

95 GeV Higgs boson and spontaneous CP -violation at the finite temperature

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The ATLAS and CMS collaborations reported a diphoton excess in the invariant mass distribution around the 95.4 GeV with a local significance of 3.1σ . Moreover, there is another 2.3σ local excess in the $b\bar{b}$ final state at large electron-positron collider in the same mass region. A plausible solution is that the Higgs sector is extended to include an additional Higgs boson with a mass of 95.4 GeV. We study a complex singlet scalar extension of the two-Higgs-doublet model in which the 95.4 GeV Higgs is from the mixing of three CP -even Higgs fields. In addition, the extended Higgs potential can achieve spontaneous CP -violation at the finite temperature and restore CP symmetry at the present temperature of the Universe. We find that the model can simultaneously explain the baryon asymmetry of the Universe, the diphoton and $b\bar{b}$ excesses around the 95.4 GeV while satisfying various relevant constraints including the experiments of collider and electric dipole moment.

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I. INTRODUCTION

The baryon asymmetry of the Universe (BAU) is a fundamental question in particle physics and cosmology. By the observation based on the big bang nucleosynthesis, the BAU is [1]

$$Y_B \equiv \rho_B/s = (8.2 - 9.2) \times 10^{-11}, \quad (1)$$

where s is the entropy density, and ρ_B is the baryon number density. The three necessary Sakharov ingredients for generating such an asymmetry dynamically: baryon number violation, C and CP violations, and a departure from thermal equilibrium [2]. The electroweak baryogenesis (EWBG) [3,4] is a promising and attractive mechanism of explaining the BAU. In order to achieve the EWBG, the SM need be extended to produce sufficient large CP -violation and a strongly first-order electroweak phase transition (PT), such as the singlet extension of SM (see for example [5–20]) and the two-Higgs-doublet model (2HDM) (see for example [21–41]).

The explicit CP -violation interactions in the lagrangian required by the EWBG can be severely constrained by the negative results in the electric dipole moment (EDM) searches for electrons [42]. The spontaneous CP -violation at finite temperature offers a valid mechanism for solving the problem. In the mechanism, no explicit CP -violating

interactions are added to the Lagrangian. Because of thermal corrections to the effective potential, the vacuum structure of the theory changes at finite temperature, and can dynamically induce sufficient large CP -violation to obtain the observed BAU via electroweak PT. However, as the temperature decreases, the observed electroweak vacuum is obtained and the CP symmetry is restored at the present temperature. Thus, there are no new CP -violating interactions beyond the Cabibbo-Kobayashi-Maskawa (CKM) mechanism at the present temperature, and no new corrections to the EDM. As a result, the stringent EDM experimental constraints will be naturally avoided. The novel mechanism has been realized in the singlet scalar extension of the SM [14,15] in which one need introduce a high dimension effective operator, the singlet pseudoscalar extension of 2HDM [43,44], and the complex singlet scalar extension of 2HDM respecting a discrete dark CP symmetry in which the dark matter is simultaneously explained [45].

Based on full Run 2 dataset, the CMS collaboration published the latest result of searching for low-mass diphoton, and confirmed a previously released excess at $m_{\gamma\gamma} = 95.4$ GeV with a local significance of 2.9σ [46]. The ATLAS collaboration also presented the result of searching for diphoton signals in the mass range from 66 to 110 GeV based on their full Run 2 dataset [47]. The ATLAS search found a diphoton excess with a local significance of 1.7σ at the same mass value as the one that was previously reported by CMS. Neglecting possible correlations, the combined signal strength corresponds to a 3.1σ local excess [48],

$$\mu_{\gamma\gamma}^{\text{exp}} = \mu_{\gamma\gamma}^{\text{ATLAS+CMS}} = \frac{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}{\sigma_{\text{SM}}(pp \rightarrow h_{95.4}^{\text{SM}} \rightarrow \gamma\gamma)} = 0.24_{-0.08}^{+0.09}, \quad (2)$$

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where ϕ is a non-SM scalar with a mass of 95.4 GeV, and $h_{95.4}^{\text{SM}}$ is a hypothetical SM-like Higgs with the same mass. Interestingly, there is a longstanding 2.3σ local excess in the $e^+e^- \rightarrow Z(\phi \rightarrow b\bar{b})$ searches at large electron-positron collider in the same mass region [49],

$$\mu_{b\bar{b}}^{\text{exp}} = 0.117 \pm 0.057. \quad (3)$$

There are numerous discussions on the excesses in the new physics models, for example Refs. [48,50–96]. In this paper, we propose a complex singlet scalar extension of 2HDM (2HDMS) to accommodate the diphoton and $b\bar{b}$ excesses around 95.4 GeV. Meanwhile, we discuss the spontaneous CP -violation at the finite temperature and possibility of explanation on the BAU.

The structure of this paper is outlined as follows. In Sec. II, we introduce some characteristic features of the 2HDMS. In Secs. III and IV, we discuss relevant theoretical and experimental constraints, and explain the diphoton and $b\bar{b}$ excesses around the 95.4 GeV. In Secs. V and VI, we discuss the PTs of the Universe and the possibility of explaining the BAU, respectively. Finally, we draw our conclusion in Sec. VII.

II. THE MODEL

The SM is augmented with a second Higgs doublet Φ_2 and a complex singlet S ,

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \frac{(v_1 + \rho_1 + i\eta_1)}{\sqrt{2}} \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \frac{(v_2 + \rho_2 + i\eta_2)}{\sqrt{2}} \end{pmatrix}, \quad S = \frac{(\rho_s + v_s + i\eta_s)}{\sqrt{2}}, \quad (4)$$

where v_1 , v_2 , and v_s are the real vacuum expectation values (VEVs) acquired by the fields Φ_1 , Φ_2 , and S , respectively. We define the ratio of the two VEVs as $\tan\beta \equiv v_2/v_1$.

The scalar potential is given by

$$\begin{aligned} V = & m_{11}^2(\Phi_1^\dagger\Phi_1) + m_{22}^2(\Phi_2^\dagger\Phi_2) + \frac{\lambda_1}{2}(\Phi_1^\dagger\Phi_1)^2 + \frac{\lambda_2}{2}(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\ & + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \left[\frac{\lambda_5}{2}(\Phi_1^\dagger\Phi_2)^2 + \lambda_6(\Phi_1^\dagger\Phi_1)(\Phi_1^\dagger\Phi_2) + \lambda_7(\Phi_2^\dagger\Phi_2)(\Phi_1^\dagger\Phi_2) + \text{H.c.} \right] \\ & + m_s^2 S S^* + \left[\frac{m_s'^2}{2} S S + \text{H.c.} \right] + \left[\frac{\mu_s}{6} S^3 + \frac{\mu_s'}{2} S S S^* + \text{H.c.} \right] \\ & \left[\frac{\lambda_1''}{24} S^4 + \frac{\lambda_2''}{6} S^2 S S^* + \text{H.c.} \right] + \frac{\lambda_3''}{4} (S S^*)^2 + S S^* [\lambda_1' \Phi_1^\dagger \Phi_1 + \lambda_2' \Phi_2^\dagger \Phi_2] + [S^2 (\lambda_4' \Phi_1^\dagger \Phi_1 + \lambda_5' \Phi_2^\dagger \Phi_2) + \text{H.c.}] \\ & + [(\lambda_6' S S^* + \lambda_7' S S + \lambda_8' S^* S^*) \Phi_2^\dagger \Phi_1 + \text{H.c.}] + \left[-m_{12}^2 \Phi_2^\dagger \Phi_1 + \frac{\mu}{2} (S - S^*) \Phi_1^\dagger \Phi_2 + \text{H.c.} \right] \end{aligned} \quad (5)$$

We focus on the CP -conservation case in which all the coupling coefficients and mass are real. For simplicity, we take $\lambda_6 = \lambda_7 = \lambda_6' = \lambda_7' = \lambda_8' = 0$ and $\mu_s = \mu_s'$ in the following discussions. A softly-broken Z_2 symmetry under which $\Phi_1 \rightarrow \Phi_1$, $S \rightarrow S$ and $\Phi_2 \rightarrow -\Phi_2$, can be imposed to achieve $\lambda_6 = \lambda_7 = \lambda_6' = \lambda_7' = \lambda_8' = 0$.

