

Phase transitions and gravitational waves in a non-Abelian vector dark matter scenario

Nico Benincasa^{1,*}, Luigi Delle Rose^{2,3,†}, Luca Panizzi^{2,3,‡}, Maimoona Razzaq^{2,3,§} and Savio Urzetta^{2,||}

¹*School of Physics, University of Electronic Science and Technology of China, 611731 Chengdu, China*

²*Dipartimento di Fisica, Università della Calabria, Arcavacata di Rende, I-87036, Cosenza, Italy*

³*INFN, Gruppo Collegato di Cosenza, Arcavacata di Rende, I-87036, Cosenza, Italy*



(Received 7 July 2025; accepted 4 October 2025; published 7 November 2025)

We study a scenario where the Standard Model is extended by a $SU(2)$ gauge group in the dark sector. The three associated dark gauge bosons are stabilized via a custodial symmetry triggered by an additional dark $SU(2)$ scalar doublet, thus making them viable dark-matter candidates. After considering the most recent constraints for this model, we analyze the phase transition dynamics and compute the power spectrum of the resulting stochastic gravitational-wave background. Finally, we find regions of the parameter space yielding the observed dark-matter relic density while also leading to a strong enough phase transition with an associated gravitational-wave signal reaching the sensitivity of a future space-based gravitational-wave detector, such as LISA, DECIGO, BBO, TianQin, or Taiji.

DOI: [10.1103/xd9f-bq4p](https://doi.org/10.1103/xd9f-bq4p)

I. INTRODUCTION

The possibility to test models of new physics in complementary ways allows to effectively narrow down the possible theories which aim to solve the open problems of the Standard Model (SM) of particle physics. For many years, this complementarity involved a combination of different astrophysical and collider observations. For a decade, the observation of gravitational waves (GWs) [1,2] added a precious source of information for this kind of phenomenological approach. Indeed, scenarios of new physics usually predict new particles which, throughout the history of the Universe, might have contributed to phenomena leading to the generation of GWs, possibly in the range of observation of current or future experiments. In particular, it has been demonstrated that a first-order phase transition (FOPT) in the early Universe leads to a stochastic gravitational-wave background [3,4] which could be detected today by future space-based gravitational-wave observatories such as LISA [5], DECIGO [6], BBO [7], TianQin [8], or Taiji [9]. Within the frame of the SM, the

electroweak phase transition (PT) is a crossover [10], hence the necessity to add beyond-SM ingredients to make it first order.

In this paper we consider a theoretical model which predicts dark matter (DM) candidates and at the same time can lead to FOPTs, allowing to probe its parameter space through GWs experiments. In this model, introduced in [11] and further studied in [12–31], the DM candidates appear in the form of massive gauge bosons of a non-Abelian $SU(2)$ group in the dark sector. Their stability is ensured by a custodial symmetry in the dark scalar sector which generates gauge boson masses through a mechanism of spontaneous symmetry breaking. The scalar doublet under the new $SU(2)$ can interact with the Higgs doublet via a quartic coupling, and the new physical scalar which is generated after the symmetry breaking mixes with the Higgs boson. In this way, the dark sector can communicate with the SM other than through gravity, potentially leading to observable signals.

Our analysis consists of the following steps. We will first identify the allowed regions of the four-dimensional parameter space of the model by combining theoretical constraints, observables from the Higgs sector of the SM and its precision tests in the electroweak sector, and multiple astrophysical constraints, updating previous analyses using the most recent experimental data. Subsequently, we will study the scalar sector of the model to determine the region of parameter space which can induce FOPTs during cosmological evolution, with the generation of GWs [32]. We will then evaluate whether the GWs can be observable by future experiments, and in which frequency range.

*Contact author: nico.benincasa@kbfi.ee

†Contact author: luigi.dellerose@unical.it

‡Contact author: luca.panizzi@unical.it

§Contact author: maimoona.razzaq@unical.it

||Contact author: sav.urze98@gmail.com

This paper is structured as follows: in Sec. II we provide a short description of the vector DM model from a dark $SU(2)$ gauge group; in Sec. III A the astrophysical and cosmological constraints are described and quantitatively discussed; in Sec. III B the theoretical and collider constraints are analyzed; in Sec. III C the quantum and finite-temperature corrections to the scalar potential as well as the cosmic phase-transition parameters needed to compute the power spectrum of the subsequent generated stochastic gravitational-wave background are provided; in Sec. IV we show the combination of all constraints and the allowed regions of parameter space, and determine the subregions where a first-order phase transition is possible, and finally compare the spectrum of predicted GWs with the reach of different future space-based gravitational-wave detectors.

II. THE $SU(2)_D$ MODEL

For our study, we consider a minimal SM extension, introduced in [11], providing non-Abelian gauge boson DM candidates whose stability is guaranteed by a custodial symmetry built in the new scalar sector, for which we will describe here the main features.

In this model, the SM gauge sector is extended with a new $SU(2)_D$ gauge group, where the D suffix stands for “dark,” under which all SM particles transform as singlets. The three gauge bosons V_D^a acquire mass through a mechanism of spontaneous symmetry breaking, in which a new scalar boson Φ_D , doublet under $SU(2)_D$, acquires a vacuum expectation value (VEV) v_D . The gauge bosons become then degenerate with mass $m_{V_D} = g_D v_D/2$, where g_D is the coupling constant of $SU(2)_D$. The Φ_D is singlet under the SM gauge group, and therefore its only possible interaction with the SM is through a quartic portal coupling with the Higgs boson. The Lagrangian of the new-physics part of the model is

$$\mathcal{L} = -\frac{1}{4}F_{D\mu\nu}F_D^{\mu\nu} + (D_\mu\Phi_D)^\dagger(D^\mu\Phi_D) + \mu_D^2(\Phi_D^\dagger\Phi_D) - \lambda_D(\Phi_D^\dagger\Phi_D)^2 - \lambda_{HD}(\Phi_D^\dagger\Phi_D)(\Phi^\dagger\Phi), \quad (2.1)$$

where $F_D^{\mu\nu}$ is the $SU(2)_D$ field strength; the covariant derivative is $D_\mu = \partial_\mu - ig_D T_D^a V_D^a$, with T_D the dark isospin. This Lagrangian is invariant under an $SO(4)$ symmetry, which is broken to the custodial $SO(3)$ symmetry by the VEV of Φ_D , ensuring the stability of the V_D^a gauge bosons, and making them viable dark matter candidates. An enlarged fermionic sector could break such symmetry, lift the degeneration between the dark vector bosons at loop level, and allow for purely SM decay modes for one of the V_D^a [28,33].

Out of the 8 degrees of freedom in the scalar sector (the components of the Higgs and dark scalar doublets), six are Goldstone bosons which give mass to the electroweak gauge bosons and to the three V_D^a of $SU(2)_D$.

The remaining two correspond to physical states, the Higgs boson h and a further CP -even scalar H_D . These states arise from the mixing between ϕ and ϕ_D , appearing in the neutral components of the Higgs and dark doublet, respectively, $\Phi_{-1/2} = \frac{1}{\sqrt{2}}(\phi + v)$ and $\Phi_{D,-1/2} = \frac{1}{\sqrt{2}}(\phi_D + v_D)$ in unitary gauge, where $v \simeq 246$ GeV is the SM Higgs VEV. The scalar mass matrix in the vacuum (v, v_D) , $\begin{pmatrix} \lambda_H v^2 & \lambda_{HD} v v_D \\ \lambda_{HD} v v_D & \lambda_D v_D^2 \end{pmatrix}$ is diagonalized by the rotation matrix $\begin{pmatrix} \cos\theta_S & \sin\theta_S \\ -\sin\theta_S & \cos\theta_S \end{pmatrix}$, where θ_S is the scalar mixing angle and λ_H are the usual Higgs self-quartic coupling. In the absence of further new particles, mixing between EW and dark gauge sectors is forbidden because the scalar doublets transform under only one of the two gauge groups.

The full model contains four free parameters from the new sector. Following [28], we choose to identify them with the dark gauge coupling, the V_D mass, the H_D mass, and the scalar mixing angle,

$$g_D, m_{V_D}, m_{H_D}, \theta_S, \quad (2.2)$$

which can be easily associated to physical observables. The conversion map between these parameters and the Lagrangian ones is the following:

$$\mu_H^2 = \frac{1}{2} \left(m_h^2 \cos^2\theta_S + m_{H_D}^2 \sin^2\theta_S + \frac{1}{2} \frac{g}{g_D} \frac{m_{V_D}}{m_W} (m_{H_D}^2 - m_h^2) \sin 2\theta_S \right), \quad (2.3)$$

$$\mu_D^2 = \frac{1}{2} \left(m_h^2 \sin^2\theta_S + m_{H_D}^2 \cos^2\theta_S + \frac{1}{2} \frac{g}{g_D} \frac{m_{V_D}}{m_W} (m_{H_D}^2 - m_h^2) \sin 2\theta_S \right), \quad (2.4)$$

$$\lambda_H = \frac{g^2}{8m_W^2} (m_h^2 \cos^2\theta_S + m_{H_D}^2 \sin^2\theta_S), \quad (2.5)$$

$$\lambda_D = \frac{g_D^2}{8m_{V_D}^2} (m_h^2 \sin^2\theta_S + m_{H_D}^2 \cos^2\theta_S), \quad (2.6)$$

$$\lambda_{HD} = \frac{gg_D}{8m_W m_{V_D}} (m_{H_D}^2 - m_h^2) \sin 2\theta_S, \quad (2.7)$$

with the Higgs mass parameter μ_H , the Higgs mass $m_h \simeq 125$ GeV, the $SU(2)_L$ gauge coupling g , and the $W^\pm$ boson mass m_W .

III. TESTING THE MODEL

This model is a minimal extension of the SM, yet it can be constrained from multiple different directions. In the following, we will first consider the theoretical bounds coming from ensuring that the scalar potential is bounded

from below and from perturbative unitarity; then we will discuss cosmological and astrophysical bounds associated with the presence of a dark matter candidate; finally, we will discuss how collider observables constrain modifications in the scalar sector.

From the numerical point of view, we have performed random scans on the four free parameters of the model: m_{V_D} and m_{H_D} logarithmically from 10 MeV to 100 TeV, g_D also logarithmically from 10^{-5} to 4π , and $\cos\theta_S$ linearly from 0 to 1. Given the vast range of the parameter space, we have also performed dedicated scans selecting only points satisfying one or all of the constraints discussed in the following sections, in order to increase the density of valid points for plotting purposes. Furthermore, we will be almost always showing results in the $\{m_{V_D}, m_{H_D}\}$ plane, to allow for a direct comparison of various constraints in the same projection of the parameter space, and also to avoid presenting a too large number of figures, which might disrupt the flow of the discussion. Other projections for further investigation can be easily produced with the information provided below.

