

Intrinsic properties of large CP violation in the complex two-Higgs-doublet model

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We investigate the parameter space supporting large CP violation (CPV) in the complex two-Higgs-doublet model with softly broken Z_2 symmetry, where the 125 GeV Higgs boson is identified as the lightest neutral Higgs boson H_1 . Through a comprehensive global scan of type-I and type-II models under theoretical, collider, and eEDM constraints, we identify distinct structures that facilitate large CPV. In type-I, gauge-sector CPV is maximized when the 125 GeV Higgs boson is nearly degenerate with a second neutral scalar. For the ensemble of physically viable points, the predicted eEDM values typically exceed 10^{-31} e cm, placing the model largely within the sensitivity of next-generation experiments. Conversely, type-II models strongly suppress gauge-sector CPV while allowing for nearly maximal CPV in the Yukawa sector. Destructive interference among various contributions allows for $|d_e|$ values as low as $\mathcal{O}(10^{-35})$ e cm, resulting in no phenomenologically relevant lower bound. Finally, we uncover the phenomenon of “hidden CPV” in the near-alignment limit, characterized by CP -violating mixing between the heavy neutral Higgs bosons governed by the angle α_3 . We demonstrate that this hidden CPV can be experimentally probed at future colliders via CP -violating Yukawa interactions of H_2 and H_3 , as well as the robust H_2 - H_3 - Z coupling.

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I. INTRODUCTION

The discovery of the 125 GeV Higgs boson [1,2] has firmly established the mechanism of electroweak symmetry breaking within the Standard Model (SM), yet it has simultaneously intensified the search for new scalar dynamics. A compelling motivation for this search arises from the observed matter–antimatter asymmetry of the Universe [3], whose dynamical generation requires the three Sakharov conditions: baryon number violation, C and CP violation (CPV), and a departure from thermal equilibrium [4].

While the SM satisfies these conditions in principle, it fails in two critical aspects. First, for a 125 GeV Higgs boson, the electroweak phase transition is a smooth

crossover [5–7], failing to provide the required departure from thermal equilibrium. Second, the CPV arising from the Cabibbo-Kobayashi-Maskawa (CKM) matrix is far too small—by roughly ten orders of magnitude—to account for the observed baryon-to-photon ratio [8–10]. These shortcomings point to the necessity of new sources of CPV and additional scalar degrees of freedom near the electroweak scale, making extended Higgs sectors with intrinsic CPV a particularly attractive and well-motivated framework for electroweak baryogenesis.

Among such extensions, the complex two-Higgs-doublet model (C2HDM) with a softly broken Z_2 symmetry stands out as a crucial benchmark [11–14]. In this framework, a single independent CP -violating phase in the scalar potential mixes the three neutral Higgs mass eigenstates H_1 , H_2 , and H_3 , ordered as $M_{H_1} < M_{H_2} < M_{H_3}$. Depending on the Z_2 charge assignments of the right-handed fermions, four Yukawa types emerge—type-I, type-II, type-X, and type-Y. This minimal extension provides all the necessary ingredients for electroweak baryogenesis [13] while remaining predictive and testable through correlated signatures in collider searches [15–29], Higgs precision measurements [27,28,30,31], and electric dipole moment (EDM) experiments [29,31–37].

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Yet, the very predictivity of the C2HDM leads to strong phenomenological tensions. In particular, the nonobservation of the electron EDM (eEDM) [29,38–40] imposes exceptionally stringent constraints on CP violation. The most recent JILA measurement, $|d_e| < 4.1 \times 10^{-30} e \text{ cm}$ (90% CL) [41], improves upon the previous ACME bound, $|d_e| < 1.1 \times 10^{-29} e \text{ cm}$ (90% CL) [42]. This dramatic tightening of the eEDM limit severely restricts the amount of CP violation in the C2HDM, confining the allowed parameter space to highly specific, nontrivial regions.

A primary candidate for such a restricted parameter space is the Higgs alignment regime, which is strongly motivated by the observed SM-like couplings of the 125 GeV Higgs boson. Such alignment may arise naturally from global symmetries of the scalar potential—a framework known as natural Higgs alignment (NHAL) [43]—where the exact symmetric limit typically precludes explicit CP violation. In these scenarios, sizable CP violation can only be attained through the soft or explicit breaking of the underlying NHAL symmetries [44,45]. In addition, the type-I and type-II Yukawa structures are also motivated by symmetry arguments [46,47].

Since the exact Higgs alignment limit is associated with CP conservation, and our focus is on the phenomenology of sizable CP violation, we do not assume such a symmetry-enforced setup. By identifying the 125 GeV state as the lightest neutral Higgs boson and stochastically sampling the remaining degrees of freedom, we determine the viable parameter regions that support sizable CP violation. This random-scan methodology provides direct access to the Higgs-sector interactions responsible for CP violation and their associated collider phenomenology.

While a substantial body of research has already explored collider observables across various benchmark channels [25,31,48–50], these works often focus on specific final states. What remains missing is a systematic characterization of the *intrinsic structure*—namely, the specific mass hierarchies and theoretical patterns—that allows the C2HDM to support large or even maximal CP violation while remaining consistent with the full suite of theoretical and experimental constraints, including the stringent eEDM bounds.

To rigorously analyze this structure, we quantify the magnitude of CP violation using two normalized measures: ξ_V for the gauge sector [51] and ζ_t for the fermionic sector [46]. These metrics allow us to map the landscape of CP violation across the viable parameter space. After isolating all points with sizable ξ_V and ζ_t , we address the following fundamental questions:

- (i) Does the requirement of large CP violation favor specific mass hierarchies, and how do the resulting mass spectra distinguish type-I from type-II scenarios?
- (ii) How are the gauge- and Yukawa-sector CPV measures correlated with the eEDM, and what

characteristic ranges of $|d_e|$ are realized within the viable parameter space?

- (iii) Can sizable CP violation manifest in the near-alignment regime, where it decouples from the 125 GeV Higgs boson?

Our results reveal that type-I and type-II follow two qualitatively distinct pathways to large CP violation when H_1 is identified as the 125 GeV Higgs boson. In type-I, both the gauge- and Yukawa-sector CPV measures can reach significant magnitudes—with ξ_V rising to ~ 0.4 and ζ_t to ~ 0.8 —but only in regions where the 125 GeV Higgs boson is nearly degenerate with H_2 . In contrast, type-II yields a strongly suppressed gauge-sector CPV measure ($\xi_V \lesssim 10^{-3}$), while still allowing for sizable Yukawa CPV with $\zeta_t \sim 0.7$ and $\zeta_b \sim 0.8$. In this case, the maximal values of ζ_t and ζ_b lead to nearly maximal CP violation within the heavy scalar sector involving H_2 and H_3 .

Finally, addressing the last question leads us to uncover a notable phenomenon emerging near the Higgs alignment limit: *hidden CP violation*. We demonstrate that CP -violating mixing between H_2 and H_3 becomes increasingly *maximal* as the Yukawa-sector CPV measures grow.