The minimization conditions of the scalar potential give

$$\begin{aligned} m_{11}^2 = & m_{12}^2 t_\beta - \frac{v^2}{2} (\lambda_1 c_\beta^2 + \lambda_{345} s_\beta^2) - \frac{v_s^2}{2} \lambda_1' - v_s^2 \lambda_4', \\ m_{22}^2 = & m_{12}^2 / t_\beta - \frac{v^2}{2} (\lambda_2 s_\beta^2 + \lambda_{345} c_\beta^2) - \frac{v_s^2}{2} \lambda_2' - v_s^2 \lambda_5', \\ m_s^2 = & -m_s'^2 - \frac{v_s^2}{12} \lambda_1'' - \frac{v_s^2}{3} \lambda_2'' - \frac{v_s^2}{4} v_s^2 \lambda_3'' - \sqrt{2} v_s \mu_s \\ & - \frac{1}{2} v^2 (\lambda_1' + 2\lambda_4') c_\beta^2 - \frac{1}{2} v^2 (\lambda_2' + 2\lambda_5') s_\beta^2, \end{aligned} \quad (6)$$

where $t_\beta \equiv \tan\beta$, $s_\beta \equiv \sin\beta$, $c_\beta \equiv \cos\beta$, and $\lambda_{345} = \lambda_3 + \lambda_4 + \lambda_5$.

After spontaneous symmetry breaking, the mass eigenstates are obtained from the original fields by the rotation matrices,

$$(h_1, h_2, h_3) = (\rho_1, \rho_2, \rho_s) R^T, \quad (7)$$

$$(A, X, G) = (\eta_1, \eta_2, \eta_s) R_A^T, \quad (8)$$

$$\begin{pmatrix} G^\pm \\ H^\pm \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} \phi_1^\pm \\ \phi_2^\pm \end{pmatrix}, \quad (9)$$

with

$$R = \begin{pmatrix} c_1 c_2 & s_1 c_2 & s_2 \\ -s_1 c_3 - c_1 s_2 s_3 & c_1 c_3 - s_1 s_2 s_3 & c_2 s_3 \\ s_1 s_3 - c_1 s_2 c_3 & -s_1 s_2 c_3 - c_1 s_3 & c_2 c_3 \end{pmatrix}, \quad R^A = \begin{pmatrix} -s_\beta c_4 & c_\beta c_4 & s_4 \\ s_\beta s_4 & -c_\beta s_4 & c_4 \\ c_\beta & s_\beta & 0 \end{pmatrix}. \quad (10)$$

The shorthand notations $s_1 \equiv \sin \alpha_1$, $s_2 \equiv \sin \alpha_2$, $s_3 \equiv \sin \alpha_3$, and $s_4 \equiv \sin \alpha_4$. The G^0 and $G^\pm$ are Goldstones, and they are absorbed by gauge bosons Z and $W^\pm$. The remaining physical states includes three CP -even states $h_{1,2,3}$, two pseudoscalars A and X , and a pair of charged scalar $H^\pm$.

Taking the scalar masses and mixing angles as the input parameters, one can express the coupling constants in the Higgs potential as

$$\begin{aligned} \mu &= \frac{\sqrt{2}c_4 s_4 (m_X^2 - m_A^2)}{v}, \\ \lambda_1 &= \frac{m_{h1}^2 R_{11}^2 + m_{h2}^2 R_{21}^2 + m_{h3}^2 R_{31}^2 - m_{12}^2 t_\beta}{v^2 c_\beta^2}, \\ \lambda_2 &= \frac{m_{h1}^2 R_{12}^2 + m_{h2}^2 R_{22}^2 + m_{h3}^2 R_{32}^2 - m_{12}^2 t_\beta^{-1}}{v^2 s_\beta^2}, \\ \lambda_3 &= \frac{m_{h1}^2 R_{11} R_{12} + m_{h2}^2 R_{21} R_{22} + m_{h3}^2 R_{31} R_{32} + 2m_{H^\pm}^2 s_\beta c_\beta - m_{12}^2}{v^2 s_\beta c_\beta}, \\ \lambda_4 &= \frac{(m_A^2 c_4^2 + m_X^2 s_4^2 - 2m_{H^\pm}^2) s_\beta c_\beta + m_{12}^2}{v^2 s_\beta c_\beta}, \\ \lambda_5 &= \frac{-(m_A^2 c_4^2 + m_X^2 s_4^2) s_\beta c_\beta + m_{12}^2}{v^2 s_\beta c_\beta}, \\ \lambda'_4 &= -\frac{3c_4^2 m_X^2 + 6m_S^2 + 3m_A^2 s_4^2 + 6\lambda'_5 s_\beta^2 v^2 + 3\sqrt{2}\mu_s v_s + (\lambda'_1 + \lambda'_2) v_s^2}{6v^2 c_\beta^2}, \\ \lambda'_1 &= -2\lambda'_4 + \frac{(m_{h1}^2 R_{11} R_{13} + m_{h2}^2 R_{21} R_{23} + m_{h3}^2 R_{31} R_{33})}{v v_s c_\beta}, \\ \lambda'_2 &= -2\lambda'_5 + \frac{(m_{h1}^2 R_{12} R_{13} + m_{h2}^2 R_{22} R_{23} + m_{h3}^2 R_{32} R_{33})}{v v_s s_\beta}, \\ \lambda''_3 &= \frac{6m_{h1}^2 R_{13}^2 + 6m_{h2}^2 R_{23}^2 + 6m_{h3}^2 R_{33}^2 - v_s (6\sqrt{2}\mu_s + \lambda'_1 v_s + 4\lambda'_2 v_s)}{3v_s^2}. \end{aligned} \quad (11)$$

The general Yukawa interactions are written as

$$\begin{aligned} -\mathcal{L} &= Y_{u2} \bar{Q}_L \tilde{\Phi}_2 u_R + Y_{d2} \bar{Q}_L \Phi_2 d_R + Y_{\ell 2} \bar{L}_L \Phi_2 e_R \\ &+ Y_{u1} \bar{Q}_L \tilde{\Phi}_1 u_R + Y_{d1} \bar{Q}_L \Phi_1 d_R + Y_{\ell 1} \bar{L}_L \Phi_1 e_R + \text{H.c.}, \end{aligned} \quad (12)$$

where $Q_L^T = (u_L, d_L)$, $L_L^T = (\nu_L, l_L)$, $\tilde{\Phi}_{1,2} = i\tau_2 \Phi_{1,2}^*$, and $Y_{u1,2}$, $Y_{d1,2}$ and $Y_{\ell 1,2}$ are 3×3 matrices in family space. We assume the Yukawa coupling matrices to be aligned to avoid the tree-level flavor changing neutral current [97,98],

$$\begin{aligned} (Y_{u1})_{ii} &= \frac{\sqrt{2}m_{ui}}{v} \rho_{1u}, & (Y_{u2})_{ii} &= \frac{\sqrt{2}m_{ui}}{v} \rho_{2u}, \\ (Y_{\ell 1})_{ii} &= \frac{\sqrt{2}m_{\ell i}}{v} \rho_{1\ell}, & (Y_{\ell 2})_{ii} &= \frac{\sqrt{2}m_{\ell i}}{v} \rho_{2\ell}, \\ (X_{d1})_{ii} &= \frac{\sqrt{2}m_{di}}{v} \rho_{1d}, & (X_{d2})_{ii} &= \frac{\sqrt{2}m_{di}}{v} \rho_{2d}, \end{aligned} \quad (13)$$

where all the off-diagonal elements are zero, and $\rho_{1f} = (c_\beta - s_\beta \kappa_f)$ and $\rho_{2f} = (s_\beta + c_\beta \kappa_f)$ with $f = u, d, \ell$. κ_f are new free parameters, which determine the Yukawa coupling matrices of up-type quark, down-type quark, and lepton. $X_{d1,2} = V_{dL}^\dagger Y_{d1,2} V_{dR}$ with $V_{dL} \equiv V_{\text{CKM}}$, where $V_{dL,R}$ are unitary matrices which transform the interaction

eigenstates to the mass eigenstates for the left-handed and right-handed down-type quark fields. The Yukawa interactions will explicitly break the Z_2 symmetry mentioned above.

