

Forbidden dark matter assisted by a first-order phase transition and the associated gravitational waves

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(Received 27 April 2026; accepted 1 September 2026; published 23 September 2026)

We propose a simple yet testable framework for light fermion dark matter (DM) with mass in the MeV–GeV range, charged under a dark $U(1)_D$ gauge symmetry. Before symmetry-breaking, DM annihilates excessively into massless dark gauge bosons, resulting in an underabundant relic and hence creating severe tension with cosmic microwave background (CMB) and indirect detection constraints. This is naturally remedied by a strong first-order phase transition (FOPT) in the dark sector, triggered by a scalar field Φ , which gives rise to masses for the dark gauge boson (X_D) and the physical scalar (ϕ). To achieve the correct relic abundance while evading indirect detection constraints, the s -wave annihilation channel $\chi\bar{\chi} \rightarrow X_D\phi$ is strictly maintained in the kinematically forbidden regime. Crucially, due to the $U(1)_D$ gauge structure, the usual $\chi\bar{\chi} \rightarrow \phi\phi$ annihilation is strictly absent at tree level and only arises at the one-loop level. Within this framework, we explore two distinct possibilities. When the loop-induced $\chi\bar{\chi} \rightarrow \phi\phi$ channel is also kinematically forbidden, it predominantly determines the relic density. Conversely, when the $\chi\bar{\chi} \rightarrow \phi\phi$ channel is kinematically allowed, a rich interplay emerges between the tree-level forbidden $\chi\bar{\chi} \rightarrow X_D\phi$ process and the loop-level allowed $\chi\bar{\chi} \rightarrow \phi\phi$ process, with the dominant contribution being dictated sensitively by the DM mass and gauge coupling. The late-time annihilations are highly suppressed in both of these channels, either by the energetic threshold of the forbidden channel or the p -wave velocity suppression of the loop-induced $\chi\bar{\chi} \rightarrow \phi\phi$ final-state process, rendering the scenario entirely safe from CMB and indirect search bounds. Moreover, the same symmetry-breaking phase transition is strongly first-order, producing a stochastic gravitational-wave background that could be probed by upcoming space-based interferometers and pulsar timing arrays. We demonstrate that achieving the observed DM abundance tightly correlates the DM mass with the nucleation temperature of the phase transition. Thus, this setup links the DM relic abundance, dark-sector dynamics, and gravitational-wave signals, offering complementary paths to discovery in both terrestrial and cosmological observations.

DOI: [10.1103/f9jp-vvft](https://doi.org/10.1103/f9jp-vvft)

I. INTRODUCTION

Understanding the identity of dark matter (DM), which constitutes approximately 85% of the matter content of the Universe, remains one of the foremost challenges in modern cosmology. Robust gravitational evidence from galaxy rotation curves, gravitational lensing, and cluster dynamics [1–3], together with precision measurements of the cosmic microwave background (CMB) [4,5], indicate

that roughly a quarter of the total energy density of the Universe is composed of nonluminous and nonbaryonic matter, whose microscopic nature remains unknown.

For several decades, Weakly Interacting Massive Particles (WIMPs) were regarded as the leading DM candidate, motivated by their natural appearance in many extensions of the Standard Model (SM) and their ability to reproduce the observed relic abundance via thermal freeze-out. However, as increasingly sensitive direct detection experiments have continued to report null results, the viability of the canonical WIMP paradigm has been placed under growing tension [6,7]. In particular, the absence of statistically significant signals across multiple generations of direct detection experiments has led to increasingly stringent exclusion limits on WIMP–nucleon scattering cross-sections. This mounting experimental pressure has motivated the exploration of alternative DM candidates, including light DM with masses $m_{\text{DM}} \lesssim 10$ GeV [8–20].

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Such sub-GeV to GeV candidates are now being actively probed by a new generation of low-threshold experiments (e.g. SuperCDMS, SENSEI, CRESST, DarkSide, DAMIC), which search for nuclear or electron recoils induced by light DM scattering [21–28].

However, even light DM produced via conventional thermal freeze-out is subject to severe cosmological and astrophysical constraints. For DM masses below the electroweak scale, thermal freeze-out generically implies sizable annihilation rates that persist well after chemical decoupling. Such late-time annihilations inject energy into the primordial plasma, and are therefore tightly constrained by observations of the CMB as well as by indirect detection experiments [29]. In particular, CMB measurements strongly limit energy injection during the epoch of recombination, excluding standard s -wave annihilating DM with masses $m_{\text{DM}} \lesssim \mathcal{O}(10)$ GeV for visible final states [4,29]. Complementary constraints arise from indirect searches for annihilation products in gamma rays and cosmic rays, such as those performed by Fermi-LAT and AMS-02, which rule out the canonical thermal annihilation cross-section for DM masses below ~ 20 GeV [30,31]. Taken together, these results effectively exclude the simplest realizations of light thermal DM.

A viable alternative is the forbidden DM scenarios [32], wherein the dominant annihilation channel is kinematically suppressed at late times. In a forbidden model, DM annihilates into slightly heavier states during freeze-out, so that annihilation shuts off exponentially as the temperature drops, evading CMB bounds [32,33]. For instance, one can arrange DM to annihilate into heavier dark-sector bosons or scalars, rather than directly into SM particles. However, if the final states are SM fields, the DM mass and couplings remain highly constrained by laboratory and astrophysical limits. We therefore consider a secluded dark sector in which the DM χ annihilates primarily into a dark gauge boson X_D and a dark Higgs scalar ϕ [which breaks the dark $U(1)$ symmetry]. This hidden-sector forbidden annihilation can yield the correct relic density while allowing sizable DM–SM couplings (via gauge kinetic mixing or other portals) that keep the model phenomenologically accessible. A key feature of this framework is its reliance on dark gauge symmetry breaking, which shapes the DM phenomenology. Before the phase transition, both the dark gauge boson X_D and the scalar ϕ are effectively massless. In this symmetric phase, the DM annihilates excessively into these massless states. If these couplings are large enough to be phenomenologically interesting, this process would completely deplete the relic abundance while remaining highly active at late times, creating severe tension with CMB and indirect detection constraints. A strongly first-order phase transition (FOPT) can resolve this issue. By dynamically generating masses for both X_D and ϕ , the FOPT introduces a kinematic threshold that alters the available annihilation phase space. Moreover, the sizable

couplings required to drive this strong FOPT naturally generate a stochastic gravitational-wave (GW) background that lies within the sensitivity of present and future GW detectors [34–39].

Following symmetry breaking, the relic abundance is governed by an interesting interplay of annihilation channels. Because the strong FOPT dynamics naturally dictate a hierarchical mass spectrum ($M_{X_D} \gg M_\phi$), the $\chi\bar{\chi} \rightarrow X_D X_D$ channel is severely phase-space suppressed. Consequently, the tree-level, s -wave process $\chi\bar{\chi} \rightarrow X_D \phi$ naturally dominates the thermal cross-section. To achieve the correct relic density while avoiding late-time constraints, we deliberately tune the mass splitting such that $2M_\chi < M_{X_D} + M_\phi$. This restricts the dominant s -wave channel to the thermally assisted “forbidden” regime during freeze-out. Furthermore, due to the specific $U(1)_D$ gauge charge assignments in our model, the $\chi\bar{\chi} \rightarrow \phi\phi$ annihilation is strictly absent at tree level; it manifests only as a one-loop process (e.g., via a triangle diagram and box diagram with gauge boson and DM running in the loop) that is intrinsically p -wave suppressed.

With the dominant tree-level $\chi\bar{\chi} \rightarrow X_D \phi$ process strictly residing in the forbidden regime, the relic density determination bifurcates into two distinct phenomenological cases. In the first case, the loop-induced $\chi\bar{\chi} \rightarrow \phi\phi$ channel is also kinematically forbidden, leaving the relic abundance exclusively determined by the forbidden $\chi\bar{\chi} \rightarrow \phi\phi$ process even if it suffers from loop effect and p -wave suppression. In the second case, the $\chi\bar{\chi} \rightarrow \phi\phi$ channel is kinematically allowed, leading to a rich interplay between the heavily Boltzmann-suppressed tree-level $\chi\bar{\chi} \rightarrow X_D \phi$ process and the phase-space open (but loop- and p -wave suppressed) $\chi\bar{\chi} \rightarrow \phi\phi$ process. The channel that ultimately dictates the relic density depends sensitively on the DM mass and the gauge coupling. Crucially, in all cases, late-time annihilations are highly suppressed, either by the energetic thresholds of the forbidden mechanism or by the velocity dependence of the p -wave loop process, thus keeping the scenario completely safe from CMB and indirect search.

In summary, this scenario not only achieves the correct thermal relic density but also offers multiple experimental avenues, including direct detection of DM, searches for new forces, as well as gravitational-wave astronomy. Moreover, the model predicts a striking correlation between the DM mass and the gravitational-wave spectrum: heavier DM requires a higher-scale phase transition to maintain the forbidden mechanism, shifting the GW peak frequency to higher values. This robust model prediction provides a powerful consistency check, linking the microscopic properties of DM to distinct signatures in gravitational-wave detectors.

The rest of the paper is organized as follows. In Sec. II, we discuss the forbidden DM. The minimal model is introduced in Sec. III, followed by the F first-order phase

transition, which is discussed in Sec. IV. We discuss the forbidden DM from the first-order phase transition in Sec. V. We discuss other phenomenological constraints in Sec. VI. Section VII discusses gravitational-wave production from the first-order phase transition. We finally conclude in Sec. VIII.

II. FORBIDDEN DARK MATTER

We consider a scenario where a Dirac fermionic DM χ couples to dark sector states $\mathcal{D}_1$ and $\mathcal{D}_2$ through an effective interaction term $\mathcal{L}_{\text{int}} \supset g_{\text{eff}} \bar{\chi} \chi \mathcal{D}_1 \mathcal{D}_2$. The relic abundance is determined through the annihilation process $\chi \bar{\chi} \rightarrow \mathcal{D}_1 \mathcal{D}_2$, with a thermally averaged cross-section of the order $\langle \sigma v \rangle \sim g_{\text{eff}}^2 / M_\chi^2$, where we assume that $\mathcal{D}_1$ and $\mathcal{D}_2$ are massless or have negligible mass compared to χ . This can reproduce the observed DM abundance for weak-scale masses and a sizable coupling. For lighter DM below the weak scale, however, the annihilation cross-section becomes too large, driving the relic density to underabundant values unless $g_{\text{eff}} \ll 1$. An attractive alternative is to consider annihilation into slightly heavier final states, the so-called forbidden channels. In this case, the annihilation is kinematically forbidden at zero temperature but proceeds in the early Universe thanks to the Boltzmann tail, leading to an exponentially suppressed cross-section at late times. This allows light DM to achieve the correct relic density even with order-one couplings [32,33].