The remainder of this paper is organized as follows. In Sec. II, we briefly review the structure of the C2HDM and define the CPV measures ξ_V and ζ_t . Section III describes our parameter scan and discusses the characteristics of the viable parameter space with large CPV measures. In Sec. IV, we present the future prospects and observability of the large CPV parameter space, including eEDM prospects, Higgs signal strengths in the near-degenerate type-I scenario, and the CP mixing in the top and τ couplings in type-II. Additionally, we study the hidden CP violation in the near-alignment limit. We conclude in Sec. V.

II. BRIEF REVIEW OF C2HDM

A. Scalar potential, vacuum structure, and neutral scalar mixing

The 2HDM extends the SM by introducing two complex $SU(2)_L$ scalar doublet fields, Φ_1 and Φ_2 , both with hypercharge $Y = 1/2$ under the $Q = T_3 + Y$ convention [12]. These fields are expressed as

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \frac{v_1 + \rho_1 + i\eta_1}{\sqrt{2}} \end{pmatrix} \quad \text{and} \quad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \frac{v_2 + \rho_2 + i\eta_2}{\sqrt{2}} \end{pmatrix}, \quad (1)$$

where v_1 and v_2 are the *real* vacuum expectation values (VEVs) of Φ_1 and Φ_2 , respectively, assuming no spontaneous CP violation. Their ratio defines the key model parameter $t_\beta = v_2/v_1$. For notational simplicity, we use the abbreviations $s_x = \sin x$, $c_x = \cos x$, and $t_x = \tan x$. The combined VEV, $v = \sqrt{v_1^2 + v_2^2} \approx 246 \text{ GeV}$, induces spontaneous electroweak symmetry breaking.

To eliminate tree-level flavor-changing neutral currents (FCNCs), a discrete Z_2 symmetry is imposed, under which

the fields transform as $\Phi_1 \rightarrow \Phi_1$ and $\Phi_2 \rightarrow -\Phi_2$ [47,52]. The scalar potential with a softly broken Z_2 symmetry is given by

$$\begin{aligned} V_\Phi(\Phi_1, \Phi_2) &= m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - (m_{12}^2 \Phi_1^\dagger \Phi_2 + \text{H.c.}) \\ &+ \frac{1}{2} \lambda_1 (\Phi_1^\dagger \Phi_1)^2 + \frac{1}{2} \lambda_2 (\Phi_2^\dagger \Phi_2)^2 + \lambda_3 (\Phi_1^\dagger \Phi_1) (\Phi_2^\dagger \Phi_2) \\ &+ \lambda_4 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1) + \frac{1}{2} [\lambda_5 (\Phi_1^\dagger \Phi_2)^2 + \text{H.c.}], \end{aligned} \quad (2)$$

where the m_{12}^2 term acts as the soft breaking parameter. Due to the Hermiticity of the Lagrangian, only the parameters m_{12}^2 and λ_5 can be complex.

The minimization conditions for the potential are

$$m_{11}^2 v_1 + \frac{\lambda_1}{2} v_1^3 + \frac{\lambda_{345}}{2} v_1 v_2^2 = \text{Re}(m_{12}^2) v_2, \quad (3)$$

$$m_{22}^2 v_2 + \frac{\lambda_2}{2} v_2^3 + \frac{\lambda_{345}}{2} v_1^2 v_2 = \text{Re}(m_{12}^2) v_1, \quad (4)$$

$$2 \text{Im}(m_{12}^2) = v_1 v_2 \text{Im}(\lambda_5), \quad (5)$$

where $\lambda_{345} \equiv \lambda_3 + \lambda_4 + \text{Re}(\lambda_5)$. The condition in Eq. (5) relates two imaginary parts, leaving a single independent CP phase. To accommodate explicit CP violation, we further assume $\phi(\lambda_5) \neq 2\phi(m_{12}^2)$ [53], ensuring that the two phases cannot be simultaneously removed by rephasing Φ_1 and Φ_2 . This framework constitutes the C2HDM.

In the C2HDM, the three neutral Higgs bosons do not possess definite CP quantum numbers. The mass eigenstates are denoted as H_1 , H_2 , and H_3 , with the mass hierarchy $M_{H_1} < M_{H_2} < M_{H_3}$. While the observed 125 GeV Higgs boson could correspond to any of these states, we focus on the “normal scenario” where the lightest state is identified with the observed particle

$$\text{Normal scenario: } M_{H_1} = 125 \text{ GeV.} \quad (6)$$

The Higgs basis $\{\mathcal{H}_1, \mathcal{H}_2\}$ is particularly advantageous as it aligns the electroweak symmetry-breaking VEV with a single doublet (defined here as $\mathcal{H}_1$) [53–57]. The relation between $\{\mathcal{H}_1, \mathcal{H}_2\}$ and $\{\Phi_1, \Phi_2\}$ is given by the rotation

$$\begin{pmatrix} \mathcal{H}_1 \\ \mathcal{H}_2 \end{pmatrix} = \begin{pmatrix} c_\beta & s_\beta \\ -s_\beta & c_\beta \end{pmatrix} \begin{pmatrix} \Phi_1 \\ \Phi_2 \end{pmatrix}, \quad (7)$$

which leads to the component decomposition

$$\begin{aligned} \mathcal{H}_1 &= \begin{pmatrix} G^\pm \\ \frac{1}{\sqrt{2}}(v + H^0 + iG^0) \end{pmatrix} \quad \text{and} \\ \mathcal{H}_2 &= \begin{pmatrix} H^\pm \\ \frac{1}{\sqrt{2}}(R_2 + iI_2) \end{pmatrix}. \end{aligned} \quad (8)$$

Here, H^0 represents the SM Higgs interaction eigenstate before any mixing with the additional neutral scalar components R_2 and I_2 . The Goldstone bosons $G^\pm$ and G^0 reside entirely in $\mathcal{H}_1$ and are absorbed by the $W^\pm$ and Z gauge bosons, respectively, whereas the physical charged Higgs eigenstates $H^\pm$ originate from $\mathcal{H}_2$.

The three neutral mass eigenstates H_i ($i = 1, 2, 3$) arise from the mixing of the two neutral CP -even components [$\rho_{1,2}$ in Eq. (1)] and the neutral CP -odd component [I_2 in Eq. (8)]. The mass-squared matrix of the neutral Higgs bosons, $\mathcal{M}^2$, is defined through

$$V_\Phi \supset \frac{1}{2} \begin{pmatrix} \rho_1 & \rho_2 & I_2 \end{pmatrix} \mathcal{M}^2 \begin{pmatrix} \rho_1 \\ \rho_2 \\ I_2 \end{pmatrix}, \quad (9)$$

where

$$\mathcal{M}^2 = \begin{pmatrix} \bar{m}^2 s_\beta^2 + \lambda_1 v^2 c_\beta^2 & -(\bar{m}^2 - \lambda_{345} v^2) s_\beta c_\beta & -\frac{1}{2} \text{Im}(\lambda_5) v^2 s_\beta \\ -(\bar{m}^2 - \lambda_{345} v^2) s_\beta c_\beta & \bar{m}^2 c_\beta^2 + \lambda_2 v^2 s_\beta^2 & -\frac{1}{2} \text{Im}(\lambda_5) v^2 c_\beta \\ -\frac{1}{2} \text{Im}(\lambda_5) v^2 s_\beta & -\frac{1}{2} \text{Im}(\lambda_5) v^2 c_\beta & \bar{m}^2 - \text{Re}(\lambda_5) v^2 \end{pmatrix}, \quad (10)$$

with $\bar{m}^2 = \text{Re}(m_{12}^2)/(s_\beta c_\beta)$.

