

Primordial black hole formation and multimessenger signals in a complex singlet extension of the standard model

Fa Peng Huang[✉], Chikako Idegawa^{✉,*} and Aidi Yang

*MOE Key Laboratory of TianQin Mission, TianQin Research Center for Gravitational Physics
and School of Physics and Astronomy, Frontiers Science Center for TianQin,
Gravitational Wave Research Center of CNSA, Sun Yat-sen University (Zhuhai Campus),
Zhuhai 519082, China*



(Received 6 November 2025; accepted 13 February 2026; published 4 March 2026)

We investigate the formation of primordial black holes (PBHs) induced by a first-order electroweak phase transition in a realistic renormalizable framework, the complex singlet extension of the Standard Model. We perform a quantitative analysis of the PBH abundance and identify parameter regions consistent with current microlensing constraints. Furthermore, we show that the same parameter space predicts observable stochastic gravitational waves within the sensitivities of future space-based detectors, as well as a sizable deviation in the Higgs triple coupling that can be probed at future lepton colliders. Our results highlight a comprehensive multimessenger framework in which PBHs, gravitational waves, and collider observations can jointly test the dynamics of a strongly first-order electroweak phase transition in the early Universe.

DOI: [10.1103/cxh3-bpl6](https://doi.org/10.1103/cxh3-bpl6)

I. INTRODUCTION

Primordial black holes (PBHs) are hypothetical black holes formed in the early Universe, long before the formation of stars and galaxies [1–3]. They can serve as unique probes of the early Universe, providing insights into high-energy physics phenomena that cannot be tested in laboratory experiments. A wide variety of mechanisms for the formation of PBHs have been proposed. Over the past few decades, one extensively studied class of models involves scenarios where large primordial quantum fluctuations are generated at small scales in some inflation models [4–7]. In such frameworks, the enhanced curvature perturbations can collapse gravitationally upon horizon reentry, giving rise to PBHs [3]. More recently, increasing attention has been devoted to alternative mechanisms in which PBHs are produced during cosmological first-order phase transitions in the early Universe. In these scenarios, the specified dynamical processes associated with the phase transition can generate regions of overdensity that may gravitationally collapse into PBHs [8–23]. These studies open up a new window for connecting PBH formation with

the microphysics of the early Universe, providing a complementary probe of new physics models.

Among the various mechanisms proposed for PBH formation during a strong first-order phase transition, a particularly well-motivated scenario involves a nonuniform nucleation rate in time. Certain regions of the Universe can undergo bubble nucleation earlier or later than others. During such a transition, true-vacuum bubbles nucleate stochastically through thermal tunneling, and different Hubble patches complete the transition at slightly different times [8,24]. As a result, regions where bubble nucleation occurs later, referred to as delayed patches, remain in the false vacuum longer and thus retain a larger vacuum energy density. Since the false-vacuum energy does not redshift while the radiation energy density decreases as the Universe expands, the delayed-decayed regions gradually become overdense compared to the surrounding space [9,25]. Once the energy-density contrast between the normal-decayed and delayed-decayed regions exceeds a critical threshold, the region collapses into a PBH upon horizon reentry. The PBH fraction f_{PBH} is determined by the rare probability that a Hubble patch remains in the false vacuum up to such a delayed time, making f_{PBH} extremely sensitive to the detailed dynamics of the phase transition [15,18]. This delayed vacuum decay mechanism therefore connects the microscopic physics of the phase transition to macroscopic cosmological observables.

While the possibility of PBH formation from a first-order phase transition has been studied extensively, most of the existing analyses have been performed in a

*Contact author: idegawa@mail.sysu.edu.cn

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

model-independent framework or within phenomenological toy models that parametrize the thermal potential. In particular, studies focusing on the electroweak phase transition (EWPT) have often adopted such approaches, since the properties of the EWPT, such as its strength and critical temperature, depend sensitively on the underlying model [26–28], and a full numerical treatment of the finite-temperature potential is computationally demanding. As a result, the dynamics of PBH formation have typically been discussed in a general manner without specifying a concrete model. A few studies have explored PBH formation within specific particle-physics models such as Refs. [14,15,18,19,21–23,29–32]. In this work, we instead focus on a realistic renormalizable setup, a complex singlet extension of the Standard Model (CxSM) [33–48], and examine the possibility of PBH formation in this framework, together with its phenomenological implications for gravitational waves (GWs) and collider observables.

As a well-motivated and minimal extension of the Standard Model (SM), the CxSM introduces a gauge-singlet complex scalar field S that couples to the Higgs doublet. This model simultaneously accommodates a viable dark matter (DM) candidate and allows for a strong first-order EWPT [42,43]. In particular, we focus on the so-called degenerate-scalar scenario, in which the additional scalar has a mass nearly degenerate with the observed Higgs boson [49,50]. This setup is known to evade current DM direct-detection bounds [51] due to destructive interference between multiple scalar-mediated amplitudes, while remaining consistent with collider constraints [52,53].

The existence of DM provides clear evidence for physics beyond the SM; however, no conclusive DM signal has been observed so far. While this null result is often interpreted in terms of decoupling scenarios, such as very heavy new states or extremely suppressed couplings to the SM particles, the steadily improving experimental sensitivity without corresponding discoveries motivates the exploration of qualitatively different possibilities. In contrast to such decoupling-based interpretations, the degenerate-scalar realization of the CxSM provides a qualitatively different possibility: the nonobservation of DM can be naturally understood without invoking decoupling from the SM.

It should also be emphasized that, in the parameter region relevant for a strong first-order EWPT in the CxSM, the relic abundance of the pseudoscalar DM component is typically much smaller than the observed DM density. This feature has been explicitly demonstrated in Ref. [42], where it was shown that the pseudoscalar DM in the degenerate-scalar scenario cannot account for the full dark matter abundance. A concrete realization of a two-component DM scenario within the CxSM, in which the missing relic density is supplemented by an additional fermion DM candidate, was subsequently studied in Ref. [47].

Therefore, in the present work, we do not require the pseudoscalar component to explain the entire DM abundance. While PBHs can in principle contribute to the DM abundance, we do not perform a detailed quantitative assessment of their contribution in this work. Instead, PBH formation is mainly used as a cosmological probe of the dynamics of the EWPT. We note, however, that PBHs could potentially complement other DM components and jointly account for the observed DM relic abundance.

In this work, we investigate the possibility of PBH formation induced by a first-order EWPT in the CxSM. We explore how the dynamics of the EWPT, determined by the scalar potential parameters, affect the resulting PBH abundance, and we compare the predicted fractions with recent microlensing constraints from HSC and OGLE [54–56]. In addition to PBH production, we also study two complementary phenomenological consequences of a strong first-order EWPT: the generation of stochastic GWs [57,58], and the deviation of the Higgs triple coupling from its SM prediction [59,60]. Through this multimessenger approach, we demonstrate that the same underlying phase transition can simultaneously lead to observable signals in cosmology, GW astronomy, and collider experiments.

This paper is organized as follows: In Sec. II, we review the formalism of PBH formation arising from a first-order EWPT. In Sec. III, we introduce the CxSM, define the relevant parameters, and discuss the degenerate-scalar scenario. The numerical results for PBH formation are presented in Sec. IV, where we analyze how the model parameters affect the transition dynamics and compare the predicted PBH abundances with current microlensing constraints from HSC and OGLE. In Sec. V, we investigate the GW signatures of the EWPT and their correlations with collider observables such as the Higgs triple coupling deviation. Finally, Sec. VI summarizes our findings and discusses the implications of this comprehensive multimessenger framework for probing the dynamics of the EWPT.

II. PRIMORDIAL BLACK HOLE FORMATION FROM ELECTROWEAK PHASE TRANSITION

In various extensions of the SM, a first-order EWPT plays an important role in the early Universe, such as those that lead to GW production and PBH formation. In particular, a first-order EWPT is often referred to as a strong first-order when the sphaleron process becomes inefficient inside the broken phase, which is approximately characterized by

$$\frac{v_C}{T_C} > 1. \quad (1)$$

T_C denotes the critical temperature at which the effective potential develops two degenerate minima, and v_C is the vacuum expectation value (VEV) of the Higgs doublet

at T_C . We here outline the procedure used to evaluate the PBH abundance generated from such a transition, following the framework developed in Refs. [15,18,22,61]. First, the decay rate of the false vacuum is given by [62]

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

where $S(T) = S_3(T)/T$, and $S_3(T)$ is the three-dimensional Euclidean action of a critical bubble. The nucleation temperature T_N is defined by the condition

$$\Gamma(T_N)H^{-4}(T_N) = 1. \quad (3)$$

The spatial average of the false vacuum fraction at time t is given by

$$F(t) = \exp \left[-\frac{4\pi}{3} \int_{t_R}^t dt' \Gamma(t') a^3(t') r^3(t, t') \right], \quad (4)$$

where $r(t, t')$ denotes the comoving radius of the true vacuum bubble, defined as

$$r(t, t') \equiv \int_{t'}^t \frac{d\tilde{t}}{a(\tilde{t})}. \quad (5)$$

The time variable can be converted into temperature using the Hubble parameter through

$$\frac{dt}{dT} = -\frac{1}{TH(T)}. \quad (6)$$

Then, we obtain

$$a(t')r(t, t') = a(T') \int_T^{T'} \frac{d\tilde{T}}{T' H(\tilde{T}) a(T')} = \frac{1}{T'} \int_T^{T'} \frac{d\tilde{T}}{H(\tilde{T})}. \quad (7)$$