The couplings of the neutral Higgs bosons with respect to the SM are given by

$$\begin{aligned}
y_V^{h_1} &= c_2 c_{\beta 1}, & y_f^{h_1} &= c_2 (c_{\beta 1} - s_{\beta 1} \kappa_f), \\
y_V^{h_2} &\simeq |s_2| s_{\beta 13} + \frac{c_2^2}{2} c_3 s_{\beta 1}, & y_f^{h_2} &\simeq |s_2| (s_{\beta 13} + c_{\beta 13} \kappa_f) + \frac{c_2^2}{2} c_3 (s_{\beta 1} + c_{\beta 1} \kappa_f), \\
y_V^{h_3} &\simeq |s_2| c_{\beta 13} - \frac{c_2^2}{2} c_3 s_{\beta 1}, & y_f^{h_3} &\simeq |s_2| (c_{\beta 13} - s_{\beta 13} \kappa_f) - \frac{c_2^2}{2} c_3 (s_{\beta 1} + c_{\beta 1} \kappa_f), \\
y_V^A &= 0, & y_A^f &= -i\kappa_f (\text{for } u) c_4, & y_f^A &= i\kappa_f c_4 (\text{for } d, \ell), \\
y_V^X &= 0, & y_X^f &= i\kappa_f (\text{for } u) s_4, & y_f^X &= -i\kappa_f s_4 (\text{for } d, \ell),
\end{aligned} \tag{14}$$

where V denotes Z or W . The shorthand notations $c_{\beta 1} \equiv \cos(\beta - \alpha_1)$, $s_{\beta 1} \equiv \sin(\beta - \alpha_1)$, $c_{\beta 13} \equiv \cos(\beta - \alpha_1 - \text{sgn}(s_2)\alpha_3)$, and $s_{\beta 13} \equiv \sin(\beta - \alpha_1 - \text{sgn}(s_2)\alpha_3)$. We take the approximation of $s_2 \simeq \text{sgn}(s_2)(1 - \frac{c_2^2}{2})$ for the expressions in the second and third lines of Eq. (14).

The Yukawa couplings of the charged Higgs are

$$\mathcal{L}_Y = -\frac{\sqrt{2}}{v} H^+ \{ \bar{u}_i [\kappa_d (V_{\text{CKM}})_{ij} m_{dj} P_R - \kappa_u m_{ui} (V_{\text{CKM}})_{ij} P_L] d_j + \kappa_\ell \bar{\nu} m_\ell P_R \ell \} + \text{H.c.}, \tag{15}$$

where $i, j = 1, 2, 3$ are the index of generation.

III. RELEVANT THEORETICAL AND EXPERIMENTAL CONSTRAINTS

In our calculations, we consider the following theoretical and experimental constraints:

(1) The signal data of the 125 GeV Higgs. We take h_2 as the observed 125 GeV state, and apply HiggsTools [99] to calculate the total χ^2 of the latest LHC rate measurements of the observed Higgs boson at 125 GeV. HiggsTools incorporates the codes HiggsSignals [99,100] and HiggsBounds [101,102]. We pay particular attention to the surviving samples with $\chi^2 - \chi_{\text{SM}}^2 < 6.18$, where χ_{SM}^2 denotes the χ^2 value of the SM. These samples are favored

compared to the SM fit result within 2σ confidence level (assuming two-dimensional parameter estimations).

(2) The searches for extra Higgses at the collider and the flavor observables. We employ HiggsTools to pick out the samples which meet the 95% confidence level exclusion limits from searches for the additional Higgs at the collider. SuperIso-v4.1 [103] is used to check the compatibility with $B \rightarrow X_s \gamma$ transitions.

(3) Vacuum stability. We impose the conditions that the vacuum should be a global minimum of the potential, and the potential remains positive when the field values approach infinity, namely boundedness from below. For large values of the fields, the quadratic and cubic terms can be neglected, and the quartic part $V_{4-\text{min}}$ in Eq. (5) is written in matrix form in the basis $B = (\Phi_1^\dagger \Phi_1, \Phi_2^\dagger \Phi_2, \rho_S^2, \eta_S^2)^T$,

$$\begin{aligned}
V_{4-\text{min}} &= C^T \frac{1}{2} \underbrace{\begin{pmatrix} \lambda_1 & \lambda_3 + \Delta & \lambda'_1 + 2\lambda'_4 & \lambda'_1 - 2\lambda'_4 \\ \lambda_3 + \Delta & \lambda_2 & \lambda'_2 + 2\lambda'_5 & \lambda'_2 - 2\lambda'_5 \\ \lambda'_1 + 2\lambda'_4 & \lambda'_2 + 2\lambda'_5 & \frac{\lambda''_1 + 3\lambda''_3 + 4\lambda''_2}{6} & \frac{-\lambda''_1 + \lambda''_3}{2} \\ \lambda'_1 - 2\lambda'_4 & \lambda'_2 - 2\lambda'_5 & \frac{-\lambda''_1 + \lambda''_3}{2} & \frac{\lambda''_1 + 3\lambda''_3 - 4\lambda''_2}{6} \end{pmatrix}}_A C \\
&= \frac{1}{2} C^T A C,
\end{aligned} \tag{16}$$

with $\Delta = 0$ for $\lambda_4 \geq |\lambda_5|$ and $\Delta = \lambda_4 - |\lambda_5|$ for $\lambda_4 < |\lambda_5|$. Boundedness from below demands A is a copositive matrix. Deleting the i th row and column of A , $i = 1, 2, 3, 4$, and four symmetric 3×3 matrices are remained. The copositivity of the symmetric order 3 matrix B with entries b_{ij} , $i, j = 1, 2, 3$ demands [75,104]

$$\begin{aligned}
b_{11} &\geq 0, & b_{22} &\geq 0, & b_{33} &\geq 0, \\
\bar{b}_{12} &= b_{12} + \sqrt{b_{11}b_{22}} \geq 0, \\
\bar{b}_{13} &= b_{13} + \sqrt{b_{11}b_{33}} \geq 0, \\
\bar{b}_{23} &= b_{23} + \sqrt{b_{22}b_{33}} \geq 0, \\
&\sqrt{b_{11}b_{22}b_{33}} + b_{12}\sqrt{b_{33}} + b_{13}\sqrt{b_{22}} + b_{23}\sqrt{b_{11}} \\
&+ \sqrt{2\bar{b}_{12}\bar{b}_{13}\bar{b}_{23}} \geq 0.
\end{aligned} \tag{17}$$

Besides, the copositivity of matrix A requires $\det(A) \geq 0 \vee (\text{adj}A)_{ij} < 0$, for some i, j . The adjugate of A is the transpose of the cofactor matrix of A , $(\text{adj}A)_{ij} = (-1)^{i+j}M_{ji}$. The M_{ij} is the determinant of the submatrix which is obtained by deleting row i and column j of A .

(4) Tree-level perturbative unitarity. We require that the eigenvalues of the $2 \rightarrow 2$ scalar-scalar scattering matrix are below 8π , which is achieved in SPheno-v4.0.5 [105] employing SARAH-SPheno [106].

(5) The oblique parameters. The oblique parameters can impose constraints on the Higgs mass spectrum in the model. For points being in agreement with the experimental observation, it was required that the values of the S and the T parameter are within the 2σ ranges, corresponding to $\chi^2 = 6.18$ for two degrees of freedom. The fit results are given in Ref. [1],

$$S = 0.00 \pm 0.07, \quad T = 0.05 \pm 0.06, \tag{18}$$

with the correlation coefficients, $\rho_{ST} = 0.92$.

In the model, the S and T are approximately calculated by [107,108]

$$\begin{aligned}
S &= \frac{1}{\pi m_Z^2} \left[\sum_{i=1,2,3} (F_S(m_Z^2, m_{h_i}^2, m_A^2)(R_{11}^A R_{i1} + R_{12}^A R_{i2})^2 \right. \\
&\quad + F_S(m_Z^2, m_{h_i}^2, m_X^2)(R_{21}^A R_{i1} + R_{22}^A R_{i2})^2 + F_S(m_Z^2, m_Z^2, m_{h_i}^2)(c_\beta R_{i1} + s_\beta R_{i2})^2 \\
&\quad \left. - m_Z^2 F_{S0}(m_Z^2, m_Z^2, m_{h_i}^2)(c_\beta R_{i1} + s_\beta R_{i2})^2 - F_S(m_Z^2, m_{H^\pm}^2, m_{H^\pm}^2) - F_S(m_Z^2, m_Z^2, m_{\text{ref}}^2) + m_Z^2 F_{S0}(m_Z^2, m_Z^2, m_{\text{ref}}^2) \right], \\
T &= \frac{1}{16\pi m_W^2 s_W^2} \left[\sum_{i=1,2,3} (-F_T(m_{h_i}^2, m_A^2)(R_{11}^A R_{i1} + R_{12}^A R_{i2})^2 - F_T(m_{h_i}^2, m_X^2)(R_{21}^A R_{i1} + R_{22}^A R_{i2})^2 + F_T(m_{H^\pm}^2, m_{h_i}^2)(c_\beta R_{i2} - s_\beta R_{i1})^2 \right. \\
&\quad + 3F_T(m_Z^2, m_{h_i}^2)(c_\beta R_{i1} + s_\beta R_{i2})^2 - 3F_T(m_W^2, m_{h_i}^2)(c_\beta R_{i1} + s_\beta R_{i2})^2 \\
&\quad \left. + F_T(m_{H^\pm}^2, m_A^2)(c_\beta R_{12}^A - s_\beta R_{11}^A)^2 + F_T(m_{H^\pm}^2, m_X^2)(c_\beta R_{22}^A - s_\beta R_{21}^A)^2 - 3F_T(m_Z^2, m_{\text{ref}}^2) + 3F_T(m_W^2, m_{\text{ref}}^2) \right],
\end{aligned} \tag{19}$$