The orthogonal mixing matrix $\mathbf{R}$ diagonalizes $\mathcal{M}^2$, yielding mass eigenstates $\{H_a, H_b, H_c\}$

$$\mathcal{M}^2 = \mathbf{R}^T \begin{pmatrix} m_{H_a}^2 & 0 & 0 \\ 0 & m_{H_b}^2 & 0 \\ 0 & 0 & m_{H_c}^2 \end{pmatrix} \mathbf{R}. \quad (11)$$

The matrix $\mathbf{R}$ is parametrized by three mixing angles $\alpha_{1,2,3}$ as

$$\mathbf{R} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{\alpha_3} & s_{\alpha_3} \\ 0 & -s_{\alpha_3} & c_{\alpha_3} \end{pmatrix} \begin{pmatrix} c_{\alpha_2} & 0 & s_{\alpha_2} \\ 0 & 1 & 0 \\ -s_{\alpha_2} & 0 & c_{\alpha_2} \end{pmatrix} \begin{pmatrix} c_{\alpha_1} & s_{\alpha_1} & 0 \\ -s_{\alpha_1} & c_{\alpha_1} & 0 \\ 0 & 0 & 1 \end{pmatrix} \\ = \begin{pmatrix} c_{\alpha_1}c_{\alpha_2} & s_{\alpha_1}c_{\alpha_2} & s_{\alpha_2} \\ -c_{\alpha_1}s_{\alpha_2}s_{\alpha_3} - s_{\alpha_1}c_{\alpha_3} & c_{\alpha_1}c_{\alpha_3} - s_{\alpha_1}s_{\alpha_2}s_{\alpha_3} & c_{\alpha_2}s_{\alpha_3} \\ -c_{\alpha_1}s_{\alpha_2}c_{\alpha_3} + s_{\alpha_1}s_{\alpha_3} & -c_{\alpha_1}s_{\alpha_3} - s_{\alpha_1}s_{\alpha_2}c_{\alpha_3} & c_{\alpha_2}c_{\alpha_3} \end{pmatrix}. \quad (12)$$

The mixing angles lie in the range $\alpha_{1,2,3} \in [-\pi/2, \pi/2]$. Since the hierarchy among m_{H_a} , m_{H_b} , and m_{H_c} is not imposed *a priori*, the states must be reordered according to their masses to define the mass eigenstates $\{H_1, H_2, H_3\}$, satisfying the standard hierarchy $M_{H_1} < M_{H_2} < M_{H_3}$. In summary, the mass eigenstates (H_1, H_2, H_3) are related to the states (ρ_1, ρ_2, I_2) via

$$\begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix} = \mathbf{R} \begin{pmatrix} \rho_1 \\ \rho_2 \\ I_2 \end{pmatrix}. \quad (13)$$

An interesting relation emerges between the mass parameters and the mixing angles. The explicit form of the mass matrix $\mathcal{M}^2$ in Eq. (10) leads to the following identity:

$$\mathbf{T} = \begin{pmatrix} c_{\beta-\alpha_1}c_{\alpha_2} & -c_{\beta-\alpha_1}s_{\alpha_2}s_{\alpha_3} + s_{\beta-\alpha_1}c_{\alpha_3} & -c_{\beta-\alpha_1}s_{\alpha_2}c_{\alpha_3} - s_{\beta-\alpha_1}s_{\alpha_3} \\ -s_{\beta-\alpha_1}c_{\alpha_2} & c_{\beta-\alpha_1}c_{\alpha_3} + s_{\beta-\alpha_1}s_{\alpha_2}s_{\alpha_3} & -c_{\beta-\alpha_1}s_{\alpha_3} + s_{\beta-\alpha_1}s_{\alpha_2}c_{\alpha_3} \\ s_{\alpha_2} & c_{\alpha_2}s_{\alpha_3} & c_{\alpha_2}c_{\alpha_3} \end{pmatrix} \quad (17)$$

Finally, we identify *CP*-conserving limits corresponding to scenarios in which one state becomes pure *CP*-odd. Given the identification $M_{H_1} = 125$ GeV, this decoupling occurs in two distinct cases [58]

- (i) H_2 is *CP*-odd: This is realized by setting $\alpha_2 = 0$ and $\alpha_3 = \pm\pi/2$, yielding

$$\mathbf{T}^T|_{H_2: CP\text{-odd}} = \begin{pmatrix} c_{\beta-\alpha_1} & -s_{\beta-\alpha_1} & 0 \\ 0 & 0 & \pm 1 \\ \mp s_{\beta-\alpha_1} & \mp c_{\beta-\alpha_1} & 0 \end{pmatrix}. \quad (18)$$

- (ii) H_3 is *CP*-odd: This is realized by setting $\alpha_2 = \alpha_3 = 0$, yielding

$$\mathbf{T}^T|_{H_3: CP\text{-odd}} = \begin{pmatrix} c_{\beta-\alpha_1} & -s_{\beta-\alpha_1} & 0 \\ s_{\beta-\alpha_1} & c_{\beta-\alpha_1} & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (19)$$

$$\mathcal{M}_{13}^2 = t_\beta \mathcal{M}_{23}^2. \quad (14)$$

Using the relation between $\mathcal{M}^2$ and the mixing matrix $\mathbf{R}$ in Eq. (11), the mass of the third neutral scalar, m_{H_c} , can be expressed as [48]

$$m_{H_c}^2 = \frac{m_{H_a}^2 R_{a3}(R_{a2}t_\beta - R_{a1}) + m_{H_b}^2 R_{b3}(R_{b2}t_\beta - R_{b1})}{R_{c3}(R_{c1} - R_{c2}t_\beta)}. \quad (15)$$

As a derived quantity, m_{H_c} may emerge as the lightest, intermediate, or heaviest state, depending on the choice of input parameters. The states $\{H_a, H_b, H_c\}$ are then reordered accordingly to define the mass eigenstates $\{H_1, H_2, H_3\}$.