In the forbidden regime, the thermally averaged annihilation cross-section can be related to the reverse process by detailed balance,

$$\langle \sigma_{\chi \bar{\chi} \rightarrow \mathcal{D}_1 \mathcal{D}_2} v \rangle = \frac{n_{\mathcal{D}_1}^{\text{eq}} n_{\mathcal{D}_2}^{\text{eq}}}{(n_\chi^{\text{eq}})^2} \langle \sigma_{\mathcal{D}_1 \mathcal{D}_2 \rightarrow \chi \bar{\chi}} v \rangle, \quad (1)$$

where n_i^{eq} denotes the equilibrium number density of species i , which for nonrelativistic particles is given by $n_i^{\text{eq}} = g_i (M_i T / 2\pi)^{3/2} e^{-M_i/T}$. Defining $r_i = M_{\mathcal{D}_i} / M_\chi$, $\Delta = (M_{\mathcal{D}_1} + M_{\mathcal{D}_2} - 2M_\chi) / 2M_\chi$ and $x = M_\chi / T$, the ratio of equilibrium densities yield

$$\frac{n_{\mathcal{D}_1}^{\text{eq}} n_{\mathcal{D}_2}^{\text{eq}}}{(n_\chi^{\text{eq}})^2} \simeq \frac{g_{\mathcal{D}_1} g_{\mathcal{D}_2}}{g_\chi^2} (r_1 r_2)^{3/2} e^{-2\Delta x},$$

where the exponential factor $e^{-2\Delta x}$ encodes the kinematic suppression from the heavier final states. Including the phase-space factor for the reverse process, the thermally averaged forbidden annihilation cross-section can be written schematically as

$$\langle \sigma_{\chi \bar{\chi} \rightarrow \mathcal{D}_1 \mathcal{D}_2} v \rangle = \frac{g_{\mathcal{D}_1} g_{\mathcal{D}_2}}{g_\chi^2} \mathcal{F}(r_1, r_2) e^{-2\Delta x} \langle \sigma_{\mathcal{D}_1 \mathcal{D}_2 \rightarrow \chi \bar{\chi}} v \rangle, \quad (2)$$

where $\mathcal{F}(r_1, r_2)$ is a kinematic function, which for nearly degenerate final states $M_{\mathcal{D}_1} \simeq M_{\mathcal{D}_2}$ reduces to the familiar

$(1 + \Delta)^{3/2}$ factor in the literature. The key point is that the annihilation rate is exponentially suppressed, so that the resulting relic abundance scales approximately as $\Omega_\chi h^2 \propto e^{2\Delta x_f}$ with $x_f = M_\chi / T_f$ the freeze-out parameter. As a result, even relatively light DM with sizable couplings can reproduce the observed relic density by appropriately tuning the mass splitting Δ .

III. THE MINIMAL MODEL

We now introduce a minimal setup that realizes the forbidden DM scenario triggered by a first-order phase transition in the dark sector. The model contains a Dirac fermion χ with charge +1 under a $U(1)_D$ gauge symmetry and a complex scalar singlet $\Phi = (\phi + i\eta) / \sqrt{2}$ with charge -1. The relevant dark sector Lagrangian is given by

$$\mathcal{L}_{U(1)_D} = \bar{\chi} (i\gamma^\mu \mathcal{D}_\mu - M_\chi) \chi - \frac{1}{4} X_{\mu\nu} X^{\mu\nu} - \frac{\epsilon}{2} B_{\mu\nu} X^{\mu\nu} + (\mathcal{D}_\mu \Phi)^* (\mathcal{D}_\mu \Phi) + \mu_\Phi^2 \Phi^* \Phi - \lambda_\Phi (\Phi^* \Phi)^2, \quad (3)$$

where the covariant derivative is defined as $\mathcal{D}_\mu = \partial_\mu + iQ_D g_D (X_D)_\mu$ with Q_D being the charge of the particle under $U(1)_D$ symmetry, and ϵ is the kinetic mixing parameter between the dark gauge field $(X_D)_\mu$ and the hypercharge gauge boson B_μ . After spontaneous breaking of the $U(1)_D$ symmetry, the scalar can be parameterized around its vacuum expectation value (VEV) as

$$\Phi = \frac{\phi + v_\phi}{\sqrt{2}}, \quad (4)$$

where v_ϕ denotes the VEV of the dark scalar field. The symmetry breaking generates masses for the dark gauge boson and the physical scalar ϕ as

$$M_{X_D} = g_D v_\phi, \quad M_\phi = \sqrt{2\lambda_\Phi} v_\phi. \quad (5)$$

A pictorial summary of the proposed scenario is shown in Fig. 1, which captures the essential thermal history of the dark sector in this framework. In the high-temperature symmetric phase, the dark scalar field sits at the origin, $v_\phi = 0$, so that both the dark gauge boson and the scalar remain massless and the forbidden annihilation channel is not operative. As the temperature drops below the nucleation temperature T_n , the dark sector undergoes a first-order phase transition, and the scalar acquires a nonzero vacuum expectation value, $v_\phi \neq 0$, thereby generating masses for both the dark gauge boson and the scalar itself.

In the broken phase, the mass hierarchy is arranged such that $M_\chi < (M_{X_D} + M_\phi) / 2$, rendering the process $\chi \bar{\chi} \rightarrow X_D \phi$ a forbidden channel at late times while still allowing it to proceed efficiently in the early Universe due to the thermal population of the heavier final states. This

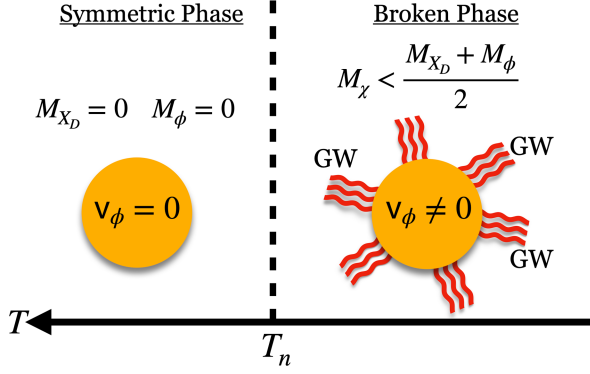


FIG. 1. Pictorial representation of our model, which gives rise to gravitational waves via a first-order phase transition and correct relic of DM through the forbidden channel.

realizes the forbidden DM mechanism discussed above and enables the correct relic abundance for light DM with sizable couplings while simultaneously producing a stochastic gravitational-wave background from the strong first-order phase transition. Taken together, this setup tightly links the origin of the DM relic density to testable signals in upcoming gravitational-wave experiments and dark sector searches. Prior to analyzing the DM phenomenology, we first outline the dynamics of the dark sector phase transition and determine the regions of parameter space that yield a strongly first-order transition.

IV. FIRST-ORDER PHASE TRANSITION

As emphasized above, the same dark sector dynamics that generate the correct relic abundance through a forbidden channel are facilitated by a strong first-order phase transition, providing a complementary source of a stochastic gravitational-wave background. In this section, we focus on the thermal evolution of the scalar field and identify regions of parameter space in which $U(1)_D$ symmetry breaking proceeds via a strongly first-order phase transition, suitable for generating potentially observable gravitational-wave signals.

The scalar potential receives thermal corrections at finite temperature, and it can be written as

$$V_{\text{eff}}(\phi, T) = V_{\text{tree}}(\phi) + V_{\text{CW}}(\phi) + V_{\text{ct}}(\phi) + V_T(\phi, T) + V_{\text{daisy}}(\phi, T), \quad (6)$$

where $V_{\text{tree}}(\phi)$ is the tree-level potential, $V_{\text{CW}}(\phi)$ is the Coleman-Weinberg one-loop correction, $V_{\text{ct}}(\phi)$ is the counterterm, $V_T(\phi)$ is the finite temperature contribution, and V_{daisy} accounts for the daisy correction. The explicit expressions of these terms are given in Appendix C. At high temperatures, the effective potential is minimized at $\phi = 0$. As the Universe cools and the temperature drops, the shape of the potential evolves, and a second, energetically favorable minimum at $\phi \neq 0$ can appear, separated

from the origin by a potential barrier. In this situation, the system undergoes a FOPT, in which bubbles of the true vacuum nucleate within the metastable false vacuum and subsequently expand.

Such an FOPT can efficiently produce gravitational waves via three main processes: bubble wall collisions, sound waves in the plasma, and magnetohydrodynamic turbulence in the plasma. The corresponding spectra and their combination into the total gravitational-wave signal are discussed in detail in Appendix D. For phenomenological purposes, the gravitational-wave signal is largely characterized by two key parameters evaluated at the nucleation temperature T_n : the strength of the transition, quantified by the ratio of the released vacuum energy to the radiation energy density, and the inverse duration of the transition in units of the Hubble rate. These are conventionally expressed as

$$\alpha(T_n) = \frac{\rho_{\text{vac}}(T_n)}{\rho_{\text{rad}}(T_n)}, \quad (7)$$

where

$$\rho_{\text{vac}}(T) = V_{\text{eff}}(\phi_{\text{false}}, T) - V_{\text{eff}}(\phi_{\text{true}}, T) - T \frac{\partial}{\partial T} [V_{\text{eff}}(\phi_{\text{false}}, T) - V_{\text{eff}}(\phi_{\text{true}}, T)], \quad (8)$$

and

$$\rho_{\text{rad}}(T) = \frac{\pi^2}{30} g_* T^4, \quad (9)$$

with g_* denoting the number of relativistic degrees of freedom (d.o.f.) at $T = T_n$. The parameter $\alpha(T_n)$ thus measures how strongly the phase-transition energy budget departs from that of a purely radiation-dominated Universe. The timescale of the transition is encoded in the quantity β , defined as the inverse duration of the phase transition and commonly expressed in units of the Hubble rate at nucleation,

$$\frac{\beta}{\mathcal{H}(T_n)} \simeq T_n \frac{d}{dT} \left(\frac{S_3}{T} \right) \Big|_{T=T_n}. \quad (10)$$

where S_3 is the three-dimensional Euclidean action of the critical bubble configuration. A smaller value of $\beta/\mathcal{H}(T_n)$ corresponds to a slower, more prolonged transition and typically enhances the gravitational-wave signal, while larger values indicate a faster transition with a more suppressed signal.

In Fig. 2, we show the region of parameter space in the g_D - λ_ϕ plane that leads to a FOPT for a fixed dark scalar VEV $v_\phi = 100$ MeV. The color code indicates the corresponding value of $\alpha(T_n)$, providing a measure of the transition strength. This plot highlights the range of gauge

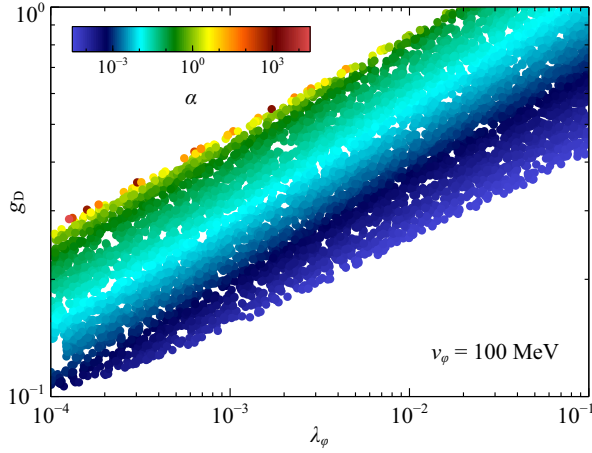


FIG. 2. Allowed regions in the dark sector gauge coupling (g_D) versus scalar self-coupling (λ_ϕ) plane for a fixed scalar VEV $v_\phi = 100$ MeV with the color scale representing the strength of the first-order phase transition (quantified by α), with larger values indicating stronger transitions.

and scalar couplings for which the dark sector undergoes a sufficiently strong FOPT to generate potentially detectable gravitational waves.