A. Cosmological and astrophysical bounds

The model parameters can be constrained by requiring that the relic density of the dark gauge bosons does not exceed the measured value by the Planck experiment [34], with a precision smaller than 1%:

$$\Omega_{\text{DM}} h^2 = 0.120 \pm 0.001, \quad (3.1)$$

with h the dimensionless Hubble constant. The annihilation processes which determine the relic density of the DM candidates in this model are of two types: those involving exclusively SM particles and/or H_D in the final state and those in which two V_D^a annihilate into one V_D^a and a scalar [11]. Representative annihilation processes of the two kinds are shown in Fig. 1.

Furthermore, the model can be tested against experimental data from different direct detection (DD) experiments, considering that the DM states can interact with nuclei through the Higgs and H_D boson.

Both relic density and direct detection constraints have been computed using `micrOMEGAs6` [35], which allows testing the signal against data from multiple experiments, such as LZ [36], PICO-60 [37], XENON1T [38], and PandaX-4T [39], for DM masses ranging from few GeV to 10 TeV, or DarkSide-50 [40] and CRESST-III [41]. These experiments are sensitive to DM masses down to a few hundreds MeV, which is contained in the range of our scan. The annihilation of DM into h and H_D after they concentrate in the center of the Sun and the Earth also allows to set bounds from the neutrinos arising from the subsequent decays of the scalars. The predicted neutrino fluxes can be computed and tested against IceCube (IC) data through `micrOMEGAs` [42].

Indirect detection (ID) constraints are evaluated through MADDM [43] using data from gamma-ray lines and gamma-ray continuum from the Fermi-LAT analysis of dwarf spheroidal galaxies [44] by imposing that the signal is compatible with exclusion limits at 95% CL.

All of these constraints identify different allowed regions in the parameter space, shown in Fig. 2 in the $\{m_{V_D}, m_{H_D}\}$ projection, where in all panels we impose that relic density cannot be overabundant. The region delimited by $m_{V_D} \lesssim m_h/2$ is therefore always excluded by the fact that it predicts a too high relic abundance for any choice of g_D and θ_S .

In the whole remaining parameter space, configurations of parameters which give a DM relic density compatible with the Planck observed value [34] within 3σ can be found. A few regions tend to accumulate solutions in our uniform scan: the areas around $m_h = m_{H_D}$, $m_{V_D} = m_h/2$, and $m_{H_D} = 2m_{V_D}$. All of these areas lead to enhancement in the annihilation or coannihilation cross sections due to interferences and resonance contributions. The constraints from DD, ID, and IC exclude large patches of the parameter space, with the direct detection constraints being by far the strongest. The areas which are not excluded are in all cases the same which feature resonances and large interferences between processes, but also those for which the set of experiments for which data are available do not have sensitivity, such as the region for very small m_{V_D} for the DD constraints. The higher density in the top-right corner

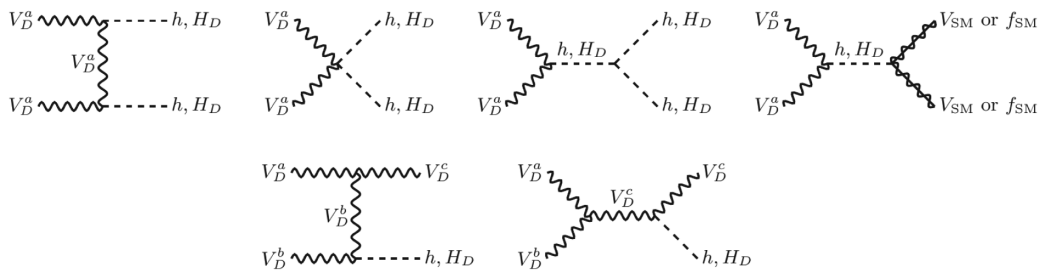


FIG. 1. Annihilation processes, involving SM particles or H_D in the final states (top row) or V_D and scalar particles in the final state (bottom row). V_{SM} and f_{SM} respectively stand for massive vector and fermion SM particles.

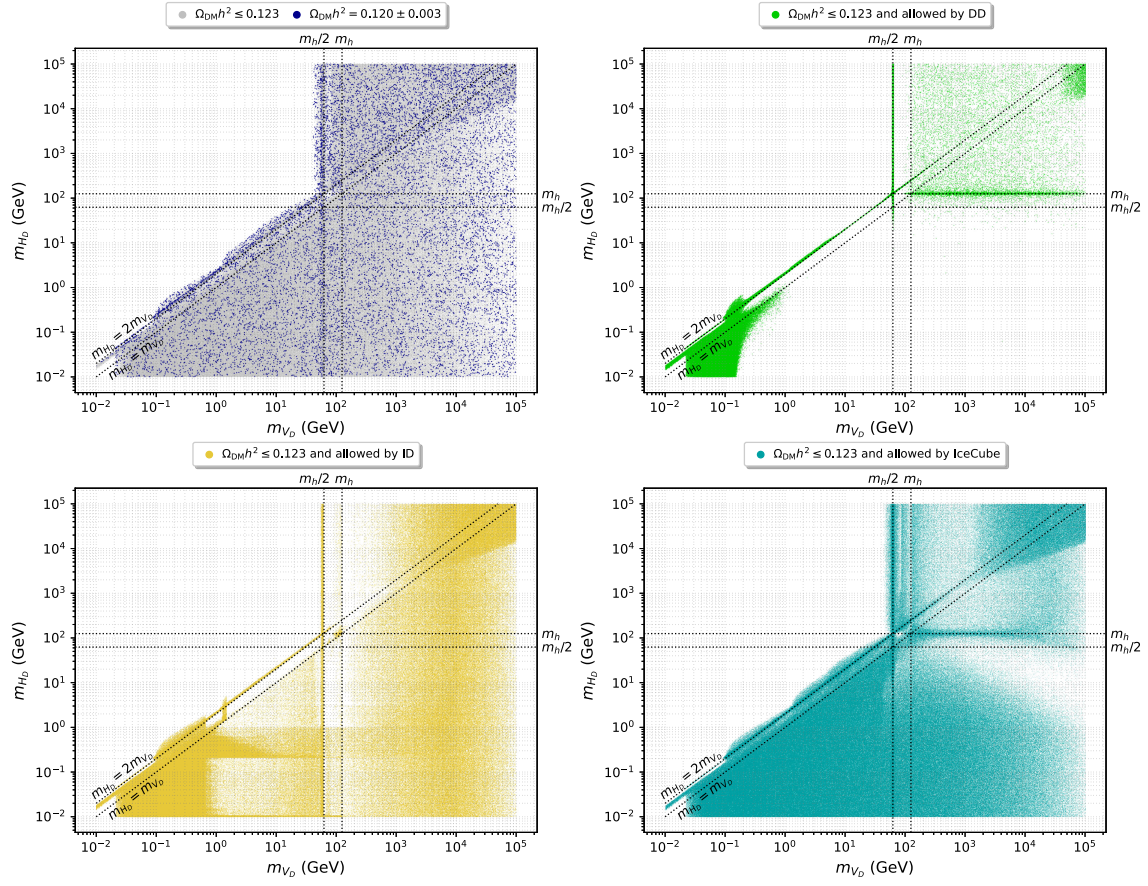


FIG. 2. Areas in the $\{m_{V_D}, m_{H_D}\}$ projection of the parameter space which are allowed by astrophysical and cosmological constraints: underabundant and exact relic density (top left), underabundant and allowed by direct detection experiments (top right), underabundant and allowed by indirect detection from Fermi-LAT (bottom left), and underabundant and allowed by neutrino fluxes measured at IceCube (bottom right).

of all panels is due to the fact that the region with high V_D and H_D masses leads to multiple contributions to the reduction of the DM relic density and all the considered experiments lose sensitivity in that region of the mass parameter space.

It is illustrative to describe in the following some of the main features which lead to patches of exclusions for the different observables.

Considering direct detection, for sufficiently small mixing angle θ_S , the contribution from the diagram involving SM states in Fig. 1 is suppressed. DM annihilation and DD cross section are thus uncorrelated in that limit. Then, in principle, for small enough mixing angle, it should be possible to find points yielding the observed relic abundance while evading DD constraints for the whole region $m_{H_D} < m_{V_D}$, including the empty region $m_{H_D} < \min(m_{V_D}, m_h/2)$ in the upper right panel of Fig. 2 [45]. However, for such a low mixing, the H_D decay into SM particles is more and more suppressed and the relic abundance quickly increases above the measured experimental value. For the same reason, if we just consider the allowed points from DD, when g_D becomes small, the cross

section would be too small almost anywhere to be detected, with the exception of a region where the V_D mass is between ~ 10 GeV and $m_h/2$, and H_D lighter than ~ 2 GeV, where the sensitivity of LZ peaks [36] and the lower limit on our g_D scan is still giving a large enough cross section which is excluded. However, by combining DD limits with the requirement of having underabundant relic density, all of the points in the light H_D area of the plot become excluded because the coupling needed to achieve the correct relic density would give a too high cross section for direct detection, with the exception of the region with very light DM which is outside the sensitivity of all experiments, and for the region with high V_D and H_D masses (top-right area of the plot).

For what concerns indirect detection, below $m_{V_D} \sim 500$ MeV the photon emitted by DM annihilation is too soft regardless of the channel, Fermi-LAT loses all sensitivity, and all points are allowed. To understand the feature appearing when m_{H_D} is between 200 MeV and 1 GeV is physical, we refer to the first panel, third row, of Figs. 3 and 4 in [43], where the photon spectrum for different annihilation channels is computed for two different DM masses.

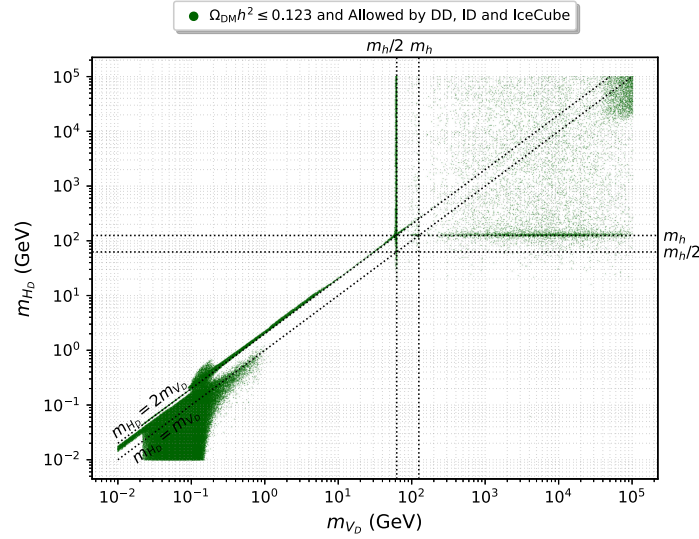


FIG. 3. Combination of all astrophysical and cosmological constraints in the $\{m_{V_D}, m_{H_D}\}$ projection.