Accordingly, the false vacuum fraction can be rewritten as a function of temperature:

$$F(T) = \exp \left[-\frac{4\pi}{3} \int_T^{T_R} \frac{dT' \Gamma(T')}{T' H(T')} a^3(T') r^3(T, T') \right], \quad (8)$$

where T_R is the reference temperature that depends on the relevant region. The initial value of $F(T)$ is set differently for the normal-decayed and the delayed-decayed regions [63]:

$$F(T) = \exp \left[-12\pi (M_{\text{Pl}} \xi_g)^4 \int_T^{T_C} \frac{dT' \Gamma(T')}{T'^6} \left(\frac{1}{T} - \frac{1}{T'} \right)^3 \right], \quad (9)$$

(normal-decayed region),

$$F(T) = \exp \left[-12\pi (M_{\text{Pl}} \xi_g)^4 \int_T^{T_d} \frac{dT' \Gamma(T')}{T'^6} \left(\frac{1}{T} - \frac{1}{T'} \right)^3 \right],$$

(delayed-decayed region), (10)

where $M_{\text{Pl}} = 2.435 \times 10^{18}$ GeV and $\xi_g = \sqrt{30/(\pi^2 g)}$, with $g = 108.75$. For the normal-decayed region, we use the critical temperature T_C as the reference temperature, while T_d , corresponding to the delayed-decayed region, is determined later to satisfy the PBH formation condition.

During a first-order EWPT, the Universe contains regions remaining in the false vacuum and regions that have transitioned to the true vacuum. The difference between the two regions is given by

$$\Delta V = V_{\text{false}} - V_{\text{true}}. \quad (11)$$

The total energy density of the Universe can be expressed as

$$\rho(T) = F(T) V_{\text{false}} + (1 - F(T)) V_{\text{true}}. \quad (12)$$

Choosing $V_{\text{true}} = 0$ as the reference, the vacuum energy density becomes

$$\rho_v(T) = F(T) \Delta V. \quad (13)$$

The evolution of the radiation energy density is then described by

$$\frac{d\rho_r(T)}{dT} = \frac{4\rho_r(T)}{T} - \frac{d\rho_v(T)}{dT}, \quad (14)$$

and the Hubble parameter is determined by

$$H^2 = \left(\frac{\dot{a}}{a} \right)^2 = \frac{1}{3M_{\text{Pl}}^2} (\rho_v(T) + \rho_r(T)), \quad (15)$$

where $\dot{a}$ denotes the derivative of the scale factor with respect to the temperature T . By substituting the above energy densities into Eq. (15), we can iteratively solve for the Hubble parameter and the scale factor until convergence.

For PBH formation, the relevant quantity is the contrast in the total energy density between the delayed- and normal-decayed regions:

$$\delta \equiv \frac{\rho_{\text{delay}}}{\rho_{\text{normal}}} - 1 = \frac{\rho_{r,\text{delay}} + \rho_{v,\text{delay}}}{\rho_{r,\text{normal}} + \rho_{v,\text{normal}}} - 1. \quad (16)$$

In this work, we adopt $\delta_c = 0.45$ as the threshold value for PBH formation [64], such that $\delta \geq \delta_c$ must be satisfied.¹ The temperature T_d for the delayed region is chosen such

¹For studies discussing the threshold value δ_c , see, for example, Refs. [22,65–67].

that the peak of δ in Eq. (16) reaches 0.45. The temperature T_d is lower than the nucleation temperature T_N , but if it becomes much lower, the probability of remaining in the false vacuum state rapidly decreases.

The resulting PBH abundance produced by the EWPT is given by [15]

$$f_{\text{PBH}}^{\text{EW}} \simeq 1.49 \times 10^{11} \left(\frac{0.25}{\Omega_{\text{CDM}}} \right) \left(\frac{T_{\text{PBH}}}{100 \text{ GeV}} \right) P(t_N), \quad (17)$$

where

$$P(t_N) = \exp \left[-\frac{4\pi}{3} \int_{t_c}^{t_d} \frac{a_{\text{delay}}^3(t)}{a_{\text{delay}}^3(t_{\text{PBH}})} \frac{1}{H_{\text{delay}}^3(t_{\text{PBH}})} \Gamma_{\text{delay}}(t) dt \right]. \quad (18)$$

Here, Ω_{CDM} denotes the present cold DM density normalized by the total energy density, and $T_{\text{PBH}}(t_{\text{PBH}})$ represents the temperature (time) at PBH formation. The PBH mass is evaluated as

$$M_{\text{PBH}} = \gamma \frac{4\pi}{3} H_{\text{delay}}^{-3}(t_{\text{PBH}}) (\rho_{r,\text{delay}} + \rho_{v,\text{delay}}), \quad (19)$$

where γ denotes the fraction of the horizon mass that collapses into a PBH, and we adopt $\gamma = 0.2$ in our analysis [3,18,68].

As mentioned in Sec. I, previous studies on PBH formation from the EWPT have often been carried out in a model-independent framework or by employing phenomenological toy models that parametrize the thermal potential. In this work, we instead focus on a well-motivated renormalizable extension of the SM—namely, the CxSM—and investigate the possibility of PBH formation, as well as the testability of the relevant parameters through future GW and collider experiments.

III. THE COMPLEX SINGLET EXTENDED HIGGS MODEL

A. Model introduction

The CxSM is an extension of the SM that introduces a complex SU(2) gauge-singlet scalar field [33]. In our study, we consider the following scalar potential:

$$V_0(\Phi, S) = \frac{m^2}{2} \Phi^\dagger \Phi + \frac{\lambda}{4} (\Phi^\dagger \Phi)^2 + \frac{\delta_2}{2} \Phi^\dagger \Phi |S|^2 + \frac{b_2}{2} |S|^2 + \frac{d_2}{4} |S|^4 + \left(a_1 S + \frac{b_1}{4} S^2 + \text{H.c.} \right), \quad (20)$$

where a global U(1) symmetry of S is softly broken by the a_1 and b_1 terms. In the following, all the couplings in Eq. (20) are assumed to be real. When the linear term in S is absent, the potential possesses a Z_2 symmetry ($S \rightarrow -S$). If the singlet S acquires a VEV, this Z_2 symmetry is

spontaneously broken, leading to the domain-wall problem [49]. To avoid this issue, we include the linear term of S in the potential Eq. (20), which explicitly breaks the Z_2 symmetry and thus prevents the formation of domain walls. Although several U(1)-breaking terms are allowed in the potential, not all of them are necessary to achieve a strong first-order EWPT and viable DM. Therefore, we adopt a minimal set of operators that is closed under renormalization.

We parametrize the scalar fields as

$$\Phi = \begin{pmatrix} G^+ \\ \frac{1}{\sqrt{2}}(v + h + iG^0) \end{pmatrix}, \quad (21)$$

$$S = \frac{1}{\sqrt{2}}(v_S + s + i\chi), \quad (22)$$

where $v(\simeq 246.22 \text{ GeV})$ and v_S denote the VEVs of Φ and S , respectively. The Nambu-Goldstone bosons G^+ and G^0 are absorbed by the W and Z bosons, respectively, after the electroweak symmetry breaking. Since we assume that all parameters in Eq. (20) are real, the scalar potential is invariant under the CP transformation ($S \rightarrow S^*$). As a result, the real and imaginary components of S do not mix, and the stability of χ is ensured, making it a scalar DM candidate.

The first derivatives of V_0 with respect to h and s are, respectively, given by

$$\frac{1}{v} \left\langle \frac{\partial V_0}{\partial h} \right\rangle = \frac{m^2}{2} + \frac{\lambda}{4} v^2 + \frac{\delta_2}{4} v_S^2 = 0, \quad (23)$$

$$\frac{1}{v_S} \left\langle \frac{\partial V_0}{\partial s} \right\rangle = \frac{b_2}{2} + \frac{\delta_2}{4} v^2 + \frac{d_2}{4} v_S^2 + \frac{\sqrt{2}a_1}{v_S} + \frac{b_1}{2} = 0, \quad (24)$$

where $\langle \dots \rangle$ indicates that all fluctuation fields are set to zero. Note that a nonzero value of v_S is enforced by $a_1 \neq 0$.