V. FORBIDDEN DARK MATTER FROM FIRST-ORDER PHASE TRANSITION

Having identified regions of the dark sector parameter space where the $U(1)_D$ symmetry breaking proceeds via a strong first-order phase transition, we now turn to the DM phenomenology of the same setup. In particular, we analyze how the phase transition reshapes the available annihilation channels and enables the forbidden process to control the relic abundance of χ .

Before detailing the cross-sections, it is crucial to explicitly define the kinematically and dynamically allowed annihilation channels dictated by the $U(1)_D$ gauge structure. Given the gauge charge assignments of the dark matter ($Q_\chi = +1$) and the dark scalar ($Q_\phi = -1$), a Yukawa interaction of the form $\bar{\chi}\Phi\chi$ is forbidden in the Lagrangian by gauge invariance; consequently, no tree-level $\bar{\chi}\chi\phi$ vertex exists. Furthermore, expanding the scalar kinetic term $|D_\mu\Phi|^2$ after spontaneous symmetry breaking generates a dimensionful $X_D X_D \phi$ vertex proportional to $g_D^2 v_\phi$, but does not yield an $X_D \phi \phi$ vertex for the real physical scalar ϕ . In the simultaneous absence of $\bar{\chi}\chi\phi$ and $X_D \phi \phi$ vertices, the process $\chi\bar{\chi} \rightarrow \phi\phi$ cannot physically occur at tree level, even within the parameter space where $M_\chi > M_\phi$. While the $\chi\bar{\chi} \rightarrow \phi\phi$ annihilation is strictly absent at tree level, it is generated at the one-loop level via box and triangle diagrams involving internal dark fermion and dark gauge boson as shown in Fig. 3. The effective interaction generated by integrating out the gauge bosons takes the form

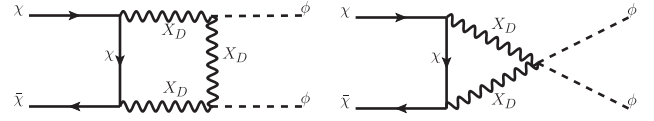


FIG. 3. Feynman diagrams for $\bar{\chi}\chi \rightarrow \phi\phi$ channel arising at one loop.

$$\mathcal{L}_{\text{eff}} \simeq \frac{12g_D^4}{16\pi^2} \frac{M_\chi}{M_{X_D}^2} \bar{\chi}\chi\phi^2. \quad (11)$$

Since parity conservation forbids the s -wave annihilation of a fermion pair into identical scalars, this loop-induced $\phi\phi$ channel strictly proceeds via p -wave and suffers a velocity suppression (v^2) at freeze-out and later times.

Consequently, in our setup, four annihilation channels can in principle contribute to setting the DM abundance:

$$\bar{\chi}\chi \rightarrow X_D X_D, \bar{\chi}\chi \rightarrow X_D \phi, \bar{\chi}\chi \rightarrow \phi\phi, \bar{\chi}\chi \rightarrow \bar{f}f, \quad (12)$$

where f denotes SM fermions. The process $\bar{\chi}\chi \rightarrow \bar{f}f$ is induced by the kinetic mixing between the dark gauge boson X_D and the hypercharge gauge boson, but its contribution to the relic density can be made negligible by taking the kinetic mixing parameter ϵ sufficiently small. In this regime, the DM abundance is primarily determined by annihilations into dark-sector states.

As discussed in Sec. IV, and evident from Fig. 2, a strong first-order phase transition in the dark sector dictates $g_D \gg \lambda_\phi$, which naturally enforces a strict mass hierarchy $M_{X_D} \gg M_\phi$ at fixed v_ϕ . Consequently, the phase-space suppression for $\bar{\chi}\chi \rightarrow X_D X_D$ is more severe than for the mixed final state $\bar{\chi}\chi \rightarrow X_D \phi$. Thus, the tree-level, s -wave process $\bar{\chi}\chi \rightarrow X_D \phi$ typically dominates the annihilation rate in the broken phase. To realize the forbidden DM mechanism, we restrict our parameter space such that this dominant channel is always kinematically forbidden at zero temperature:

$$2M_\chi < M_{X_D} + M_\phi. \quad (13)$$

This channel thus becomes active in the early Universe solely due to the high-energy Boltzmann tail of the thermal distribution. While $X_D \phi$ remains universally forbidden in our analysis, the loop-induced $\bar{\chi}\chi \rightarrow \phi\phi$ channel can be either allowed ($M_\chi > M_\phi$) or forbidden ($M_\chi < M_\phi$) depending on the specific scalar mass. In the following subsections, we analyze these two distinct phenomenological regimes and delineate the parameter spaces where the relic density is governed by the intricate interplay between tree-level exponential suppression and loop-level p -wave suppression.

A. Dark matter relic when $\bar{\chi}\chi \rightarrow X_D \phi$ is forbidden but $\bar{\chi}\chi \rightarrow \phi\phi$ is allowed

We first discuss the scenario in which the loop-induced $\bar{\chi}\chi \rightarrow \phi\phi$ process is kinematically allowed while the tree-level $\bar{\chi}\chi \rightarrow X_D \phi$ channel remains strictly forbidden. We

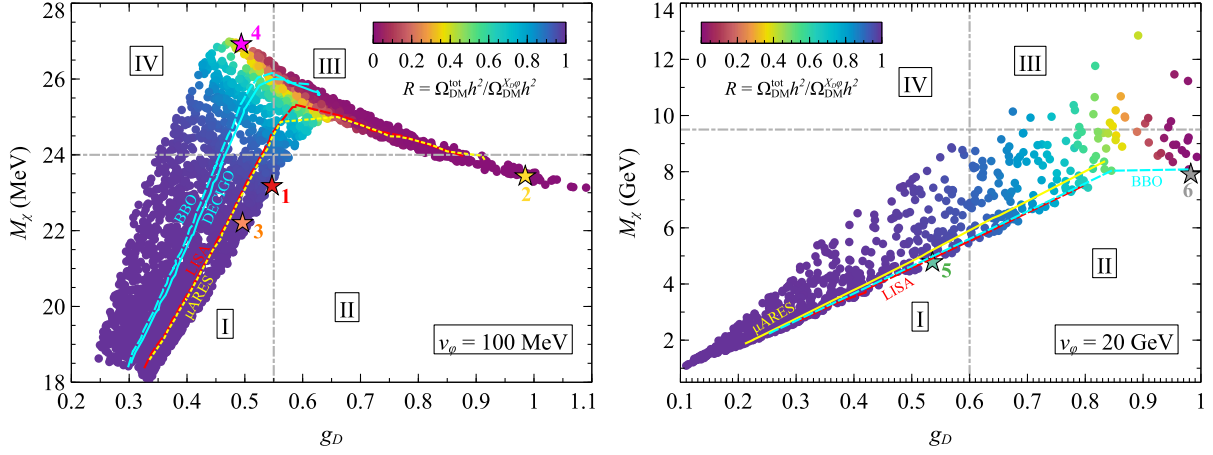


FIG. 4. Left: correct DM relic parameter space in the plane of M_χ versus g_D for fixed VEV of ϕ , $v_\phi = 100$ MeV. The same points also give rise to a first-order phase transition. The color code represents the ratio R . The SNR contours are shown for BBO [40] (cyan dashed), DECIGO [41] (cyan solid), LISA [42] (red solid), and μ ARES [43] (yellow dotted) for $\mathcal{S} = 10$ with $\tau_{\text{obs}} = 1$ year. See Sec. VII for more details on SNR. Right: the same with $v_\phi = 20$ GeV.

define the mass splitting parameter as $\Delta_{X_D\phi} = \frac{M_{X_D} + M_\phi - 2M_\chi}{2M_\chi}$. The DM yield $Y_\chi \equiv n_\chi/s$ evolves according to the Boltzmann equation:

$$\frac{dY_\chi}{dz} = -\frac{sz}{\mathcal{H}(M_\chi)} (\langle\sigma v\rangle_{\bar{\chi}\chi \rightarrow X_D\phi}^f + \langle\sigma v\rangle_{\bar{\chi}\chi \rightarrow \phi\phi}^a) \times (Y_\chi^2 - (Y_\chi^{\text{eq}})^2). \quad (14)$$

where $z = M_\chi/T$ is the dimensionless parameter used to track the evolution of temperature, Y_χ^{eq} is the equilibrium abundance of χ i.e. $Y_\chi^{\text{eq}}(z) = 0.116 \frac{g_\chi}{g_*} x^2 K_2(z)$, g_* is the effective number of relativistic d.o.f. at temperature T , and $K_2(z)$ is the modified Bessel function of the second kind. The Hubble parameter and entropy density are respectively given by $\mathcal{H}(M_\chi) = 1.66\sqrt{g_*} \frac{M_\chi^2}{M_{\text{Pl}}}$, $s = \frac{2\pi^2}{45} g_* M_\chi^3 z^{-3}$, where $M_{\text{Pl}} = 1.22 \times 10^{19}$ GeV is the Planck mass. The thermally averaged cross-section for the forbidden channel is related to the allowed reverse process via detailed balance as

$$\langle\sigma v\rangle_{\bar{\chi}\chi \rightarrow X_D\phi}^f = \frac{n_{X_D}^{\text{eq}} n_\phi^{\text{eq}}}{(n_\chi^{\text{eq}})^2} \langle\sigma v\rangle_{X_D\phi \rightarrow \bar{\chi}\chi}^a. \quad (15)$$

The expression for thermally averaged cross-section is given in Appendix B.

