained by LHC data from invisible Higgs branching ratio and Higgs signal strength bounds. In addition to that, the scalar potential in these models is inevitably CP conserving.

If one were to abandon the idea of DM, the 2HDM scalar potential could, in general, accommodate CP violation. However, such new sources of CP violation modify the SM-Higgs couplings and contribute to the electric dipole moments (EDMs) of the neutron, electron, and certain atomic nuclei [17] and are, therefore, highly constrained by experiment [18–21]. The purely scalar singlet extension of SM, however, is CP conserving regardless of an apparent phase in the VEVs or the parameters of the potential. Since a scalar singlet does not have any gauge interactions, there is a large degree of arbitrariness in the definition of action of charge conjugation on the singlet which renders the model CP conserving [22]. One could summarize this discussion in the following list where S represents a singlet and ϕ a doublet scalar. The superscripts Z_2^+ and Z_2^- stand for even and odd Z_2 charges, respectively:

- (i) $\{\phi_1^{Z_2^+}, \phi_2^{Z_2^+}, S^{Z_2^-}\}$ with vacuum alignment $(v, 0)$: very constrained DM, no CP violation.
- (ii) $\{\phi_1^{Z_2^+}, \phi_2^{Z_2^-}\}$ with vacuum alignment (v_1, v_2) : very constrained CP violation, no DM.

To accommodate CP violation and a viable DM candidate one needs to go beyond the simple doublet or singlet extensions of the SM. If the active sector is extended to accommodate CP violation and the solitary inert sector is to provide the DM candidate, one runs into the same limitations as in the 2HDM, the IDM, and the Higgs portal model [23,24]. Extending the inert sector to accommodate both DM and new sources of CP violation, on the other hand, relieves the model from these constraints.

Such dark/inert sources of CP violation were introduced for the first time in [25] and studied further in [26–29] in a 3HDM framework where the CP mixed dark/inert scalars interact with the SM particles through Higgs and the gauge bosons. It was shown that dark CP violation provides a handle on the otherwise fixed gauge-scalar couplings and opens up a large region of the parameter space to

accommodate a CP -violating DM candidate in agreement with cosmological and collider experiments. Moreover, since the inert sector is protected from directly coupling to the SM particles, the dark CP violation receives no limitation from the EDM experiments. Again, the following list aims to summarize this discussion:

- (i) $\{\phi_1^{Z_2^+}, \phi_2^{Z_2^+}, S^{Z_2^-}\}$ with $(v_1, v_2, 0)$: very constrained DM, very constrained CP violation.
- (ii) $\{\phi_1^{Z_2^+}, \phi_2^{Z_2^-}, \phi_3^{Z_2^-}\}$ with $(v, 0, 0)$: DM and CP violation.

It is important to note that when the extended inert sector contains a doublet and a singlet scalar, as opposed to two doublet scalars, the amount of CP violation is inevitably reduced due to the presence of the singlet [30]. In addition to that, there are fewer coannihilation channels present for the DM candidate. As a result, the model is more constrained and less favorable to the 3HDM counterpart.

Here, we present a novel mechanism in which the CP -violating dark particles only interact with the SM through the gauge bosons, primarily the Z boson. Such Z -portal dark CP violation is realized in the regions of the parameter space where the Higgs portal couplings are of order 10^{-4} or less, and as a result the Higgs-mediated (co) annihilation processes are subdominant and have negligible contributions to the DM relic density. We show that such Z -portal CP violating DM can still thermalize and satisfy all experimental and observational bounds and discuss the implications of such phenomena for electroweak baryogenesis.

The layout of the remainder of the paper is as follows. In Sec. II, we present the details of the scalar potential and the theoretical and experimental limits on its parameters. In Sec. III, we construct and justify our benchmark scenarios. In Sec. IV, we show the effect of dark CP violation on the production and annihilation of DM, and in Sec. V, we conclude and present the outlook for our future studies.

II. THE EXTENDED SCALAR SECTOR

A scalar potential extended by Higgs doublets, which is symmetric under a group G of phase rotations, can be written as the sum of two parts: V_0 with terms symmetric under any phase rotation, and V_G with terms symmetric under G [31,32]. As a result, a Z_2 -symmetric 3HDM can be written as

$$\begin{aligned}
 V_{3\text{HDM}} &= V_0 + V_{Z_2}, \\
 V_0 &= -\mu_1^2(\phi_1^\dagger\phi_1) - \mu_2^2(\phi_2^\dagger\phi_2) - \mu_3^2(\phi_3^\dagger\phi_3) + \lambda_{11}(\phi_1^\dagger\phi_1)^2 + \lambda_{22}(\phi_2^\dagger\phi_2)^2 + \lambda_{33}(\phi_3^\dagger\phi_3)^2 + \lambda_{12}(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) \\
 &\quad + \lambda_{23}(\phi_2^\dagger\phi_2)(\phi_3^\dagger\phi_3) + \lambda_{31}(\phi_3^\dagger\phi_3)(\phi_1^\dagger\phi_1) + \lambda'_{12}(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) + \lambda'_{23}(\phi_2^\dagger\phi_3)(\phi_3^\dagger\phi_2) + \lambda'_{31}(\phi_3^\dagger\phi_1)(\phi_1^\dagger\phi_3), \\
 V_{Z_2} &= -\mu_{12}^2(\phi_1^\dagger\phi_2) + \lambda_1(\phi_1^\dagger\phi_2)^2 + \lambda_2(\phi_2^\dagger\phi_3)^2 + \lambda_3(\phi_3^\dagger\phi_1)^2 + \text{H.c.},
 \end{aligned} \tag{1}$$

where the three Higgs doublets, ϕ_1, ϕ_2, ϕ_3 , transform under the Z_2 group, respectively, as

$$g_{Z_2} = \text{diag}(-1, -1, +1). \quad (2)$$

Note that adding the four Z_2 -respecting terms in V_{Z_2} ensures that our 3HDM potential is symmetric only under this Z_2 group [31] and no larger symmetry. There exist other Z_2 -respecting terms $(\phi_1^\dagger \phi_2)(\phi_1^\dagger \phi_1)$, $(\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_2)$, $(\phi_3^\dagger \phi_1)(\phi_2^\dagger \phi_3)$, and $(\phi_1^\dagger \phi_2)(\phi_3^\dagger \phi_3)$, whose inclusion does not add to the phenomenology of the model.¹ We have, therefore, set the coefficients of these terms to zero for simplicity.