Furthermore, the physical mass basis is related to the Higgs basis via the mixing matrix $\mathbf{T}$ [12]

$$\begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix} \equiv \mathbf{T}^T \begin{pmatrix} H^0 \\ R_2 \\ I_2 \end{pmatrix}. \quad (16)$$

Combining Eqs. (7) and (13), the explicit form of $\mathbf{T}$ is

B. Neutral Higgs couplings and global measures of *CP* violation

The interactions of the neutral Higgs bosons with the gauge and fermionic sectors are central to the phenomenology of the C2HDM.¹ The interaction vertex $H_i V_\mu V_\nu$ (where $V = W^\pm, Z$) is parametrized as

$$ic(H_i VV)g_{\mu\nu}g_V m_V, \quad (20)$$

where $g_W = g$ and $g_Z = g/c_W$, with $c_W = \cos\theta_W$ denoting the cosine of the Weinberg angle. The gauge coupling modifiers are given in terms of the mixing matrix elements by

$$c(H_i VV) = T_{1i}. \quad (21)$$

Since current Higgs precision data favor a positive sign for the gauge coupling modifier κ_V [59–61], we restrict

¹For the complete set of C2HDM couplings, we refer the reader to Appendix A of Ref. [48].

our analysis to the case where

$$c(H_1 VV) > 0. \quad (22)$$

The Yukawa Lagrangian for a fermion field ψ_f with mass m_f is parametrized by

$$\mathcal{L}_Y = - \sum_{i=1}^3 \frac{m_f}{v} \bar{\psi}_f [c^e(H_i f f) + i c^o(H_i f f) \gamma_5] \psi_f H_i, \quad (23)$$

where $c^e(H_i f f)$ and $c^o(H_i f f)$ denote the CP -even and CP -odd Yukawa coupling modifiers, respectively. For type-I and type-II models, the modifiers for the top quark, bottom quark, and τ lepton take the following forms:

$$\begin{aligned} \text{Type-I: } c^e(H_i t \bar{t}) &= c^e(H_i b \bar{b}) = c^e(H_i \tau \bar{\tau}) = \frac{R_{i2}}{s_\beta}, \\ c^o(H_i t \bar{t}) &= -c^o(H_i b \bar{b}) = -c^o(H_i \tau \bar{\tau}) = -\frac{R_{i3}}{t_\beta}, \end{aligned} \quad (24)$$

$$\begin{aligned} \text{Type-II: } c^e(H_i t \bar{t}) &= \frac{R_{i2}}{s_\beta}, \\ c^e(H_i b \bar{b}) &= c^e(H_i \tau \bar{\tau}) = \frac{R_{i1}}{c_\beta}, \\ c^o(H_i t \bar{t}) &= -\frac{R_{i3}}{t_\beta}, \\ c^o(H_i b \bar{b}) &= c^o(H_i \tau \bar{\tau}) = -t_\beta R_{i3}. \end{aligned} \quad (25)$$

To quantify the extent of CP violation in the scalar sector, various measures have been proposed and studied extensively [46,51,62,63]. In this work, we adopt two normalized global CPV measures: ξ_V for the gauge sector and $\zeta_{t,b}$ for the Yukawa sector. By definition, both measures satisfy $\xi_V \in [0, 1]$ and $\zeta_{t,b} \in [0, 1]$.

The gauge-sector measure ξ_V is defined as [51]

$$\xi_V = 27 [c(H_1 VV)c(H_2 VV)c(H_3 VV)]^2 = 27 \prod_{i=1}^3 T_{1i}^2. \quad (26)$$

For the Yukawa sector, we introduce two normalized sum variables [46]

$$\zeta_{23} = 2 \sum_{i=1}^3 (R_{i2} R_{i3})^2, \quad \zeta_{13} = 2 \sum_{i=1}^3 (R_{i1} R_{i3})^2. \quad (27)$$

Given the distinct structures of the CP -even and CP -odd Yukawa couplings in type-I and type-II models [see Eqs. (24) and (25)], the normalized CPV measures for the top quark, b quark, and τ lepton are identified as follows:

$$\begin{cases} \text{Type-I:} & \zeta_t = \zeta_b = \zeta_\tau = \zeta_{23} \\ \text{Type-II:} & \zeta_t = \zeta_{23}, \quad \zeta_b = \zeta_\tau = \zeta_{13} \end{cases} \quad (28)$$

C. Higgs alignment limit

Before presenting our numerical results, we investigate the Higgs alignment limit, a critical regime where the lightest neutral scalar, H_1 , exhibits gauge and Yukawa couplings identical to those of the SM Higgs boson. In the C2HDM, the exact Higgs alignment limit is realized when $\alpha_1 = \beta$ and $\alpha_2 = 0$. Under these conditions, the mixing matrix $\mathbf{T}$ simplifies to

$$\mathbf{T}|_{\alpha_1=\beta, \alpha_2=0} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{\alpha_3} & -s_{\alpha_3} \\ 0 & s_{\alpha_3} & c_{\alpha_3} \end{pmatrix}. \quad (29)$$

A primary implication of the Higgs alignment limit is the recovery of CP conservation in the scalar potential. This is manifest in the behavior of the imaginary part of the quartic coupling λ_5 [58]

$$\begin{aligned} \text{Im } \lambda_5 &= -\frac{1}{c_\beta s_\beta v^2} \{ c_\beta [c_{\alpha_1} c_{\alpha_2} s_{\alpha_2} M_{H_1}^2 - c_{\alpha_2} s_{\alpha_3} (c_{\alpha_1} s_{\alpha_2} s_{\alpha_3} + s_{\alpha_1} c_{\alpha_3}) M_{H_2}^2 \\ &\quad + c_{\alpha_2} c_{\alpha_3} (s_{\alpha_1} s_{\alpha_3} - c_{\alpha_1} s_{\alpha_2} c_{\alpha_3}) M_{H_3}^2] + s_\beta [s_{\alpha_1} c_{\alpha_2} s_{\alpha_2} M_{H_1}^2 \\ &\quad + c_{\alpha_2} s_{\alpha_3} (c_{\alpha_1} c_{\alpha_3} - s_{\alpha_1} s_{\alpha_2} s_{\alpha_3}) M_{H_2}^2 - c_{\alpha_2} c_{\alpha_3} (c_{\alpha_1} s_{\alpha_3} + s_{\alpha_1} s_{\alpha_2} c_{\alpha_3}) M_{H_3}^2] \}, \end{aligned} \quad (30)$$

which vanishes identically upon substituting the alignment conditions $\alpha_1 = \beta$ and $\alpha_2 = 0$. This symmetry requirement forces the CP -violating off-diagonal elements of the mass matrix to vanish—specifically $\mathcal{M}_{13}^2 = \mathcal{M}_{23}^2 = 0$ in Eq. (10)—thereby decoupling the CP -odd state from the CP -even sector. Furthermore, the condition $\text{Im } \lambda_5 = 0$ leads to the indeterminacy of M_{H_c} in Eq. (15) as the exact alignment limit is approached.

In the exact Higgs alignment limit, the mixing angle α_3 is further constrained by the requirement of CP conservation. Applying the alignment conditions $\alpha_1 = \beta$ and $\alpha_2 = 0$ to the mixing matrix in Eq. (13), the heavy neutral states can be expressed as

$$H_2 = c_{\alpha_3} (-s_\beta \rho_1 + c_\beta \rho_2) + s_{\alpha_3} I_2, \quad (31)$$

$$H_3 = -s_{\alpha_3}(-s_\beta \rho_1 + c_\beta \rho_2) + c_{\alpha_3} I_2, \quad (32)$$

where $\rho_{1,2}$ are CP -even components and I_2 is CP -odd. Since CP conservation requires H_2 and H_3 to be pure CP eigenstates, α_3 must take the values 0 or $\pm\pi/2$.