Below 200 MeV, the main decay channels are into up and down quarks and, more suppressed, into electrons; while the hadronic decays will lead to too soft photons because it is kinematically not possible to produce pions at tree level, the suppressed electron decays can produce enough photons only if the coupling strength is large enough, which explains the low number of allowed points in that region, especially for light DM masses. Above $m_{H_D} \sim 200$ MeV, two channels open for H_D decays: muons and strange quarks. The strange quarks will produce photons through hadronization and pion production, but their energy is too low to be detected. For muons, the number of photons arising from this channel is much lower than other channels and, as the DM mass decreases, the sensitivity of Fermi-LAT also decreases. As m_{H_D} increases various kinematic thresholds for hadron production are crossed, leading to a large number of photons and a larger exclusion rate.

Let us finally move to the constraints from IceCube. The sensitivity of the experiment is limited to the neutrino energy range 39 GeV to 4×10^5 GeV (see [42] and references therein). Indeed, we have verified that below a DM mass of 39 GeV there is no exclusion from this observable. The regions with lower density of allowed points in the higher m_{V_D} region are due to an interplay of effects which enter the neutrino flux at Earth, described by Eq. (4) in [42]: they involve the capture rate in the Sun, the annihilation cross section of the DM in the sun, and the BRs of this process into SM final states (which depend on the mass of H_D); for example, decays into W^+W^- , ZZ , and $t\bar{t}$, when allowed, produce neutrinos with high energy within the sensitivity reach of IceCube, and the masses of V_D and H_D determine the annihilation cross section, and therefore the number of produced neutrinos.

The combination of astrophysical and cosmological bounds, shown in Fig. 3, leaves only the resonant areas,

the area delimited by $m_{V_D} \gtrsim m_h$ and $m_{H_D} \gtrsim m_h/2$, and the region with $m_{V_D} \lesssim 100$ MeV where DD becomes insensitive.

B. Theoretical and collider constraints

Requiring that the scalar potential is bounded from below leads to specific conditions on the couplings of the scalar fields defined in Eq. (2.1):

$$\lambda_H, \lambda_D \geq 0, \quad \lambda_{HD} \geq -2\sqrt{\lambda_H \lambda_D}. \quad (3.2)$$

Perturbative unitarity (PU) requires that the scattering amplitudes for two-particle processes do not grow uncontrollably with energy. This imposes constraints on the combinations of coupling constants derived from the eigenvalues of the scattering matrix [46]:

$$\text{PU} = |(a_+, a_-, b_+, b_-, \lambda_{HD})| < 8\pi, \quad (3.3)$$

where the parameters a_+ , a_- , b_+ , and b_- are defined in terms of the couplings of the Higgs fields as follows:

$$\begin{aligned} a_{\pm} &= 3(\lambda_H + \lambda_D) \pm \sqrt{9(\lambda_H - \lambda_D)^2 + 4\lambda_{HD}^2}, \\ b_{\pm} &= \lambda_H + \lambda_D \pm |\lambda_H - \lambda_D|. \end{aligned} \quad (3.4)$$

Electroweak precision observables (EWPO) serve as a valuable method for parametrizing the impact of new physics on electroweak measurements. These observables are typically represented by the Peskin-Takeuchi parameters S , T , and U [47,48]. The contributions of our model to the S , T , and U parameters are [49]

$$\bar{T} = \frac{3}{16\pi s_W^2} \left[\cos^2 \theta_S \left(f_1 \left(\frac{m_h^2}{m_W^2} \right) - \frac{1}{c_W^2} f_1 \left(\frac{m_h^2}{m_Z^2} \right) \right) + \sin^2 \theta_S \left(f_1 \left(\frac{m_{H_D}^2}{m_W^2} \right) - \frac{1}{c_W^2} f_1 \left(\frac{m_{H_D}^2}{m_Z^2} \right) \right) - \left(f_1 \left(\frac{m_h^2}{m_W^2} \right) - \frac{1}{c_W^2} f_1 \left(\frac{m_h^2}{m_Z^2} \right) \right) \right], \quad (3.5)$$

$$\bar{S} = \frac{1}{2\pi} \left[\cos^2 \theta_S f_2 \left(\frac{m_h^2}{m_W^2} \right) + \sin^2 \theta_S f_2 \left(\frac{m_{H_D}^2}{m_W^2} \right) - f_2 \left(\frac{m_h^2}{m_Z^2} \right) \right], \quad (3.6)$$

$$\bar{U} = \frac{1}{2\pi} \left[\cos^2 \theta_S f_2 \left(\frac{m_h^2}{m_W^2} \right) + \sin^2 \theta_S f_2 \left(\frac{m_{H_D}^2}{m_W^2} \right) - f_2 \left(\frac{m_h^2}{m_Z^2} \right) \right] - \bar{S}, \quad (3.7)$$

where $f_1(x)$ and $f_2(x)$ are loop functions defined by [50]:

$$f_1(x) = \frac{x \log x}{x-1}, \quad (3.8)$$

$$f_2(x) = \begin{cases} \frac{1}{12} \left[-2x^2 + 9x + (x-3)(x^2 - 4x + 12) + \frac{1-x}{x} f_1(x) \right. \\ \left. + 2\sqrt{(4-x)x}(x^2 - 4x + 12) \arctan \left(\sqrt{\frac{4-x}{x}} \right) \right], & 0 < x < 4, \\ \frac{1}{12} \left[-2x^2 + 9x + (x-3)(x^2 - 4x + 12) + \frac{1-x}{x} f_1(x) \right. \\ \left. + \sqrt{(x-4)x}(x^2 - 4x + 12) \log \left(\frac{x - \sqrt{(x-4)x}}{x + \sqrt{(x-4)x}} \right) \right], & x \geq 4. \end{cases} \quad (3.9)$$

To constrain the parameter space of our model, the total values of S , T , and U must be within the confidence level of 95% with respect to the measured ones [51].

Finally, we checked the consistency of the parameters of our model by exploiting constraints from the Large Hadron Collider (LHC). This has been done with the HiggsTools software (HiggsBounds and HiggsSignals) [52], which enforces bounds from direct searches for scalar particles, as well as the compatibility of the predictions of the Higgs coupling modification in our model with the corresponding measurements at the LHC. The scalar mixing angle θ_S is strongly constrained by the signal strength measurement to $|\theta| \lesssim O(0.1)$, suppressing large deviations from SM-like behavior, unless the new scalar is highly degenerate with the SM Higgs.

Figure 4 displays the allowed regions in the m_{V_D}, m_{H_D} plane after imposing individual theoretical and collider constraints on our model parameter space, where one can see that PU and EWPO strongly limit the highest values achievable for the H_D mass [$m_{H_D} \lesssim O(10^4 \text{ GeV})$].

It is illustrative here to describe the horizontal bands appearing in the HiggsBounds results, which might look as artifacts, while in fact they represent well identified physical regions. These bands are horizontal because the dominant parameters for the determination of the collider constraints on the production and decay of the scalars of the

theory are the H_D mass and the scalar mixing angle, while the DM mass plays a subdominant role. Such bands are then related to the sensitivity regions of the various searches (both LEP and LHC) implemented in HiggsBounds and tend to exclude more strongly a lighter H_D , especially if the Higgs boson can decay into it, and allow heavier H_D masses. The DM mass plays a more relevant role for searches looking for invisible decays, visible in this projection as different point densities for fixed H_D mass, and depending on the search there can be sensitivity in the region where $m_{V_D} < m_h/2$, where the Higgs boson has a further invisible decay which can be excluded by dedicated LHC searches (see, e.g., [53,54]).

Figure 5 presents the combined effect of all theoretical and collider constraints, showing the surviving parameter points in green on the (m_{V_D}, m_{H_D}) plane (left) and on the $(m_H, \cos \theta_S)$ plane (right). In the latter it is clearly visible how the mixing angle is tightly constrained unless H_D is almost degenerate with the Higgs boson, and the sensitivity areas of the aforementioned collider searches.

C. Phase transitions and gravitational waves

1. Loop corrections to scalar potential

The tree-level potential is corrected at one-loop order in perturbation theory by the Coleman-Weinberg potential,

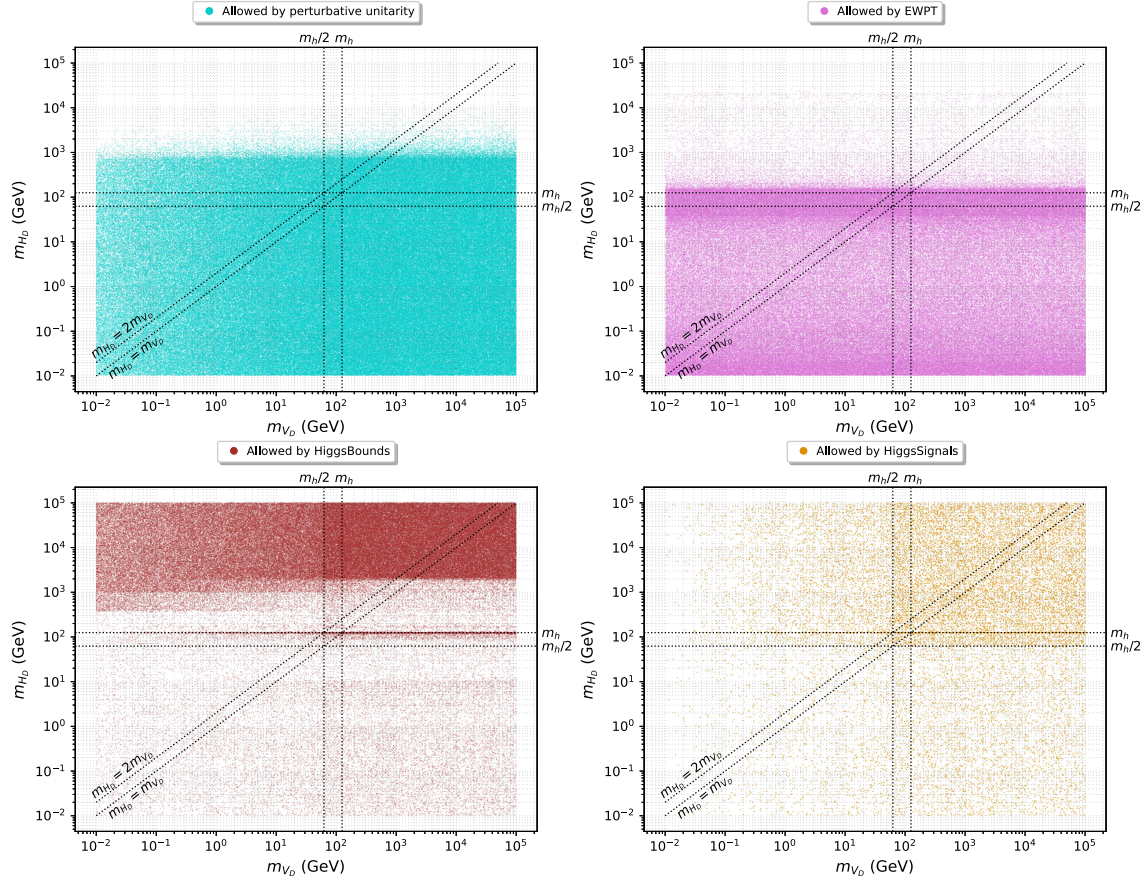


FIG. 4. Allowed areas in the $\{m_{V_D}, m_{H_D}\}$ projection of the $SU(2)_D$ model parameter space allowed by theoretical and collider constraints. From top left to bottom right: perturbative unitarity, EW precision tests, HiggsBounds and HiggsSignals.