The composition of the doublets is as follows:

$$\phi_1 = \begin{pmatrix} H_1^+ \\ \frac{H_1 + iA_1}{\sqrt{2}} \end{pmatrix}, \quad \phi_2 = \begin{pmatrix} H_2^+ \\ \frac{H_2 + iA_2}{\sqrt{2}} \end{pmatrix}, \quad \phi_3 = \begin{pmatrix} G^+ \\ \frac{v + h + iG^0}{\sqrt{2}} \end{pmatrix}, \quad (3)$$

where ϕ_1 and ϕ_2 are the Z_2 -odd *inert* doublets, $\langle \phi_1 \rangle = \langle \phi_2 \rangle = 0$, and ϕ_3 is the one Z_2 -even *active* doublet, $\langle \phi_3 \rangle = v/\sqrt{2} \neq 0$, which plays the role of the SM Higgs doublet, with h being the SM Higgs boson and $G^\pm, G^0$ the would-be Goldstone bosons. Note that the Z_2 charges assigned to each doublet are according to the Z_2 generator in Eq. (2): odd- Z_2 charge to the inert doublets, ϕ_1 and ϕ_2 , and even- Z_2 charge to the active doublet, ϕ_3 . Therefore, the symmetry of the potential is respected by the vacuum $(0, 0, v/\sqrt{2})$.

The CP -even scalar h contained in the active doublet ϕ_3 has the tree-level couplings of the SM Higgs boson. Thus CP violation is only introduced in the *inert* sector which is forbidden from mixing with the *active* sector by the conservation of the Z_2 symmetry. As a result, the amount of CP violation is not limited by EDMs. This phenomenon of unlimited dark CP violation was first introduced in [25]. The DM candidate, which is the lightest particle among the CP -mixed neutral fields from the inert doublets,² is indeed stable due to the unbroken Z_2 symmetry.

A. Explicit CP violation

The parameters of the phase invariant part, V_0 , are real by construction. We introduce explicit CP violation through complex parameters, $\mu_{12}^2, \lambda_1, \lambda_2, \lambda_3$ in the potential in Eq. (1). Let us emphasize, however, that λ_1 (and other

dark sector parameters $\lambda_{11}, \lambda_{22}, \lambda_{12}, \lambda'_{12}$) only concern inert scalars self-interactions and do not influence tree-level DM and collider phenomenology of the model. These parameters are only constrained through perturbative unitarity and positivity of V bounds and do not play any role in our tree-level analysis. Their value has, therefore, been set to 0.1.

The parameters which are phenomenologically relevant are μ_3^2, λ_{33} which are fixed by the Higgs mass, and $\mu_1^2, \mu_2^2, \mu_{12}^2, \lambda_{31}, \lambda_{23}, \lambda'_{31}, \lambda'_{23}, \lambda_2, \lambda_3$ which appear in inert scalar masses and scalar couplings. The latter nine parameters are, in principle, independent. However, for simplicity, here we study the model in the *dark democracy* limit [25,28,33,34], where

$$\mu_1^2 = \mu_2^2, \quad \lambda_3 = \lambda_2, \quad \lambda_{31} = \lambda_{23}, \quad \lambda'_{31} = \lambda'_{23}. \quad (4)$$

The model is still explicitly CP violating, provided $(\lambda_{22} - \lambda_{11})[\lambda_1(\mu_{12}^2)^2 - \lambda_1^*(\mu_{12}^2)^2] \neq 0$ [35,36] after imposing the dark democracy limit.³ Note that one can redefine the doublets as

$$\begin{cases} \phi_1 \rightarrow \phi_1 e^{i\theta_{12}/2} \\ \phi_2 \rightarrow \phi_2 e^{-i\theta_{12}/2} \\ \phi_3 \rightarrow \phi_3 \end{cases} \Rightarrow \begin{cases} |\mu_{12}^2| e^{i\theta_{12}} \rightarrow |\mu_{12}^2| \\ |\lambda_2| e^{i\theta_2} \rightarrow |\lambda_2| e^{i(\theta_2 + \theta_{12})} \end{cases}. \quad (5)$$

As a result, the only relevant CP -violating parameter⁴ in the dark democracy limit is $\theta_2 + \theta_{12}$, the “shifted” phase of the λ_2 coupling, which is referred to as θ_{CPV} throughout the paper.

B. The mass spectrum

The point $(0, 0, \frac{v}{\sqrt{2}})$ is the minimum of the potential when $v^2 = \frac{\mu_3^2}{\lambda_{33}}$. The only active doublet, ϕ_3 , plays the role of the SM Higgs doublet, with $G^0, G^\pm$ as the massless Goldstone bosons, and h as the SM-like Higgs with $m_h^2 = 2\mu_3^2 = 2\lambda_{33}v^2 = (125 \text{ GeV})^2$.

1. The charged inert states

The charged mass-squared matrix in the $(H_1^\pm, H_2^\pm)$ basis is calculated to be

$$\mathcal{M}_C^2 = \begin{pmatrix} -\mu_2^2 + \frac{1}{2}\lambda_{23}v^2 & -\mu_{12}^2 \\ -\mu_{12}^2 & -\mu_2^2 + \frac{1}{2}\lambda_{23}v^2 \end{pmatrix}, \quad (6)$$

¹The terms $(\phi_1^\dagger \phi_2)(\phi_1^\dagger \phi_1)$, $(\phi_1^\dagger \phi_2)(\phi_2^\dagger \phi_2)$ only appear in inert scalar self-interactions and do not affect our tree-level analysis. Including the terms $(\phi_3^\dagger \phi_1)(\phi_2^\dagger \phi_3)$, $(\phi_1^\dagger \phi_2)(\phi_3^\dagger \phi_3)$ resembles itself as a shift in the value of the $\lambda_{2,3}$ parameters. Our numerical analysis confirms that removing these terms does not lose the generality of the model.

²We avoid regions of parameter space where one of the charged inert scalars is the lightest inert particle.

³For a 3HDM to be CP -violating at least one CP -odd invariant has to be nonzero. Therefore, the nonvanishing CP -odd invariant condition provided here is a sufficient condition but not necessary since any other nonzero CP -odd invariant ensures CP violation in the model.