In summary, the exact Higgs alignment limit within the C2HDM with softly broken Z_2 symmetry is characterized by the following simultaneous conditions:

$$\begin{aligned} \text{Higgs alignment: } \alpha_1 &= \beta, & \alpha_2 &= 0, \\ \text{Im } \lambda_5 &= 0, & \alpha_3 &\in \{0, \pm\pi/2\}. \end{aligned} \quad (33)$$

To quantify small but finite deviations from this limit, we introduce the normalized alignment measure

$$\delta_{H\text{-align}} = \sqrt{\frac{s_{\beta-\alpha_1}^2 + s_{\alpha_2}^2}{2}}, \quad (34)$$

where $\delta_{H\text{-align}} = 0$ corresponds to the case of exact alignment.

In the exact Higgs alignment limit, the coupling modifiers for H_1 recover their SM values: $c(H_1 VV) = 1$, $c^e(H_1 ff) = 1$, and $c^o(H_1 ff) = 0$. For the heavy scalars $H_{2,3}$, the gauge coupling modifiers vanish, $c(H_{2,3} VV) = 0$, while the Yukawa coupling modifiers are given by

Higgs alignment limit ($\alpha_1 = \beta, \alpha_2 = 0$):

$$\text{Type-I: } \begin{cases} c^e(H_2 ff) = \frac{c_{\alpha_3}}{t_\beta}, & c^o(H_2 tt) = -c^o(H_2 bb) = -c^o(H_2 \tau\tau) = -\frac{s_{\alpha_3}}{t_\beta}, \\ c^e(H_3 ff) = -\frac{s_{\alpha_3}}{t_\beta}, & c^o(H_3 tt) = -c^o(H_3 bb) = -c^o(H_3 \tau\tau) = -\frac{c_{\alpha_3}}{t_\beta}, \end{cases} \quad (35)$$

$$\text{Type-II: } \begin{cases} c^e(H_2 tt) = \frac{c_{\alpha_3}}{t_\beta}, & c^e(H_2 bb) = c^e(H_2 \tau\tau) = -t_\beta c_{\alpha_3}, \\ c^o(H_2 tt) = -\frac{s_{\alpha_3}}{t_\beta}, & c^o(H_2 bb) = c^o(H_2 \tau\tau) = -t_\beta s_{\alpha_3}, \\ c^e(H_3 tt) = -\frac{s_{\alpha_3}}{t_\beta}, & c^e(H_3 bb) = c^e(H_3 \tau\tau) = t_\beta s_{\alpha_3}, \\ c^o(H_3 tt) = -\frac{c_{\alpha_3}}{t_\beta}, & c^o(H_3 bb) = c^o(H_3 \tau\tau) = -t_\beta c_{\alpha_3}. \end{cases} \quad (36)$$

We explicitly retain the α_3 dependence in these expressions, as it plays a central role in the emergence of hidden CP violation discussed in Sec. IV D.

III. COMPREHENSIVE SCAN

In this section, we describe the comprehensive global scan performed to identify C2HDM parameter regions capable of supporting large CP violation. We begin by detailing our scanning methodology, including the choice of input parameters and the rigorous set of theoretical and experimental constraints—most notably the eEDM—applied to ensure phenomenological viability. We then characterize the intrinsic properties of the surviving parameter space, revealing that type-I and type-II models realize large CP violation through fundamentally different phenomenological patterns. For type-I, we identify a unique near-degenerate mass hierarchy that is essential for generating large CP violation in the gauge sector. For type-II, we show that large CP violation is restricted to the Yukawa sector of heavy scalars.

A. Parameter setup and scanning methodology

We investigate the parameter space of the C2HDM with a softly broken Z_2 symmetry. The scalar potential in Eq. (2) depends on the parameters $\{m_{11}^2, m_{22}^2, m_{12}^2, \lambda_{1,\dots,5}\}$, where m_{12}^2 and λ_5 can, in general, be complex. A global $U(1)$ rephasing of the Higgs doublets allows us to remove one of these two phases, leaving a single physical CP -violating phase $\arg(\lambda_5(m_{12}^2)^{-2})$. Consequently, the model is characterized by nine independent real parameters.

A physical basis for these nine input parameters consists of

$$\{m_{H_a}, m_{H_b}, m_{H^\pm}, \alpha_1, \alpha_2, \alpha_3, t_\beta, \text{Re}(m_{12}^2), v\}, \quad (37)$$

where m_{H_a} and m_{H_b} denote the masses of two neutral scalars.

To assess the theoretical consistency and phenomenological viability of the parameter points, we employ the public code SCANNERS v2.0.0 [64,65], which fully implements the C2HDM. For efficient sampling, particularly to ensure consistency with Higgs precision data, SCANNERS utilizes an alternative input parametrization focused on the properties of one neutral Higgs boson, identified here as the observed 125 GeV state ($H_a \equiv H_1$) [65]

$$\{m_{H_a} = 125 \text{ GeV}, m_{H_b}, m_{H^\pm}, t_\beta, \text{Re}(m_{12}^2), c^2(H_a VV), |c(H_a t\bar{t})|^2, \text{sign}(R_{a3}), R_{b3}\}. \quad (38)$$

Here, $c(H_a VV)$ is the effective $H_a VV$ coupling normalized to the SM value [see Eq. (21)], and $|c(H_a t\bar{t})|^2 = (c^e(H_a t\bar{t}))^2 + (c^o(H_a t\bar{t}))^2$ parametrizes the coupling strength to top quarks.

We perform a random scan over the following parameter ranges

$$\begin{aligned} m_{H_b} &\in [130, 2000] \text{ GeV}, & M_{H^\pm} &\in [90, 2000] \text{ GeV}, \\ t_\beta &\in [0.5, 50], & \text{Re}(m_{12}^2) &\in [0, 2 \times 10^6] \text{ GeV}^2, \\ \text{sign}(R_{a3}) &\in \{\pm 1\}, & R_{b3} &\in [-1, 1], \\ c^2(H_a VV) &\in [0.7, 1.0], & |c(H_a t\bar{t})|^2 &\in [0.6, 1.6]. \end{aligned} \quad (39)$$

We restrict our scan to the mass range below 2 TeV to focus on the parameter space most relevant for current and near-future collider searches, where the additional Higgs bosons remain within the direct reach of the High-Luminosity LHC and potential future Higgs factories.