The mass matrix of (h, s) is expressed as

$$\mathcal{M}_S^2 = \begin{pmatrix} \lambda v^2/2 & \delta_2 v v_S/2 \\ \delta_2 v v_S/2 & \Lambda^2 \end{pmatrix}, \quad \Lambda^2 \equiv \frac{d_2}{2} v_S^2 - \sqrt{2} \frac{a_1}{v_S}. \quad (25)$$

The mass matrix in Eq. (25) is diagonalized by an orthogonal matrix $O(\theta)$ as

$$O(\theta)^\top \mathcal{M}_S^2 O(\theta) = \begin{pmatrix} m_{h_1}^2 & 0 \\ 0 & m_{h_2}^2 \end{pmatrix}, \quad O(\theta) = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix}, \quad (26)$$

where θ denotes the mixing angle. The mass eigenstates (h_1, h_2) are related to the gauge eigenstates (h, s) through the mixing matrix as

$$\begin{pmatrix} h \\ s \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} h_1 \\ h_2 \end{pmatrix}. \quad (27)$$

We emphasize that the limit $\theta \rightarrow 0$ corresponds to the SM-like limit $(h_1 \rightarrow h, h_2 \rightarrow s)$. The mass eigenvalues are obtained as

$$m_{h_1, h_2}^2 = \frac{1}{2} \left(\frac{\lambda}{2} v^2 + \Lambda^2 \mp \frac{\frac{\lambda}{2} v^2 - \Lambda^2}{\cos 2\theta} \right) \quad (28)$$

$$= \frac{1}{2} \left(\frac{\lambda}{2} v^2 + \Lambda^2 \mp \sqrt{\left(\frac{\lambda}{2} v^2 - \Lambda^2 \right)^2 + 4 \left(\frac{\delta_2}{2} v v_S \right)^2} \right), \quad (29)$$

$$\cos 2\theta = \frac{\frac{\lambda}{2} v^2 - \Lambda^2}{m_{h_1}^2 - m_{h_2}^2}. \quad (30)$$

We identify h_1 with the Higgs boson observed in the LHC experiments—i.e., $m_{h_1} = 125$ GeV. The χ mass is determined by the soft breaking terms a_1 and b_1 as

$$\begin{aligned} m_\chi^2 &= \frac{b_2}{2} - \frac{b_1}{2} + \frac{\delta_2}{4} v^2 + \frac{d_2}{4} v_S^2 \\ &= -\frac{\sqrt{2} a_1}{v_S} - b_1, \end{aligned} \quad (31)$$

where the tadpole condition [Eq. (24)] has been used in the second equality.

For later convenience, we summarize the relations between the input and output parameters. Here, we take $\{v, m_{h_1}, m_{h_2}, \theta, a_1, v_S, m_\chi\}$ as input parameters, while the Lagrangian parameters $\{m^2, b_2, \lambda, d_2, \delta_2, b_1\}$ can be expressed in terms of the inputs. Among these, m^2 and b_2 are eliminated using the tadpole conditions in Eqs. (23) and (24):

$$m^2 = -\frac{\lambda}{2} v^2 - \frac{\delta_2}{2} v_S^2, \quad (32)$$

$$b_2 = -\frac{\delta_2}{2} v^2 - \frac{d_2}{2} v_S^2 - \sqrt{2} \frac{a_1}{v_S} - b_1. \quad (33)$$

The remaining four Lagrangian parameters are then written as

$$\lambda = \frac{2}{v^2} (m_{h_1}^2 \cos^2 \theta + m_{h_2}^2 \sin^2 \theta), \quad (34)$$

$$\delta_2 = \frac{1}{v v_S} (m_{h_1}^2 - m_{h_2}^2) \sin 2\theta, \quad (35)$$

$$d_2 = 2 \left(\frac{m_{h_1}}{v_S} \right)^2 \sin^2 \theta + 2 \left(\frac{m_{h_2}}{v_S} \right)^2 \cos^2 \theta + 2\sqrt{2} \frac{a_1}{v_S^3}, \quad (36)$$

$$b_1 = -m_\chi^2 - \frac{\sqrt{2}}{v_S} a_1. \quad (37)$$

Theoretical constraints on the quartic couplings in the scalar potential are summarized below. To ensure that the potential is bounded from below, the following conditions must hold²:

$$\lambda > 0, \quad d_2 > 0. \quad (38)$$

Furthermore, the couplings λ and d_2 are constrained by perturbative unitarity [49]:

$$\lambda < \frac{16\pi}{3}, \quad d_2 < \frac{16\pi}{3}. \quad (39)$$

In addition, requiring the eigenvalues of the mass matrix in Eq. (29) to be positive leads to the following tree-level stability condition [33]:

$$\frac{2\lambda\Lambda^2}{v_S^2} = \lambda \left(d_2 - \frac{2\sqrt{2}a_1}{v_S^3} \right) > \delta_2^2. \quad (40)$$

B. Degenerate scalar scenario

Recent DM direct-detection experiments have provided stringent upper limits on the spin-independent DM-nucleon scattering cross section [51], which strongly constrain various Higgs-portal-type DM models. In the CxSM, the scattering process of the DM χ with nucleons is mediated by h_1 and h_2 . It has been pointed out [49] that the DM-quark scattering amplitude can be significantly suppressed when the masses of h_1 and h_2 are nearly degenerate. This so-called degenerate scalar scenario naturally evades the current direct-detection bounds without introducing additional symmetry structures.

In this work, although we do not focus on the detailed DM phenomenology, the degenerate scalar region remains of particular interest, because it is simultaneously consistent with both direct-detection constraints and collider searches. The coupling of h_1 (h_2) to the SM fermions is scaled by $\cos \theta$ ($-\sin \theta$):

$$\mathcal{L}_Y = \frac{m_q}{v} \bar{q} q (h_1 \cos \theta - h_2 \sin \theta), \quad (41)$$

where m_q denotes a mass of the quark q . The couplings of h_1 and h_2 to the SM gauge bosons are modified in the same

²For $\delta_2 < 0$, an additional condition $\lambda d_2 > \delta_2^2$ is required. In this work, we assume $\delta_2 > 0$.

way. Consequently, the partial decay widths of h_1 and h_2 into an SM final state X are given by

$$\Gamma_{h_1 \rightarrow XX} = \cos^2 \theta \Gamma_{h \rightarrow XX}^{\text{SM}}(m_{h_1}), \quad (42)$$

$$\Gamma_{h_2 \rightarrow XX} = \sin^2 \theta \Gamma_{h \rightarrow XX}^{\text{SM}}(m_{h_2}), \quad (43)$$

where $\Gamma_{h \rightarrow XX}^{\text{SM}}(m_{h_{1,2}})$ denotes the corresponding SM Higgs partial decay width evaluated at $m_{h_{1,2}}$.

When the two scalar masses are nearly degenerate ($m_{h_1} \simeq m_{h_2}$), their production and decay processes cannot be experimentally distinguished, and only the combined signal is observed:

$$\Gamma_{h_1 \rightarrow XX} + \Gamma_{h_2 \rightarrow XX} \simeq \Gamma_{h \rightarrow XX}^{\text{SM}}(m_h). \quad (44)$$

Therefore, the total Higgs signal strength in the degenerate limit coincides with that of the SM regardless of the mixing angle θ . Such a parameter region is thus phenomenologically motivated, being compatible with current Higgs measurements while remaining consistent with DM direct-detection constraints.

IV. NUMERICAL ANALYSIS OF PBH FORMATION

Before presenting the numerical results for PBH formation, we first discuss the characteristics of the EWPT in the CxSM. The thermal evolution of the vacuum in this model proceeds from the singlet-like phase ($\langle \Phi \rangle, \langle S \rangle$) = $(0, v'_S)$ to the electroweak symmetry-broken phase (v, v_S) , where $\langle \Phi \rangle$ and $\langle S \rangle$ denote the VEVs of the SM Higgs doublet and the singlet scalar, respectively. It has been shown in previous studies that the tree-level contributions play an essential role in determining the nature of the EWPT in the CxSM [42]. In particular, the parameters δ_2 and d_2 are the key factors that control the strength of the transition. Using the tree-level potential and thermal masses, the critical temperature T_C and the Higgs VEV v_C at T_C can be approximately expressed as follows³:

$$\begin{aligned} T_C &\simeq \sqrt{\frac{1}{2\Sigma_\Phi} \left(-m^2 - \frac{(v'_{SC})^2}{2} \delta_2 \right)}, \\ v_C &\simeq \sqrt{\frac{2\delta_2 (v'_{SC})^2}{\lambda} \left(1 - \frac{v_{SC}}{v'_{SC}} \right)}, \end{aligned} \quad (45)$$

³Although the cubic terms in the fields arising from thermal boson loops serve as the source of the EWPT, we focus on the tree-level potential supplemented by the thermal masses (quadratic terms in the fields), which dominate at high temperature, in order to illustrate that the tree-level potential plays an essential role in forming the potential barrier between vacua. In the numerical analysis presented later, we employ the full one-loop effective potential.

where all quantities evaluated at T_C are denoted with the subscript C . It is found that a large δ_2 is favored to realize Eq. (1). Σ_Φ and Σ_S represent the two-point self-energies of the fields Φ and S , respectively, which are related to the thermal mass corrections. (Their explicit forms can be found in Ref. [42].) On the other hand, v'_S is obtained from the tadpole condition at $(0, v'_S)$ as

$$(v'_{SC})^3 + \frac{2(b_1 + b_2 + 2\Sigma_S)}{d_2} v'_{SC} + \frac{4\sqrt{2}a_1}{d_2} = 0. \quad (46)$$

When real solutions exist, v'_S is scaled by $1/\sqrt{d_2}$, implying that a smaller d_2 is required to obtain a larger v_C . However, by comparing the potential energies of the two vacua and imposing the condition that the potential takes its global minimum at (v, v_S) , the case with $\delta_2 \gg 1$ and $d_2 \ll 1$ is excluded. An appropriate choice for realizing a strong first-order EWPT is $\delta_2 = \mathcal{O}(1)$ and $d_2 = \mathcal{O}(1)$. As seen from Eq. (35), in the Higgs-degenerate region, a small v_S and a large mixing angle θ are favored in order to realize a moderately large δ_2 . Moreover, as seen from Eq. (36), d_2 depends inversely on v_S^2 (and partly on v_S^3 through the a_1 term). With the small value of v_S required to realize an appropriate δ_2 , the value of d_2 becomes highly sensitive to small variations in a_1 and m_{h_2} .