⁴Recall that, as a dark parameter, λ_1 and its phase are not relevant for our tree-level analysis.

with the two physical charged states, $S_{1,2}^\pm$, as eigenstates

$$S_1^\pm = \frac{H_1^\pm + H_2^\pm}{\sqrt{2}}, \quad S_2^\pm = \frac{H_1^\pm - H_2^\pm}{\sqrt{2}}. \quad (7)$$

The masses of the physical charged scalars, eigenvalues of the $\mathcal{M}_C^2$ matrix, are calculated to be

$$m_{S_1^\pm}^2 = -\mu_2^2 - \mu_{12}^2 + \frac{1}{2}\lambda_{23}v^2, \quad m_{S_2^\pm}^2 = -\mu_2^2 + \mu_{12}^2 + \frac{1}{2}\lambda_{23}v^2, \quad (8)$$

where we take $\mu_{12}^2 > 0$ and, therefore, $m_{S_1^\pm} < m_{S_2^\pm}$.

2. The neutral inert states

The neutral mass-squared matrix in the (H_1, H_2, A_1, A_2) basis is calculated to be

$$\mathcal{M}_N^2 = \frac{1}{4} \begin{pmatrix} \Lambda_c^+ & -2\mu_{12}^2 & -\Lambda_s & 0 \\ -2\mu_{12}^2 & \Lambda_c^+ & 0 & \Lambda_s \\ -\Lambda_s & 0 & \Lambda_c^- & -2\mu_{12}^2 \\ 0 & \Lambda_s & -2\mu_{12}^2 & \Lambda_c^- \end{pmatrix}, \quad (9)$$

where

$$\Lambda_s = 2\lambda_2 \sin \theta_{\text{CPV}} v^2 \quad \text{and} \quad \Lambda_c^\pm = -2\mu_2^2 + (\lambda_{23} + \lambda'_{23} \pm 2\lambda_2 \cos \theta_{\text{CPV}}) v^2, \quad (10)$$

with the four CP -mixed neutral scalars, $S_{1,2,3,4}$, as eigenstates

$$S_1 = \frac{\alpha H_1 - A_1 + \alpha H_2 + A_2}{\sqrt{2}\sqrt{\alpha^2 + 1}}, \quad S_2 = \frac{H_1 + \alpha A_1 + H_2 - \alpha A_2}{\sqrt{2}\sqrt{\alpha^2 + 1}}, \\ S_3 = \frac{\beta H_1 + A_1 - \beta H_2 + A_2}{\sqrt{2}\sqrt{\beta^2 + 1}}, \quad S_4 = \frac{-H_1 + \beta A_1 + H_2 + \beta A_2}{\sqrt{2}\sqrt{\beta^2 + 1}}, \quad (11)$$

where α and β are defined as

$$\alpha = \frac{-\mu_{12}^2 + v^2|\lambda_2| \cos \theta_{\text{CPV}} - \Lambda^-}{v^2|\lambda_2| \sin \theta_{\text{CPV}}}, \quad \beta = \frac{-\mu_{12}^2 - v^2|\lambda_2| \cos \theta_{\text{CPV}} + \Lambda^+}{v^2|\lambda_2| \sin \theta_{\text{CPV}}}, \quad (12)$$

and $\Lambda^\mp$ as

$$\Lambda^\mp = \sqrt{(\mu_{12}^2)^2 + v^4|\lambda_2|^2 \mp 2v^2\mu_{12}^2|\lambda_2| \cos \theta_{\text{CPV}}}. \quad (13)$$

The masses of the physical neutral CP -mixed inert scalars, eigenvalues of the $\mathcal{M}_N^2$ matrix, are calculated to be

$$m_{S_{1,2}}^2 = -\mu_2^2 + \frac{v^2}{2}(\lambda'_{23} + \lambda_{23}) \mp \Lambda^-, \quad m_{S_{3,4}}^2 = -\mu_2^2 + \frac{v^2}{2}(\lambda'_{23} + \lambda_{23}) \mp \Lambda^+. \quad (14)$$

We take S_1 to be lightest inert particle and the DM candidate. Throughout the paper, the notations S_1 and DM will be used interchangeably.

3. The span of θ_{CPV}

To reproduce the results of [33,34] with $\lambda_2 < 0$, we require⁵

$$\frac{\pi}{2} < \theta_{\text{CPV}} < \frac{3\pi}{2} \Rightarrow m_{S_1} < m_{S_2}, m_{S_3}, m_{S_4}. \quad (15)$$

Consequently, in the first and fourth quadrants S_3 becomes the lightest neutral inert particle,

$$\left. \begin{array}{l} 0 < \theta_{\text{CPV}} < \frac{\pi}{2} \\ \frac{3\pi}{2} < \theta_{\text{CPV}} < 2\pi \end{array} \right\} \Rightarrow m_{S_3} < m_{S_1}, m_{S_2}, m_{S_4}. \quad (16)$$

At $\theta_{\text{CPV}} = \frac{\pi}{2}, \frac{3\pi}{2}$ where $\Lambda^+ = \Lambda^-$, a mass degeneracy between neutral inert particles occurs,

$$\theta_{\text{CPV}} = \frac{\pi}{2}, \frac{3\pi}{2} \Rightarrow \begin{cases} m_{S_1} = m_{S_3}, \\ m_{S_2} = m_{S_4}. \end{cases} \quad (17)$$

The model is reduced to the CP -conserving limit when $\theta_{\text{CPV}} = 0, \pi$ which renders $S_{1,3}$ to CP -even and $S_{2,4}$ to CP -odd particles,

$$\theta_{\text{CPV}} = 0, \pi \Rightarrow \text{CP-conserving limit: } \begin{cases} S_{1,3} = \frac{H_1 \pm H_2}{\sqrt{2}}, \\ S_{2,4} = \frac{A_1 \pm A_2}{\sqrt{2}}. \end{cases} \quad (18)$$

We take all other parameters of the potential to be positive.

4. The input parameters of the model

As independent input parameters of the model, we take

$$m_{S_1}, \quad m_{S_2}, \quad m_{S_1^\pm}, \quad m_{S_2^\pm}, \quad \theta_{\text{CPV}}, \quad g_{h\text{DM}}, \quad (19)$$

where $g_{h\text{DM}} \equiv g_{S_1 S_1 h}$ is the Higgs-DM coupling, with the relevant terms in the Lagrangian appearing as

$$\mathcal{L} \supset g_{ZS_i S_j} Z_\mu (S_i \partial^\mu S_j - S_j \partial^\mu S_i) + \frac{v}{2} g_{S_i S_i h} h S_i^2 \\ + v g_{S_i S_j h} h S_i S_j + v g_{S_i^\pm S_j^\mp h} h S_i^\pm S_j^\mp. \quad (20)$$

⁵Note that for $\lambda_2 > 0$, the above arguments hold provided the neutral inert particles are relabeled as $S_1 \leftrightarrow S_3$ and $S_2 \leftrightarrow S_4$.