Each generated point is subjected to a comprehensive set of constraints implemented within SCANNERS, as outlined below

- (i) *Theoretical requirements*: We require the scalar potential to be bounded from below [66], ensure perturbative unitarity by constraining the eigenvalues of the $2 \rightarrow 2$ scalar scattering matrix to be less than 8π , and enforce the stability of the electroweak vacuum against decay to deeper minima [67–69].
- (ii) *Electroweak and flavor constraints*: Parameter points must be consistent with electroweak precision data, specifically the oblique parameters (S , T , U) [70], compatible with the global fit values from the Particle Data Group at 95% CL [71]. Constraints from flavor physics are applied at 95% CL, including bounds from $B \rightarrow X_s \gamma$, neutral meson mixing, and various rare B -meson decays ($B_{d,s} \rightarrow \mu^+ \mu^-$, $B \rightarrow \tau \nu$, $B_s \rightarrow \phi \gamma$) [72–77].
- (iii) *Higgs collider data*: Agreement with Higgs data is evaluated using the HIGGSTOOLS framework (version 1.1.3) [78], which comprises three sublibraries: HIGGSPREDICTIONS, HIGGSSIGNALS [79], and HIGGSBOUNDS [80]. HIGGSSIGNALS tests the compatibility of the 125 GeV Higgs boson (H_1) with LHC signal strength measurements, requiring the deviation $\Delta\chi^2 = \chi_{\text{C2HDM}}^2 - \chi_{\text{SM}}^2$ to lie within the 2σ region of the SM best-fit value. HIGGSBOUNDS checks that the predicted signal rates for additional Higgs bosons (H_2 , H_3 , $H^\pm$) are not excluded by direct searches at the LEP, Tevatron, or LHC. Since the constraint on the CP -mixing angle $\alpha_\tau^{H_1}$, defined as $\alpha_\tau^{H_1} = \arctan(c^o(H_1 \tau \tau) / c^e(H_1 \tau \tau))$, is not natively implemented in SCANNERS, we explicitly impose the

ATLAS bound derived from $h_{125} \rightarrow \tau^+ \tau^-$ decays. We require consistency with the ATLAS measurement $\alpha_\tau^{H_1} = 9^\circ \pm 34^\circ$ (95% CL) [81].

- (iv) *eEDM*: The predicted eEDM, which primarily arises from two-loop Barr–Zee diagrams [82] and has been extensively studied in two-Higgs-doublet models [83,84], must satisfy the stringent 90% CL upper limit $|d_e| < 4.1 \times 10^{-30} e \text{ cm}$ set by the JILA collaboration [41].

In evaluating consistency with Higgs precision data, we also include near-degenerate scenarios where the H_2 mass is close to 125 GeV. To avoid numerical singularities associated with exact mass degeneracy, we impose the following mass conditions:

$$\begin{aligned} M_{H_1} &= 125 \text{ GeV}, & M_{H_2} &\geq 126 \text{ GeV}, \\ M_{H_3} &\geq 130 \text{ GeV}. \end{aligned} \quad (40)$$

Through extensive random scans, we collected 1.1×10^5 viable points for type-I and an equal number for type-II, all of which satisfy the full set of constraints including the mass conditions in Eq. (40). The eEDM constraint proves to be the most restrictive factor; only about 1% of the points satisfying all other requirements survive this bound in both type-I and type-II, thereby significantly shaping the allowed parameter space.

We now clarify the treatment of near-degenerate scenarios involving the observed state and the second neutral scalar ($M_{H_1} \approx M_{H_2}$) in HIGGSSIGNALS [85], which requires careful consideration in the signal-strength evaluation. HIGGSSIGNALS compares theoretical predictions to experimental measurements by summing the signal strengths (μ_i) of all neutral scalars whose masses fall within the effective mass window of a given channel, $m_h \pm \Delta m$. However, the widths of these mass windows differ substantially between two categories of analysis channels.

The first category comprises the *mass-sensitive* channels, namely the high-resolution $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ$ analyses based on Refs. [86,87]. These channels employ exceptionally narrow mass windows, approximately $[124.74, 125.76] \text{ GeV}$ and $[124.37, 125.81] \text{ GeV}$, respectively. HIGGSSIGNALS includes only the contributions from neutral Higgs bosons whose masses fall strictly within these windows.² Our mass condition of $M_{H_2} \geq 126 \text{ GeV}$

²This restriction is necessary because the experimental inputs for these channels [86,87] are provided as signal-strength fits fixed at the 125 GeV peak, rather than as model-independent invariant-mass spectra. Consequently, a simultaneous two-state fit to overlapping H_1 and H_2 contributions is not feasible, making the use of predefined narrow windows the only statistically consistent approach.

ensures that H_2 lies outside these high-resolution windows. Consequently, H_2 does not contribute to the predicted signal strengths of these dominant precision channels, which drive the global χ^2 fit.

In contrast, the remaining *mass-insensitive* channels—covering almost all other Higgs decay modes such as $H \rightarrow WW$, $H \rightarrow \tau\tau$, $H \rightarrow b\bar{b}$, and $H \rightarrow \mu\mu$ (and lower-resolution measurements of $\gamma\gamma$ and ZZ)—employ substantially wider mass windows. In these analyses, the second Higgs state H_2 often falls within the corresponding mass window, and HIGGSIGNALS appropriately accounts for its presence by summing its contribution to the total predicted signal strength

$$\mu_i^{\text{pred}} = \mu_i(H_1) + \mu_i(H_2). \quad (41)$$

This distinction is central to the viability of the near-degenerate scenario. The overall χ^2 is dominated by the high-resolution channels, which possess much smaller experimental uncertainties $\delta\mu_i^{\text{obs}}$ and thus carry the greatest statistical weight in the global fit [79].³ By construction, H_2 with $M_{H_2} > 126$ GeV is excluded from the dominant, mass-sensitive channels [86,87]. Consequently, although it contributes to the less precise, mass-insensitive channels, its impact on the global χ^2 remains marginal. As shown in the following analysis, this treatment allows near-degenerate scenarios to remain phenomenologically viable, particularly within the type-I C2HDM.

B. Characteristics of the large CP violation parameter space in type-I

Having detailed our random scanning methodology in the previous subsection, we now analyze the intrinsic characteristics of the viable parameter points for type-I. Our analysis places particular focus on regions exhibiting sizable CP violation, as quantified by the gauge and Yukawa CPV measures, ξ_V and ζ_t .

We begin by mapping the viable type-I parameter space in the (M_{H_2}, M_{H_3}) plane, as illustrated in Fig. 1. The color coding reflects the hierarchy of constraints: beige points satisfy standard theoretical and experimental bounds [constraints (i)–(iii)], while gray points represent the subset that additionally satisfies the eEDM bound [constraints (i)–(iv)]. The profound impact of the eEDM constraint is immediately apparent; the EDM bound excludes approximately 99% of the parameter points that were otherwise allowed by collider and flavor physics. Interestingly, the surviving (gray) points exhibit a preference for near-degenerate H_2 and H_3 masses when $M_{H_2} \gtrsim 250$ GeV, and permit the

decoupling limit where the BSM scalar masses can reach up to the 2 TeV scan limit.

To systematically explore the phenomenology of large CP violation, we introduce the benchmark condition $\xi_V > 0.1$, which is highlighted by the colored points in Fig. 1. While numerically arbitrary, this threshold serves as a practical tool to isolate regions with non-negligible CP -violating effects. Furthermore, this criterion underscores a fundamental difference between type-I and type-II: ξ_V can reach values up to ~ 0.4 in type-I but remains suppressed below $\mathcal{O}(10^{-3})$ in type-II. Quantitatively, $\xi_V = 0.1$ corresponds to a geometric mean of the gauge couplings $[\prod_i c(H_i VV)^2]^{1/3} \approx 0.155$, representing a physically meaningful level of CP mixing in the gauge sector.