where

$$\begin{aligned}
F_T(a, b) &= \frac{1}{2}(a + b) - \frac{ab}{a - b} \log\left(\frac{a}{b}\right), & F_S(a, b, c) &= B_{22}(a, b, c) - B_{22}(0, b, c), \\
F_{S0}(a, b, c) &= B_0(a, b, c) - B_0(0, b, c)
\end{aligned} \tag{20}$$

with

$$\begin{aligned}
B_{22}(a, b, c) &= \frac{1}{4} \left[b + c - \frac{1}{3}a \right] - \frac{1}{2} \int_0^1 dx X \log(X - i\epsilon), \\
B_0(a, b, c) &= - \int_0^1 dx \log(X - i\epsilon), \\
X &= bx + c(1 - x) - ax(1 - x).
\end{aligned} \tag{21}$$

IV. EXCESSES AROUND 95.4 GeV

The assumed origins of the diphoton and $b\bar{b}$ excesses around the 95.4 GeV are the resonant productions of the lightest CP -even Higgs h_1 in the model. In the narrow width approximation, these signal strengths can be expressed as follows:

$$\mu_{b\bar{b}} = \frac{\sigma_{2\text{HDMS}}(e^+e^- \rightarrow Zh_1)}{\sigma_{\text{SM}}(e^+e^- \rightarrow Zh_{95.4}^{\text{SM}})} \times \frac{\text{BR}_{2\text{HDMS}}(h_1 \rightarrow b\bar{b})}{\text{BR}_{\text{SM}}(h_{95.4}^{\text{SM}} \rightarrow b\bar{b})} = |y_V^{h_1}|^2 \frac{\text{BR}_{2\text{HDMS}}(h_1 \rightarrow b\bar{b})}{\text{BR}_{\text{SM}}(h_{95.4}^{\text{SM}} \rightarrow b\bar{b})} \tag{22}$$

$$\begin{aligned}\mu_{\gamma\gamma} &= \frac{\sigma_{2\text{HDMS}}(gg \rightarrow h_1)}{\sigma_{\text{SM}}(gg \rightarrow h_{95.4}^{\text{SM}})} \times \frac{\text{BR}_{2\text{HDMS}}(h_1 \rightarrow \gamma\gamma)}{\text{BR}_{\text{SM}}(h_{95.4}^{\text{SM}} \rightarrow \gamma\gamma)} \\ &\simeq |y_u^{h_1}|^2 \frac{\text{BR}_{2\text{HDMS}}(h_1 \rightarrow \gamma\gamma)}{\text{BR}_{\text{SM}}(h_{95.4}^{\text{SM}} \rightarrow \gamma\gamma)}\end{aligned}\quad (23)$$

In our calculations, we take the mixing parameters and Yukawa parameters as follows:

$$\begin{aligned}0 \leq c_2 \leq 0.4, \quad s_{\beta 13} &= 1.0, \quad 0.5 \leq \tan \beta \leq 2, \\ 0 \leq s_{\beta 1} \leq 1.0, \quad 0.2 \leq s_4 &\leq 0.7, \\ \kappa_u = 0, \quad \kappa_\ell = 0, \quad -5.0 \leq \kappa_d &\leq 5.0.\end{aligned}\quad (24)$$

For $s_{\beta 13} = 1.0$ and $c_2 \rightarrow 0$, the couplings between the h_2 and SM particles approach to those of SM, which is favored by the signal data of the 125 GeV Higgs. In order to weaken the bounds on the extra Higgses from the direct searches at the LHC and indirect searches via the flavor physics data, we choose $\kappa_u = \kappa_\ell = 0$ and $-5.0 \leq \kappa_d \leq 5.0$. The nonzero κ_d can make $y_d^{h_1}$ to be smaller than $y_V^{h_1}$ and $y_u^{h_1}$, and suppress

the total width of h_1 significantly. As a result, $\text{Br}(h_1 \rightarrow \gamma\gamma)$ is sizably enhanced.

In Figs. 1 and 2 we respectively display the parameter points explaining the $b\bar{b}$ and diphoton excesses at the 95.4 GeV within 1σ and 2σ ranges without imposing any other constraints. From Fig. 1, it can be observed that c_2 is imposed a lower bound for a given $c_{\beta 1}$ since the signal strength $\mu_{b\bar{b}}$ is directly proportional to $|y_{h_1}^V|^2$. The Yukawa coupling of bottom quark to h_1 is proportional to $c_{\beta 1} - s_{\beta 1}\kappa_d$, and therefore a relative large value of $|c_{\beta 1} - s_{\beta 1}\kappa_d|$ is favored to explain the $b\bar{b}$ excess, especially for the explanation within 1σ range. The allowed κ_d and $c_{\beta 1}$ are displayed in the right panel of Fig. 1. Different from the case of the $b\bar{b}$ excess, the explanation of diphoton excess at the 95.4 GeV favors a relative small $|c_{\beta 1} - s_{\beta 1}\kappa_d|$, which can suppress the total width of h_1 , and then enhance $\text{Br}(h_1 \rightarrow \gamma\gamma)$. As a result, a positive κ_d is favored for the case of $c_{\beta 1} > 0$ and $s_{\beta 1} > 0$. When $c_{\beta 1}$ approaches to 1.0, a negative κ_d is allowed to explain the diphoton excess with 2σ range, as shown in the right panel of Fig. 2.

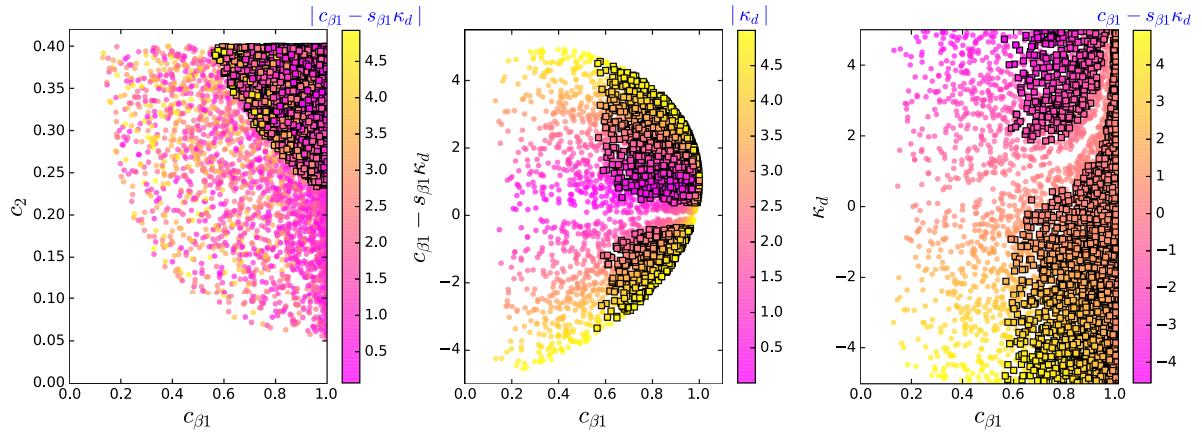


FIG. 1. The samples accommodating the $b\bar{b}$ excess at the 95.4 GeV. The excess is explained within 1σ and 2σ ranges for the squares and circles, respectively.