Solving the Boltzmann equation yields the late-time value of Y_χ , from which the present relic abundance follows as

$$\Omega_{\text{DM}} h^2 \simeq 0.118 \left(\frac{Y_\chi}{4.2 \times 10^{-10}} \right) \left(\frac{M_\chi}{1 \text{ GeV}} \right). \quad (16)$$

In the left panel of Fig. 4, we map the correct relic parameter space in the $M_\chi - g_D$ plane for a fixed scalar VEV: $v_\phi = 100$ MeV. The color code represents the ratio

$$R = \Omega_{\text{DM}}^{\text{tot}} h^2 / \Omega_{\text{DM}}^{X_D\phi} h^2, \quad (17)$$

where $\Omega_{\text{DM}}^{\text{tot}} h^2$ is obtained by solving Eq. (14) while $\Omega_{\text{DM}}^{X_D\phi} h^2$ is obtained from the same Eq. (14) but keeping the $\bar{\chi}\chi \rightarrow \phi\phi$ channel switched off. We note that R quantifies the relative importance of the two channels. R approaching unity indicates that the heavily Boltzmann-suppressed, yet strongly coupled s -wave $\bar{\chi}\chi \rightarrow X_D\phi$ channel dictates the relic density. Conversely, $R \ll 1$ indicates that the exponential penalty on $\bar{\chi}\chi \rightarrow X_D\phi$ is too severe, allowing the phase-space open (but loop and p -wave suppressed) $\bar{\chi}\chi \rightarrow \phi\phi$ channel to take over. As the DM mass M_χ increases, the cross-section for the allowed $\bar{\chi}\chi \rightarrow \phi\phi$ channel inherently grows due to the mass insertion factor $(M_\chi/M_{X_D})^2$ embedded in its effective coupling. At the same time, maintaining the correct relic density and FOPT conditions across the parameter space often requires a larger scalar mass M_ϕ . This larger M_ϕ increases the required mass splitting $\Delta_{X_D\phi}$, which inflicts a severe exponential penalty on the $\bar{\chi}\chi \rightarrow X_D\phi$ channel until the $\bar{\chi}\chi \rightarrow \phi\phi$ process inevitably dominates.

To illustrate this dynamical crossover, we select four benchmark points (BP1–BP4) listed in Table I. Comparing BP1 (red star) and BP2 (yellow star), we nearly fix the DM mass but significantly increase M_{X_D} (and thus g_D). In the left panel of Fig. 5 (BP1), the allowed $\bar{\chi}\chi \rightarrow \phi\phi$ channel (blue curve) freezes out early due to its loop suppression, while the forbidden $\bar{\chi}\chi \rightarrow X_D\phi$ channel (red curve) remains active longer, solely determining the relic abundance. However, for BP2 (right panel), the increased M_{X_D} leads to a much larger mass splitting ($\Delta_{X_D\phi}(\text{BP2}) \approx 1.53$ vs $\Delta_{X_D\phi}(\text{BP1}) \approx 0.61$). This inflicts a massive exponential penalty on the $\bar{\chi}\chi \rightarrow X_D\phi$ channel, causing it to decouple early. Consequently, the relic abundance for BP2 is entirely

TABLE I. Benchmark points for satisfying correct DM relic via $\bar{\chi}\chi \rightarrow X_D\phi(f)$, $\bar{\chi}\chi \rightarrow \phi\phi(a)$ channels, and giving first-order phase transition. $\Delta_{X_D\phi}(BP0) = 0.4091$, $\Delta_{X_D\phi}(BP1) = 0.60842$, $\Delta_{X_D\phi}(BP2) = 1.53066$, $\Delta_{X_D\phi}(BP3) = 0.2544036$, $\Delta_{X_D\phi}(BP4) = 0.411547441420$.

BPs	M_χ (MeV)	M_{X_D} (MeV)	M_ϕ (MeV)	λ_ϕ	v_ϕ (MeV)	g_D	T_c (MeV)	T_n (MeV)	α_*	$\beta_*/\mathcal{H}_*$
BP0	42.326	112.059	7.222	0.000104	500	0.2241	52.0094	35.432	0.1332	8880.491
BP1	23.1786	54.6895	5.840255	0.0017054	100	0.546895	17.41913	1.516388	839.4085	334.49253
BP2	23.4358	98.49544	20.120645	0.02024	100	0.984954	35.27115	24.25950	0.16288	658.2282
BP3	22.2034	49.60992	6.094203	0.00185696	100	0.496099	20.237876	13.308489	0.204574	1809.8969
BP4	26.919842	49.387419	26.60985	0.03540420	100	0.493874198	75.08041	74.9224079	0.00041969	139918.089
BP5	4778.9589	10713.3738	1164.1505	0.001694	20000	0.5356686	3564.6272	1143.8917	0.86177	849.7935
BP6	7912.7921	19649.45091	3751.20211	0.017589	20000	0.98247254	6535.093	3046.63690	0.152777	191.2885

governed by the g_D^8 -enhanced $\bar{\chi}\chi \rightarrow \phi\phi$ process. Conversely, examining BP3 (orange star) and BP4 (magenta star) in Fig. 6, we keep M_{X_D} nearly fixed while increasing M_χ . In BP3, $\bar{\chi}\chi \rightarrow X_D\phi$ dominates. For BP4, the simultaneous requirement of the correct relic density dictates a notably heavier dark scalar ($M_\phi \approx 26.6$ MeV compared to 6.1 MeV in BP3). This larger scalar mass drives up the mass-splitting parameter ($\Delta_{X_D\phi}(BP4) \approx 0.41$ compared to $\Delta_{X_D\phi}(BP3) \approx 0.25$), heavily suppressing the $\bar{\chi}\chi \rightarrow X_D\phi$ channel. Concurrently, the larger M_χ boosts the $(M_\chi/M_{X_D}^2)^2$ mass-insertion factor, causing the $\bar{\chi}\chi \rightarrow \phi\phi$ channel to definitively govern the freeze-out. This same qualitative interplay is observed for a higher symmetry-breaking scale ($v_\phi = 20$ GeV) in the right panel of Fig. 4, represented by BP5 and BP6.

To explicitly connect this particle physics parameter space with upcoming macroscopic cosmological probes, we have additionally overlaid gravitational-wave signal-to-noise ratio (SNR) ($\mathcal{S}$) contours for various future observatories in Fig. 4. The contours represent a detection threshold of $\mathcal{S} = 10$ with an observation time of $\tau_{\text{obs}} = 1$ year for BBO [40] (cyan dashed), DECIGO [41] (cyan solid), LISA [42] (red solid), and μARES [43]

(yellow dotted). These contours visually demonstrate the gravitational-wave detectability prospects directly alongside the valid dark matter relic parameter space; a detailed discussion of the exact SNR calculation and the corresponding phase transition dynamics is found in Sec. VII.

B. Dark matter relic when both $\bar{\chi}\chi \rightarrow X_D\phi$ and $\bar{\chi}\chi \rightarrow \phi\phi$ are forbidden

We now examine the regime where $M_\chi < M_\phi$, rendering both the $\bar{\chi}\chi \rightarrow X_D\phi$ and the loop-induced $\bar{\chi}\chi \rightarrow \phi\phi$ channels kinematically forbidden at zero temperature. We define the mass splitting for the $\bar{\chi}\chi \rightarrow \phi\phi$ channel as $\Delta_{\phi\phi} = \frac{M_\phi - M_\chi}{M_\chi}$. The relic density is obtained by solving the Boltzmann equation:

$$\frac{dY_\chi}{dz} = -\frac{sZ}{\mathcal{H}(M_\chi)} (\langle\sigma v\rangle_{\bar{\chi}\chi \rightarrow X_D\phi}^f + \langle\sigma v\rangle_{\bar{\chi}\chi \rightarrow \phi\phi}^f) \times (Y_\chi^2 - (Y_\chi^{\text{eq}})^2), \quad (18)$$

where both forbidden cross-sections are detailed and balanced from their respective allowed processes:

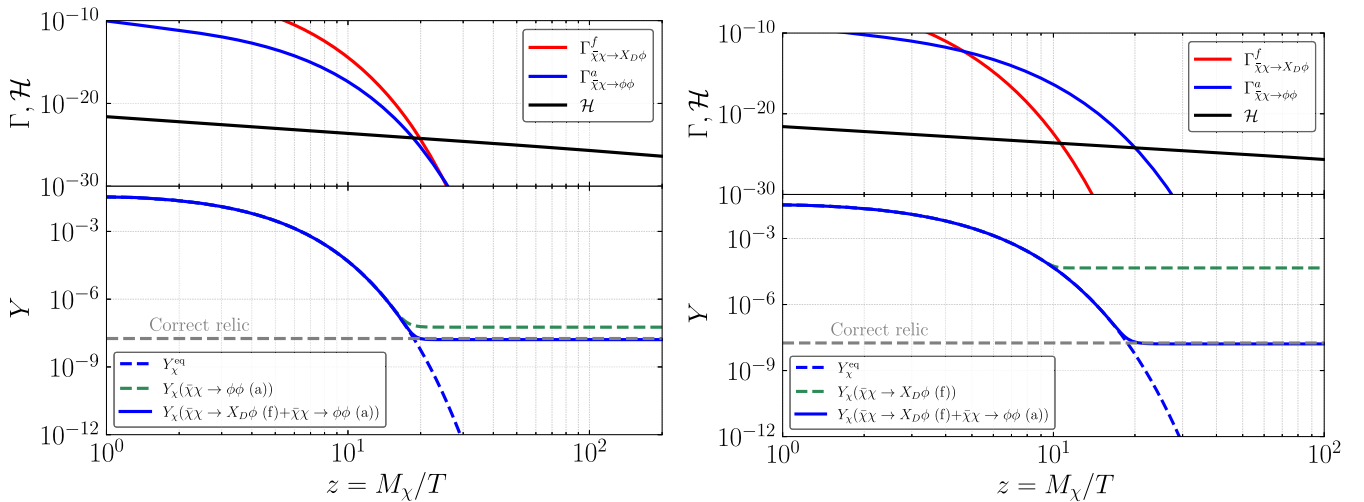


FIG. 5. Left: interaction rate comparison with Hubble (top panel) and cosmological evolution of the DM (bottom panel) for BP1 from Table I. Right: same for BP2 from Table I.