The parameters of the model can be written in terms of the physical observables in Eq. (19):

$$\begin{aligned}
\mu_{12}^2 &= \frac{1}{2}(m_{S_2^\pm}^2 - m_{S_1^\pm}^2), \\
\lambda_{23} &= \frac{1}{v^2}(2\mu_2^2 + m_{S_2^\pm}^2 + m_{S_1^\pm}^2), \\
\lambda'_{23} &= \frac{1}{v^2}(m_{S_2}^2 + m_{S_1}^2 - m_{S_2^\pm}^2 - m_{S_1^\pm}^2), \\
|\lambda_2| &= \frac{1}{v^2} \left[\mu_{12}^2 \cos \theta_{\text{CPV}} + \frac{1}{4} \sqrt{(2\mu_{12}^2 \cos \theta_{\text{CPV}})^2 + (m_{S_2}^2 - m_{S_1}^2)^2 - (m_{S_2^\pm}^2 - m_{S_1^\pm}^2)^2} \right], \\
\mu_2^2 &= \frac{v^2}{2} g_{h\text{DM}} - \frac{v^2}{\alpha^2 + 1} (2\alpha \sin \theta_{\text{CPV}} + (\alpha^2 - 1) \cos \theta_{\text{CPV}}) |\lambda_2| - \frac{1}{2}(m_{S_2}^2 + m_{S_1}^2).
\end{aligned} \tag{21}$$

Using these relations, all other relevant masses and couplings can be expressed in terms of the independent input parameters of the model, for example, the masses of other inert scalars:

$$\begin{aligned}
m_{S_{3,4}}^2 &= m_{S_1}^2 + \Lambda^\mp \mp \Lambda^+, \quad \text{where} \\
(\Lambda^+)^2 &= (\Lambda^-)^2 + 4v^2 \mu_{12}^2 |\lambda_2| \cos \theta_{\text{CPV}}.
\end{aligned} \tag{22}$$

The neutral scalar-gauge couplings are derived to be

$$\begin{aligned}
|g_{ZS_1S_3}| &= |g_{ZS_2S_4}| = g_Z^{\text{CPC}} \left(\frac{\alpha + \beta}{\sqrt{\alpha^2 + 1} \sqrt{\beta^2 + 1}} \right), \\
|g_{ZS_1S_4}| &= |g_{ZS_2S_3}| = g_Z^{\text{CPC}} \left(\frac{\alpha\beta - 1}{\sqrt{\alpha^2 + 1} \sqrt{\beta^2 + 1}} \right),
\end{aligned} \tag{23}$$

where $g_Z^{\text{CPC}} = \frac{e}{2c_{\theta_W} s_{\theta_W}}$ is the $|g_{ZS_iS_j}|$ value in the CP -conserving limit, with e the elementary charge and c_{θ_W} , s_{θ_W} the sine and cosine of the weak mixing angle. Note that

$$g_{ZS_1S_3}^2 + g_{ZS_1S_4}^2 = (g_Z^{\text{CPC}})^2, \quad g_{ZS_2S_3}^2 + g_{ZS_2S_4}^2 = (g_Z^{\text{CPC}})^2, \tag{24}$$

where $g_{ZS_1S_2} = g_{ZS_3S_4} = 0$ in the dark democracy limit. Similarly, the charged scalar-gauge couplings are derived to be

$$\begin{aligned}
|g_{W^\pm S_1^\mp S_1}| &= |g_{W^\pm S_2^\mp S_2}| = g_W^{\text{CPC}} \left(\frac{\alpha}{\sqrt{\alpha^2 + 1}} \right), \\
|g_{W^\pm S_1^\mp S_2}| &= |g_{W^\pm S_2^\mp S_1}| = g_W^{\text{CPC}} \left(\frac{1}{\sqrt{\alpha^2 + 1}} \right), \\
|g_{W^\pm S_1^\mp S_3}| &= |g_{W^\pm S_2^\mp S_4}| = g_W^{\text{CPC}} \left(\frac{1}{\sqrt{\beta^2 + 1}} \right), \\
|g_{W^\pm S_1^\mp S_4}| &= |g_{W^\pm S_2^\mp S_3}| = g_W^{\text{CPC}} \left(\frac{\beta}{\sqrt{\beta^2 + 1}} \right),
\end{aligned}$$

where $g_W^{\text{CPC}} = \frac{e}{s_{\theta_W}}$ is the $|g_{W^\pm S_i^\mp S_j}|$ value in the CP -conserving limit. Note that, unlike the CP -conserving

limit, the strength of gauge-scalar interactions depend on the parameters α and β in Eq. (12), which in turn depend on m_{S_i} in the presence of CP violation.

C. Constraints on the parameter space

The parameter space of the model is constrained by theoretical, observational, and experimental bounds which are satisfied in all our benchmark scenarios to follow:

- (1) To satisfy theoretical constraints, we require the potential to be bounded from below and for the Hessian to be positive definite [25], using the conservative sufficient limits of

$$\begin{aligned}
\lambda_{ii} &> 0, \quad \lambda_{ij} + \lambda'_{ij} > -2\sqrt{\lambda_{ii}\lambda_{jj}}, \\
|\lambda_{1,2,3}| &< |\lambda_{ii}|, |\lambda_{ij}|, |\lambda'_{ij}|, \quad i \neq j = 1, 2, 3.
\end{aligned} \tag{25}$$

We take all couplings to be $|\lambda_i| \leq 4\pi$ in accordance with perturbative unitarity limits.

- (2) Parametrized by the electroweak (EW) oblique parameters S , T , U [37–40], inert particles S_i , $S_i^\pm$ may introduce important radiative corrections to gauge boson propagators. We impose a 2σ agreement with EW precision observables (EWPOs) at 95% confidence level (C.L.) [41],

$$S = 0.05 \pm 0.11, \quad T = 0.09 \pm 0.13, \quad U = 0.01 \pm 0.11. \tag{26}$$

Similar to the 2HDM, this condition requires each charged state to be close in mass with a neutral state, in the dark sector [42].