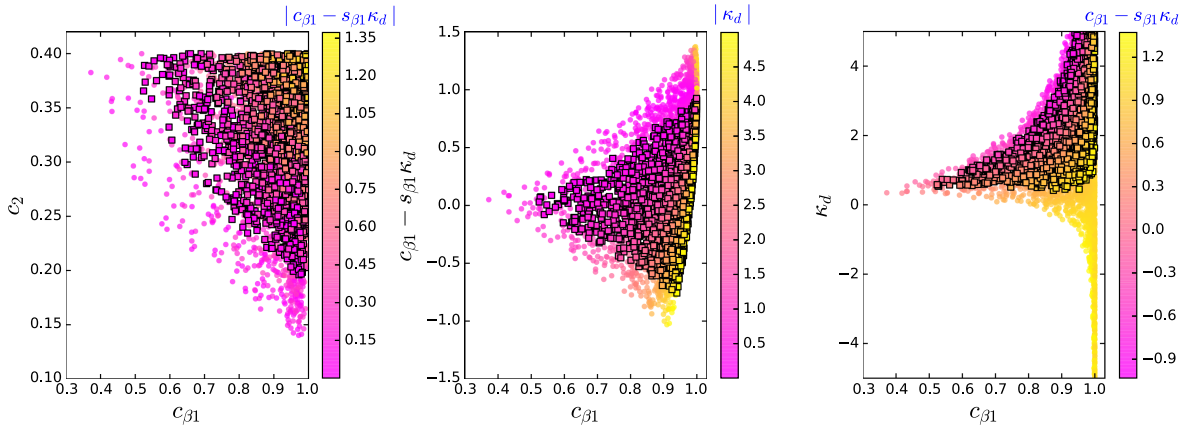


FIG. 2. Same as the Fig. 1, but for the diphoton excess at the 95.4 GeV.

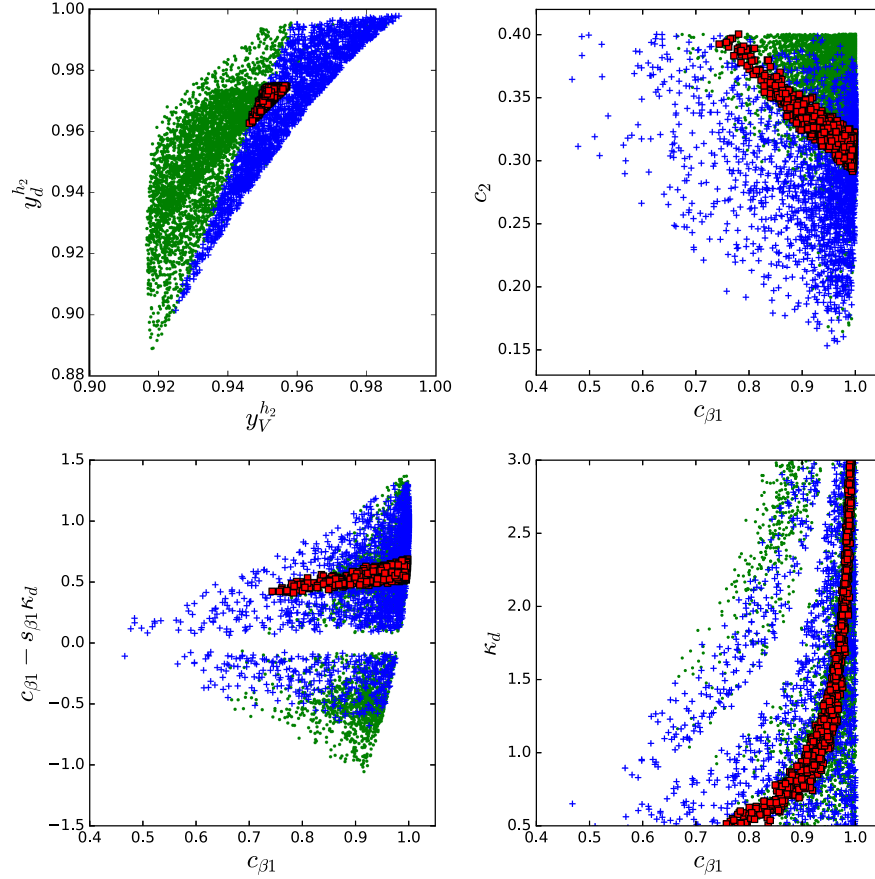


FIG. 3. The circles (green) are excluded by the constraints of “pre-95” while the pluses (blue) and squares (red) are allowed. The $b\bar{b}$ and diphoton excesses are simultaneously explained within 1σ and 2σ ranges for squares (red) and the pluses (blue).

After imposing the constraints of “pre-95” (denoting theoretical constraints, the oblique parameters, the signal data of the 125 GeV Higgs, the searches for extra Higgses at the LHC and the flavor observables), we show the parameter points explaining the diphoton and $b\bar{b}$ excesses at the 95.4 GeV in Fig. 3. We find that the whole parameter space chosen in the Eq. (24) can be consistent with the measurement of $B \rightarrow X_s \gamma$. The choice of $\kappa_u = 0$ can suppress the top Yukawa couplings to H , A , X , and $H^\pm$, and reduce the production cross sections of these extra Higgses at the LHC. As a result, the exclusion bounds of searches for additional Higgs bosons at the LHC can be markedly weakened. The signal data of the 125 GeV Higgs impose stringent constraints on $y_V^{h_2}$ and $y_d^{h_2}$, and exclude many parameter points explaining the $b\bar{b}$ and diphoton excesses. In the region of $0.15 < c_2 < 0.4$ and $0.45 < c_{\beta 1} < 1.0$, the model can simultaneously explain the $b\bar{b}$ and diphoton excesses at the 95.4 GeV within 2σ ranges. The $|\kappa_d|$ is favored to decrease with $c_{\beta 1}$, and $0 < \kappa_d < 1$ is required for the case of $c_{\beta 1} < 0.6$. There is a small region explaining both $b\bar{b}$ and diphoton excesses within 1σ range, in which c_2 , $c_{\beta 1}$, and κ_d have strong dependencies.

V. PHASE TRANSITION

To examine the PTs we first study the effective scalar potential of the model at finite temperature. The neutral elements of Φ_1 and Φ_2 are parametrized as $\frac{\varphi_1}{\sqrt{2}}$ and $\frac{\varphi_2 + i\varphi_3}{\sqrt{2}}$, and the singlet scalar field is parametrized as $\frac{\chi + i\eta_s}{\sqrt{2}}$. It is plausible to choose the imaginary part of the neutral component of Φ_1 to be zero because the potential of Eq. (5) only depends on the relative phase of the neutral components of Φ_1 and Φ_2 .

The full effective potential at finite temperature is gauge-dependent [109,110], which is composed of four parts: the tree level potential, the Coleman-Weinberg term [111], the finite temperature corrections [112] and the resummed daisy corrections [113,114]. In this paper we consider the high-temperature approximation of effective potential, which keeps only the tree-level potential and the thermal mass terms in the high-temperature expansion. Therefore, the effective potential is gauge invariant, and it does not depend on the resummation scheme and the renormalization scheme. The high-temperature approximation of effective potential is written as