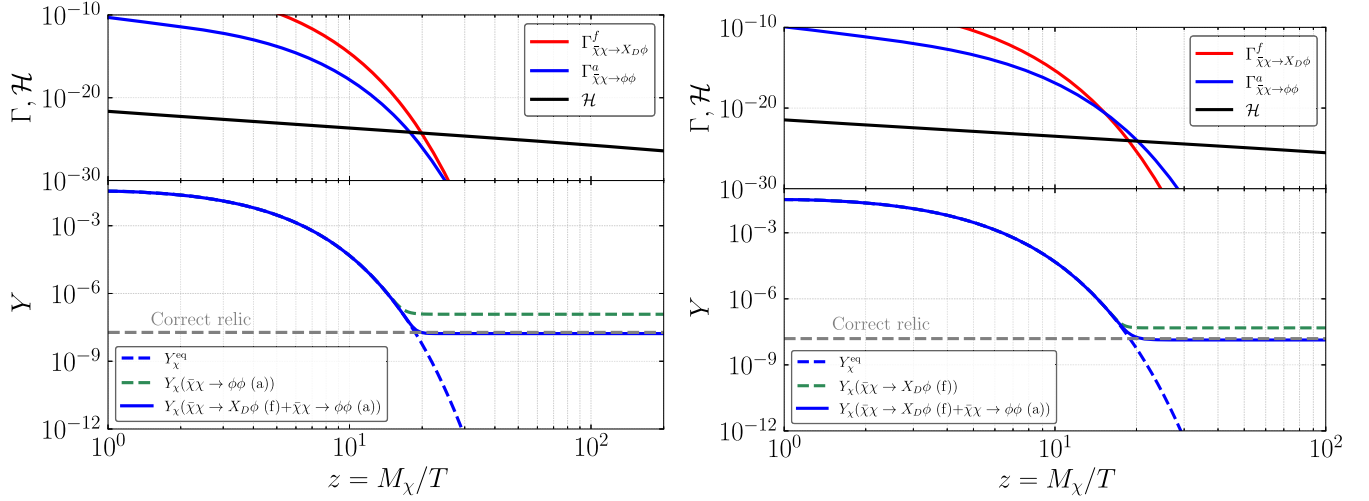


FIG. 6. Left: interaction rate comparison with Hubble (top panel) and cosmological evolution of the DM (bottom panel) for BP3 from Table I. Right: same for BP4 from Table I.

$$\begin{aligned} \langle \sigma v \rangle_{\tilde{\chi}\tilde{\chi} \rightarrow X_D \phi}^f &= \frac{n_{X_D}^{\text{eq}} n_\phi^{\text{eq}}}{(n_\chi^{\text{eq}})^2} \langle \sigma v \rangle_{X_D \phi \rightarrow \tilde{\chi}\tilde{\chi}}^a, \\ \langle \sigma v \rangle_{\tilde{\chi}\tilde{\chi} \rightarrow \phi\phi}^f &= \frac{(n_\phi^{\text{eq}})^2}{(n_\chi^{\text{eq}})^2} \langle \sigma v \rangle_{\phi\phi \rightarrow \tilde{\chi}\tilde{\chi}}^a. \end{aligned} \quad (19)$$

In this completely forbidden regime, the relic abundance depends on a competition between distinct suppressions. Because the FOPT demands $M_{X_D} \gg M_\phi$, the kinematic threshold for the $\tilde{\chi}\chi \rightarrow X_D \phi$ channel is much larger than the threshold for the $\tilde{\chi}\chi \rightarrow \phi\phi$ channel. Thus, the $\tilde{\chi}\chi \rightarrow X_D \phi$ channel suffers a significantly larger exponential Boltzmann penalty compared to the milder penalty on the $\tilde{\chi}\chi \rightarrow \phi\phi$ channel. When this disparity in Boltzmann suppression overcomes the intrinsic loop and p -wave penalties of the $\tilde{\chi}\chi \rightarrow \phi\phi$ process, the $\tilde{\chi}\chi \rightarrow \phi\phi$ channel

dominates the freeze-out. In the left panel of Fig. 7, we show the DM relic abundance as a function of $\Delta_{\phi\phi}$ for BP7–BP10. The right panel explicitly illustrates this dynamic for BP9: the milder exponential penalty allows the $\tilde{\chi}\chi \rightarrow \phi\phi$ channel to remain in equilibrium longer, cleanly determining the final DM relic abundance.

Finally, in Fig. 8, we present the viable parameter space in the ϵ – M_{X_D} plane for selected benchmarks across both regimes i.e. BP2, BP5, BP8, and BP9 (top left, top right, bottom left, and bottom right, respectively), as defined in Tables I and II. For each benchmark, the shaded and hatched regions summarize current constraints from DM–electron and DM–nucleon scattering experiments, as well as bounds from accelerator and astrophysical probes, while the vertical black dashed line indicates the value of M_{X_D} that reproduces the observed relic abundance.

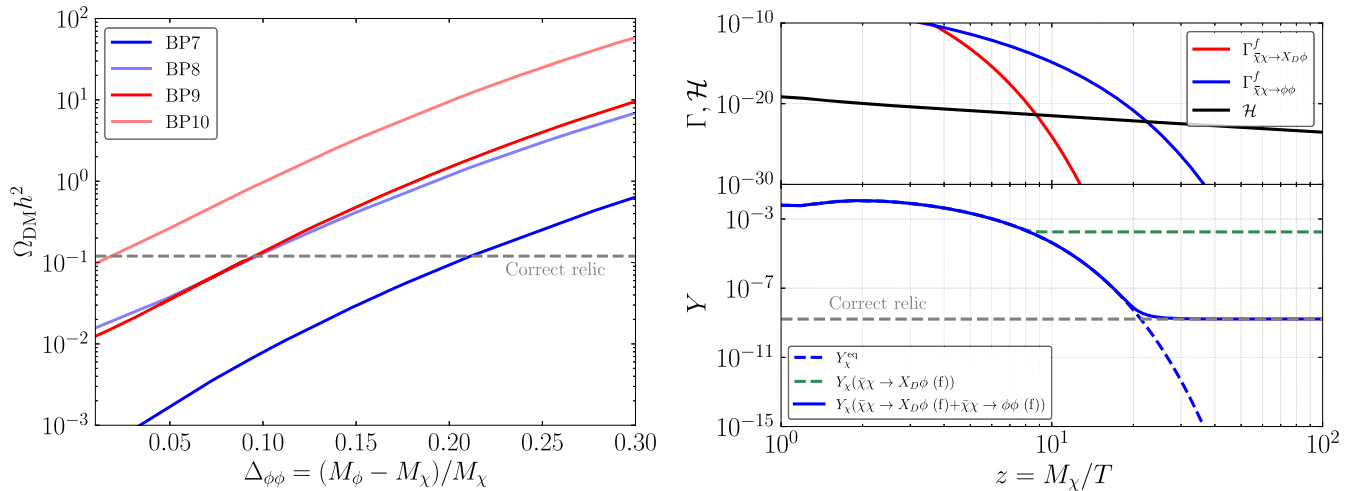


FIG. 7. Left: dark matter relic density $\Omega_{\text{DM}} h^2$ as a function of $\Delta_{\phi\phi}$ for the benchmark points in Table II. Right: interaction rate comparison with Hubble (top panel) and cosmological evolution of the DM (bottom panel) for BP9 from Table II.

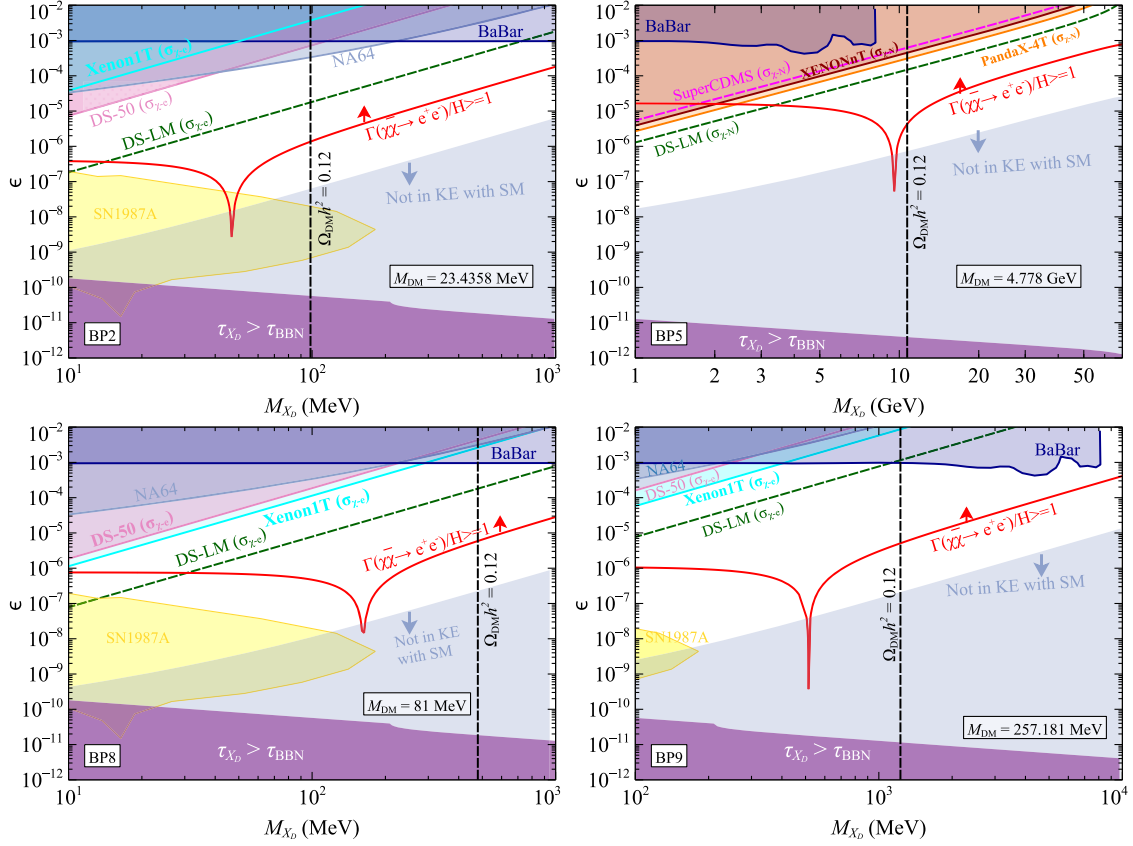


FIG. 8. Parameter space in the ϵ - M_{X_D} plane for *BP2*, *BP5*, *BP8*, and *BP9* (top left, top right, bottom left, bottom right respectively), as defined in Tables I and II. Shaded regions indicate possible constraints; the vertical black dashed line shows the value of M_{X_D} that yields the correct dark matter relic abundance for each benchmark.

The light blue region corresponds to parameter choices where χ fails to maintain kinetic equilibrium with the SM bath, primarily determined by the t -channel process $\chi e^- \rightarrow \chi e^-$. The region above the red solid curve satisfies $\Gamma_{\chi\bar{\chi} \rightarrow e^+e^-}/H \gtrsim 1$ at $z = 10$, implying that annihilations into e^+e^- substantially modify the relic density and the simple forbidden-channel picture no longer applies. The purple band near $\epsilon \sim 10^{-11}$ is excluded by big bang nucleosynthesis (BBN) considerations, since in this regime the dark gauge boson X_D becomes too long-lived and disrupts standard nucleosynthesis.

In Fig. 8, current exclusion limits from DM- e scattering experiments, specifically DS-50 [44] and XENON1T [21], are depicted as pink and cyan shaded regions, respectively.

Correspondingly, direct-detection exclusion limits from the DM-nucleon experiments XENONnT [45] and PandaX-4T [46] are indicated by the dark red and orange shaded regions. To highlight the future testability of the allowed parameter space, projected sensitivities from upcoming experiments are overlaid as nonshaded lines: the sensitivity of DS-LM [47] is represented by green dashed lines, while the projected reach of SuperCDMS [48] is shown as a magenta dashed line. Additional constraints from NA64 [49] (light blue shaded), BABAR [50] (dark blue shaded), and SN1987A [51–53] (light orange shaded) are included as well. Similar features and trends are visible across all four benchmark panels, illustrating that the interplay between relic-density requirements, kinetic equilibrium,

TABLE II. Benchmark points for satisfying correct DM relic via $\bar{\chi}\chi \rightarrow X_D\phi(f)$, $\bar{\chi}\chi \rightarrow \phi\phi(f)$ channels, and giving first-order phase transition. $\Delta_{\phi\phi}(BP7) = 0.21$, $\Delta_{\phi\phi}(BP8) = 0.097139$, $\Delta_{\phi\phi}(BP9) = 0.0947368$, $\Delta_{\phi\phi}(BP10) = 0.0138817$.