Remarkably, requiring large gauge-sector CP violation ($\xi_V > 0.1$) imposes severe restrictions on the masses of the heavy neutral scalars. Imposing the large CPV condition $\xi_V > 0.1$ selects approximately 8.8% of the viable (gray) parameter points. However, these surviving points are not distributed uniformly; rather, they are strictly confined to the low-mass region, as evident from the left panel of Fig. 1. The magnified view in the right panel reveals that $\xi_V > 0.1$ requires $M_{H_2} \lesssim 159$ GeV and $M_{H_3} \lesssim 197$ GeV.

Turning to the charged Higgs sector, Fig. 2 maps the viable parameter points in the $(M_{H^\pm}, t_\beta)$ plane. Unlike the neutral scalar masses, $M_{H^\pm}$ and t_β are not significantly restricted by EDM measurements alone, as evidenced by the considerable overlap between the beige and gray points. However, imposing the large CPV condition $\xi_V > 0.1$ significantly narrows the allowed space to $t_\beta > 5.1$ and $M_{H^\pm} < 229.4$ GeV. It is encouraging that large CPV measures systematically favor low $M_{H^\pm}$, that is readily accessible to direct searches at high-energy colliders.

Having identified the specific mass regions favored by large CP violation, we now examine the underlying correlations among the individual Higgs-vector-boson couplings, $c(H_i VV)$, which directly dictate the collider phenomenology. In Fig. 3, we present the correlations of $c(H_3 VV)$ with $c(H_1 VV)$ (left panel) and with $c(H_2 VV)$ (right panel).

In the left panel, we observe that the viable (gray) points for $c(H_3 VV)$ are distributed broadly and symmetrically around zero, while $c(H_1 VV)$ ranges between approximately 0.84 and 1, consistent with an SM-like 125 GeV Higgs. The Higgs alignment limit, where $c(H_1 VV) = 1$ and $c(H_3 VV) = 0$, remains viable, corresponding to the decoupling of H_3 from the gauge sector. However, the large ξ_V region (colored points) necessitates a sizable deviation from alignment: specifically, $|c(H_3 VV)| \gtrsim 0.14$ and $c(H_1 VV) \lesssim 0.935$. Furthermore, as ξ_V increases, both $|c(H_3 VV)|$ and $|1 - c(H_1 VV)|$ increase. Consequently, future precision measurements of the Higgs couplings will critically test this scenario; for example, if $c(H_1 VV)$ is constrained to be > 0.95 , the CPV measure ξ_V cannot exceed 0.1.

³Note that the full χ^2 evaluation in HIGGSIGNALS uses the covariance matrix of correlated experimental and theoretical uncertainties [79].

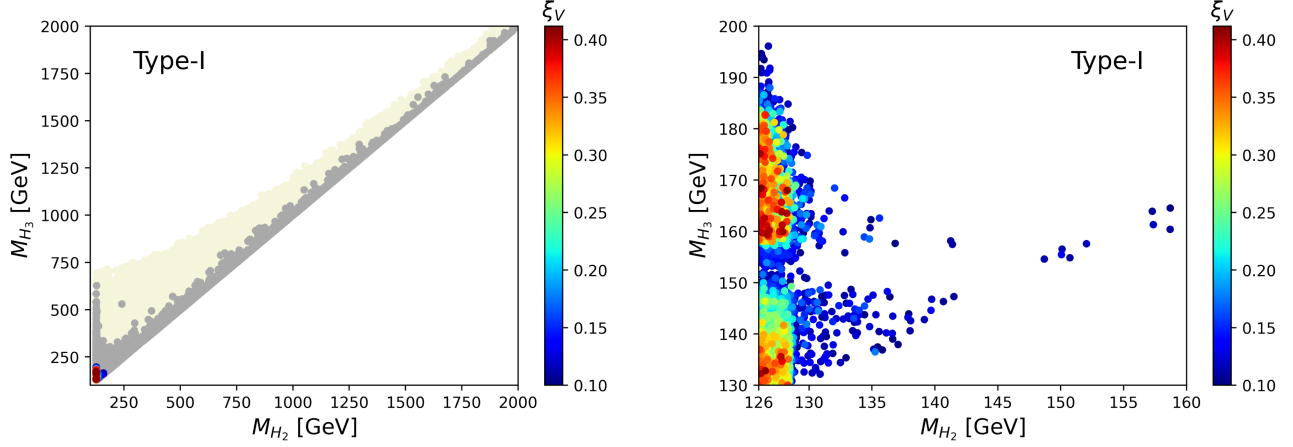


FIG. 1. Viable parameter points in the (M_{H_2}, M_{H_3}) plane for the type-I C2HDM, with the color code of ξ_V . Beige points satisfy theoretical requirements and constraints from electroweak precision data, flavor physics, and Higgs collider searches. Gray points represent the subset additionally satisfying the eEDM bound. Colored points highlight the region with large CP violation ($\xi_V > 0.1$), layered by the value of ξ_V , with higher values plotted on top. The right panel provides a magnified view focused on the large ξ_V region.

The right panel of Fig. 3 illustrates the correlation between $c(H_3VV)$ and $c(H_2VV)$. In contrast to the symmetric distribution of $c(H_3VV)$, the distribution of $c(H_2VV)$ is notably asymmetric, exhibiting a clear preference for positive values. Crucially, in the parameter space with large ξ_V , only $c(H_2VV) \gtrsim 0.2$ is permitted. We also observe that ξ_V increases concurrently with both $c(H_2VV)$ and $|c(H_3VV)|$, confirming that ξ_V serves as a robust representative measure for CP violation in the gauge sector.

We next shift our focus to CP violation in the Yukawa sector, quantified by ζ_t . Figure 4 displays the distribution of viable points in the (M_{H_2}, M_{H_3}) and $(M_{H^\pm}, t_\beta)$ planes, color-coded by ζ_t . Similar to the gauge CPV case, the subset with $\zeta_t > 0.1$ constitutes roughly 10% of the viable parameter space. However, a key distinction is that ζ_t can reach significantly higher values with $\max(\zeta_t) \approx 0.836$,

compared to ξ_V . This difference arises because ζ_t is a sum variable, whereas ξ_V is multiplicative; for instance, if H_3 is purely CP -odd, ξ_V vanishes, yet ζ_t remains nonzero provided $\beta \neq \alpha_1$.

Furthermore, $\zeta_t > 0.1$ allows for the decoupling limit where H_2 and H_3 are heavy and nearly degenerate (up to the 2 TeV scan limit)—a region strictly excluded by the $\xi_V > 0.1$ condition. Nevertheless, the highest ζ_t values still correlate with lighter BSM masses. The inset in the left panel of Fig. 4 clearly shows that $\zeta_t > 0.6$ requires $M_{H_2} < 129$ GeV, again confining the model to a scenario where H_2 is nearly degenerate with H_1 .

The right panel of Fig. 4 reveals two distinct clusters in the $(M_{H^\pm}, t_\beta)$ plane for large ζ_t . The first cluster features a lighter charged Higgs ($M_{H^\pm} \lesssim 250$ GeV) and spans a broad range of $t_\beta \approx 2$ –50, whereas the second cluster extends to the heavy $M_{H^\pm}$ limit (up to 2 TeV) but is realized only at low t_β ($t_\beta \lesssim 3$). This latter restriction arises because the type-I Yukawa couplings scale approximately as $1/t_\beta$; consequently, small t_β values are necessary to prevent strong suppression of the couplings in the heavy mass regime. Notably, the most extreme CPV values ($\zeta_t > 0.8$) are predominantly found within the lighter charged Higgs cluster.