- (3) The contribution of the inert scalars to the total decay width of the EW gauge bosons constrains the masses of the inert scalars to be [43]

$$\begin{aligned}
m_{S_i} + m_{S_{1,2}^\pm} &\geq m_{W^\pm}, \quad m_{S_i} + m_{S_j} \geq m_Z, \\
2m_{S_{1,2}^\pm} &\geq m_Z, \quad i, j = 1, 2, 3, 4.
\end{aligned} \tag{27}$$

- (4) Nonobservation of charged scalars puts a model-independent lower bound on their mass [44–46] and an upper bound on their lifetime [47] to be

$$\begin{aligned} m_{S_{1,2}^\pm} &\geq 70 \text{ GeV}, \\ \tau_{S_{1,2}^\pm} &\leq 10^{-7} \text{ s} \quad \Rightarrow \quad \Gamma_{S_{1,2}^\pm}^{\text{tot}} \geq 6.58 \times 10^{-18} \text{ GeV}, \end{aligned} \quad (28)$$

to guarantee their decay within the detector. In all our benchmark scenarios, the mass of both charged scalars is above 95 GeV, and their decay width, primarily to $S_i^\pm \rightarrow S_j W^\pm$, is of the order of 10^{-1} GeV, which is well within limits.

- (5) Any model introducing new decay channels for the SM-Higgs boson is constrained by an upper limit on the Higgs total decay width, $\Gamma_{\text{tot}}^h \leq 9 \text{ MeV}$ [48], and Higgs signal strengths [49–51]. In our model, the SM-like Higgs could decay to a pair of inert scalars, provided $m_{S_i} + m_{S_j} < m_h$ and $S_{i,j}$ are long-lived enough ($\tau \geq 10^{-7} \text{ s}$). As a result, $S_{i,j}$ will not decay inside the detector and therefore contribute to the Higgs *invisible* decay, $h \rightarrow S_i S_j$, with a branching ratio of

$$\text{BR}(h \rightarrow S_i S_j) = \frac{\sum_{i,j} \Gamma(h \rightarrow S_i S_j)}{\Gamma_h^{\text{SM}} + \sum_{i,j} \Gamma(h \rightarrow S_i S_j)}, \quad (29)$$

with

$$\begin{aligned} \Gamma(h \rightarrow S_i S_j) &= \frac{g_{hS_i S_j}^2 v^2}{32\pi m_h^3} ((m_h^2 - (m_{S_i} + m_{S_j})^2) \\ &\quad \times (m_h^2 - (m_{S_i} - m_{S_j})^2))^{1/2}, \end{aligned} \quad (30)$$

which sets strong limits on the Higgs-inert couplings. Moreover, the partial decay $\Gamma(h \rightarrow \gamma\gamma)$ receives contributions from the inert charged scalars. The combined ATLAS and CMS Run I results for Higgs to $\gamma\gamma$ signal strength require $\mu_{\gamma\gamma} = 1.14_{-0.36}^{+0.38}$ [49]. In Run II, ATLAS reports $\mu_{\gamma\gamma} = 0.99_{-0.14}^{+0.14}$ [50], and CMS reports $\mu_{\gamma\gamma} = 1.18_{-0.14}^{+0.17}$ [51] with both of which we are in 2σ agreement.

- (6) Reinterpretation of LEP 2 and LHC Run I searches for supersymmetric (SUSY) particles (mainly sneutrinos and sleptons) for the IDM [44,52] excludes the region of parameter space where the following conditions are simultaneously satisfied ($i, j = 2, \dots, 4$):

$$m_H \leq 80 \text{ GeV}, \quad m_A \leq 100 \text{ GeV}, \quad \Delta m(H, A) \geq 8 \text{ GeV}. \quad (31)$$

We take these limits into account for our DM candidate paired with any other neutral scalar. We also check the validity of our benchmark scenarios against LHC searches for new particles in accordance with the analysis for the IDM [53].

- (7) DM relic density measurements from the Planck experiment [12],

$$\Omega_{\text{DM}} h^2 = 0.1197 \pm 0.0022, \quad (32)$$

require the relic abundance of the DM candidate to lie within these bounds if it constitutes 100% of DM in the universe.

A DM candidate with $\Omega_{\text{DM}} h^2$ smaller than the observed value is allowed; however, an additional DM candidate is needed to complement the missing relic density. Regions of the parameter space corresponding to values of $\Omega_{\text{DM}} h^2$ larger than the Planck upper limit are excluded.

We impose a 3σ agreement with the observation on the relic abundance of our DM candidate, S_1 .

- (8) The latest XENON1T results for DM direct detection experiments [54] and FermiLAT results for indirect detection searches [55] do not constrain the model any further. Having set the Higgs portal couplings to zero in our benchmark scenarios, the largest direct detection cross section is $\sigma_{\text{DM-N}} \approx 10^{-14} \text{ pb}$ and the largest indirect detection cross section is $\langle v\sigma \rangle \approx 10^{-32} \text{ cm}^3/\text{s}$, both of which are well below the limits [56].

III. DM ABUNDANCE AND THE SELECTION OF BENCHMARKS

The relic abundance of the DM candidate, S_1 , after freeze-out is given by the solution of the Boltzmann equation,

$$\frac{dn_{S_1}}{dt} = -3Hn_{S_1} - \langle \sigma_{\text{eff}} v \rangle [(n_{S_1})^2 - (n_{S_1}^{\text{eq}})^2], \quad (33)$$

where n_{S_1} is the number density of the S_1 particle, H is the Hubble parameter, and $n_{S_1}^{\text{eq}}$ is the number density of S_1 at equilibrium. The thermally averaged effective (co)annihilation cross section, $\langle \sigma_{\text{eff}} v \rangle$, receives contributions from all relevant annihilation processes of any $S_i S_j$ pair into SM particles, so that

$$\begin{aligned} \langle \sigma_{\text{eff}} v \rangle &= \sum_{i,j} \langle \sigma_{ij} v_{ij} \rangle \frac{n_{S_i}^{\text{eq}} n_{S_j}^{\text{eq}}}{n_{S_1}^{\text{eq}} n_{S_1}^{\text{eq}}}, \quad \text{where} \\ \frac{n_{S_i}^{\text{eq}}}{n_{S_1}^{\text{eq}}} &\sim \exp\left(-\frac{m_{S_i} - m_{S_1}}{T}\right). \end{aligned} \quad (34)$$

However, only processes with the $S_i - S_1$ mass splitting comparable to the thermal bath temperature T provide a sizable contribution.

A common feature of nonminimal Higgs DM models is that in a large region of the parameter space the most important process for DM annihilation is through the $S_1 S_1 \rightarrow h_{\text{SM}} \rightarrow f\bar{f}$ channel whose efficiency depends on both the DM mass and the Higgs-DM coupling. In the region where $m_{\text{DM}} < m_h/2$, generally one requires a large Higgs-DM coupling in order to produce relic density in agreement with Eq. (32). However, such large Higgs-DM coupling leads to large direct detection and indirect detection cross sections and significant deviations from SM-Higgs coupling measurements, which are ruled out by experimental and observational data. On the other hand, a small Higgs-DM coupling fails to annihilate the DM candidate effectively and leads to the overclosure of the universe. This is where coannihilation processes play an important role as they can contribute to changes in the DM relic density.