which in the Landau gauge and in the $\overline{\text{MS}}$ renormalization scheme reads as [55]

$$V_1(\phi, \phi_D) = \frac{1}{64\pi^2} \sum_i n_i m_i^4 \left(\ln \frac{m_i^2}{\mu^2} - C_i \right), \quad (3.10)$$

where the sum ranges over the fields $i = W^\pm, Z, t, \phi, \phi_D, G^j, G_D^j, V_D, j \in \{0, +, -\}$, with G and G_D the SM and dark Goldstone bosons respectively (as usual, we retain only the top quark contribution and neglect all of the other lighter fermions). The numbers of bosonic and fermionic

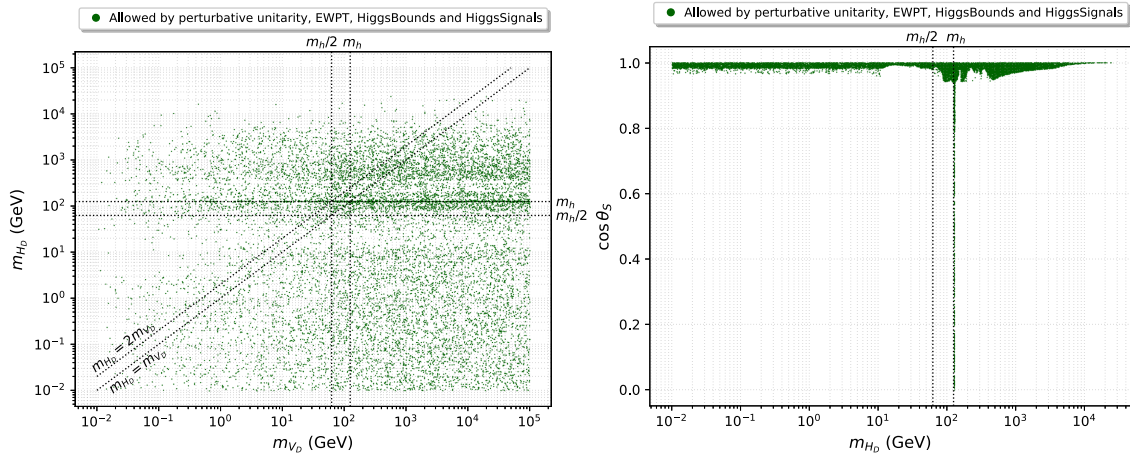


FIG. 5. Combination of all collider and theoretical constraints in the $\{m_{V_D}, m_{H_D}\}$ (left) and $\{m_{H_D}, \cos \theta_S\}$ (right) projections.

degrees of freedom are counted by n_i and explicitly given by $n_{W^\pm} = 6, n_Z = 3, n_{V_D} = 9, n_t = -12, n_\phi = n_{\phi_D} = n_{G^i} = n_{G_D^i} = 1$. The constants C_i are peculiar to the renormalization scheme and in the $\overline{\text{MS}}$ are $C_i = 3/2$ for scalars, fermions, and $C_i = 5/6$ for vector bosons. Finally, the renormalization scale can be chosen to be $\mu = v$.

The field-dependent masses in the scalar sector correspond to the eigenvalues of the mass matrix,

$$M_S^2 = \begin{pmatrix} -\mu_H^2 + 3\lambda_H\phi^2 + \frac{\lambda_{HD}}{2}\phi_D^2 & \lambda_{HD}\phi\phi_D \\ \lambda_{HD}\phi\phi_D & -\mu_D^2 + 3\lambda_D\phi_D^2 + \frac{\lambda_{HD}}{2}\phi^2 \end{pmatrix}, \quad (3.11)$$

and to the Goldstone field-dependent masses,

$$\begin{aligned} m_G^2 &= -\mu_H^2 + \lambda_H\phi^2 + \frac{\lambda_{HD}}{2}\phi_D^2, \\ m_{G_D}^2 &= -\mu_D^2 + \lambda_D\phi_D^2 + \frac{\lambda_{HD}}{2}\phi^2, \end{aligned} \quad (3.12)$$

while the field-dependent top quark and gauge boson masses are given by

$$\begin{aligned} m_t^2 &= \frac{y_t^2}{2}\phi^2, & m_{W^\pm}^2 &= \frac{g^2}{4}\phi^2, \\ m_Z^2 &= \frac{g^2 + g_D^2}{4}\phi^2, & m_{V_D}^2 &= \frac{g_D^2}{4}\phi_D^2, \end{aligned} \quad (3.13)$$

where y_t , g , and g_D are the top Yukawa coupling, the $U(1)_Y$, the $SU(2)_L$, and the $SU(2)_D$ gauge couplings, respectively.

To change the renormalization prescription one can add finite counterterms to the Coleman-Weinberg potential. Indeed, it is often convenient to directly use physical parameters as input values. For this purpose, one can enforce the VEVs, the scalar masses, and the portal coupling to their tree-level values by including a counterterm potential,

$$\begin{aligned} V_{\text{ct}}(\phi, \phi_D) &= \delta\mu_H^2\phi^2 + \delta\lambda_H\phi^4 + \delta\mu_D^2\phi_D^2 + \delta\lambda_D\phi_D^4 \\ &+ \delta\lambda_{HD}\phi^2\phi_D^2, \end{aligned} \quad (3.14)$$

in which the coefficients are fixed by the following renormalization conditions:

$$\left. \frac{\partial V_{\text{ct}}}{\partial \phi} \right|_{\text{VEV}} = - \left. \frac{\partial V_1}{\partial \phi} \right|_{\text{VEV}}, \quad (3.15)$$

$$\left. \frac{\partial V_{\text{ct}}}{\partial \phi_D} \right|_{\text{VEV}} = - \left. \frac{\partial V_1}{\partial \phi_D} \right|_{\text{VEV}}, \quad (3.16)$$

$$\begin{aligned} \left. \frac{\partial^2 V_{\text{ct}}}{\partial \phi^2} \right|_{\text{VEV}} &= - \left(\left. \frac{\partial^2 V_1}{\partial \phi^2} \right|_{G^i=0} + \frac{1}{32\pi^2} \sum_i \left(\frac{\partial m_{G^i}^2}{\partial \phi} \right)^2 \ln \frac{m_{\phi_{\text{IR}}}^2}{\mu^2} \right) \Big|_{\text{VEV}}, \end{aligned} \quad (3.17)$$

$$\begin{aligned} \left. \frac{\partial^2 V_{\text{ct}}}{\partial \phi_D^2} \right|_{\text{VEV}} &= - \left(\left. \frac{\partial^2 V_1}{\partial \phi_D^2} \right|_{G_D^i=0} + \frac{1}{32\pi^2} \sum_i \left(\frac{\partial m_{G_D^i}^2}{\partial \phi_D} \right)^2 \ln \frac{m_{\phi_{\text{IR}}}^2}{\mu^2} \right) \Big|_{\text{VEV}}, \end{aligned} \quad (3.18)$$

$$\left. \frac{\partial^2 V_{\text{ct}}}{\partial \phi \partial \phi_D} \right|_{\text{VEV}} = - \left. \frac{\partial^2 V_1}{\partial \phi \partial \phi_D} \right|_{\text{VEV}}, \quad (3.19)$$

where VEV corresponds to the vacuum $(\phi, \phi_D) = (v, v_D)$ and the infrared regulators $m_{\phi_{\text{IR}}}^2$ and $m_{\phi_D_{\text{IR}}}^2$ prevent the singularities arising from the Goldstone boson contributions. Following [56], we adopt the choice $m_{\phi_{\text{IR}}}^2 = m_h^2$ and $m_{\phi_D_{\text{IR}}}^2 = m_{H_D}^2$. The counterterms are explicitly given by

$$\delta\lambda_H = \frac{1}{8v^3} (\partial_\phi V_1 - \partial_{\phi^2}^2 V_1 v) \Big|_{\text{VEV}}, \quad (3.20)$$

$$\delta\lambda_D = \frac{1}{8v_D^3} (\partial_{\phi_D} V_1 - \partial_{\phi_D^2}^2 V_1 v_D) \Big|_{\text{VEV}}, \quad (3.21)$$

$$\delta\lambda_{HD} = - \left. \frac{\partial_{\phi\phi_D}^2 V_1}{4vv_D} \right|_{\text{VEV}}, \quad (3.22)$$

$$\delta\mu_H^2 = \frac{1}{4v} (\partial_{\phi^2}^2 V_1 v - 3\partial_\phi V_1 + \partial_{\phi\phi_D}^2 V_1 v_D) \Big|_{\text{VEV}}, \quad (3.23)$$

$$\delta\mu_D^2 = \frac{1}{4v_D} (\partial_{\phi_D^2}^2 V_1 v_D - 3\partial_{\phi_D} V_1 + \partial_{\phi\phi_D}^2 V_1 v) \Big|_{\text{VEV}}, \quad (3.24)$$

where the derivatives are evaluated at the tree-level vacuum.

The one-loop finite-temperature contributions to the effective potential are given by [57]

$$V_T(\phi, \phi_D) = \frac{T^4}{2\pi} \sum_i n_i J_{\text{B/F}} \left(\frac{m_i^2}{T^2} \right), \quad (3.25)$$

where the thermal bosonic and fermionic functions $J_{\text{B/F}}$ are defined as

$$J_{\text{B/F}}(y^2) = \int_0^\infty dx x^2 \ln \left(1 \mp e^{-\sqrt{x^2 + y^2}} \right). \quad (3.26)$$

It is customary to deal with IR divergences arising from the zero bosonic Matsubara mode in the high-temperature limit by considering the so-called daisy resummation [58].

This procedure amounts to adding the leading-order thermal correction or Debye mass $\Pi_i T^2$ to the field-dependent mass m_i^2 in V_1 and V_T [59]. For the scalar sector it consists in the shifts

$$\begin{aligned}\mu_H^2 &\rightarrow \mu_H^2 + c_H T^2, \\ \mu_D^2 &\rightarrow \mu_D^2 + c_D T^2,\end{aligned}\quad (3.27)$$

where

$$c_H = \frac{1}{16}(g'^2 + 3g^2 + 4y_t^2) + \frac{\lambda_H}{2} + \frac{\lambda_{HD}}{6}, \quad (3.28)$$

$$c_D = \frac{3}{16}g_D^2 + \frac{\lambda_D}{2} + \frac{\lambda_{HD}}{6}, \quad (3.29)$$

with the SM contribution taken from [60].