BPs	M_χ (MeV)	M_{X_D} (MeV)	M_ϕ (MeV)	λ_Φ	v_ϕ (MeV)	g_D	T_c (MeV)	T_n (MeV)	α_*	$\beta_*/\mathcal{H}_*$
BP7	30.53	145.41	36.96	0.06829	100	1.4541	47.3661	24.4982	0.39091	95.04275
BP8	81	479.225	88.869	0.015796	500	0.95845	158.0208	71.6957	0.7169459	187.9065
BP9	257.181	1226.57	281.545	0.039634	1000	1.22657	408.497	241.9205	0.086404	269.9877
BP10	1768.0408	6761.662	1792.5843	0.06426717	5000	1.35233	2453.449	1869.49	0.0186588	704.74155

and the various experimental and cosmological bounds consistently carves out a narrow but well-defined region in ϵ - M_{X_D} space.

VI. OTHER RELEVANT CONSTRAINTS

In addition to relic density, direct detection, and other requirements discussed in the previous section, the scenario is subject to several cosmological and collider bounds that further shape the viable parameter space. In this section, we discuss the limits from BBN on the lifetime of the scalar ϕ , and bounds on invisible Higgs decays from LHC measurements.

A. BBN constraint from Φ decay

The light scalar, ϕ , can decay into the SM fermions after mixing with the SM Higgs boson. The partial decay width into a fermion pair is given as

$$\Gamma_{\phi \rightarrow ff} = \sin^2 \theta \left(\frac{m_f}{v_0} \right)^2 \frac{M_\phi}{8\pi} \left(1 - \frac{4m_f^2}{M_\phi^2} \right)^{\frac{3}{2}}, \quad (20)$$

where m_f is the mass of the fermion, $\sin \theta$ is the mixing angle between the SM Higgs and scalar ϕ , and v_0 is the VEV of SM Higgs. The corresponding lifetime is given by

$$\tau_\phi = (\Gamma_{\phi \rightarrow ff})^{-1}. \quad (21)$$

The singlet scalar must decay before the onset of BBN; otherwise, its late-time decay would disrupt the successful predictions of primordial nucleosynthesis.

B. Constraint from Higgs invisible decay

If the dark sector particles X_D and ϕ are lighter than $M_h/2$, the SM Higgs boson can decay into pairs of these particles, contributing to its invisible decay width. Recent LHC measurements constrain the invisible branching ratio to $\mathcal{B}r(h \rightarrow \text{inv}) \leq 0.145$ [54]. The partial widths for Higgs decays into ϕ and X_D pairs are

$$\Gamma_{h \rightarrow \phi\phi} = \frac{\lambda_{H\Phi}^2 v_h^2}{8\pi M_h} \sqrt{1 - \frac{4M_\phi^2}{M_h^2}}, \quad (22)$$

$$\begin{aligned} \Gamma_{h \rightarrow X_D X_D} &= \frac{g_D^2 M_h^3 \sin^2 \theta}{8\pi M_{X_D}^2} \sqrt{1 - \frac{4M_{X_D}^2}{M_h^2}} \\ &\times \left(1 - \frac{4M_{X_D}^2}{M_h^2} + \frac{12M_{X_D}^4}{M_h^4} \right), \end{aligned} \quad (23)$$

where $\lambda_{H\Phi}$ is the Higgs–portal coupling defined in Eq. (A7) in Appendix A.

In Fig. 9, we show the combined BBN and Higgs invisible decay constraints in the $\sin \theta$ - M_ϕ plane. The gray shaded region corresponds to $\tau_\phi > \tau_{\text{BBN}} = 1$ s and is

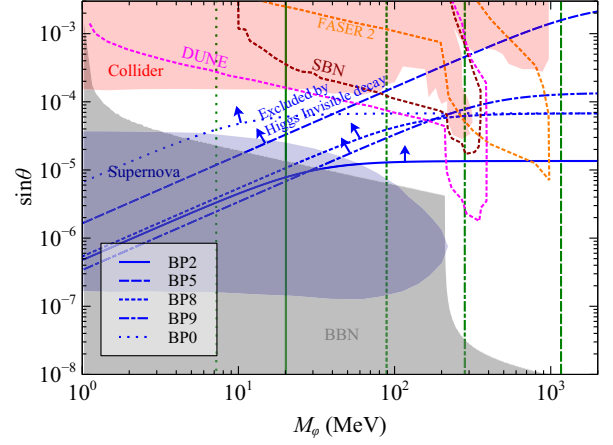


FIG. 9. Allowed parameter space from BBN and Higgs invisible decay constraints in the $\sin \theta$ - M_ϕ plane. The gray shaded region is excluded because $\tau_\phi > \tau_{\text{BBN}}$. The regions above the blue lines are excluded by Higgs invisible decay bounds. The constraint from supernova [55] is shown with blue shaded region. The collider limits are shown with red shaded region [56–58]. The future sensitivities from DUNE [59], SBN [60], and FASER 2 [61] are shown with different colored dotted lines.

therefore excluded by BBN. The constraints from supernova [55] and colliders [56–58] are shown with blue and red shaded regions, respectively. The future sensitivities of DUNE [59], SBN [60], and FASER 2 [61] are shown by magenta, dark red, and orange dotted lines, respectively. The blue curves represent the Higgs invisible decay bounds for the five benchmark points (BP0, BP2, BP5, BP8, and BP9) in Tables I and II. For each benchmark, we fix g_D and M_{X_D} to their benchmark values and vary $\{M_\phi, \sin \theta\}$, such that the region below each blue line satisfies $\mathcal{B}r(h \rightarrow \text{inv}) \leq 0.145$, while the region above is excluded. The values of M_ϕ corresponding to BP0, BP2, BP5, BP8, and BP9 are indicated by green spaced-dotted, solid, dashed, dotted, and dashed-dotted vertical lines, respectively. Evaluating the viability of our benchmark points against this comprehensive set of limits, we find that BP1, BP3, BP4, and BP7 are firmly excluded by the combined BBN and Higgs invisible decay constraints. Similarly, BP2 is ruled out by both the Higgs invisible decay bounds and supernova constraints. Furthermore, although BP9 safely evades the BBN and Higgs limits, it is ultimately excluded by existing collider bounds. Consequently, only BP0, BP5, BP6, BP8, and BP10 successfully survive all current experimental and cosmological constraints, establishing them as highly viable targets for future planned facilities.

VII. GRAVITATIONAL WAVE FROM FIRST-ORDER PHASE TRANSITION

Having established the viability of our benchmark points against various constraints, we now turn to the gravitational-wave signals generated by the dark-sector phase

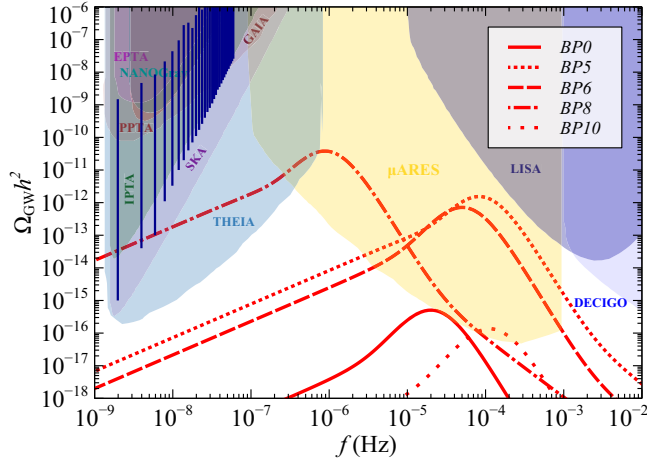


FIG. 10. Gravitational-wave spectrum for the five benchmark points listed in Tables I and II. The colored shaded regions indicate the projected sensitivities of different gravitational-wave detectors.

transition. The stochastic gravitational-wave background arising from a strong first-order phase transition is shaped by contributions from bubble collisions, sound waves in the plasma, and magnetohydrodynamic turbulence, as detailed in Appendix D. Here, we compute the GW spectrum for the five benchmark points (BP0, BP5, BP6, BP8, and BP10) that are allowed from all constraints, listed in Tables I and II, and compare it against the projected sensitivities of existing and planned detectors.

In Fig. 10, we present the predicted stochastic gravitational-wave spectra expressed as the GW energy density amplitude $\Omega_{\text{GW}} h^2$ versus frequency f for the selected benchmark points. For context, we overlay the current and projected sensitivities of a wide array of GW observatories, including NANOGrav [62,63], EPTA [64], PPTA [65], IPTA [66], SKA [67], LISA [42], THEIA [68], GAIA [68], DECIGO [41], and μ ARES [43]. The vertical blue lines denote the frequency and amplitude range currently favored by the NANOGrav 15-year dataset. The properties of these GW signals depend sensitively on the underlying dark-sector dynamics. For instance, the MeV-scale DM mass of BP8 dictates a nucleation temperature T_n that is correspondingly in the MeV range (cf. Table II). This phase transition scale naturally positions the resulting GW signal within the nHz to μ Hz frequency bands, placing it within the observational reach of pulsar timing arrays (PTAs), SKA, μ ARES, and astrometric observatories like THEIA. Notably, the low-frequency tail of this spectrum marginally intersects the parameter space indicated by the NANOGrav observations [62], demonstrating that such MeV-scale transitions can begin to probe regions of current observational interest. Conversely, while BP0 also features an MeV-scale DM mass and T_n where the relic density is dominantly decided by the $X_D \phi$ channel, its phase transition is characterized by a substantially larger inverse

duration parameter ($\beta_*/\mathcal{H}_*$). This rapid transition severely suppresses the overall GW amplitude and shifts the peak frequency to $\sim 2 \times 10^{-5}$ Hz, placing it only at the very edge of the projected μ ARES sensitivity. Finally, moving to the heavier mass regime, the GeV-scale DM masses characterizing BP5, BP6, and BP10 necessitate significantly higher phase-transition scales ($T_n \sim \mathcal{O}(\text{GeV})$). Consequently, their GW peak frequencies are shifted substantially upward to $\sim 10^{-4}$ Hz, bringing them into the optimal sensitivity windows of μ ARES and LISA.