$$\begin{aligned}
V_{\text{eff}}(\varphi_1, \varphi_2, \varphi_3, \chi, \eta_s) = & \frac{1}{2}(m_{11}^2 + \Pi_{\varphi_1})\varphi_1^2 + \frac{1}{2}(m_{22}^2 + \Pi_{\varphi_2})(\varphi_2^2 + \varphi_3^2) + \frac{1}{2}(m_S^2 + m_S'^2 + \Pi_\chi)\chi^2 \\
& + \frac{1}{2}(m_S^2 - m_S'^2 + \Pi_{\eta_s})\eta_s^2 + \frac{1}{8}(\lambda_1\varphi_1^4 + \lambda_2\varphi_2^4 + \lambda_3\varphi_3^4) + \frac{1}{4}\lambda_{345}\varphi_1^2\varphi_2^2 + \frac{1}{4}\bar{\lambda}_{345}\varphi_1^2\varphi_3^2 \\
& + \frac{\lambda_2}{4}\varphi_3^2\varphi_2^2 - m_{12}^2\varphi_1\varphi_2 - \frac{\mu}{\sqrt{2}}\varphi_3\eta_s\varphi_1 + \frac{\lambda_1'}{4}(\chi^2 + \eta_s^2)\varphi_1^2 + \frac{\lambda_2'}{4}(\chi^2 + \eta_s^2)(\varphi_2^2 + \varphi_3^2) \\
& + \frac{\lambda_4'}{2}(\chi^2 - \eta_s^2)\varphi_1^2 + \frac{\lambda_5'}{2}(\chi^2 - \eta_s^2)(\varphi_3^2 + \varphi_2^2) + \left(\frac{\lambda_1''}{48} + \frac{\lambda_3''}{16}\right)(\chi^4 + \eta_s^4) + \frac{1}{8}(\lambda_3'' - \lambda_1'')\chi^2\eta_s^2 \\
& + \frac{\lambda_2''}{12}(\chi^4 - \eta_s^4) + \frac{\sqrt{2}}{3}\mu_s\chi^3, \\
\Pi_{\varphi_1} = & \left[\frac{9g^2}{2} + \frac{3g'^2}{2} + 6\lambda_1 + 4\lambda_3 + 2\lambda_4 + 2\lambda_1' + 6y_t^2(c_\beta - s_\beta\kappa_u)^2 + 6y_b^2(c_\beta - s_\beta\kappa_d)^2 \right] \frac{T^2}{24}, \\
\Pi_{\varphi_2} = & \left[\frac{9g^2}{2} + \frac{3g'^2}{2} + 6\lambda_2 + 4\lambda_3 + 2\lambda_4 + 2\lambda_2' + 6y_t^2(s_\beta + c_\beta\kappa_u)^2 + 6y_b^2(s_\beta + c_\beta\kappa_d)^2 \right] \frac{T^2}{24}, \\
\Pi_{\varphi_3} = & \Pi_{\varphi_2}, \\
\Pi_\chi = & [4\lambda_1' + 4\lambda_2' + 2\lambda_2'' + 2\lambda_3'' + 8\lambda_4' + 8\lambda_5'] \frac{T^2}{24}, \\
\Pi_{\eta_s} = & [4\lambda_1' + 4\lambda_2' - 2\lambda_2'' + 2\lambda_3'' - 8\lambda_4' - 8\lambda_5'] \frac{T^2}{24},
\end{aligned} \tag{25}$$

where $\bar{\lambda}_{345} = \lambda_3 + \lambda_4 - \lambda_5$, $y_t = \frac{\sqrt{2}m_t}{v}$, $y_b = \frac{\sqrt{2}m_b}{v}$, and Π_i denotes the thermal mass terms of the field i .

In a first-order PT, the bubble nucleation rate per unit volume at the finite temperature is given by [115–117]

$$\Gamma \approx A(T)e^{-S_3/T}, \tag{26}$$

where $A(T) \sim T^4$ is a prefactor and S_3 is a three-dimensional Euclidian action,

$$S_3 = 4\pi \int_0^\infty dr r^2 \left[\sum_{i=1}^5 \frac{1}{2} \left(\frac{d\phi_i}{dr} \right)^2 + V_{\text{eff}} \right]. \tag{27}$$

The $\phi_{i=1,2,3,4,5}$ denotes φ_1 , φ_2 , φ_3 , χ , and η_s , which are determined by bounce equations [118]

$$\frac{d^2\phi_i}{dr^2} + \frac{2}{r} \frac{d\phi_i}{dr} = \frac{\partial V_{\text{eff}}}{\partial \phi_i}, \quad (i = 1, 2, 3, 4, 5). \tag{28}$$

We employ FindBounce to solve the bounce equations, and obtain the bubble wall VEV profiles [119]. At the nucleation temperature T_n , the thermal tunneling probability for bubble nucleation per horizon volume and per horizon time is of order one, and this condition is roughly converted to the relation $\frac{S_3(T)}{T}|_{T=T_n} \approx 140$.

The CP symmetry is conserved both at very high temperature and the present temperature, and it is only spontaneously broken at the finite temperature. Therefore, the Universe need undergo multistep PTs. It is noted that

the effective potential V_{eff} in Eq. (25) has a Z_2 symmetry: $\varphi_3 \rightarrow -\varphi_3$ and $\eta_s \rightarrow -\eta_s$. Therefore, there will not be a bias between transitions to $(\langle\varphi_1\rangle, \langle\varphi_2\rangle, \langle\varphi_3\rangle, \langle\chi\rangle, \langle\eta_s\rangle)$ and $(\langle\varphi_1\rangle, \langle\varphi_2\rangle, -\langle\varphi_3\rangle, \langle\chi\rangle, -\langle\eta_s\rangle)$ from the origin $(0, 0, 0, 0, 0)$ GeV. The two kinds of bubbles will create baryon asymmetry of opposite signs. As a result, the net BAU averaged over the entire Universe will be zero resulting in a null BAU. The issue can be solved in a special pattern of three-step PTs. The first step is the Z_2 -breaking PT, and the second step is a strongly first-order electroweak PT during which the BAU is generated. Finally, through the third-step PT the observed vacuum is obtained and the CP symmetry is recovered at the present temperature. During the first Z_2 -breaking PT, different patches of the Universe would end up in either $+\langle\eta_s\rangle$ or $-\langle\eta_s\rangle$ vacuum, forming a network of domains. One may add a Z_2 symmetry breaking term, $-i\mu_3(S - S^*)^3$, which leads to a potential difference between the vacua with $\pm\langle\eta_s\rangle$, ΔV . The amount of bias ΔV needed for the deeper minimum to be the only vacuum remaining when the next electroweak PT transition happens is very tiny, $\Delta V/T^4 > 10^{-16}$ [5,120]. Despite the μ_3 term breaks the CP symmetry explicitly, the CP -conservation is still a safe approximation in the model because of a negligible value of μ_3 .

We apply CosmoTransitions to analyze the PTs [121]. After imposing the constraints of “pre-95” and requirement of explaining both $b\bar{b}$ and diphoton excesses within 2σ ranges, we find two different patterns of the three-step PTs with the characteristics mentioned above, and take two benchmark points BP1 and BP2 to discuss the PTs in

TABLE I. Input parameters for the BP1 and the BP2.

	$m_{h_1}(\text{GeV})$	$m_{h_2}(\text{GeV})$	$m_{h_3} = m_{H^\pm}(\text{GeV})$	$m_A(\text{GeV})$	$m_X(\text{GeV})$	$m_{12}^2(\text{GeV})^2$	$m_s^2(\text{GeV})^2$
BP1	95.4	125.0	451.42	80.99	412.39	5314.36	4720.09
BP2	95.4	125.0	570.22	70.40	275.10	2482.07	5067.11

	t_β	$s_{\beta 13}$	c_2	$s_{\beta 1}$	s_4	λ'_5	λ''_1	λ''_2	κ_d	$\kappa_u = \kappa_\ell$	$\mu_s(\text{GeV})$	$v_s(\text{GeV})$	$\mu_{b\bar{b}}$	$\mu_{\gamma\gamma}$
BP1	1.147	1.0	0.277	0.382	0.317	-1.429	3.908	3.166	0.954	0.0	73.48	-50.11	0.0489	0.118
BP2	0.950	1.0	0.247	0.498	0.426	-1.024	4.384	2.324	0.891	0.0	101.63	-67.05	0.0280	0.104

detailed. The key input parameters are displayed in Table I. Figures 4 and 5 exhibit the phase histories for the BP1 and BP2. For the BP1, at a very high temperature, the minimum of the potential is at the origin since the contributions of the thermal mass terms to the effective potential are proportional to T^2 . When the temperature decreases to $T = 83.15$ GeV, a second-order PT starts during which η_s develops a nonzero VEV and the other four

fields still remain zero. At $T = 65.81$ GeV, a strongly first-order electroweak PT happens which breaks electroweak symmetry, $(0, 0, 0, 0, 68.27) \text{ GeV} \rightarrow (68.88, 66.0, 53.86, -40.83, 25.47) \text{ GeV}$ for $(\langle\varphi_1\rangle, \langle\varphi_2\rangle, \langle\varphi_3\rangle, \langle\chi\rangle, \langle\eta_s\rangle)$. The strength of PT ($\xi = \frac{\sqrt{\varphi_1^2 + \varphi_2^2 + \varphi_3^2}}{T}$) is 1.66, and the BAU is generated via the EWBG mechanism. At $T = 56.59$ GeV, another second-order PT takes place. Next, the vacuum evolves along the final phase, and ultimately ends in the