The thermal corrections also affect the longitudinal components of the gauge bosons. The mass of the W_L is simply obtained by shifting $m_W^2(\phi)$ with the Debye mass $\Pi_{W_L} T^2 = \frac{11}{6}g^2 T^2$, while those of the neutral EW bosons, Z_L and γ_L , are given by

$$m_{Z_L}^2 = \frac{1}{2} \left[m_Z^2(\phi) + \frac{11}{6} \frac{g^2}{\cos^2 \theta_W} T^2 + \Delta \right], \quad (3.30)$$

$$m_{\gamma_L}^2 = \frac{1}{2} \left[m_Z^2(\phi) + \frac{11}{6} \frac{g^2}{\cos^2 \theta_W} T^2 - \Delta \right] \quad (3.31)$$

with

$$\Delta^2 = \left(\frac{\phi^2}{4} + \frac{11}{6} T^2 \right)^2 (g'^2 - g^2)^2 + \frac{g'^2 g^2}{4} \phi^4. \quad (3.32)$$

These relations above are obtained from the diagonalization of the thermally corrected mass matrix of the (W_3, B) gauge fields. Finally, the contribution to the longitudinal part of the dark gauge boson V_{D_L} is $\Pi_{V_{D_L}} T^2 = \frac{5}{6}g_D^2 T^2$ [22].

2. The phase transition

A couple of parameters capture most of the information related to a first-order phase transition which proceeds through the nucleation of bubbles of true vacuum. The strength of the transition is characterized by the vacuum energy $\Delta\epsilon$ released during the nucleation process. The corresponding dimensionless parameter, α , is obtained by normalizing this vacuum energy to the energy density of the plasma in the symmetric phase ρ_{rad} [61]:

$$\alpha \equiv \frac{\Delta\epsilon}{\rho_{\text{rad}}}\bigg|_{T=T_*}, \quad \Delta\epsilon \equiv \epsilon|_{\text{false vacuum}} - \epsilon|_{\text{true vacuum}} \quad (3.33)$$

with $\epsilon = V_{\text{eff}} - \frac{T}{4} \frac{\partial V_{\text{eff}}}{\partial T}$, T_* the temperature at which GWs are generated ($T_* \simeq T_n$, where T_n is the nucleation temperature), and $\rho_{\text{rad}} = \frac{\pi^2}{30} g_{\text{eff}}(T) T^4$, where $g_{\text{eff}}(T)$ is the effective number of relativistic degrees of freedom at the

temperature T . The inverse time duration of the phase transition, β , is usually normalized by the Hubble parameter $H_* \equiv H(T_*)$ [62]:

$$\frac{\beta}{H_*} = T \frac{d(S_3/T)}{dT} \bigg|_{T=T_*}, \quad (3.34)$$

where S_3 is the three-dimensional Euclidean action computed for an $O(3)$ -symmetric critical bubble. A large β/H_* means that we can safely neglect the Hubble expansion of the Universe during the phase-transition process.

3. The gravitational waves

A single bubble cannot generate GWs, because its quadrupole moment vanishes due to its spherical symmetry. One of the key ingredients for a stochastic GW background to be generated is the collisions between these bubbles of true vacuum which clearly break the aforementioned spherical symmetry [63]. For bubbles evolving in a thermal plasma, additional contributions to the GW signal come from the metric perturbation induced by the perturbed plasma: sound waves [64,65], and magnetohydrodynamic (MHD) turbulence [66]. In the former case, subsonic and supersonic bubbles, respectively, create compression waves beyond the bubble wall and rarefaction waves behind it. When acoustic waves from different bubbles overlap, the induced shear stress of the metric generates gravitational waves. Finally, bubble collision is also responsible for the turbulent motion of the fully ionized plasma and these MHD turbulences source a GW background.

The power spectrum $h^2 \Omega_{\text{GW}}$ of the stochastic GW background from a cosmic phase transition thus arises from three contributions [67]: $h^2 \Omega_{\text{GW}} \simeq h^2 \Omega_{\text{col}} + h^2 \Omega_{\text{sw}} + h^2 \Omega_{\text{turb}}$. In our model, we find that the dominant contribution is given by sound waves [68–70]:

$$h^2 \Omega_{\text{GW}}(f) \simeq h^2 \Omega_{\text{sw}}(f) = h^2 \Omega_{\text{sw}}^{\text{peak}} S_{\text{sw}}(f), \quad (3.35)$$

with

$$\begin{aligned}h^2 \Omega_{\text{sw}}^{\text{peak}} &= 1.23 \times 10^{-6} \left(\frac{H_*}{\beta} \right) \left(\frac{\kappa_{\text{sw}} \alpha}{1 + \alpha} \right)^2 \left(\frac{100}{g_*} \right)^{1/3} v_w \Upsilon, \\ S_{\text{sw}}(f) &= \left(\frac{f}{f_{\text{sw}}} \right)^3 \left(\frac{7}{4 + 3(f/f_{\text{sw}})^2} \right)^{7/2},\end{aligned}\quad (3.36)$$

where κ_{sw} is the efficiency factor for the conversion of the vacuum energy into the bulk motion of the plasma and f_{sw} is the sound-wave peak frequency.¹ The suppression factor for a radiation-dominated Universe is defined as [73,74]

¹Note that the perturbative treatment of cosmic phase transitions may suffer from uncertainties due to gauge or renormalization-scale dependence that could potentially lower the resulting GW signal strength [71,72].

$$\Upsilon = 1 - \frac{1}{\sqrt{2\tau_{\text{sh}}H + 1}}, \quad (3.37)$$

where τ_{sh} is the time after which sound waves do not source the GW production and where $\Upsilon \simeq \tau_{\text{sh}}H$ when $\tau_{\text{sh}}H \ll 1$, thus when this time is much smaller than a Hubble time.

IV. COMBINED RESULTS AND DISCUSSION

The combination of all theoretical, cosmological, astrophysical, and collider bounds is presented in the $\{m_{V_D}, m_{H_D}\}$ projection of the parameter space in Fig. 6, where we show both points which lead to an underabundant DM relic density (meaning that other DM components must exist) and those for which the gauge boson DM of the $SU(2)_D$ model is sufficient to explain the entire observed abundance. Notice that the difference in the density of allowed but just underabundant points and allowed points with exact relic density is just meant to represent graphically the proportion of the latter with respect to the former. Indeed, by (largely) increasing the number of scanned points, it would be possible to fill the allowed areas with points with exact relic density, but this would be computationally extremely expensive and beyond the scope of this analysis.

The distribution of points precisely identifies specific regions of the parameter space: they span a range delimited by $m_h/2 \lesssim m_{H_D} \lesssim m_{V_D}$ and cluster on three narrow regions characterized by $m_{H_D} \simeq m_h$ or $m_{V_D} \simeq m_h/2$ or $m_{H_D} \simeq 2m_{V_D}$. The $m_{V_D} \simeq m_h/2$ or $m_{H_D} \simeq 2m_{V_D}$ values make the DM annihilation processes with a resonant scalar in the s -channel effective, thus reducing the relic density even

for small g_D (within the scanned range). When the two scalars are almost degenerate, on the other hand, their mixing can be large, and therefore both scalars can decay to the same channels as the SM Higgs boson.

Next, we analyze the phase transition dynamics for each point in Fig. 6 (in addition to a refined scan, see below) via the COSMOTRANSITIONS package [75]. Our results are shown in Fig. 7, which is a projection of the scanned parameter space. It indicates which values of the DM mass m_{V_D} and the dark Higgs mass m_{H_D} lead to a strong FOPT. We can identify three main regions, as in Fig. 6. One corresponds to $m_{V_D} \simeq m_h/2$, while another one is identified by the condition $m_{V_D} \simeq m_{H_D}/2$. Points with symmetry nonrestoration are found in the latter region: the minimum of the potential is at $(0, v'_D)$ at high temperature, instead of $(0,0)$ because the coefficient c_D in Eq. (3.29) is negative, therefore preventing the symmetry to be restored as the temperature increases. Finally, the last region is broader and corresponds to $m_{V_D} \in [4 \times 10^2, 10^3]$ GeV, with $|m_{H_D} - m_h| \lesssim m_h/2$ GeV. We find that the strong FOPTs are obtained in this region and that some of these points lead to the observed DM relic density. We then make a refined scan in this strong FOPT region to find more such points. The combined results of the general and refined scan are shown in the upper-left panel in Fig. 7. We also show each PT step contribution in the remaining panels. We then immediately see that the strongest phase transitions are obtained for the PT step $\phi \rightarrow \phi\phi_D$, from the double-step FOPT $O \rightarrow \phi \rightarrow \phi\phi_D$, where in this notation O corresponds to the origin $(0,0)$, ϕ corresponds to a phase along the ϕ axis,

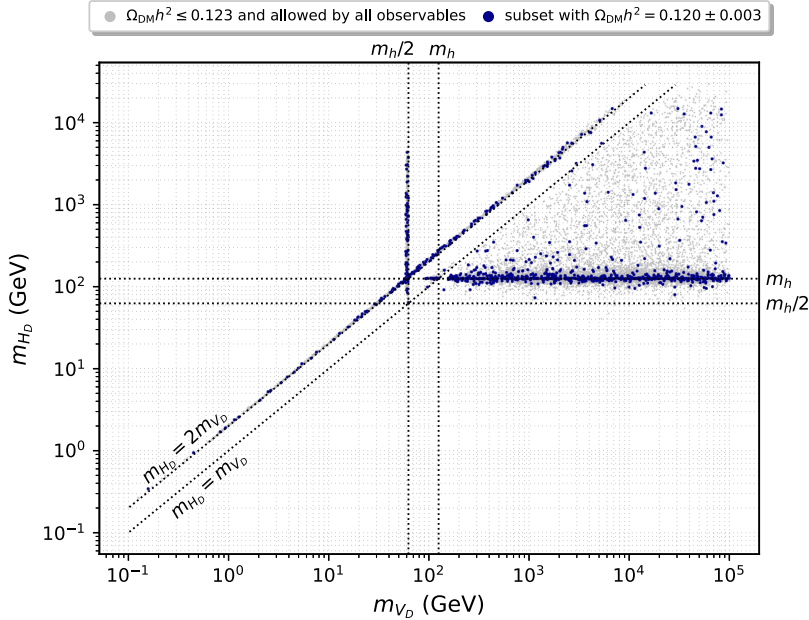


FIG. 6. Points in the $\{m_{V_D}, m_{H_D}\}$ projection of the parameter space with underabundant relic density and satisfying all theoretical, collider, and astrophysical constraints. Highlighted in dark blue color are the points which reconstruct the measured relic density within 3σ .