In models with extended dark sectors, in addition to the standard Higgs mediated annihilation channels of DM, there exists the possibility of coannihilation with heavier states, provided they are close in mass [25,27,28,33,34]. The relevance of this effect depends not only on the DM mass and the mass splittings but also on the strength of the standard DM annihilation channel.

It is worth pointing out that in the IDM, where by construction CP violation is not allowed, the only coannihilation process is through the Z -mediated $HA \rightarrow Z \rightarrow f\bar{f}$ channel whose subdominant effect fails to revive the model in the low mass region. Extending the inert sector, as shown in [27,33,34] in the CP -conserving limit, opens up several coannihilation channels, both Higgs-mediated $H_1 H_2 \rightarrow h \rightarrow f\bar{f}$ and Z -mediated $H_1 A_{1,2} \rightarrow Z \rightarrow f\bar{f}$. However, their collective contribution to DM coannihilation is not sufficient and one still needs a nonzero Higgs-DM coupling to satisfy relic density bounds. Introducing CP violation in the extended dark sector [25,28,57] opens up many coannihilation channels through the Higgs and Z bosons, $S_i S_j \rightarrow h/Z \rightarrow f\bar{f}$, which can significantly affect the DM phenomenology. In fact, the Z -mediated coannihilations can be strong enough to relieve the model of the need for any Higgs-mediated (co)annihilation processes.

To show the effect of Z portal CP violation on the abundance of DM, we set the Higgs-DM coupling to zero, $g_{h\text{DM}} = 0$, thereby removing the main DM annihilation process, $S_1 S_1 \rightarrow h \rightarrow f\bar{f}$. All other $S_i S_j h$ vertex coefficients are also reduced to a point where their resulting coannihilation processes have negligible contributions to the DM relic density. So, the only communication between the dark sector and the visible sector is through the gauge bosons $W^\pm$ and Z .

It is important to note that the phenomenon of dark CP violation is not realizable in purely scalar singlet extensions of the SM. An extended dark sector with a doublet plus a singlet could accommodate dark CP violation; however, the presence of the singlet dilutes the CP -violating effects,

since it has no couplings to SM gauge bosons. We would like to point out that this is the reason Ref. [30] does not find points leading to 100% of DM relic density in their low mass region. Furthermore, the effects of dark CP violation, which is through the ZZZ observable [58], in their model is considerably smaller [28,29].

In the region of the parameter space where Higgs portal interactions are negligible ($g_{h\text{DM}} \approx 10^{-4}$), the total DM annihilation cross section receives contributions from the following:

(i) *DM annihilation processes:*

$$\begin{aligned} S_1 S_1 &\rightarrow VV, & S_1 S_1 &\rightarrow VV^* \rightarrow Vff', \\ S_1 S_1 &\rightarrow V^* V^* \rightarrow ff'ff', \end{aligned} \quad (35)$$

where V is any of the SM gauge bosons. In the $m_{\text{DM}} < m_W$ region, the processes with off-shell gauge bosons dominate over the ones with on-shell gauge bosons.

(ii) *DM coannihilation processes:*

$$S_1 S_{2,3,4} \rightarrow Z^* \rightarrow f\bar{f}, \quad S_1 S_{1,2}^\pm \rightarrow W^{\pm*} \rightarrow ff', \quad (36)$$

where the coannihilating dark scalars are up to 20% heavier than the DM candidate.

(iii) *(Co)annihilation of other dark states:*

$$\begin{aligned} S_i S_i &\rightarrow VV, & S_i S_i &\rightarrow VV^* \rightarrow Vff', \\ S_i S_i &\rightarrow V^* V^* \rightarrow ff'ff', & S_i S_j &\rightarrow V^* \rightarrow ff', \end{aligned} \quad (37)$$

where $S_i \neq S_j$ are any of the dark scalars $S_{2,3,4}$, $S_{1,2}^\pm$ which are all close in mass.

Taking all such processes into account, we define the following benchmark scenarios with distinct DM phenomenology. It is convenient to introduce the mass splittings between the DM candidate and other inert scalars as

$$\delta_{12} = m_{S_2} - m_{S_1}, \quad \delta_c = m_{S_2^\pm} - m_{S_1^\pm}, \quad \delta_{1c} = m_{S_1^\pm} - m_{S_1}. \quad (38)$$

A. Benchmarks of type 1

In agreement with reinterpreted SUSY searches in Eq. (31), we devise two benchmark scenarios in the low mass region, $45 \text{ GeV} < m_{S_1} \leq 80 \text{ GeV}$,

$$\mathbf{B}_1 \mathbf{D}_4 \mathbf{C}_1 : \delta_{12} = 4 \text{ GeV}, \quad \delta_c = 1 \text{ GeV}, \quad \delta_{1c} = 50 \text{ GeV},$$

$$\mathbf{B}_1 \mathbf{D}_8 \mathbf{C}_1 : \delta_{12} = 8 \text{ GeV}, \quad \delta_c = 1 \text{ GeV}, \quad \delta_{1c} = 50 \text{ GeV},$$

where all neutral inert particles are close in mass and are much lighter than the inert charged particles,

$$m_{S_1} \sim m_{S_3} \sim m_{S_2} \sim m_{S_4} \ll m_{S_1^\pm} \sim m_{S_2^\pm}. \quad (39)$$

If one were to ignore these reinterpreted SUSY bounds, one can also construct benchmark scenarios with larger mass splittings,

$$\mathbf{B}_1\mathbf{D}_{12}\mathbf{C}_1: \delta_{12}=12 \text{ GeV}, \quad \delta_c=1 \text{ GeV}, \quad \delta_{1c}=50 \text{ GeV},$$

$$\mathbf{B}_1\mathbf{D}_{20}\mathbf{C}_1: \delta_{12}=20 \text{ GeV}, \quad \delta_c=1 \text{ GeV}, \quad \delta_{1c}=50 \text{ GeV}.$$

With a larger δ_{12} , the neutral inert particles split into two groups, with S_1 and S_3 close in mass and lighter than S_2 and S_4 which are also close in mass, and all lighter than the charged inert scalars,

$$m_{S_1} \sim m_{S_3} \lesssim m_{S_2} \sim m_{S_4} \ll m_{S_1^\pm} \sim m_{S_2^\pm}. \quad (40)$$