Figure 1 illustrates the parameter dependence of PBH formation in the (δ_2, d_2) plane. Here, we take $v_S = 0.6$ GeV and $\theta = \pi/4$ as representative values to realize appropriate magnitudes of δ_2 and d_2 in the degenerate scalar scenario. The red band indicates the region where the PBH fraction satisfies $10^{-4} \leq f_{\text{PBH}} \leq 1$,

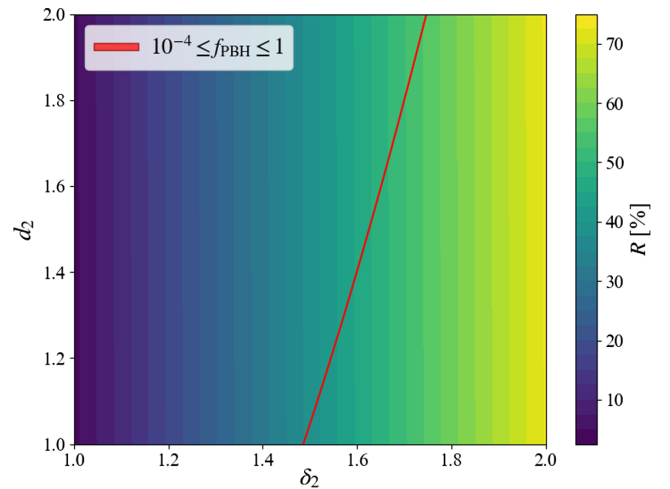


FIG. 1. Parameter dependence of the PBH fraction in the (δ_2, d_2) plane. Here we take $v_S = 0.6$ GeV and $\theta = \pi/4$ as representative values to realize appropriate magnitudes of δ_2 and d_2 in the degenerate scalar scenario. The red line represents the region where $10^{-4} \leq f_{\text{PBH}} \leq 1$. The color map shows the deviation of the Higgs triple coupling from the SM value defined as Eq. (47), which serves as an indicator of the EWPT strength.

TABLE I. Benchmark points used in the analysis. The left columns show the input parameters, while the right columns summarize the derived parameters δ_2 and d_2 computed from Eqs. (35) and (36).

Benchmark	Input parameters							Derived parameters	
	v [GeV]	m_{h_1} [GeV]	m_{h_2} [GeV]	θ [rad]	a_1 [GeV ³]	v_S [GeV]	m_χ [GeV]	δ_2	d_2
BP1	246.22	125	124	$\pi/4$	-6576.1720	0.6	62.5	1.6855	1.7442
BP2	246.22	125	124	$\pi/4$	-6576.1723	0.6	62.5	1.6855	1.7402
BP3	246.22	125	124	$\pi/4$	-6576.1724	0.6	62.5	1.6855	1.7389
BP4	246.22	125	123.999995	$\pi/4$	-6576.1720	0.6	62.5	1.6855	1.7407
BP5	246.22	125	123.999993	$\pi/4$	-6576.1720	0.6	62.5	1.6855	1.7393

corresponding to efficient PBH formation. This narrow red band, which almost appears as a single line, indicates that f_{PBH} is extremely sensitive to small variations in the parameters δ_2 and d_2 . In fact, as shown in Eqs. (17) and (18), the PBH fraction f_{PBH} is proportional to the probability of PBH formation $P(t_N)$, which is itself proportional to the decay rate Γ in Eq. (2). Since Γ is approximately given by $e^{-S_3/T}$, we eventually obtain $f_{\text{PBH}} \propto e^{-S_3/T}$. Therefore, PBH production is highly sensitive to the details of S_3/T . The upward slope reflects that a stronger first-order EWPT tends to occur for larger δ_2 and smaller d_2 , and that efficient PBH formation is realized only in a limited region of this parameter space. The color map shows the deviation of the Higgs triple coupling from the SM value, $g_{hhh}^{\text{SM}} = 3m_h^2/v$, defined as

$$R \equiv \frac{g_{h_1 h_1 h_1} - g_{hhh}^{\text{SM}}}{g_{hhh}^{\text{SM}}} \times 100[\%] \quad \text{with } \mathcal{L} = g_{h_1 h_1 h_1} h_1^3. \quad (47)$$

Although R is not directly related to PBH formation, it provides a useful indicator of the strength of the EWPT. The triple coupling of the SM-like state can be written as

$$g_{h_1 h_1 h_1} = \frac{3}{2} [\lambda v \cos^3 \theta + \delta_2 v_S \cos^2 \theta \sin \theta + \delta_2 v \cos \theta \sin^2 \theta + d_2 v_S \sin^3 \theta]. \quad (48)$$

Among these terms, the term proportional to d_2 is suppressed by the small singlet VEV v_S , which we take to be small in the Higgs-degenerate setup. As a result, the contribution of the d_2 term to $g_{h_1 h_1 h_1}$ becomes subdominant, and the dependence of the triple coupling on d_2 is hardly visible in our parameter scan.

Next, as an illustrative example, we fix the value of δ_2 and slightly vary d_2 . For clarity, instead of using the parameters (δ_2, d_2) , we reexpress the parameter space in terms of the physical quantities (m_{h_2}, a_1) , which are more directly related to observable properties of the model. This representation allows us to visualize how a small variation in m_{h_2} or a_1 can lead to a significant change in

S_3/T , reflecting the strong sensitivity of PBH formation to these parameters.

The benchmark points are summarized in Table I. We treat $(v, m_{h_1}, m_{h_2}, \theta, a_1, v_S, m_\chi)$ as input parameters, while (δ_2, d_2) are derived quantities calculated from Eqs. (35) and (36). BP1 serves as the reference point, where the value of $\delta_2 = 1.6855$ is chosen so that the model parameters reproduce the setup studied in Refs. [42,43]. Reference [49] notes that the mass difference $|m_{h_1} - m_{h_2}| \lesssim 3$ GeV has not been conclusively ruled out by LHC experiments [69]. In addition, at the ILC 250 [70], the recoil mass technique may allow for the discrimination of a 1 GeV mass difference, but resolving mass differences smaller than 1 GeV would require even higher precision. Thus, we focus on the Higgs mass differences of around 1 GeV. The DM mass m_χ is fixed at 62.5 GeV as an example, since it has a negligible impact on the dynamics of the EWPT and PBH formation. In BP2 and BP3, the parameter a_1 is varied with all other inputs fixed, whereas in BP4 and BP5, m_{h_2} is slightly varied instead. The strong sensitivity of d_2 to these small variations originates from the small value of v_S adopted in our setup. As seen from Eq. (36), d_2 scales as $1/v_S^2$ or partly $1/v_S^3$, so that even a tiny change in a_1 or m_{h_2} leads to a shift in d_2 , as shown in Table I. In contrast, the dependence of δ_2 on m_{h_2} is comparatively weak. These benchmark points are therefore suitable for probing how small variations in a_1 or m_{h_2} can influence S_3/T , and consequently, the PBH abundance.

TABLE II. Parameters characterizing the EWPT for each benchmark point.

Benchmark	EWPT parameters			
	v_C [GeV]	T_C [GeV]	v_C/T_C	T_N [GeV]
BP1	201.05	105.41	1.9073	56.442
BP2	201.20	105.31	1.9106	54.751
BP3	201.25	105.28	1.9116	54.094
BP4	201.18	105.32	1.9102	54.965
BP5	201.23	105.28	1.9113	54.283

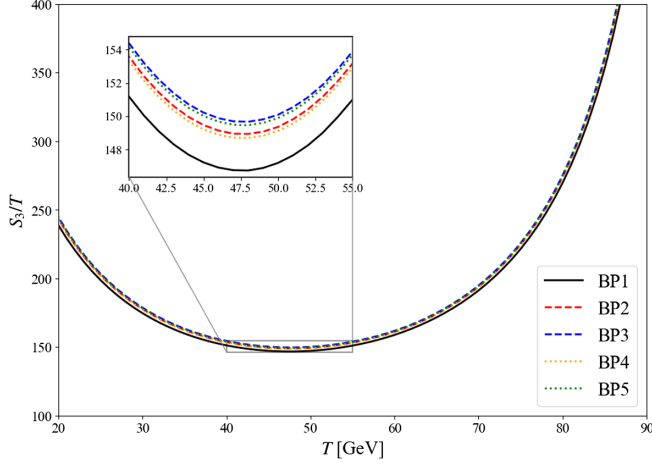


FIG. 2. Temperature dependence of S_3/T , which characterizes the strength of the EWPT. The inset shows an enlarged view around the minimum of S_3/T .