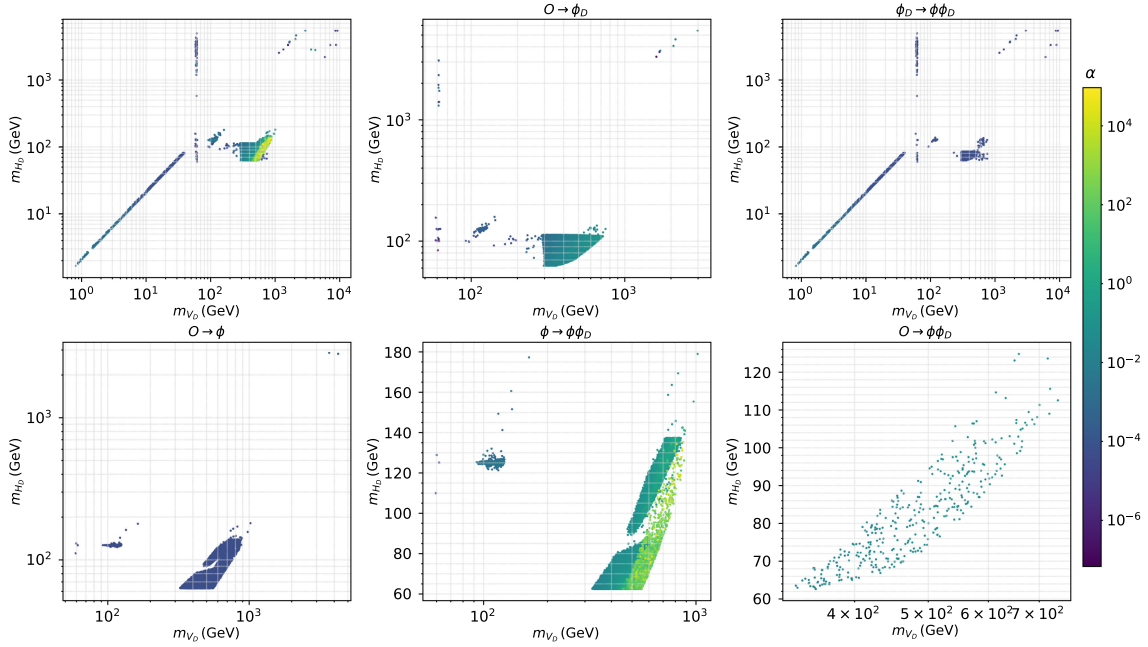


FIG. 7. Projection of the parameter space on the plane spanned by m_{V_D} and m_{H_D} . The upper-left panel shows the full results of our scan, while the other panels only show the results for a specific PT step. The color code, for each panel, represents the PT strength α .

and $(v', 0)$ and $\phi\phi_D$ correspond to a phase in the plane, (v'', v_D') . The reason why, although this strong FOPT region contains points with different PT patterns ($O \rightarrow \phi_D \rightarrow \phi\phi_D$, $O \rightarrow \phi \rightarrow \phi\phi_D$, and $O \rightarrow \phi\phi_D$, only the second one yields $\alpha \geq 1$), is the following: This is because, as in [76], PT starting from the origin are found to be weaker. Moreover, the transition $\phi_D \rightarrow \phi\phi_D$ cannot be arbitrarily strong as the value of the VEV along the ϕ axis is bounded by $v \simeq 246$ GeV. However, for the transition $\phi \rightarrow \phi\phi_D$, a large $v_D(T_n)$ necessarily leads to a large α , as the latter is proportional to $v_D(T_n)/T_n$, with $T_n \sim 10^2$ GeV. Indeed these strong FOPTs come from the point with $v_D(T=0) \in [10^3, 2 \times 10^3]$ GeV. Finally, note that the two points in the upper-right corner of the lower-left panel

give rise to a PT step along the direction $\phi \rightarrow \phi_D$. Actually the full pattern for these two points is a triple-step FOPT $O \rightarrow \phi \rightarrow \phi_D \rightarrow \phi\phi_D$, from which the PT step $\phi \rightarrow \phi_D$ yields a strong enough GW to be detected. However, these two points lead to an underabundant DM, $h^2\Omega_{\text{DM}} \sim 0.01$.

Next, in the left panel in Fig. 8, we show the correlation between the phase-transition parameters α and β/H_* . As expected, we obtain that the slower the phase transition (smaller β/H_*), the stronger (larger α) it is. Moreover, we find that in the present model, the DM abundance is not correlated to the PT strength, as we obtain the correct DM relic density for a large range of values of α . In the right panel, the relevant temperatures for the study of phase transitions, namely T_c and T_n , are displayed, with the color

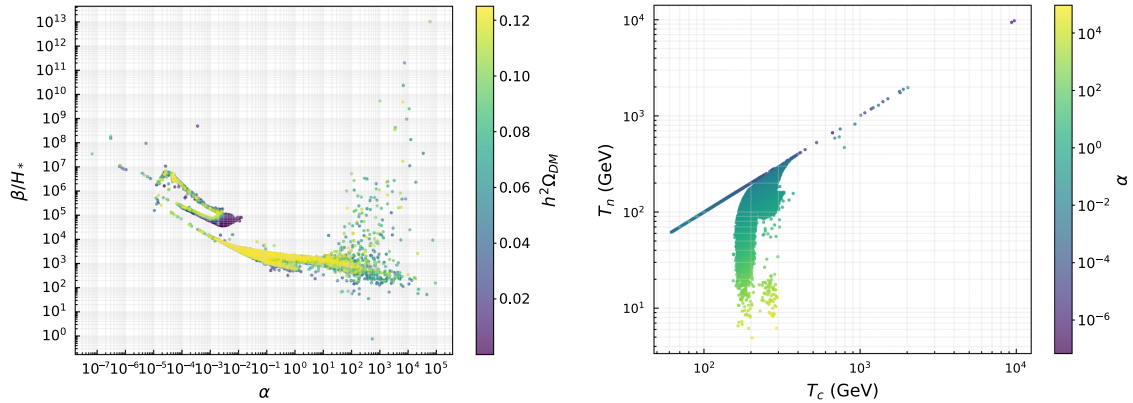


FIG. 8. The left panel shows the correlation between the PT strength α and the normalized inverse PT duration β/H_* . The color bar indicates the DM abundance. The right panel shows the correlation between the critical temperature T_c and the nucleation one T_n , with the PT strength α for the color code.

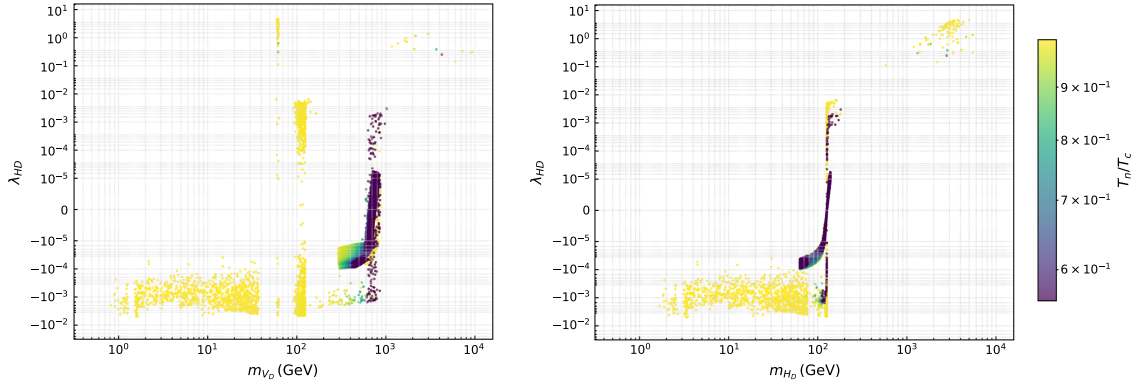


FIG. 9. The left panel shows the correlation between the portal coupling λ_{HD} and the DM mass m_{V_D} , while the right panel shows the correlation between λ_{HD} and the dark Higgs mass m_{H_D} . For both panels, the color bar shows the value of the ratio T_n/T_c .

code indicating the PT strength α . Along the oblique line, we have $T_n \simeq T_c$, which means that the PT occurs as soon as the new minimum in the potential becomes more stable than the current phase. It therefore results in a small value of α . By contrast, the more T_n deviates from the line $T_n = T_c$, the slower is the phase transition, and as shown in the left panel, the bigger the PT strength α . This shows that α and the ratio T_n/T_c are negatively correlated: the smaller T_n/T_c (larger α), the larger the amount of supercooling.

Then, in Fig. 9, we show how the portal coupling λ_{HD} is correlated to the masses m_{V_D} and m_{H_D} . We can first notice that, in most of the parameter space leading to FOPTs, $|\lambda_{\text{HD}}|$ is small: $|\lambda_{\text{HD}}| \lesssim 6 \times 10^{-3}$, and we have $m_{H_D} \lesssim 155$ GeV. Furthermore, a stronger FOPT, thus a smaller value of T_n/T_c , is preferred for $m_{V_D} \in [4 \times 10^2, 10^3]$ GeV

and a moderate mass splitting $|m_{H_D} - m_h| \lesssim m_h/2$ GeV, as already stated when describing Fig. 7. The refined scan yields points in the hockey-stick region bounded by $\lambda_{\text{HD}} \in [-10^{-4}, 2 \times 10^{-5}]$, where a lot of them both provide the observed DM relic density, as well as a strong enough FOPT to be detectable. In addition, we find that such strong FOPTs require $g_D \in [0.6, 1]$.

Figure 10 shows the peak of the GW power spectrum $h^2\Omega_{\text{GW}}^{\text{peak}}$ from FOPTs and the corresponding frequency f^{peak} of this spectrum at its peak. The darker points indicate those giving the correct DM abundance. The power-law-integrated sensitivity curves [77] of LISA, DECIGO, BBO, TianQin, and Taiji for an observation time of 4 years and signal-to-noise ratio equal to 10, are displayed and we see that a significant number of points can potentially be

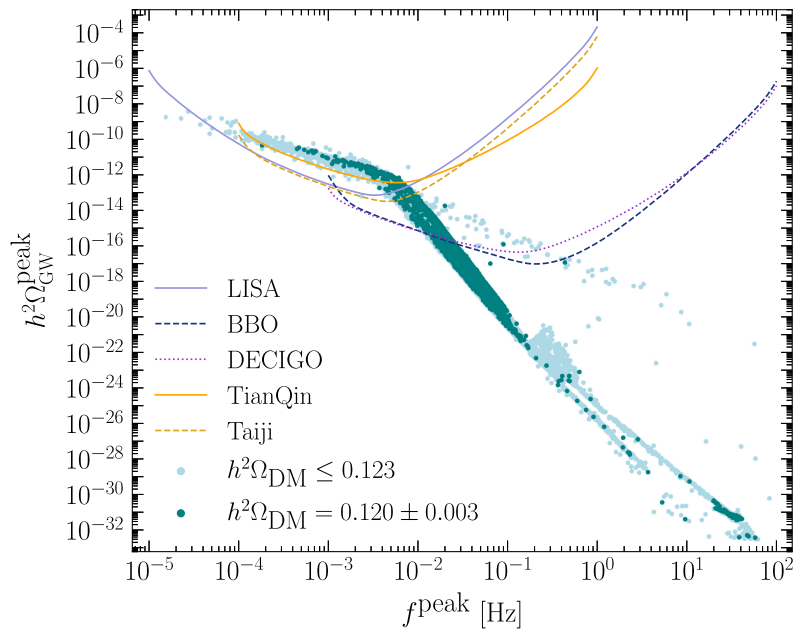


FIG. 10. Peak of the GW signal $h^2\Omega_{\text{GW}}^{\text{peak}}$ as a function of the peak frequency f^{peak} . Darker points are those yielding the observed DM relic density. The power-law integrated sensitivity curves of LISA, DECIGO, BBO, TianQin, and Taiji are also shown.