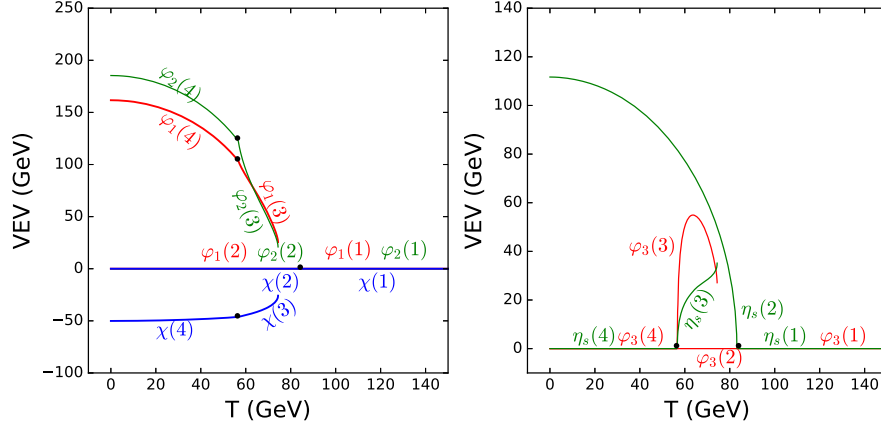


FIG. 4. Phase histories for the BP1. $\phi(i)$ denotes the VEV of ϕ at the i th phase, where $\phi = \varphi_1, \varphi_2, \chi, \varphi_3, \eta_s$. The black dot represents the occurrence of a second-order PT.

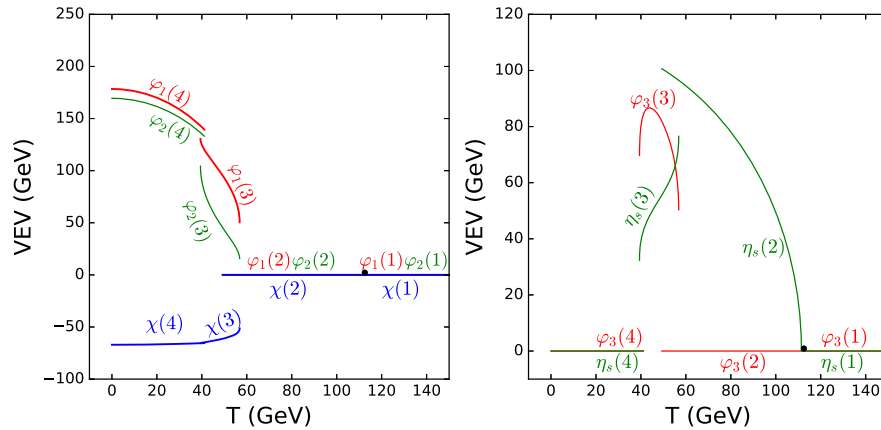


FIG. 5. Same as Fig. 1, but for the BP2.

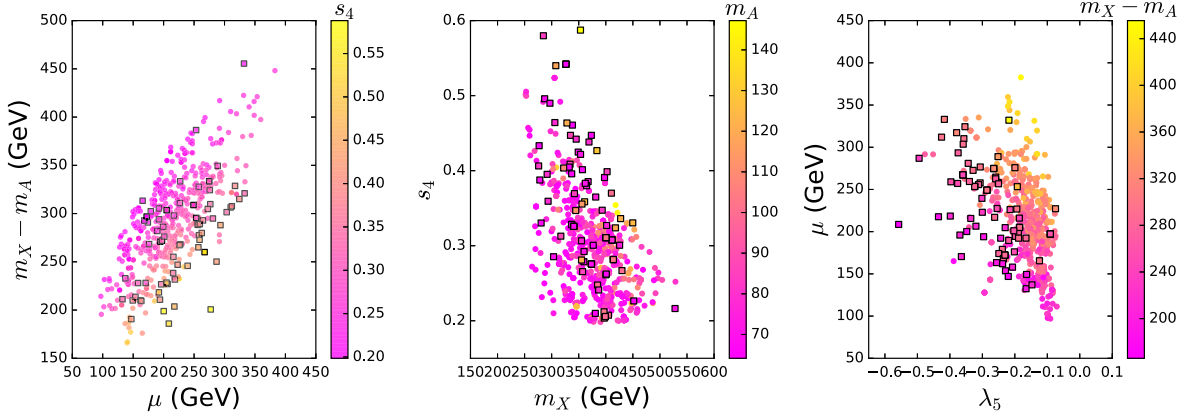


FIG. 6. The scattering plots achieving the three-step PTs with the characteristics mentioned in the text. The third step is respectively first-order PT and second-order PT for the squares and circles.

observed values at $T = 0$ GeV while the CP symmetry is restored. Meanwhile, $\xi > 1$ is always remained so that the BAU is not washed out by the sphaleron process [122].

For the BP2, the first step is second-order PT and the second step is strongly first-order PT, which are similar to those of BP1. However, the third step is first-order PT during which the observed vacuum is obtained and the CP symmetry is restored.

There are no explicit CP -violating interactions in the scalar potential in Eq. (5). Due to the thermal corrections to the effective potential in Eq. (25), the η_s field first acquires a nonzero VEV at finite temperature. Naturally, the mass term and interaction terms of the S field will play an important role. In addition, the μ term [see Eqs. (5) and (25)] can lead to a close correlation between $\langle \varphi_3 \rangle$ and $\langle \eta_s \rangle$ of the potential minimum, which plays a key role in the φ_3 field acquiring a nonzero VEV as the temperature evolves. Thus, an imaginary part of the top quark mass is induced by the interaction of $t\bar{t}(\varphi_2 + i\varphi_3)$.

In Fig. 6 we display some parameter points achieving the three-step PTs. From Fig. 6, we find that the three-step PTs satisfying our requirements favor an appropriate value of μ . As a result, according to Eq. (11), m_A and m_X is required to have a large mass splitting, and the mass splitting tends to increase with a decrease of s_4 . Also m_X is favored to have a large value for a small s_4 . Besides, the PT, whose third step is first order, tends to be achieved in the region of relatively large $|\lambda_5|$.

In BP1 and BP2, the 70–80 GeV lightest CP -odd scalar boson A is predicted. Because of $\kappa_u = \kappa_\ell = 0$, $\kappa_d = 0.954$ (0.891) for BP1 and BP2, the A coupling to the top quark is absent according to Eq. (14), leading to a negligible production cross section via gluon-gluon fusion process at the LHC. The A can be mainly produced via the $pp \rightarrow Ab\bar{b}$ process at the LHC, and then A decays into $b\bar{b}$. However, it is challenging to isolate potential signals from the overwhelming QCD background, and several strategies need be improved, such as precise event selection with

optimized kinematic cuts, and advanced jet substructure techniques to distinguish signal-like jets from QCD jets. Additionally, the machine learning-based analyses can enhance signal-to-background discrimination.

VI. BARYOGENESIS

We take the WKB approximation method to calculate the CP -violating source terms and chemical potentials transport equations of particle species in the wall frame with a radial coordinate z [23,123,124]. When a top quark penetrates the bubble wall, it acquires a complex mass as a function of z ,

$$m_t(z) = \frac{y_t}{\sqrt{2}} \sqrt{(c_\beta \varphi_1(z) + s_\beta \varphi_2(z))^2 + s_\beta^2 \varphi_3^2(z)} e^{i\Theta_t(z)} \quad (29)$$

with

$$\begin{aligned} \Theta_t(z) &= \Theta_Z(z) + \arctan \frac{s_\beta \varphi_3(z)}{c_\beta \varphi_1(z) + s_\beta \varphi_2(z)}, \\ \partial_z \Theta_Z(z) &= - \frac{\varphi_2^2(z) + \varphi_3^2(z)}{\varphi_1^2(z) + \varphi_2^2(z) + \varphi_3^2(z)} \partial_z \Theta(z), \\ \Theta(z) &= \arctan \frac{\varphi_3(z)}{\varphi_2(z)}. \end{aligned} \quad (30)$$

Because the imaginary part of the neutral element of Φ_1 is chosen to be zero in our discussions, the nonvanishing Z_μ field produce an additional CP -violating force acting on the top quark, which is removed by a local axial transformation of top quark, reintroducing an additional overall phase $\Theta_Z(z)$ into m_t [28].

The complex mass of the top quark drives the transport equations, which include effects of the strong sphaleron process (Γ_{ss}) [23,125], W-scattering (Γ_W) [23,126], the top Yukawa interaction (Γ_y) [23,126], the top helicity flips (Γ_M) [23,126], and the Higgs number violation (Γ_h) [23,126].