B. Benchmarks of type 2

In the low mass region, $45 \text{ GeV} < m_{S_1} \leq 80 \text{ GeV}$, in agreement with reinterpreted SUSY searches in Eq. (31), we define

$$\mathbf{B}_2\mathbf{D}_{55}\mathbf{C}_1: \delta_{12}=55 \text{ GeV}, \quad \delta_c=1 \text{ GeV}, \quad \delta_{1c}=50 \text{ GeV},$$

$$\mathbf{B}_2\mathbf{D}_{55}\mathbf{C}_{15}: \delta_{12}=55 \text{ GeV}, \quad \delta_c=15 \text{ GeV}, \quad \delta_{1c}=50 \text{ GeV},$$

where only one neutral inert particle is close in mass with the DM candidate,

$$m_{S_1} \sim m_{S_3} \ll m_{S_2} \sim m_{S_4} \sim m_{S_1^\pm} \sim m_{S_2^\pm}. \quad (41)$$

C. Benchmarks of type 3

In the heavy mass region $m_{S_1} \geq 80 \text{ GeV}$, the reinterpreted SUSY bounds in Eq. (31) do not apply any more, so any δ_{12} mass splitting is allowed. Moreover, the charged particles are allowed to be close in mass with the DM. We define three benchmark scenarios:

$$\mathbf{B}_3\mathbf{D}_5\mathbf{C}_1: \delta_{12}=5 \text{ GeV}, \quad \delta_c=1 \text{ GeV}, \quad \delta_{1c}=1 \text{ GeV},$$

$$\text{where } m_{S_1} \sim m_{S_3} \sim m_{S_2} \sim m_{S_4} \sim m_{S_1^\pm} \sim m_{S_2^\pm}, \quad (42)$$

$$\mathbf{B}_3\mathbf{D}_{55}\mathbf{C}_1: \delta_{12}=55 \text{ GeV}, \quad \delta_c=1 \text{ GeV}, \quad \delta_{1c}=1 \text{ GeV},$$

$$\text{where } m_{S_1} \sim m_{S_3} \sim m_{S_1^\pm} \sim m_{S_2^\pm} \ll m_{S_2} \sim m_{S_4}, \quad (43)$$

FIG. 1. The relevant and non-negligible ZS_iS_j couplings in type 1 scenarios for an exemplary m_{S_1} of 57 GeV.

FIG. 2. The relic abundance of S_1 for different DM masses in type 1 benchmark scenarios. The horizontal red band shows the Planck observation limit on the abundance of DM.

$$\mathbf{B}_3\mathbf{D}_{55}\mathbf{C}_{22}: \delta_{12} = 55 \text{ GeV}, \quad \delta_c = 22 \text{ GeV}, \quad \delta_{1c} = 1 \text{ GeV},$$

$$\text{where } m_{S_1} \sim m_{S_3} \sim m_{S_1^\pm} \ll m_{S_2^\pm} \sim m_{S_2} \sim m_{S_4}. \quad (44)$$

IV. THE EFFECT OF DARK CP VIOLATION ON THE ABUNDANCE

A. Benchmarks of type 1

In benchmark scenarios of type 1, all neutral inert particles are relatively close in mass and are much lighter than the inert charged particles. Here the main coannihilation channel is through the $S_i S_j \rightarrow Z^* \rightarrow f \bar{f}$ processes. With very small $S_1 - S_2$ mass splitting, in $\mathbf{B}_1\mathbf{D}_4\mathbf{C}_1$ and $\mathbf{B}_1\mathbf{D}_8\mathbf{C}_1$ scenarios, the coannihilation of DM with other neutral scalars is so strong that DM is underproduced irrespective of the CP -violating angle. A larger $S_1 - S_2$ mass splitting, as in $\mathbf{B}_1\mathbf{D}_{12}\mathbf{C}_1$ and $\mathbf{B}_1\mathbf{D}_{20}\mathbf{C}_1$ scenarios, weakens these coannihilation processes and leads to a larger relic abundance of S_1 . The efficiency of the coannihilation process also depends on the strength of the $ZS_i S_j$ coupling. Figure 1 shows the strength of the relevant and non-negligible $ZS_i S_j$ couplings in all four scenarios for an exemplary m_{S_1} of 57 GeV. As expected the $g_{ZS_1 S_3}$ coupling vanishes at $\theta_{CPV} = \pi$, i.e., in the CP -conserving limit which is

where S_1 and S_3 are reduced to two CP -even particles as discussed in Eq. (18).

Note that as δ_{12} increases from $\mathbf{B}_1\mathbf{D}_4\mathbf{C}_1$ to $\mathbf{B}_1\mathbf{D}_{20}\mathbf{C}_1$ scenarios which reduces the coannihilation probability of S_1 with other neutral dark particles, also the coupling of the main coannihilation channel, $g_{ZS_1 S_3}$ is reduced. As a result, the DM abundance is considerably larger in the latter scenarios. Figure 2 shows the abundance of S_1 for different DM masses in all type 1 scenarios.

It is worth noting that in a given benchmark scenario as DM mass increases to values comparable with $m_{W^\pm}$ and m_Z , the $S_i S_j \rightarrow VV$ annihilation channels open up and reduce the DM number density. As a result, DM is always underproduced for $m_{S_1} \gtrsim 80$ GeV. Figure 3 shows the regions with correct abundance and underabundance of DM in $\mathbf{B}_1\mathbf{D}_{12}\mathbf{C}_1$ and $\mathbf{B}_1\mathbf{D}_{20}\mathbf{C}_1$ scenarios. At $\theta_{CPV} = \pi$, where the model is CP -conserving, the strength of the $S_i S_j Z$ gauge couplings is fixed. As a result, the intermediate mass region $54 \text{ GeV} \lesssim m_{S_1} \lesssim 70 \text{ GeV}$ is ruled out due to the overproduction of DM. Varying the CP -violating phase, θ_{CPV} , changes the strength of the $S_i S_j Z$ couplings and the $S_1 - S_3$ mass splitting, with the smallest mass splitting close to the $\pi/2$ and $3\pi/2$ boundaries. Therefore, one expects a more effective coannihilation of S_1 with S_3 and a smaller relic density as θ_{CPV} moves away from the CP -conserving limit and

FIG. 3. Regions producing 100% of DM in green and underproducing DM in blue in the $\theta_{\text{CPV}} - m_{\text{DM}}$ plane. The horizontal dashed red line represents the CP -conserving limit.

toward the maximum CP violation at the $\pi/2$ and $3\pi/2$ boundaries. Figure 3 illustrates this behavior in $B_1D_{12}C_1$ and $B_1D_{20}C_1$ scenarios. Note that in the latter scenario, due to the large δ_{12} , the intermediate mass overproduces DM regardless of the size of the $S_1 - S_3$ mass splitting.