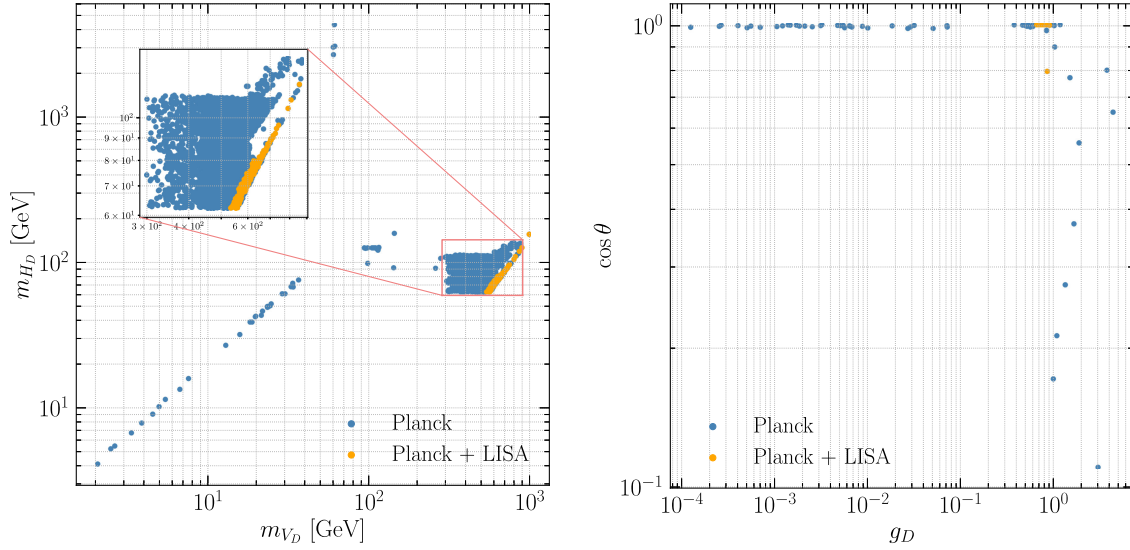


FIG. 11. Results from both the general and refined scans indicating which values of the four free parameters lead to the observed DM relic density (blue) while also being detectable by LISA (orange). The enlarged box in the left panel only contains points from the refined scan. The left panel shows a projection on the plane spanned by the dark-sector masses m_{V_D} and m_{H_D} , while the right panel shows the correlation between the remaining free parameters: g_D and $\cos \theta$.

detected by these five GW detectors. Moreover, a large part of these points also provide a DM abundance that matches the Planck observations. These dark-color points are essentially from the refined scan, which focuses on strong FOPTs. This is why points with a weaker GW signal, $h^2 \Omega_{\text{GW}}^{\text{peak}} \lesssim 10^{-22}$, arising from the more general scan, are less populated than the dark points; the refined scan mostly provides points with $h^2 \Omega_{\text{GW}}^{\text{peak}} \gtrsim 10^{-22}$.

Finally, in Fig. 11, we show the values of the four free parameters of this model that can account for the observed DM relic density, as constrained by the Planck mission. These are depicted by blue points. Those points which, in addition, lead to a GW signal reaching the sensitivity of LISA are depicted in orange. As in Figs. 7–11 presents results from both the general and refined scans combined. The latter has been identified when describing Fig. 7 and correspond to a region of the parameter space leading to strong FOPTs. In order to clearly highlight points from this refined scan, we add an enlarged box containing only points from the latter in the left panel of Fig. 11. Regarding the right panel, points from the refined scan lie in the range $g_D \in [0.6, 1]$, with $\cos \theta \simeq 1$. We can then clearly see that the most interesting points, the orange ones, mostly come from this refined scan. They are approximately obtained for $m_{V_D} \in [500, 900]$ GeV, $m_{H_D} \in [m_h/2, 125]$ GeV, $g_D \in [0.65, 0.85]$ GeV, and $\cos \theta \simeq 1$.

V. CONCLUSIONS

In this paper we have discussed the possibility to generate gravitational waves during a first-order phase transition in a scenario where the Standard Model is

extended by a new $SU(2)$ gauge group. In this scenario, none of the SM particles transform under this $SU(2)$, but the SM Higgs doublet can interact with the scalar doublet which generates the new gauge boson masses through a quadrilinear term in the scalar potential. This also induces a mixing between the Higgs boson and the physical scalar degree of freedom which is left after spontaneous symmetry breaking. The custodial symmetry of the scalar sector noticeably implies that the new massive gauge bosons are stable, and therefore are dark matter candidates. We have tested this model against the most recent experimental data from collider and astrophysical observables and determined the allowed regions of its four-dimensional parameter space. In these regions we found we could obtain the observed DM abundance and that strong first-order phase transitions are possible. We have then shown that the resulting power spectrum of the generated stochastic gravitational-wave background is well within the sensitivity range of different gravitational-wave detectors such as LISA, DECIGO, BBO, TianQin, or Taiji. We have thus shown how gravitational waves can provide a precious complementary way to probe scenarios of new physics, enhancing their discovery potential or at least narrow down their parameter space.

ACKNOWLEDGMENTS

N.B. is supported by the National Natural Science Foundation of China (Grant No. 12475105). The work of L.P. is supported by ICSC—Centro Nazionale di Ricerca in High Performance Computing, Big Data and Quantum Computing, funded by the European

Union—NextGenerationEU. The work of L. D. R. has been funded by the European Union—Next Generation EU through Research Grant No. P2022Z4P4B “SOPHYA—Sustainable Optimised PHYsics Algorithms: fundamental physics to build an advanced society” under the program PRIN 2022 PNRR of the Italian Ministero dell’Università e Ricerca (MUR) and by Research Grant No. 20227S3M3B “Bubble Dynamics in Cosmological Phase Transitions”

under the program PRIN 2022 of the Italian Ministero dell’Università e Ricerca (MUR).

DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

-
- [1] B. P. Abbott *et al.* (LIGO Scientific and Virgo Collaborations), Observation of gravitational waves from a binary black hole merger, *Phys. Rev. Lett.* **116**, 061102 (2016).
 - [2] B. P. Abbott *et al.* (LIGO Scientific and Virgo Collaborations), GW170817: Observation of gravitational waves from a binary neutron star inspiral, *Phys. Rev. Lett.* **119**, 161101 (2017).
 - [3] E. Witten, Cosmic separation of phases, *Phys. Rev. D* **30**, 272 (1984).
 - [4] C. J. Hogan, Gravitational radiation from cosmological phase transitions, *Mon. Not. R. Astron. Soc.* **218**, 629 (1986).
 - [5] P. Amaro-Seoane *et al.*, Laser interferometer space antenna, [arXiv:1702.00786](#).
 - [6] S. Kawamura *et al.*, The Japanese space gravitational wave antenna: DECIGO, *Classical Quantum Gravity* **28**, 094011 (2011).
 - [7] V. Corbin and N. J. Cornish, Detecting the cosmic gravitational wave background with the big bang observer, *Classical Quantum Gravity* **23**, 2435 (2006).
 - [8] J. Luo *et al.* (TianQin Collaborations), TianQin: A spaceborne gravitational wave detector, *Classical Quantum Gravity* **33**, 035010 (2016).
 - [9] W.-R. Hu and Y.-L. Wu, The Taiji program in space for gravitational wave physics and the nature of gravity, *Natl. Sci. Rev.* **4**, 685 (2017).
 - [10] K. Kajantie, M. Laine, K. Rummukainen, and M. E. Shaposhnikov, Is there a hot electroweak phase transition at $m_H \gtrsim m_W$?, *Phys. Rev. Lett.* **77**, 2887 (1996).
 - [11] T. Hambye, Hidden vector dark matter, *J. High Energy Phys.* **01** (2009) 028.
 - [12] T. Hambye and M. H. G. Tytgat, Confined hidden vector dark matter, *Phys. Lett. B* **683**, 39 (2010).
 - [13] F. Chen, J. M. Cline, and A. R. Frey, Nonabelian dark matter: Models and constraints, *Phys. Rev. D* **80**, 083516 (2009).
 - [14] J. L. Diaz-Cruz and E. Ma, Neutral SU(2) gauge extension of the standard model and a vector-boson dark-matter candidate, *Phys. Lett. B* **695**, 264 (2011).
 - [15] S. Bhattacharya, J. L. Diaz-Cruz, E. Ma, and D. Wegman, Dark vector-gauge-boson model, *Phys. Rev. D* **85**, 055008 (2012).
 - [16] T. Hambye and A. Strumia, Dynamical generation of the weak and dark matter scale, *Phys. Rev. D* **88**, 055022 (2013).
 - [17] S. Fraser, E. Ma, and M. Zakeri, $SU(2)_N$ model of vector dark matter with a leptonic connection, *Int. J. Mod. Phys. A* **30**, 1550018 (2015).
 - [18] C. Gross, O. Lebedev, and Y. Mambrini, Non-Abelian gauge fields as dark matter, *J. High Energy Phys.* **08** (2015) 158.
 - [19] N. Bernal, X. Chu, C. Garcia-Cely, T. Hambye, and B. Zaldivar, Production regimes for self-interacting dark matter, *J. Cosmol. Astropart. Phys.* **03** (2016) 018.
 - [20] S. Di Chiara and K. Tuominen, A minimal model for SU(N) vector dark matter, *J. High Energy Phys.* **11** (2015) 188.
 - [21] B. Barman, S. Bhattacharya, S. K. Patra, and J. Chakraborty, Non-Abelian vector boson dark matter, its unified route and signatures at the LHC, *J. Cosmol. Astropart. Phys.* **12** (2017) 021.
 - [22] I. Baldes and C. Garcia-Cely, Strong gravitational radiation from a simple dark matter model, *J. High Energy Phys.* **05** (2019) 190.
 - [23] B. Barman, S. Bhattacharya, and M. Zakeri, Non-Abelian vector boson as FIMP dark matter, *J. Cosmol. Astropart. Phys.* **02** (2020) 029.
 - [24] E. Hall, T. Konstandin, R. McGehee, H. Murayama, and G. Servant, Baryogenesis from a dark first-order phase transition, *J. High Energy Phys.* **04** (2020) 042.
 - [25] D. Marfatia and P.-Y. Tseng, Gravitational wave signals of dark matter freeze-out, *J. High Energy Phys.* **02** (2021) 022.
 - [26] J. Hisano, A. Ibarra, and R. Nagai, Direct detection of vector dark matter through electromagnetic multipoles, *J. Cosmol. Astropart. Phys.* **10** (2020) 015.
 - [27] N. Baoche, A. Ahriche, G. Faisel, and S. Nasri, Phenomenology of the hidden SU(2) vector dark matter model, *Phys. Rev. D* **104**, 075022 (2021).
 - [28] A. Belyaev, A. Deandrea, S. Moretti, L. Panizzi, D. A. Ross, and N. Thongyoi, A fermionic portal to vector dark matter from a new gauge sector, [arXiv:2204.03510](#).
 - [29] T. Abe, K. S. Babu, and A. Kaladharan, Gravitational wave signatures of a chiral fermion dark matter model, *J. Cosmol. Astropart. Phys.* **07** (2024) 048.
 - [30] M. Arteaga, A. Ghoshal, and A. Strumia, Gravitational waves and black holes from the phase transition in models of dynamical symmetry breaking, *J. Cosmol. Astropart. Phys.* **05** (2025) 029.
 - [31] M. Arteaga Tupia, A. Ghoshal, and A. Strumia, Primordial magnetogenesis from a supercooled dynamical electroweak phase transition, [arXiv:2506.16387](#).