The transport equations are given by

$$\begin{aligned}
0 &= 3v_W K_{1,t}(\partial_z \mu_{t,2}) + 3v_W K_{2,t}(\partial_z m_t^2) \mu_{t,2} + 3(\partial_z u_{t,2}) \\
&\quad - 3\Gamma_y(\mu_{t,2} + \mu_{t^c,2} + \mu_{h,2}) - 6\Gamma_M(\mu_{t,2} + \mu_{t^c,2}) - 3\Gamma_W(\mu_{t,2} - \mu_{b,2}) \\
&\quad - 3\Gamma_{ss}[(1 + 9K_{1,t})\mu_{t,2} + (1 + 9K_{1,b})\mu_{b,2} + (1 - 9K_{1,t})\mu_{t^c,2}], \\
0 &= 3v_W K_{1,t}(\partial_z \mu_{t^c,2}) + 3v_W K_{2,t}(\partial_z m_t^2) \mu_{t^c,2} + 3(\partial_z u_{t^c,2}) \\
&\quad - 3\Gamma_y(\mu_{t,2} + \mu_{b,2} + 2\mu_{t^c,2} + 2\mu_{h,2}) - 6\Gamma_M(\mu_{t,2} + \mu_{t^c,2}) \\
&\quad - 3\Gamma_{ss}[(1 + 9K_{1,t})\mu_{t,2} + (1 + 9K_{1,b})\mu_{b,2} + (1 - 9K_{1,t})\mu_{t^c,2}], \\
0 &= 3v_W K_{1,b}(\partial_z \mu_{b,2}) + 3(\partial_z u_{b,2}) - 3\Gamma_y(\mu_{b,2} + \mu_{t^c,2} + \mu_{h,2}) - 3\Gamma_W(\mu_{b,2} - \mu_{t,2}), \\
&\quad - 3\Gamma_{ss}[(1 + 9K_{1,t})\mu_{t,2} + (1 + 9K_{1,b})\mu_{b,2} + (1 - 9K_{1,t})\mu_{t^c,2}], \\
0 &= 4v_W K_{1,h}(\partial_z \mu_{h,2}) + 4(\partial_z u_{h,2}) - 3\Gamma_y(\mu_{t,2} + \mu_{b,2} + 2\mu_{t^c,2} + 2\mu_{h,2}) - 4\Gamma_h \mu_{h,2}, \\
S_t &= -3K_{4,t}(\partial_z \mu_{t,2}) + 3v_W \tilde{K}_{5,t}(\partial_z u_{t,2}) + 3v_W \tilde{K}_{6,t}(\partial_z m_t^2) u_{t,2} + 3\Gamma_t^{\text{tot}} u_{t,2}, \\
0 &= -3K_{4,b}(\partial_z \mu_{b,2}) + 3v_W \tilde{K}_{5,b}(\partial_z u_{b,2}) + 3\Gamma_b^{\text{tot}} u_{b,2}, \\
S_t &= -3K_{4,t}(\partial_z \mu_{t^c,2}) + 3v_W \tilde{K}_{5,t}(\partial_z u_{t^c,2}) + 3v_W \tilde{K}_{6,t}(\partial_z m_t^2) u_{t^c,2} + 3\Gamma_t^{\text{tot}} u_{t^c,2}, \\
0 &= -4K_{4,h}(\partial_z \mu_{h,2}) + 4v_W \tilde{K}_{5,h}(\partial_z u_{h,2}) + 4\Gamma_h^{\text{tot}} u_{h,2}.
\end{aligned} \tag{31}$$

The $\mu_{i,2}$ and $u_{i,2}$ are the second-order CP -odd chemical potential and the plasma velocity of the particle $i = t, t^c, b, h$. S_i is the source term,

$$S_i = -v_W K_{8,t} \partial_z (m_t^2 \partial_z \Theta_t) + v_W K_{9,t} (\partial_z \Theta_t) m_t^2 (\partial_z m_t^2). \tag{32}$$

The Γ_i^{tot} are the total reaction rate of the particle i [23,124], and the functions $K_{a,i}$ and $\tilde{K}_{a,i}$ ($a = 1-9$) are given in Ref. [124].

We solve the transport equations with the boundary conditions $\mu_i(z = \pm\infty) = 0$ ($i = t, t^c, b, h$), and obtain the chemical potentials μ_i and velocity perturbations u_i of each particle species. The weak sphaleron process converts the left-handed quark number into a baryon asymmetry, which is calculated as

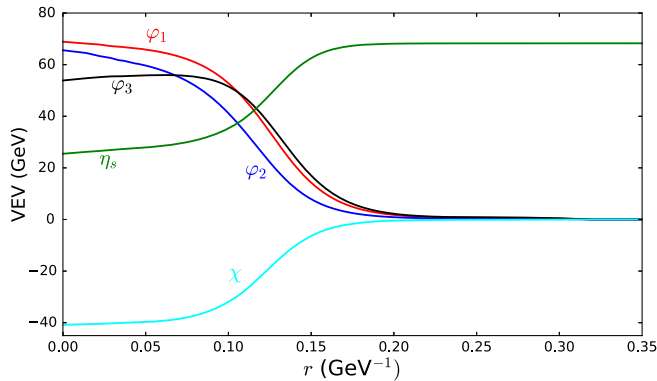


FIG. 7. The radial nucleation bubble wall VEV profiles for the second-step PT of the BP1.

$$Y_B = \frac{405\Gamma_{ws}}{4\pi^2 v_w g_* T_n} \int_0^\infty dz \mu_{B_L}(z) f_{\text{sph}}(z) \exp\left(-\frac{45\Gamma_{ws} z}{4v_w}\right). \tag{33}$$

$\Gamma_{ws} \simeq 10^{-6} T_n$ is the weak sphaleron rate inside bubble [127] and v_w is the wall velocity. μ_{B_L} is the chemical potential of left-handed quarks, and the function $f_{\text{sph}}(z) = \min(1, 2.4 \frac{T_n}{\Gamma_{ws}} e^{-40v_n(z)/T_n})$ with $v_n(z) = \sqrt{\langle \varphi_1(z) \rangle^2 + \langle \varphi_2(z) \rangle^2 + \langle \varphi_3(z) \rangle^2}$ is employed to smoothly interpolate between the sphaleron rate in the unbroken and broken phases [28].

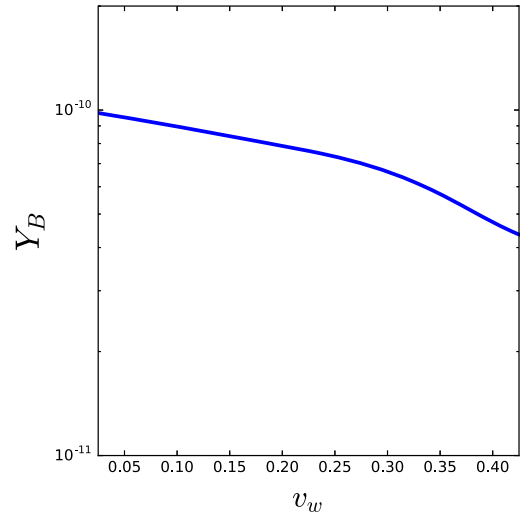


FIG. 8. Y_B as a function of v_w for the BP1.

The calculation of BAU for the BP1 depends on the bubble wall profiles for the second-step PT, and they are given in Fig. 7. The WKB method of calculating transport equations needs the condition of $L_W T_n \gg 1$, where L_W is the width of bubble wall. The $L_W T_n$ for the BP1 is approximately estimated to be 2.6. Figure 8 displays Y_B as a function of the wall velocity v_w for the BP1. Y_B tends to decrease with an increase of v_w , and $8 \times 10^{-11} < Y_B < 10 \times 10^{-11}$ for $0.025 < v_w < 0.19$.

VII. CONCLUSION

In this work we studied a complex singlet scalar extension of the 2HDM in which the 95.4 GeV Higgs is from the mixing of three CP -even Higgs fields, and there is a mixing between the two CP -odd Higgs fields. Considering relevant theoretical and experimental constraints, we found that the

model can simultaneously explain the diphoton and $b\bar{b}$ excesses at the 95.4 GeV while achieving spontaneous CP -violation at the finite temperature. Some strong dependencies among several key parameters are displayed. Finally, we picked out a benchmark point accommodating the diphoton and $b\bar{b}$ excesses simultaneously, and found that the model can explain BAU while satisfying the bound of EDM.

ACKNOWLEDGMENTS

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