The parameters relevant to the EWPT are summarized in Table II. v_C , T_C , and T_N , defined in Eq. (3), are somewhat sensitive to such small parameter changes. A strong first-order EWPT with $v_C/T_C > 1$ is realized at all benchmark points. It is also seen that a smaller d_2 generally results in a slightly stronger first-order EWPT, as reflected in the larger values of v_C/T_C . Furthermore, Fig. 2 shows the temperature dependence of S_3/T , which reflects the strength of the EWPT. All benchmark points exhibit U-shaped S_3/T curves, as typically seen in a strong first-order EWPT. The inset highlights the region near the minimum of S_3/T . Benchmark points corresponding to a stronger first-order EWPT show slightly larger values of S_3/T .

The results of the PBH analysis are summarized in Table III and visualized in Fig. 3. The quantities shown in Table III are defined in Sec. II. T_d denotes the reference temperature of the delayed-decayed region; T_{PBH} is the PBH formation temperature; f_{PBH} in Eq. (17) represents the present PBH fraction; and M_{PBH} in Eq. (19) is the PBH mass, normalized by the solar mass $M_\odot$. Figure 3 shows the predicted PBH abundance f_{PBH} as a function of the PBH

TABLE III. Parameters related to PBH formation for each benchmark point. Here, T_d is the reference temperature of the delayed-decayed region; T_{PBH} is the temperature at which PBH is produced; f_{PBH} denotes the present PBH fraction; and M_{PBH} is the PBH mass, which is normalized by the solar mass $M_\odot$.

PBH-related parameters				
Benchmark	T_d [GeV]	T_{PBH} [GeV]	f_{PBH}	$M_{\text{PBH}}/M_\odot$
BP1	44.292	29.375	3.01×10^{-214}	6.95×10^{-6}
BP2	46.581	21.364	2.09×10^{-12}	1.03×10^{-5}
BP3	47.464	15.567	5.68	1.47×10^{-5}
BP4	46.287	21.960	3.55×10^{-21}	1.00×10^{-5}
BP5	47.208	17.647	3.68×10^{-2}	1.27×10^{-5}

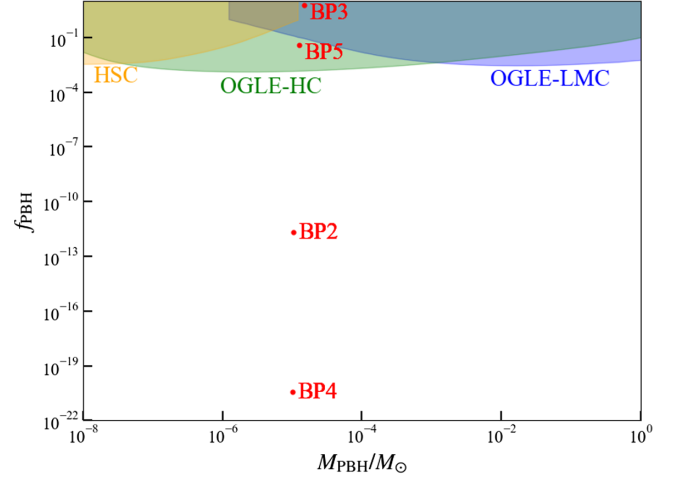


FIG. 3. Predicted PBH abundance f_{PBH} as a function of the PBH mass, normalized by the solar mass ($M_{\text{PBH}}/M_\odot$). The colored shaded regions represent current microlensing constraints: Subaru HSC [54] (HSC, orange), the long-term OGLE Large Magellanic Cloud survey (OGLE-LMC, blue) [55], and the OGLE high-cadence Magellanic Clouds survey (OGLE-HC, green) [56]. The red points indicate the benchmark scenarios BP2–BP5, while BP1 is omitted because its predicted PBH fraction is too small to appear in the plot.

mass normalized by the solar mass ($M_{\text{PBH}}/M_\odot$), together with the existing microlensing constraints from Subaru HSC [54], the long-term OGLE Large Magellanic Cloud survey (OGLE-LMC) [55], and the OGLE high-cadence Magellanic Clouds survey (OGLE-HC) [56]. The red points correspond to the benchmark points BP2–BP5, while BP1 is not visible, since its PBH fraction is too small to appear in the plot. As seen from Table III, the predicted PBH masses are all around $M_{\text{PBH}} \sim 10^{-5} M_\odot$, reflecting the similar nucleation temperatures among the benchmark points.

V. GRAVITATIONAL WAVE IN SYNERGY WITH LEPTON COLLIDER SIGNATURES

In addition to PBH formation, the first-order EWPT can generate observable GWs, providing an independent probe of the underlying dynamics of the phase transition. Together with collider measurements, such as the deviation of the Higgs triple coupling from its SM value, GWs offer a complementary way to explore the electroweak-scale physics responsible for the phase transition. In this section, we study the GW signal predicted in the present model and discuss its detectability in future space-based interferometers, as well as its correlation with collider observables. Since the favored parameter region for the PBH production could lead to obvious deviation of the Higgs triple coupling as shown in Eq. (47), future lepton colliders, such as CEPC, FCC-ee, and ILC could be able to test this scenario through loop effects.

In our analysis, we employ the formalism of Ref. [71] to evaluate the GW signals. One of the key quantities that characterize the dynamics of the phase transition is the strength parameter α , defined as the ratio of the released vacuum energy density to the radiation energy density of the plasma:

$$\alpha \equiv \frac{\Delta\rho_v}{\rho_r} = \frac{[\Delta V - T \frac{d\Delta V}{dT}]_{T=T_*}}{\frac{\pi^2}{30} g_* T_*^4}, \quad (49)$$

where ΔV is the potential energy difference between the false and true vacua, T_* denotes the characteristic temperature of the transition, and g_* is the effective number of relativistic degrees of freedom at T_* . Here, we adopt $g_* = 108.75$. A larger value of α corresponds to a stronger first-order transition.

Another important parameter characterizing the transition is the characteristic length scale, which corresponds to the typical distance between nucleated bubbles at the time of the transition. This scale is expected to carry most of the released energy, and therefore plays a crucial role in determining the GWs generated by bulk fluid motion and bubble collisions. The mean bubble separation is usually adopted as this characteristic length scale. Following

Ref. [71], the mean bubble separation, or the effective bubble radius R_* , is given by

$$R_* = n_b^{-1/3}, \quad n_b = \int_{t_c}^t dt' \frac{a(t')^3}{a(t)^3} \Gamma(t') F(t'), \quad (50)$$

where n_b denotes the number density of bubbles nucleated up to time t .

Sources of GWs arise from bubble collisions [72–77], sound waves [78–81], and turbulence induced by the percolation process [82–87]. The total GW spectrum is the sum of these three contributions:

$$h_{100}^2 \Omega_{\text{GW}}(f) = h_{100}^2 \Omega_{\text{co}}(f) + h_{100}^2 \Omega_{\text{sw}}(f) + h_{100}^2 \Omega_{\text{turb}}(f). \quad (51)$$

Here, $h_{100} = H_0/(100 \text{ km s}^{-1} \text{ Mpc}^{-1}) = 0.67$ is the dimensionless Hubble parameter from the Planck data, with H_0 being the Hubble constant today [88]. We compute the GW spectrum using the R_* -based fitting formulas of Ref. [71]. The spectral contributions from bubble collisions, sound waves, and turbulence are given, respectively, by

$$\begin{aligned} h_{100}^2 \Omega_{\text{co}}(f) &\simeq 1.67 \times 10^{-5} \left(\frac{H_* R_*}{(8\pi)^{1/3}} \right)^2 \left(\frac{\kappa_{\text{co}} \alpha}{1 + \alpha} \right)^2 \left(\frac{100}{g_*} \right)^{1/3} \frac{0.11 v_w}{0.42 + v_w^2} \frac{3.8(f/f_{\text{co}})^{2.8}}{1 + 2.8(f/f_{\text{co}})^{3.8}}, \\ h_{100}^2 \Omega_{\text{sw}}(f) &\simeq 1.64 \times 10^{-6} (H_* \tau_{\text{sw}}) (H_* R_*) \left(\frac{\kappa_v \alpha}{1 + \alpha} \right)^2 \left(\frac{100}{g_*} \right)^{1/3} \left(\frac{f}{f_{\text{sw}}} \right)^3 \left[\frac{7}{4 + 3(f/f_{\text{sw}})^2} \right]^{7/2}, \\ h_{100}^2 \Omega_{\text{turb}}(f) &\simeq 1.14 \times 10^{-4} (H_* R_*) \left(\frac{\kappa_{\text{turb}} \alpha}{1 + \alpha} \right)^{3/2} \left(\frac{100}{g_*} \right)^{1/3} \frac{(f/f_{\text{turb}})^3}{(1 + f/f_{\text{turb}})^{11/3} (1 + 8\pi f/H_*)}. \end{aligned} \quad (52)$$

Here, H_* and R_* are evaluated at the percolation temperature $T_* = T_p$, defined by $F(T_p) = 0.7$. The duration of the sound-wave source is

$$\tau_{\text{sw}} = \min \left[\frac{1}{H_*}, \frac{R_*}{\bar{U}_f} \right], \quad (53)$$

where the root-mean-square fluid velocity is approximated as $\bar{U}_f^2 \simeq \frac{3}{4} \frac{\kappa_v \alpha}{1 + \alpha}$.