- [32] P. Athron, C. Balázs, A. Fowlie, L. Morris, and L. Wu, Cosmological phase transitions: From perturbative particle physics to gravitational waves, *Prog. Part. Nucl. Phys.* **135**, 104094 (2024).
- [33] A. Belyaev, A. Deandrea, S. Moretti, L. Panizzi, and N. Thongyoi, A fermionic portal to a non-Abelian dark sector, [arXiv:2203.04681](#).
- [34] N. Aghanim *et al.* (Planck Collaborations), Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020); *Astron. Astrophys.* **652**, C4(E) (2021).
- [35] G. Alguero, G. Belanger, F. Boudjema, S. Chakraborti, A. Goudelis, S. Kraml, A. Mjallal, and A. Pukhov, micrOMEGAs6.0: N-component dark matter, *Comput. Phys. Commun.* **299**, 109133 (2024).
- [36] J. Aalbers *et al.* (LZ Collaborations), First dark matter search results from the LUX-ZEPLIN (LZ) experiment, [arXiv:2207.03764](#).
- [37] C. Amole *et al.* (PICO Collaborations), Dark matter search results from the PICO-60 C₃F₈ bubble chamber, *Phys. Rev. Lett.* **118**, 251301 (2017).
- [38] E. Aprile *et al.* (XENON Collaborations), Dark matter search results from a one ton-year exposure of XENON1T, *Phys. Rev. Lett.* **121**, 111302 (2018).
- [39] Y. Meng *et al.* (PandaX-4T Collaborations), Dark matter search results from the PandaX-4T commissioning run, *Phys. Rev. Lett.* **127**, 261802 (2021).
- [40] P. Agnes *et al.* (DarkSide Collaborations), Low-mass dark matter search with the DarkSide-50 experiment, *Phys. Rev. Lett.* **121**, 081307 (2018).
- [41] A. H. Abdelhameed *et al.* (CRESST Collaborations), First results from the CRESST-III low-mass dark matter program, *Phys. Rev. D* **100**, 102002 (2019).
- [42] G. Bélanger, J. Da Silva, T. Perrillat-Bottonet, and A. Pukhov, Limits on dark matter proton scattering from neutrino telescopes using micrOMEGAs, *J. Cosmol. Astropart. Phys.* **12** (2015) 036.
- [43] F. Ambrogio, C. Arina, M. Backovic, J. Heisig, F. Maltoni, L. Mantani, O. Mattelaer, and G. Mohlabeng, MadDM v.3.0: A comprehensive tool for dark matter studies, *Phys. Dark Universe* **24**, 100249 (2019).
- [44] A. Albert *et al.* (Fermi-LAT and DES Collaborations), Searching for dark matter annihilation in recently discovered Milky Way satellites with Fermi-LAT, *Astrophys. J.* **834**, 110 (2017).
- [45] G. Arcadi, C. Gross, O. Lebedev, S. Pokorski, and T. Toma, Evading direct dark matter detection in Higgs portal models, *Phys. Lett. B* **769**, 129 (2017).
- [46] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher, and J. P. Silva, Theory and phenomenology of two-Higgs-doublet models, *Phys. Rep.* **516**, 1 (2012).
- [47] M. E. Peskin and T. Takeuchi, Estimation of oblique electroweak corrections, *Phys. Rev. D* **46**, 381 (1992).
- [48] M. E. Peskin and T. Takeuchi, A new constraint on a strongly interacting Higgs sector, *Phys. Rev. Lett.* **65**, 964 (1990).
- [49] A. Beniwal, M. Lewicki, M. White, and A. G. Williams, Gravitational waves and electroweak baryogenesis in a global study of the extended scalar singlet model, *J. High Energy Phys.* **02** (2019) 183.
- [50] W. Grimus, L. Lavoura, O. M. Ogreid, and P. Osland, The oblique parameters in multi-Higgs-doublet models, *Nucl. Phys. B* **801**, 81 (2008).
- [51] S. Navas *et al.* (Particle Data Group), Review of particle physics, *Phys. Rev. D* **110**, 030001 (2024).
- [52] H. Bahl, T. Biekötter, S. Heinemeyer, C. Li, S. Paasch, G. Weiglein, and J. Wittbrodt, HiggsTools: BSM scalar phenomenology with new versions of HiggsBounds and HiggsSignals, *Comput. Phys. Commun.* **291**, 108803 (2023).
- [53] M. Aaboud *et al.* (ATLAS Collaborations), Search for invisible Higgs boson decays in vector boson fusion at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Phys. Lett. B* **793**, 499 (2019).
- [54] G. Aad *et al.* (ATLAS Collaborations), Combination of searches for invisible decays of the Higgs boson using 139 fb⁻¹ of proton-proton collision data at $\sqrt{s} = 13$ TeV collected with the ATLAS experiment, *Phys. Lett. B* **842**, 137963 (2023).
- [55] S. Coleman and E. Weinberg, Radiative corrections as the origin of spontaneous symmetry breaking, *Phys. Rev. D* **7**, 1888 (1973).
- [56] J. M. Cline, K. Kainulainen, and M. Trott, Electroweak baryogenesis in two Higgs doublet models and B meson anomalies, *J. High Energy Phys.* **11** (2011) 089.
- [57] L. Dolan and R. Jackiw, Symmetry behavior at finite temperature, *Phys. Rev. D* **9**, 3320 (1974).
- [58] D. J. Gross, R. D. Pisarski, and L. G. Yaffe, QCD and instantons at finite temperature, *Rev. Mod. Phys.* **53**, 43 (1981).
- [59] R. R. Parwani, Resummation in a hot scalar field theory, *Phys. Rev. D* **45**, 4695 (1992); **48**, 5965(E) (1993).
- [60] M. E. Carrington, Effective potential at finite temperature in the standard model, *Phys. Rev. D* **45**, 2933 (1992).
- [61] J. R. Espinosa, T. Konstandin, J. M. No, and G. Servant, Energy budget of cosmological first-order phase transitions, *J. Cosmol. Astropart. Phys.* **06** (2010) 028.
- [62] C. Grojean and G. Servant, Gravitational waves from phase transitions at the electroweak scale and beyond, *Phys. Rev. D* **75**, 043507 (2007).
- [63] M. Kamionkowski, A. Kosowsky, and M. S. Turner, Gravitational radiation from first order phase transitions, *Phys. Rev. D* **49**, 2837 (1994).
- [64] C. J. Hogan, Gravitational radiation from cosmological phase transitions, *Mon. Not. R. Astron. Soc.* **218**, 629 (1986).
- [65] M. Hindmarsh, S. J. Huber, K. Rummukainen, and D. J. Weir, Gravitational waves from the sound of a first order phase transition, *Phys. Rev. Lett.* **112**, 041301 (2014).
- [66] C. Caprini, R. Durrer, and G. Servant, The stochastic gravitational wave background from turbulence and magnetic fields generated by a first-order phase transition, *J. Cosmol. Astropart. Phys.* **12** (2009) 024.
- [67] C. Caprini *et al.*, Science with the space-based interferometer eLISA. II: Gravitational waves from cosmological phase transitions, *J. Cosmol. Astropart. Phys.* **04** (2016) 001.
- [68] M. Hindmarsh, S. J. Huber, K. Rummukainen, and D. J. Weir, Shape of the acoustic gravitational wave power spectrum from a first order phase transition, *Phys. Rev. D* **96**, 103520 (2017); *Phys. Rev. D* **101**, 089902(E) (2020).

- [69] C. Caprini *et al.*, Detecting gravitational waves from cosmological phase transitions with LISA: An update, *J. Cosmol. Astropart. Phys.* **03** (2020) 024.
- [70] K. Schmitz, LISA sensitivity to gravitational waves from sound waves, *Symmetry* **12**, 1477 (2020).
- [71] D. Croon, O. Gould, P. Schicho, T. V. I. Tenkanen, and G. White, Theoretical uncertainties for cosmological first-order phase transitions, *J. High Energy Phys.* **04** (2021) 055.
- [72] P. Athron, C. Balazs, A. Fowlie, L. Morris, G. White, and Y. Zhang, How arbitrary are perturbative calculations of the electroweak phase transition?, *J. High Energy Phys.* **01** (2023) 050.
- [73] H.-K. Guo, K. Sinha, D. Vagie, and G. White, Phase transitions in an expanding universe: Stochastic gravitational waves in standard and non-standard histories, *J. Cosmol. Astropart. Phys.* **01** (2021) 001.
- [74] J. Ellis, M. Lewicki, J. M. No, and V. Vaskonen, Gravitational wave energy budget in strongly supercooled phase transitions, *J. Cosmol. Astropart. Phys.* **06** (2019) 024.
- [75] C. L. Wainwright, CosmoTransitions: Computing cosmological phase transition temperatures and bubble profiles with multiple fields, *Comput. Phys. Commun.* **183**, 2006 (2012).
- [76] N. Benincasa, L. Delle Rose, K. Kannike, and L. Marzola, Multi-step phase transitions and gravitational waves in the inert doublet model, *J. Cosmol. Astropart. Phys.* **12** (2022) 025.
- [77] E. Thrane and J. D. Romano, Sensitivity curves for searches for gravitational-wave backgrounds, *Phys. Rev. D* **88**, 124032 (2013).