A distinctive feature of this scenario is the strong correlation between the DM mass M_χ and the nucleation temperature T_n of the phase transition. This connection arises naturally from the fundamental kinematic requirements of our framework: to maintain the dominant tree-level $X_D \phi$ annihilation channel in the strictly forbidden regime ($M_{X_D} + M_\phi > 2M_\chi$), the dark sector masses M_{X_D} and M_ϕ must necessarily scale upward with the DM mass M_χ . Since these masses are generated via spontaneous symmetry breaking ($M_{X_D} = g_D v_\phi$ and $M_\phi = \sqrt{2\lambda_\phi} v_\phi$), achieving heavier final states without pushing the gauge and quartic couplings into nonperturbative regimes requires a larger scalar VEV, v_ϕ , for heavier DM. Elevating the VEV inherently elevates the entire thermal scale of the dark sector. Consequently, both the critical temperature T_c and the nucleation temperature T_n generally increase with M_χ , as is evident across the benchmark points spanning the MeV to GeV regimes in Tables I and II. This robust model prediction that heavier secluded DM is intrinsically accompanied by higher-scale phase transitions and correspondingly higher-frequency GW signals provides a powerful consistency check for the scenario and establishes a clear target for near-future multifrequency GW observations.

The detection prospects of the gravitational-wave signal at a given observatory can be effectively quantified by the signal-to-noise ratio (SNR), defined as [69]

$$\mathcal{S} = \sqrt{\tau_{\text{obs}} \int_{f_{\text{min}}}^{f_{\text{max}}} df \left[\frac{\Omega_{\text{GW}} h^2}{\Omega_{\text{expt}} h^2} \right]^2}, \quad (24)$$

where τ_{obs} is the total observation time and $\Omega_{\text{expt}} h^2$ represents the noise sensitivity spectrum of the specific detector.

In Fig. 4, we overlay SNR contours corresponding to a detection threshold of $\mathcal{S} = 10$ (assuming $\tau_{\text{obs}} = 1$ year) for BBO (cyan dashed), DECIGO (cyan solid), LISA (red solid), and μ ARES (yellow dotted) across the M_χ versus g_D parameter space. For a given M_χ , the region to the right of each respective contour yields an SNR strictly greater than 10. The physical origin of this behavior is straightforward: increasing the dark gauge coupling g_D strengthens the first-order phase transition, resulting in a larger latent heat fraction α_* . This stronger FOPT subsequently drives a larger gravitational-wave amplitude $\Omega_{\text{GW}} h^2$, elevating the resulting SNR according to Eq. (24). However, we note that

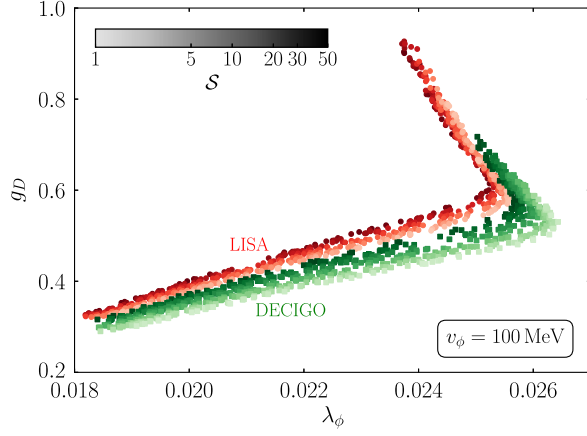


FIG. 11. FOPT parameter space consistent with DM relic abundance in the plane of g_D versus λ_ϕ . Red color represents LISA SNR, and the green color represents DECIGO SNR. The color grading denotes SNR values ranging from 1 to 50.

for DM mass, $M_\chi \gtrsim 23$ GeV, we observe the same SNR, $S = 10$, for two different g_D values for any fixed DM mass. This feature is a consequence of imposing the correct relic density criteria, as explained in Sec. V.

To further show the correlation of the phase-transition dynamics and the relic density parameter space, we isolate the FOPT parameter space in Fig. 11, mapping the exact SNR values across the g_D versus λ_ϕ plane for the points previously validated in Fig. 4. Here, we highlight parameter configurations yielding an SNR between 1 and 50, using a red color gradient for LISA and a green color gradient for DECIGO, with darker shades indicating higher SNR. As expected, a larger g_D enhances the GW amplitude and pushes the SNR toward higher values. Notably, the LISA SNR band is offset and lies at higher g_D values compared to the DECIGO band. Because LISA's projected strain sensitivity is lower than DECIGO's in this specific frequency range ($\Omega_{\text{expt}}^{\text{LISA}} h^2 > \Omega_{\text{expt}}^{\text{DECIGO}} h^2$), achieving the same SNR requires a correspondingly stronger gravitational-wave signal ($\Omega_{\text{GW}}^{\text{LISA}} h^2 > \Omega_{\text{GW}}^{\text{DECIGO}} h^2$). This strictly necessitates a stronger phase transition, thereby demanding a systematically larger gauge coupling $g_D^{\text{LISA}} > g_D^{\text{DECIGO}}$. Ultimately, the presence of these distinct, highly sensitive SNR bands highlights the robust testability of our framework. We note that the SNR bands for BBO and μARES are very similar to those of DECIGO and LISA, respectively, as can be seen from Fig. 4. Therefore, we refrain from showing BBO and μARES SNR bands in Fig. 11. Future detection of this stochastic gravitational-wave background by observatories such as LISA or DECIGO would not merely signify a cosmological phase transition but could also pinpoint the underlying microscopic gauge and scalar dynamics of the dark sector. When combined with ongoing direct detection and collider searches, these gravitational-wave signatures offer a powerful, complementary avenue for decisively probing the parameter space of this forbidden DM scenario.

VIII. CONCLUSIONS

We have presented a minimal $U(1)_D$ gauge extension of the Standard Model featuring a Dirac fermion dark-matter candidate χ and a complex singlet scalar Φ . The spontaneous breaking of the $U(1)_D$ symmetry is absolutely indispensable for the viability of the scenario: in its absence, the dark gauge boson X_D and the scalar ϕ would remain massless, leading to unsuppressed annihilations that would severely deplete the DM relic abundance and catastrophically violate late-time cosmological bounds. The symmetry-breaking elegantly resolves this by serving a dual purpose: it generates masses for both X_D and ϕ to reshape the annihilation kinematics, and it simultaneously triggers a strong first-order phase transition in the dark sector.

A defining feature of this model is the specific $U(1)_D$ charge assignment, which strictly forbids the standard $\bar{\chi}\chi \rightarrow \phi\phi$ annihilation at tree level, restricting it to a velocity-dependent p -wave loop process. To achieve the correct relic abundance, we restrict our parameter space such that the dominant tree-level s -wave annihilation, $\bar{\chi}\chi \rightarrow X_D\phi$, is strictly kinematically forbidden at zero temperature. Consequently, we find that the DM relic density is determined by two distinct regimes: (i) when the $\bar{\chi}\chi \rightarrow \phi\phi$ channel is kinematically allowed (creating a rich interplay between a forbidden tree-level process and an allowed loop-level p -wave process) and (ii) when both the $\bar{\chi}\chi \rightarrow X_D\phi$ and $\bar{\chi}\chi \rightarrow \phi\phi$ channels are kinematically forbidden.

Crucially, our scenario comfortably evades stringent constraints from CMB observations and indirect detection experiments through two distinct protective mechanisms: the energetic threshold fundamentally shuts off the forbidden $\bar{\chi}\chi \rightarrow X_D\phi$ channel at late times, while the explicit p -wave nature of the loop-induced $\bar{\chi}\chi \rightarrow \phi\phi$ channel renders its late-time cross-section vanishingly small. The benchmark points that satisfy the observed relic density also safely navigate existing direct detection limits, Big Bang Nucleosynthesis bounds, and Higgs invisible decay constraints, while remaining highly testable by near-future, low-threshold experiments.

Furthermore, the FOPT driving this dark sector mechanism produces a stochastic gravitational-wave background spanning the nHz to mHz frequency range, placing the signal within reach of existing pulsar timing arrays (NANOGrav, EPTA, PPTA, IPTA) and future detectors, including SKA, THEIA, μARES , LISA, and DECIGO. We uncover a robust correlation between the DM mass M_χ and the nucleation temperature T_n : maintaining the necessary kinematic threshold ($M_{X_D} + M_\phi > 2M_\chi$) for heavier DM requires a correspondingly larger scalar VEV, thereby elevating the phase-transition temperature scale. This dynamically shifts the GW peak frequency upward for heavier DM candidates. Taken together, this framework demonstrates that a rigorous connection between DM

freeze-out and early-Universe phase transitions can simultaneously address the relic density puzzle and yield rich, mutually reinforcing experimental signatures across multiple observational frontiers.

ACKNOWLEDGMENTS

P. K. P. would like to acknowledge the Ministry of Education, Government of India, for providing financial support for his research via the Prime Minister's Research Fellowship (PMRF) scheme. Parts of this work were developed through discussions among the authors at the *Phoenix 2025* conference, and we thank the organizers for providing a stimulating environment. S. M. acknowledges support from the IIT Goa Startup Grant [2025/SG/SM/057].

DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

APPENDIX A: SM HIGGS AND Φ MIXING

The scalar potential of the model containing the singlet scalar Φ and the SM Higgs doublet H is given by

$$V(H, \Phi) = -\mu_\Phi^2 \Phi^2 - \mu_H^2 (H^\dagger H) + \lambda_H (H^\dagger H)^2 + \lambda_\Phi \Phi^4 + \lambda_{H\Phi} (H^\dagger H) \Phi^2. \quad (\text{A1})$$

Parameterizing the scalars H and Φ as

$$H = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_h + h \end{pmatrix}, \quad \Phi = \frac{\phi + v_\phi}{\sqrt{2}}, \quad (\text{A2})$$

where the VEVs are estimated to be

$$v_h = \sqrt{\frac{2\mu_H^2 - \lambda_{H\Phi} v_\phi^2}{2\lambda_H}}, \quad v_\phi = \sqrt{\frac{2\mu_\Phi^2 - \lambda_{H\Phi} v_h^2}{2\lambda_\Phi}}.$$

The mass-squared matrix spanning the h and ϕ is given by

$$M_{h\phi}^2 = \begin{pmatrix} 2\lambda_H v_h^2 & \lambda_{H\Phi} v_h v_\phi \\ \lambda_{H\Phi} v_h v_\phi & 2\lambda_\Phi v_\phi^2 \end{pmatrix}. \quad (\text{A3})$$

Diagonalizing the above mass matrix, we get the mass eigenstates h_1 with mass $M_{h_1} = 125$ GeV and h_2 with mass M_{h_2} . The $h - \phi$ mixing angle is then given by

$$\tan 2\theta \sim \frac{\lambda_{H\Phi} v_h v_\phi}{\lambda_H v_h^2 - \lambda_\Phi v_\phi^2}. \quad (\text{A4})$$

The scalar couplings can be expressed in terms of the scalar masses and mixing angle as

$$\lambda_H = \frac{\cos^2 \theta M_{h_1}^2 + \sin^2 \theta M_{h_2}^2}{2v_h^2}, \quad (\text{A5})$$

$$\lambda_\Phi = \frac{\cos^2 \theta M_{h_2}^2 + \sin^2 \theta M_{h_1}^2}{2v_\phi^2}, \quad (\text{A6})$$

$$\lambda_{H\Phi} = \cos \theta \sin \theta \frac{M_{h_1}^2 - M_{h_2}^2}{v_h v_\phi}. \quad (\text{A7})$$