For $v_w \simeq 1$, the efficiency factors are approximated by [71,89,90]

$$\begin{aligned} \kappa_{\text{co}} &\simeq \frac{1}{1 + 0.715\alpha} \left(0.715\alpha + \frac{4}{27} \sqrt{\frac{3\alpha}{2}} \right), \\ \kappa_v &\simeq \frac{\alpha}{0.73 + 0.083\sqrt{\alpha} + \alpha}, \quad \kappa_{\text{turb}} \simeq (0.05 - 0.1) \kappa_v. \end{aligned} \quad (54)$$

In our numerical analysis, we set $v_w = 0.95$ and $\kappa_{\text{turb}} = 0.1 \kappa_v$ for illustration. The corresponding peak frequencies are

$$\begin{aligned} f_{\text{co}} &\simeq 1.65 \times 10^{-5} \text{ Hz} \frac{(8\pi)^{1/3}}{H_* R_*} \frac{0.62 v_w}{1.8 - 0.1 v_w + v_w^2} \left(\frac{T_*}{100 \text{ GeV}} \right) \left(\frac{g_*}{100} \right)^{1/6}, \\ f_{\text{sw}} &\simeq 2.6 \times 10^{-5} \text{ Hz} \frac{1}{H_* R_*} \left(\frac{T_*}{100 \text{ GeV}} \right) \left(\frac{g_*}{100} \right)^{1/6}, \\ f_{\text{turb}} &\simeq 7.91 \times 10^{-5} \text{ Hz} \frac{1}{H_* R_*} \left(\frac{T_*}{100 \text{ GeV}} \right) \left(\frac{g_*}{100} \right)^{1/6}. \end{aligned} \quad (55)$$

To quantitatively assess the detectability of the predicted GWs, we compute the signal-to-noise ratio (SNR) for each benchmark point. The SNR is evaluated as follows:

$$\text{SNR} = \sqrt{\mathcal{T} \int_{f_{\min}}^{f_{\max}} df \left(\frac{h_{100}^2 \Omega_{\text{GW}}(f)}{h_{100}^2 \Omega_{\text{sens}}(f)} \right)^2}, \quad (56)$$

where $\mathcal{T}$ denotes the observation time, and $h_{100}^2 \Omega_{\text{sens}}(f)$ represents the detector sensitivity curve [91]. In our analysis, we adopt $\mathcal{T} \simeq 1.0 \times 10^8$ s for both LISA and TianQin, corresponding to approximately 4 yr of observation time. A GW signal is considered detectable if $\text{SNR} > \text{SNR}_{\text{thr}}$, where we set the threshold value $\text{SNR}_{\text{thr}} = 5$.

The GW spectra in this work are computed using the semianalytic prescription developed in Ref. [71], which provides a practical framework for estimating GW signals from strongly supercooled first-order phase transitions based on the mean bubble separation. More refined treatments of GW production in the supercooling regime have also been discussed in the literature, such as those presented in Refs. [92,93]. In particular, Ref. [93] incorporates the distribution of bubble radii and collision times through a radius-dependent efficiency factor, showing that these effects mainly modify the overall normalization of the GW spectrum while leaving its spectral shape largely unchanged. As a result, although quantitative differences can arise among different prescriptions due to scheme dependence, they lead to consistent qualitative trends for the GW signal as a function of the phase transition parameters. Since the focus of the present work is on the correlations between PBH formation, GW observables, and collider signatures, the use of the prescription of Ref. [71] is sufficient for the purposes of our analysis.

Table IV summarizes the key parameters characterizing the GW and collider phenomenology of the benchmark points. The quantities T_p , α_p , and $(R_*)_p$ represent the percolation temperature, the strength of the phase transition, and the mean bubble separation at T_p , respectively.

Larger α_p and R_* correspond to a stronger and more prolonged phase transition, which in turn leads to more energetic GW production. The deviation of the Higgs triple coupling from the SM value R is around 50% for all benchmark points. At the one-loop level, this Higgs triple coupling can increase the SM Zh cross section ($\delta\sigma_{Zh} \equiv \sigma_{Zh}/\sigma_{Zh}^{\text{SM}} - 1$) about $\delta\sigma_{Zh} = 0.82\%$ at 240 GeV CEPC [94,95]. Under the CEPC TDR nominal operating conditions, the cross section is projected to be measured with a relative precision of 0.26% [96,97]. Recent advances in AI-assisted reconstruction and analysis suggest that this precision could be significantly improved, potentially reaching 0.1% in an aggressive scenario [96,97]. One can see that CEPC has the ability to test this scenario with high precision. Similar discussions could also be performed at ILC and FCC-ee. For ILC, the expected precision of the Zh cross section is about 0.4% [98]. For FCC-ee, the expected precision is about 0.1% [99,100]. Current Higgs data at LHC require the Higgs triple coupling modification to be within the range $(-0.4, 6.3)$. At the High-Luminosity LHC (HL-LHC), with a projected integrated luminosity of 3 ab^{-1} of high-quality data per experiment, the ATLAS and CMS Collaborations anticipate determining the Higgs triple coupling with a precision better than 30% [101]. Since this scenario can modify the Higgs triple coupling around 50%, CEPC, ILC, FCC-ee, and HL-LHC are able to test this scenario. We also list the SNRs for LISA and TianQin, computed under the assumption of four years of observation. All benchmark points yield $\text{SNR} > 5$ for LISA and comparable values for TianQin, implying that the predicted GW signals would be well within the reach of both detectors. Overall, this table illustrates that in the CxSM, a strong first-order EWPT leading to PBH formation can simultaneously give rise to sizable deviations in the Higgs self-coupling and detectable GWs, providing mutually complementary probes of the same underlying dynamics.

Figure 4 visualizes the GW spectra corresponding to the SNR results summarized in Table IV. The colored curves represent the predicted spectra for the benchmark points,

TABLE IV. Parameters relevant to the GW and collider phenomenology. Here, T_p , α_p , and $(R_*)_p$ denote the percolation temperature, the phase transition strength, and the mean bubble separation at T_p , respectively. R is the relative deviation of the Higgs triple coupling from the SM value defined in Eq. (47) and the predicted modification of the Zh production cross section, $\delta\sigma_{Zh} \equiv \sigma_{Zh}/\sigma_{Zh}^{\text{SM}} - 1$, which reflects the loop-level effect of the enhanced Higgs self-coupling. SNRs for LISA and TianQin are computed assuming an observation time of $\mathcal{T} \simeq 1.0 \times 10^8$ s. The line styles correspond to those used in Fig. 4.

GW and collider-related parameters								
Benchmark	T_p [GeV]	α_p	$(R_*)_p$ [GeV $^{-1}$]	R [%]	$\delta\sigma_{Zh}$	SNR@LISA	SNR@TianQin	GW
BP1	43.581	0.62256	7.07×10^{13}	51.25	0.82%	93.35	16.81	Solid black line
BP2	38.047	1.0108	1.61×10^{14}	51.25	0.82%	59.15	9.32	Dashed red line
BP3	35.287	1.3321	2.22×10^{14}	51.25	0.82%	47.48	6.48	Dashed blue line
BP4	38.848	0.93730	1.46×10^{14}	51.25	0.82%	63.02	10.32	Dotted orange line
BP5	36.121	1.2222	2.03×10^{14}	51.25	0.82%	50.82	7.25	Dotted green line

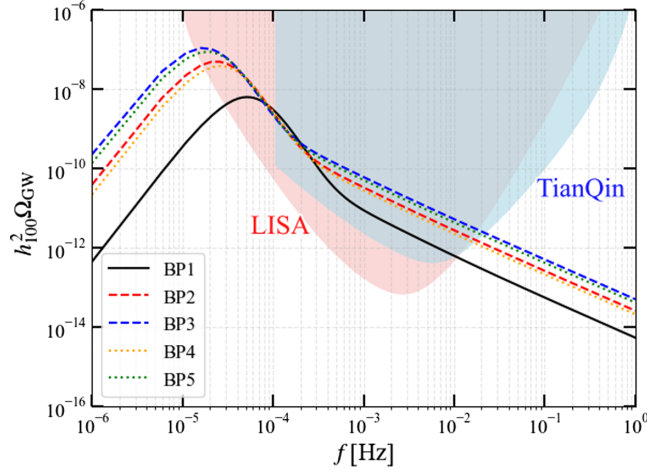


FIG. 4. The predicted GW spectra for the benchmark points BP1–BP5. The shaded regions indicate the sensitivity curves corresponding to $\text{SNR} = 5$ of future space-based GW detectors, LISA [102–104] (red) and TianQin [105,106] (blue). GW signals above these curves are detectable with $\text{SNR} > 5$.

while the shaded regions indicate the expected sensitivity ranges of future space-based detectors LISA (red) and TianQin (blue). As expected from the SNR values, the overall amplitude of the GW spectra increases with the strength of the first-order EWPT. This trend reflects the correlation between a strongly supercooled transition, efficient PBH formation, and enhanced GW production. All benchmark points yield GW spectra within the sensitivity bands of both LISA and TianQin, demonstrating that the parameter space leading to PBH formation can be simultaneously tested by upcoming GW observations.