B. Benchmarks of type 2

In both $B_2D_{55}C_1$ and $B_2D_{55}C_{15}$ scenarios, with only S_3 close in mass with the DM, there exists only one coannihilation channel, namely the $S_1 S_3 \rightarrow Z^* \rightarrow f \bar{f}$ channel, which dictates the behavior of the model in the low mass

FIG. 4. The change of the $g_{ZS_1S_3}$ coupling with respect to the CP -violating angle in the type 2 scenarios for a given mass of $m_{S_1} = 50$ GeV.

FIG. 5. The change in DM relic density for various DM masses with respect to the CP -violating angle. The horizontal red band shows the Planck observation limit on the abundance of DM.

region. As DM mass approaches the $W^\pm$, Z masses, the $S_1 S_1 \rightarrow VV$ ($V = W^\pm, Z$) annihilation channels open up and reduce the DM number density, which leads to an underproduction of DM for m_{DM} above this range, irrespective of the CP -violating angle.

Recall that the $ZS_1 S_3$ coupling is sensitive to the changes in the CP -violating angle. Figure 4 shows the absolute value of the $ZS_1 S_3$ coupling for an exemplary DM mass of 50 GeV with respect to θ_{CPV} in type 2 benchmark

scenarios. Due to the striking difference between the values of the coupling in the two benchmark scenarios, one expects a substantial difference in the DM relic density in the two models. The $B_2 D_{55} C_1$ scenario consistently overproduces DM in the mass range $45 \text{ GeV} < m_{\text{DM}} < 75 \text{ GeV}$ (except for large CP -violating angles around the Z resonance region $m_{\text{DM}} \approx m_Z/2$ where S_1 and S_3 are very close in mass). In this scenario, the $ZS_1 S_3$ coupling is so weak that it fails to coannihilate DM effectively in this mass

FIG. 6. Regions producing 100% of DM in green and underproducing DM in blue in the $\theta_{\text{CPV}} - m_{\text{DM}}$ plane. The horizontal dashed red line represents the CP -conserving limit.

FIG. 7. The relative strength of the $g_{VS_i S_j}$ coupling in the two scenarios for a given mass of $m_{S_1} = 90 \text{ GeV}$.

range. The $B_2D_{55}C_{15}$ scenario, on the other hand, has a large enough ZS_1S_3 coupling at large θ_{CPV} to satisfy the Planck limit on the DM relic density. Figure 5 confirms this behavior where DM relic density for various DM masses is shown. The $B_2D_{55}C_1$ scenario overproduces DM for masses below 75 GeV, and the $B_2D_{55}C_{15}$ scenario produces DM in agreement with the Planck limit for large CP -violating angles for this mass range. As mentioned before, both scenarios underproduce DM for larger masses when the $S_1S_1 \rightarrow VV$ annihilation channel is open.

Figure 6 shows regions where S_1 contributes to 100% of the observed DM in green, and regions where it only provides a fraction of the observed relic density in blue, in the $\theta_{CPV} - m_{DM}$ plane. The blank regions are ruled out by Planck observations as they lead to an overproduction of DM. Note that in the CP -conserving limit where $\theta_{CPV} = \pi$ the model fails to comply with the Planck observations.

C. Benchmarks of type 3

The type 3 benchmark scenarios are defined in the heavy mass region, $m_{DM} \geq 80$ GeV, where the charged inert scalars could also be close in mass with S_1 , thereby providing new coannihilation channels for the DM candidate. When studying the DM phenomenology of the model,

it is not only the annihilation and coannihilation of DM but also the (co)annihilation of other inert particles among each other that should be taken into account.

Figure 7 shows the relevant and non-negligible VS_1S_i couplings for all three type 3 benchmark scenarios where $V = W^\pm, Z$ and S_i is a neutral or charged inert particle. Due to the presence of so many coannihilation processes, type 3 scenarios underproduce DM. Note, however, that in the heavy mass region the annihilation $S_1S_1 \rightarrow VV$ is dominant whose coupling is independent of the CP -violating angle. Therefore, one expects a similar behavior in all three scenarios.

Figure 8 shows the relic abundance of the DM candidate for various DM masses. As mentioned before, all three scenarios have a similar behavior, with the $B_3D_{55}C_1$ scenario providing a slightly larger relic density in comparison to the $B_3D_5C_1$ scenario due to a larger δ_{12} mass splitting. Scenario $B_3D_{55}C_{15}$ provides only two coannihilation channels for S_1 ; however, they have larger couplings compared to the $B_3D_{55}C_1$ case, which leads to a slightly smaller relic abundance for S_1 . The reason we do not see the revival of the very heavy mass region $m_{DM} > 400$ GeV as shown in [25,34] is the absence of the Higgs mediated processes $S_iS_j \rightarrow h \rightarrow VV$, where $S_{i,j}$ is any neutral or

FIG. 8. The change DM relic density for various DM masses with respect to the CP -violating angle. The horizontal red band shows the Planck observation limit on the abundance of DM.

charged inert particle. These Higgs-mediated processes have a destructive interference with pure gauge processes $S_i S_j \rightarrow VV$, which would have revived the heavy mass region.

V. CONCLUSION AND OUTLOOK

The scalar potential is the least constrained sector of the SM which, if extended, could provide new sources of CP violation and viable DM candidates. New sources of CP violation are often limited due to their contribution to the EDMs. Introducing CP violation in the dark sector, however, is unconstrained as it has no contribution to the EDMs. On the other hand, dark sectors interacting with the visible sector through a Higgs portal are under tension, as they require a large portal coupling for efficient annihilation of DM and a small portal coupling to satisfy direct and indirect detection experiments and SM-Higgs data.

We present a novel mechanism in which the CP -violating dark particles only interact with the SM through the gauge bosons, primarily the Z boson, in the framework of a three Higgs doublet model. In the region where Higgs portal

interactions are subdominant, we show that the Z portal CP -violating DM can still thermalize and satisfy all experimental and observational data.

In the context of electroweak baryogenesis, the extended scalar sector could easily accommodate a strong first order phase transition. The efficient transfer of the unconstrained dark CP violation to the visible sector is under study and will be the subject of our future publication.

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