APPENDIX B: DM ANNIHILATION AND SCATTERING CROSS-SECTIONS

We used CalcHEP [70] for computing the cross-sections after implementing the model in LanHEP [71]. The thermally averaged cross-section for any process $12 \rightarrow 34$ is given as

$$\begin{aligned} \langle \sigma v \rangle_{12 \rightarrow 34} &= \frac{T}{8m_1^2 m_2^2 K_2(\frac{m_1}{T}) K_2(\frac{m_2}{T})} \\ &\times \int_{\max[(m_1+m_2)^2, (m_3+m_4)^2]}^{\infty} ds \sigma(s) \sqrt{s} K_1\left(\frac{\sqrt{s}}{T}\right) \\ &\times (s - \max[(m_1+m_2)^2, (m_3+m_4)^2]). \quad (\text{B1}) \end{aligned}$$

APPENDIX C: DETAILS OF FIRST-ORDER PHASE TRANSITION

The finite temperature effective potential is given as

$$V_{\text{eff}}(\phi, T) = V_{\text{tree}}(\phi) + V_{\text{CW}}(\phi) + V_{\text{ct}}(\phi) + V_T(\phi, T) + V_{\text{daisy}}(\phi, T). \quad (\text{C1})$$

The tree-level part of the potential is given as

$$V_{\text{tree}} = -\frac{1}{2} \lambda_\Phi v_\phi^2 \phi^2 + \frac{1}{4} \lambda_\Phi \phi^4. \quad (\text{C2})$$

The one-loop Coleman-Weinberg zero-temperature contribution in $\overline{\text{MS}}$ renormalization scheme is given as

$$\begin{aligned} V_{\text{CW}} &= \frac{1}{64\pi^2} m_\phi^4(\phi) \left(\log \left(\frac{m_\phi^2(\phi)}{\mu_R^2} \right) - \frac{3}{2} \right) \\ &+ \frac{3}{64\pi^2} m_{X_D}^4(\phi) \left(\log \left(\frac{m_{X_D}^2(\phi)}{\mu_R^2} \right) - \frac{5}{6} \right), \quad (\text{C3}) \end{aligned}$$

where the field-dependent masses are given as

$$m_\phi(\phi) = \sqrt{-\lambda_\Phi v^2 + 3\lambda_\Phi \phi^2}; \quad M_{Z_D}(\phi) = g_D \phi. \quad (C4)$$

The renormalization scale is fixed at $\mu_R \equiv v_\phi$. The counter-term is

$$V_{\text{ct}} = -\frac{\delta\mu_\Phi^2}{2}\phi^2 + \frac{\delta\lambda_\Phi}{4}\phi^4, \quad (C5)$$

which is derived by solving

$$\left. \frac{\partial(V_{\text{CW}} + V_{\text{ct}})}{\partial\phi} \right|_{\phi=v_\phi} = 0, \quad \left. \frac{\partial^2(V_{\text{CW}} + V_{\text{ct}})}{\partial\phi^2} \right|_{\phi=v_\phi} = 0. \quad (C6)$$

The finite temperature contribution is given as

$$V_T = \frac{T^4}{2\pi^2} J_B\left(\frac{m_\phi(\phi)}{T}\right) + \frac{T^4}{2\pi^2} J_B\left(\frac{m_\eta(\phi)}{T}\right) + 3 \frac{T^4}{2\pi^2} J_B\left(\frac{m_{X_D}(\phi)}{T}\right), \quad (C7)$$

where $m_\eta(\phi) = \sqrt{-\lambda_\phi v^2 + \lambda_\phi \phi^2}$ and

$$J_B(y) = \int_0^\infty x^2 \log(1 - e^{-\sqrt{x^2 + y^2}}). \quad (C8)$$

The daisy contribution is

$$V_{\text{daisy}} = -\frac{g_D^3 T}{12\pi} ((\phi^2 + T^2)^{3/2} - \phi^3). \quad (C9)$$

The bubble nucleation rate per unit volume at a finite temperature is given by

$$\Gamma(T) \simeq \left(\frac{S_3(T)}{2\pi T}\right)^{3/2} T^4 e^{-\frac{S_3(T)}{T}}, \quad (C10)$$

where S_3 is the three-dimensional Euclidean action given as

$$S_3(T) = \int d^3r \left[\frac{1}{2} \left(\frac{d\phi}{dr} \right)^2 + V_{\text{eff}}(\phi, T) \right], \quad (C11)$$

which is calculated by solving

$$\frac{d^2\phi}{dr^2} + \frac{2}{r} \frac{d\phi}{dr} - \frac{dV_{\text{eff}}(\phi, T)}{d\phi} = 0, \quad (C12)$$

with the following boundary conditions:

$$\left. \frac{d\phi}{dr} \right|_{r=0} = 0, \quad \phi(\infty) = \phi_{\text{false}}. \quad (C13)$$

The nucleation process begins when the bubble nucleation rate becomes comparable to the Hubble expansion rate, which is described by

$$\Gamma(T_n) \approx \mathcal{H}(T_n)^4, \quad (C14)$$

where T_n is the nucleation temperature.

APPENDIX D: GRAVITATIONAL WAVES

(i) *Bubble collision*: The peak frequency and amplitude of the GWs generated by the bubble collision are given by

$$f_{\text{peak}}^{\text{col}} = 1.65 \times 10^{-5} \text{ (Hz)} \left(\frac{g_*(T_n)}{100} \right)^{1/6} \left(\frac{T_n}{0.1 \text{ TeV}} \right) \times \frac{0.64}{2\pi} \left(\frac{\beta(T_n)}{\mathcal{H}(T_n)} \right), \quad (D1)$$

$$\Omega_{\text{col}} h^2 = 1.65 \times 10^{-5} \left(\frac{100}{g_*(T_n)} \right)^{1/3} \left(\frac{\mathcal{H}(T_n)}{\beta(T_n)} \right) \left(\frac{\kappa_\phi(T_n) \alpha(T_n)}{1 + \alpha(T_n)} \right)^2 \frac{A(a+b)^c}{\left(b \left(\frac{f}{f_{\text{peak}}^{\text{col}}} \right)^{-a/c} + a \left(\frac{f}{f_{\text{peak}}^{\text{col}}} \right)^{b/c} \right)^c}, \quad (D2)$$

with the efficiency factor

$$\kappa_\phi = \frac{1}{1 + 0.715\alpha(T_n)} \left(0.715\alpha(T_n) + \frac{4}{27} \sqrt{\frac{3\alpha(T_n)}{2}} \right), \quad (D3)$$

and $a = 1.03$, $b = 1.84$, $c = 1.45$, $A = 5.93 \times 10^{-2}$ [35,72,73].

(ii) *Sound waves*: The peak frequency and amplitude of the GWs generated by the sound waves are expressed by [35]

$$f_{\text{peak}}^{\text{sw}} = 8.9 \times 10^{-6} \text{ (Hz)} \left(\frac{g_*(T_n)}{100} \right)^{1/6} \left(\frac{1}{v_w} \right) \times \left(\frac{T_n}{0.1 \text{ TeV}} \right) \left(\frac{z_p}{10} \right) \left(\frac{\beta(T_n)}{\mathcal{H}(T_n)} \right), \quad (D4)$$

where the wall velocity $v_w \sim 1$ and $z_p \sim 10$.

$$\Omega_{\text{sw}} h^2 = 2.59 \times 10^{-6} \left(\frac{100}{g_*(T_n)} \right)^{1/3} \left(\frac{\mathcal{H}(T_n)}{\beta(T_n)} \right) v_w \Upsilon(T_n) \times \left(\frac{\kappa_{\text{sw}}(T_n) \alpha(T_n)}{1 + \alpha(T_n)} \right)^2 \frac{7^{3.5} \left(\frac{f}{f_{\text{peak}}^{\text{sw}}(T_n)} \right)^3}{\left(4 + 3 \left(\frac{f}{f_{\text{peak}}^{\text{col}}(T_n)} \right)^2 \right)^{3.5}}. \quad (D5)$$

The efficiency factor is [74]

$$\kappa_{\text{sw}} = \frac{\alpha(T_n)}{0.73 + 0.083\sqrt{\alpha(T_n)} + \alpha(T_n)}. \quad (\text{D6})$$

The suppression factor is given as

$$\Upsilon = 1 - \frac{1}{\sqrt{1 + 2\tau_{\text{sw}}(T_n)\mathcal{H}(T_n)}}, \quad (\text{D7})$$

where the lifetime of the sound wave is given as [75]

$$\tau_{\text{sw}} = \frac{R_*(T_n)}{U_f(T_n)}, \quad (\text{D8})$$

where mean bubble separation is given by

$$R_* = \frac{(8\pi)^{1/3}v_w}{\beta(T_n)}, \quad (\text{D9})$$

and the rms fluid velocity is given as

$$U_f = \sqrt{\frac{3\kappa_{\text{sw}}(T_n)\alpha(T_n)}{4(1 + \alpha(T_n))}}. \quad (\text{D10})$$

(iii) *Turbulence*: The peak frequency and the amplitude of the GWs generated by the turbulence are expressed by [35,73,76]

$$f_{\text{peak}}^{\text{turb}} = 2.7 \times 10^{-5} \text{ (Hz)} \left(\frac{g_*(T_n)}{100} \right)^{1/6} \left(\frac{1}{v_w} \right) \left(\frac{T_n}{0.1 \text{ TeV}} \right) \left(\frac{\beta(T_n)}{\mathcal{H}(T_n)} \right), \quad (\text{D11})$$

$$\Omega_{\text{turb}} h^2 = 3.35 \times 10^{-4} \left(\frac{100}{g_*(T_n)} \right)^{1/3} \left(\frac{\mathcal{H}(T_n)}{\beta(T_n)} \right) \times \left(\frac{\kappa_{\text{turb}}(T_n)\alpha(T_n)}{1 + \alpha(T_n)} \right)^2 v_w \times \frac{\left(\frac{f}{f_{\text{peak}}^{\text{sw}}(T_n)} \right)^3}{\left(1 + \left(\frac{f}{f_{\text{peak}}^{\text{col}}(T_n)} \right) \right)^{3.6} \left(1 + \frac{8\pi f}{h_*(T_n)} \right)}, \quad (\text{D12})$$

where the efficiency factor is $\kappa_{\text{turb}} = 0.1\kappa_{\text{sw}}$ [76], and the inverse of the Hubble time at the epoch of GW production, redshifted to today, is expressed as

$$h_* = 1.65 \times 10^{-5} \text{ (Hz)} \left(\frac{g_*(T_n)}{100} \right)^{1/6} \left(\frac{T_n}{0.1 \text{ TeV}} \right). \quad (\text{D13})$$

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