Beyond these specific benchmark realizations, the points shown above are chosen as representative examples to illustrate the multimessenger implications of a strong first-order EWPT in the CxSM. More generally, the underlying correlations among different observables are driven by the properties of the scalar potential. In particular, a stronger EWPT typically leads to a larger released vacuum energy, resulting in a stronger stochastic GW signal, while simultaneously inducing a larger deviation in the Higgs triple coupling. PBH formation through delayed vacuum decay is even more sensitive to the EWPT dynamics, as the PBH abundance depends exponentially on the nucleation rate and can vary dramatically across the allowed parameter space. Therefore, while the benchmark points provide concrete realizations where PBH formation, GW signals, and collider observables coexist, the complementarity of these probes is a generic feature associated with a strong first-order EWPT.

VI. CONCLUSIONS AND DISCUSSIONS

In this work, we have investigated the possibility of PBH formation induced by a first-order EWPT in a realistic

model—namely, the CxSM. This model contains a scalar DM candidate and two Higgs bosons, and it is known to be compatible with both DM direct-detection constraints and Higgs search results when the two Higgs states are nearly degenerate in mass.

We have shown that the strength of the EWPT and the resulting PBH fraction are strongly dependent on the quartic couplings δ_2 and d_2 in the scalar potential. A strong first-order EWPT, required for efficient PBH formation, is realized for $\delta_2 = \mathcal{O}(1)$ and $d_2 = \mathcal{O}(1)$, which are consistent with the tree-level vacuum stability conditions. In addition, due to the Higgs mass degeneracy, a small singlet VEV v_S and a large mixing angle θ are found to be necessary for realizing a strong first-order EWPT.

To further explore the sensitivity of PBH formation, we have analyzed benchmark points that differ only by tiny variations in a_1 and m_{h_2} . Even such small shifts induce large changes in the Euclidean action S_3/T , and hence in the PBH fraction, reflecting the double-exponential sensitivity $f_{\text{PBH}} \propto e^{-S_3/T}$. In practice, a larger f_{PBH} is obtained for stronger first-order EWPT, while milder transitions yield negligible fractions. Depending on the benchmark points, we found fractions of phenomenological interest that remain consistent with current microlensing bounds. This sensitivity also implies that theoretical and numerical choices—e.g., the thermal resummation scheme—can lead to appreciable differences in the predicted f_{PBH} ; a systematic assessment of such scheme dependence is left for future work. Related studies can be found in Ref. [107].

We have also analyzed the GW signals generated by the first-order EWPT using the R_* formalism, which accounts for the contributions from bubble collisions, sound waves, and turbulence. The resulting spectra reflect the strength of the transition, with stronger first-order EWPT producing larger GW amplitudes. All benchmark points yield SNRs exceeding 10 for LISA and comparable sensitivities for TianQin, indicating that the predicted signals are within the reach of future space-based detectors. In the same parameter region, the deviation of the Higgs triple coupling from its SM value reaches about 50% (testable at CEPC, ILC, FCC-ee, and HL-LHC) and therefore modifies the $e^+e^- \rightarrow Zh$ cross section about 0.82% at one-loop level at 240 GeV CEPC. These results suggest that both GW and collider measurements can provide complementary probes of this realistic PBH model from EWPT.

In summary, we have demonstrated that in the degenerate scalar scenario of the CxSM, a strong first-order EWPT can simultaneously lead to PBH formation, observable GW signals, and measurable deviations in the Higgs self-coupling. These results reveal the multimessenger nature of electroweak-scale physics, showing that PBH, GW, and collider observations together provide powerful and complementary means of probing their

common origin from a first-order EWPT in the early Universe.

ACKNOWLEDGMENTS

We are grateful to Masanori Tanaka and Tomo Takahashi for helpful discussions on the PBH calculations, to Siyu Jiang for valuable advice on the GW computations, and to Chen Zhou for discussions on the future sensitivity of Higgs triple coupling at HL-LHC. This work is supported by the National Natural Science Foundation of China (NNSFC) Grants No. 12475111, No. 12205387, and the

Fundamental Research Funds for the Central Universities, Sun Yat-sen University.

DATA AVAILABILITY

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

-
- [1] S. Hawking, *Mon. Not. R. Astron. Soc.* **152**, 75 (1971).
 - [2] B. J. Carr and S. W. Hawking, *Mon. Not. R. Astron. Soc.* **168**, 399 (1974).
 - [3] B. J. Carr, *Astrophys. J.* **201**, 1 (1975).
 - [4] J. Garcia-Bellido, A. D. Linde, and D. Wands, *Phys. Rev. D* **54**, 6040 (1996).
 - [5] M. Kawasaki, N. Sugiyama, and T. Yanagida, *Phys. Rev. D* **57**, 6050 (1998).
 - [6] P. Ivanov, P. Naselsky, and I. Novikov, *Phys. Rev. D* **50**, 7173 (1994).
 - [7] J. Yokoyama, *Astron. Astrophys.* **318**, 673 (1997).
 - [8] S. W. Hawking, I. G. Moss, and J. M. Stewart, *Phys. Rev. D* **26**, 2681 (1982).
 - [9] H. Kodama, M. Sasaki, and K. Sato, *Prog. Theor. Phys.* **68**, 1799 (1982).
 - [10] I. G. Moss, *Phys. Rev. D* **50**, 676 (1994).
 - [11] R. V. Konoplich, S. G. Rubin, A. S. Sakharov, and M. Y. Khlopov, *Phys. At. Nucl.* **62**, 1593 (1999).
 - [12] H. Deng and A. Vilenkin, *J. Cosmol. Astropart. Phys.* **12** (2017) 044.
 - [13] H. Deng, *J. Cosmol. Astropart. Phys.* **09** (2020) 023.
 - [14] K. Kawana and K.-P. Xie, *Phys. Lett. B* **824**, 136791 (2022).
 - [15] K. Hashino, S. Kanemura, and T. Takahashi, *Phys. Lett. B* **833**, 137261 (2022).
 - [16] Y. Gouttenoire and T. Volansky, *Phys. Rev. D* **110**, 043514 (2024).
 - [17] Y. Gouttenoire, *Phys. Lett. B* **855**, 138800 (2024).
 - [18] S. Kanemura, M. Tanaka, and K.-P. Xie, *J. High Energy Phys.* **06** (2024) 036.
 - [19] D. Gonçalves, A. Kaladharan, and Y. Wu, *Phys. Rev. D* **111**, 035009 (2025).
 - [20] S. Balaji, M. Fairbairn, and M. O. Olea-Romacho, *Phys. Rev. D* **109**, 075048 (2024).
 - [21] M. Lewicki, P. Toczec, and V. Vaskonen, *Phys. Dark Universe* **50**, 102075 (2025).
 - [22] K. Hashino, S. Kanemura, T. Takahashi, M. Tanaka, and C.-M. Yoo, *J. Cosmol. Astropart. Phys.* **09** (2025) 006.
 - [23] K. Murai, K. Sakurai, and F. Takahashi, *J. High Energy Phys.* **07** (2025) 065.
 - [24] A. H. Guth and E. J. Weinberg, *Phys. Rev. D* **23**, 876 (1981).
 - [25] M. Lewicki, P. Toczec, and V. Vaskonen, *J. High Energy Phys.* **09** (2023) 092.
 - [26] K. Kajantie, M. Laine, K. Rummukainen, and M. E. Shaposhnikov, *Phys. Rev. Lett.* **77**, 2887 (1996).
 - [27] M. Quiros, *arXiv:hep-ph/9901312*.
 - [28] M. Dine, R. G. Leigh, P. Y. Huet, A. D. Linde, and D. A. Linde, *Phys. Rev. D* **46**, 550 (1992).
 - [29] S. Balaji, J. Gonçalves, D. Marfatia, A. P. Morais, and R. Pasechnik, *J. Cosmol. Astropart. Phys.* **10** (2025) 064.
 - [30] Q.-H. Cao, M. Tanaka, J.-C. Wang, K.-P. Xie, and J.-J. Zhang, *arXiv:2505.23007*.
 - [31] M. Kierkla, N. Ramberg, P. Schicho, and D. Schmitt, *arXiv:2506.15496*.
 - [32] H.-X. Zhang, K. Hashino, H. Ishida, and S. Matsuzaki, *J. High Energy Phys.* **10** (2025) 207.
 - [33] V. Barger, P. Langacker, M. McCaskey, M. Ramsey-Musolf, and G. Shaughnessy, *Phys. Rev. D* **79**, 015018 (2009).
 - [34] V. Barger, M. McCaskey, and G. Shaughnessy, *Phys. Rev. D* **82**, 035019 (2010).
 - [35] M. Gonderinger, H. Lim, and M. J. Ramsey-Musolf, *Phys. Rev. D* **86**, 043511 (2012).
 - [36] R. Coimbra, M. O. P. Sampaio, and R. Santos, *Eur. Phys. J. C* **73**, 2428 (2013).
 - [37] M. Jiang, L. Bian, W. Huang, and J. Shu, *Phys. Rev. D* **93**, 065032 (2016).
 - [38] C.-W. Chiang, M. J. Ramsey-Musolf, and E. Senaha, *Phys. Rev. D* **97**, 015005 (2018).
 - [39] W. Cheng and L. Bian, *Phys. Rev. D* **98**, 023524 (2018).
 - [40] B. Grzadkowski and D. Huang, *J. High Energy Phys.* **08** (2018) 135.
 - [41] N. Chen, T. Li, Y. Wu, and L. Bian, *Phys. Rev. D* **101**, 075047 (2020).
 - [42] G.-C. Cho, C. Idegawa, and E. Senaha, *Phys. Lett. B* **823**, 136787 (2021).
 - [43] G.-C. Cho, C. Idegawa, and E. Senaha, *Phys. Rev. D* **106**, 115012 (2022).
 - [44] F. Egle, M. Mühlleitner, R. Santos, and J. a. Viana, *Phys. Rev. D* **106**, 095030 (2022).

- [45] G.-C. Cho, C. Idegawa, and R. Sugihara, *Phys. Lett. B* **839**, 137757 (2023).
- [46] C. Idegawa and E. Senaha, *Phys. Lett. B* **848**, 138332 (2024).
- [47] G.-C. Cho, C. Idegawa, and R. Inumiya, *Nucl. Phys. B* **1007**, 116688 (2024).
- [48] K. Funakubo and C. Idegawa, [arXiv:2506.23199](#).
- [49] S. Abe, G.-C. Cho, and K. Mawatari, *Phys. Rev. D* **104**, 035023 (2021).
- [50] G.-C. Cho and C. Idegawa, *Nucl. Phys. B* **994**, 116320 (2023).
- [51] J. Aalbers *et al.* (LZ Collaboration), *Phys. Rev. Lett.* **135**, 011802 (2025).
- [52] G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. D* **101**, 012002 (2020).
- [53] CMS Collaboration, CERN Report No. CMS-PAS-HIG-19-005, 2020.
- [54] H. Niikura *et al.*, *Nat. Astron.* **3**, 524 (2019).
- [55] P. Mróz *et al.*, *Nature (London)* **632**, 749 (2024).
- [56] P. Mróz *et al.*, *Astrophys. J. Lett.* **976**, L19 (2024).
- [57] E. Witten, *Phys. Rev. D* **30**, 272 (1984).
- [58] C. J. Hogan, *Mon. Not. R. Astron. Soc.* **218**, 629 (1986).
- [59] C. Grojean, G. Servant, and J. D. Wells, *Phys. Rev. D* **71**, 036001 (2005).
- [60] S. Kanemura, Y. Okada, and E. Senaha, *Phys. Lett. B* **606**, 361 (2005).
- [61] J. Liu, L. Bian, R.-G. Cai, Z.-K. Guo, and S.-J. Wang, *Phys. Rev. D* **105**, L021303 (2022).
- [62] A. D. Linde, *Nucl. Phys. B* **216**, 421 (1983); **B223**, 544(E) (1983).
- [63] J. Ellis, M. Lewicki, and J. M. No, *J. Cosmol. Astropart. Phys.* **04** (2019) 003.
- [64] T. Harada, C.-M. Yoo, and K. Kohri, *Phys. Rev. D* **88**, 084051 (2013); **89**, 029903(E) (2014).
- [65] I. Musco, V. De Luca, G. Franciolini, and A. Riotto, *Phys. Rev. D* **103**, 063538 (2021).
- [66] A. Escrivà, C. Germani, and R. K. Sheth, *Phys. Rev. D* **101**, 044022 (2020).
- [67] G. Franciolini, Y. Gouttenoire, and R. Jinno, [arXiv:2503.01962](#).
- [68] R.-G. Cai, Y.-S. Hao, and S.-J. Wang, *Sci. China Phys. Mech. Astron.* **67**, 290411 (2024).
- [69] V. Khachatryan *et al.* (CMS Collaboration), *Eur. Phys. J. C* **74**, 3076 (2014).
- [70] K. Fujii *et al.*, [arXiv:1710.07621](#).
- [71] X. Wang, F. P. Huang, and X. Zhang, *J. Cosmol. Astropart. Phys.* **05** (2020) 045.
- [72] A. Kosowsky, M. S. Turner, and R. Watkins, *Phys. Rev. D* **45**, 4514 (1992).
- [73] A. Kosowsky, M. S. Turner, and R. Watkins, *Phys. Rev. Lett.* **69**, 2026 (1992).
- [74] A. Kosowsky and M. S. Turner, *Phys. Rev. D* **47**, 4372 (1993).
- [75] M. Kamionkowski, A. Kosowsky, and M. S. Turner, *Phys. Rev. D* **49**, 2837 (1994).
- [76] C. Caprini, R. Durrer, and G. Servant, *Phys. Rev. D* **77**, 124015 (2008).
- [77] S. J. Huber and T. Konstandin, *J. Cosmol. Astropart. Phys.* **09** (2008) 022.
- [78] M. Hindmarsh, S. J. Huber, K. Rummukainen, and D. J. Weir, *Phys. Rev. Lett.* **112**, 041301 (2014).
- [79] J. T. Giblin, Jr. and J. B. Mertens, *J. High Energy Phys.* **12** (2013) 042.
- [80] J. T. Giblin and J. B. Mertens, *Phys. Rev. D* **90**, 023532 (2014).
- [81] M. Hindmarsh, S. J. Huber, K. Rummukainen, and D. J. Weir, *Phys. Rev. D* **92**, 123009 (2015).
- [82] C. Caprini and R. Durrer, *Phys. Rev. D* **74**, 063521 (2006).
- [83] T. Kahniashvili, A. Kosowsky, G. Gogoberidze, and Y. Maravin, *Phys. Rev. D* **78**, 043003 (2008).
- [84] T. Kahniashvili, L. Campanelli, G. Gogoberidze, Y. Maravin, and B. Ratra, *Phys. Rev. D* **78**, 123006 (2008); **79**, 109901(E) (2009).
- [85] T. Kahniashvili, L. Kisslinger, and T. Stevens, *Phys. Rev. D* **81**, 023004 (2010).
- [86] C. Caprini, R. Durrer, and G. Servant, *J. Cosmol. Astropart. Phys.* **12** (2009) 024.
- [87] P. Binetruy, A. Bohe, C. Caprini, and J.-F. Dufaux, *J. Cosmol. Astropart. Phys.* **06** (2012) 027.
- [88] N. Aghanim *et al.* (Planck), *Astron. Astrophys.* **641**, A6 (2020); **652**, C4(E) (2021).
- [89] J. R. Espinosa, T. Konstandin, J. M. No, and G. Servant, *J. Cosmol. Astropart. Phys.* **06** (2010) 028.
- [90] C. Grojean and G. Servant, *Phys. Rev. D* **75**, 043507 (2007).
- [91] K. Schmitz, *J. High Energy Phys.* **01** (2021) 097.
- [92] J. Ellis, M. Lewicki, and J. M. No, *J. Cosmol. Astropart. Phys.* **07** (2020) 050.
- [93] M. Lewicki and V. Vaskonen, *Eur. Phys. J. C* **83**, 109 (2023).
- [94] F. P. Huang, P.-H. Gu, P.-F. Yin, Z.-H. Yu, and X. Zhang, *Phys. Rev. D* **93**, 103515 (2016).
- [95] F. P. Huang, Y. Wan, D.-G. Wang, Y.-F. Cai, and X. Zhang, *Phys. Rev. D* **94**, 041702 (2016).
- [96] W. Abdallah *et al.* (CEPC Study Group), *Radiat. Detect. Technol. Methods* **8**, 1 (2024); **9**, 184(E) (2025).
- [97] X. Ai *et al.*, *Chin. Phys. C* **49**, 123108 (2025).
- [98] A. Aryshev *et al.* (ILC International Development Team), [10.2172/1873702](#) (2022).
- [99] A. Abada *et al.* (FCC Collaboration), *Eur. Phys. J. Spec. Top.* **228**, 261 (2019).
- [100] M. Benedikt *et al.* (FCC Collaboration), *Eur. Phys. J. C* **85**, 1468 (2025).
- [101] G. Aad *et al.* (ATLAS and CMS Collaborations), [arXiv:2504.00672](#).
- [102] C. Caprini *et al.*, *J. Cosmol. Astropart. Phys.* **04** (2016) 001.
- [103] P. Amaro-Seoane, H. Audley, S. Babak, J. Baker, E. Barausse, P. Bender, E. Berti, P. Binetruy, M. Born, D. Bortoluzzi *et al.*, [arXiv:1702.00786](#).
- [104] C. Caprini *et al.*, *J. Cosmol. Astropart. Phys.* **03** (2020) 024.
- [105] J. Luo *et al.* (TianQin Collaboration), *Classical Quantum Gravity* **33**, 035010 (2016).
- [106] X.-C. Hu, X.-H. Li, Y. Wang, W.-F. Feng, M.-Y. Zhou, Y.-M. Hu, S.-C. Hu, J.-W. Mei, and C.-G. Shao, *Classical Quantum Gravity* **35**, 095008 (2018).
- [107] P. Bittar, S. Roy, and C. E. M. Wagner, *J. High Energy Phys.* **12** (2025) 021.