Finally, we address the implications of large CP violation for the soft Z_2 breaking scale. Figure 5 illustrates the viable points in the $(t_\beta, |m_{12}^2|)$ plane. Even within the general viable space (gray points), $|m_{12}^2|$ deviates significantly from zero, confirming that soft Z_2 breaking is essential. The lower bound is approximately $|m_{12}^2|_{\min} \approx 53$ GeV² at $t_\beta \simeq 4$, rising as t_β deviates from this value. Conversely, the upper bound decreases with t_β , reaching $\mathcal{O}(10^6)$ GeV² for $t_\beta \lesssim 1$.

Remarkably, requiring large CP violation imposes much tighter and distinct restrictions on $|m_{12}^2|$. As shown in the left panel, the condition $\xi_V > 0.1$ favors a narrow band

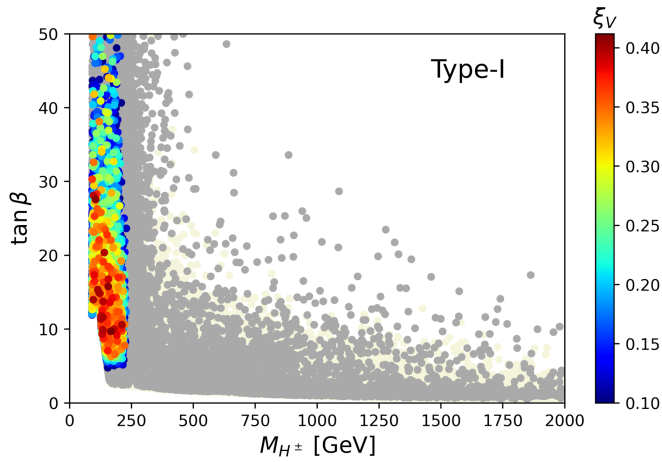


FIG. 2. Viable parameter points in the $(M_{H^\pm}, t_\beta)$ plane for the type-I C2HDM. The color bar indicates the CPV measure ξ_V . The beige, gray, and colored points follow the same convention as in Fig. 1.

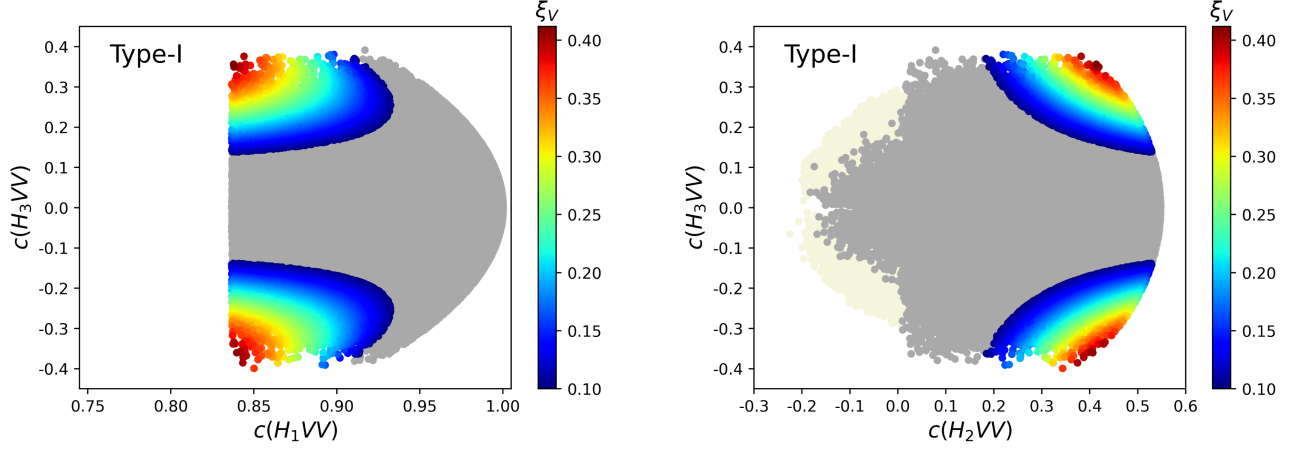


FIG. 3. Viable parameter points in the $(c(H_1VV), c(H_3VV))$ plane (left) and in the $(c(H_2VV), c(H_3VV))$ plane (right) for the type-I C2HDM. The beige, gray, and colored points follow the convention from Fig. 1, where the color bar indicates the value of the CP -violating measure ξ_V .

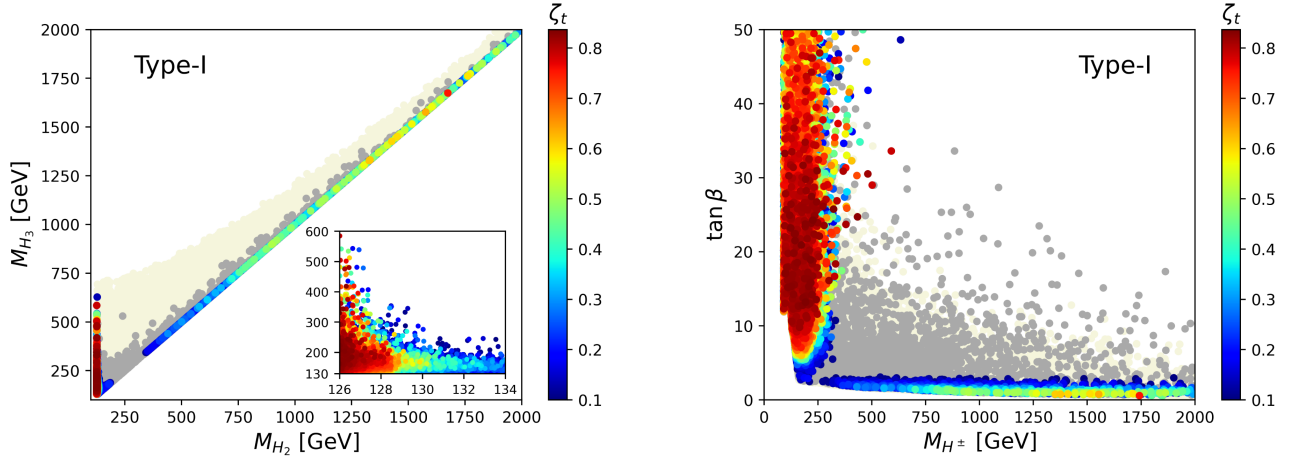


FIG. 4. Viable parameter points in the (M_{H_2}, M_{H_3}) plane (left) and the $(M_{H^\pm}, t_\beta)$ plane (right) for the type-I C2HDM. The color bar indicates the value of the CP -violating measure ζ_t . The beige and gray points follow the convention from Fig. 1, while the colored points represent the subset with $\zeta_t > 0.1$. The inset in the left panel provides a magnified view of the low-mass region.

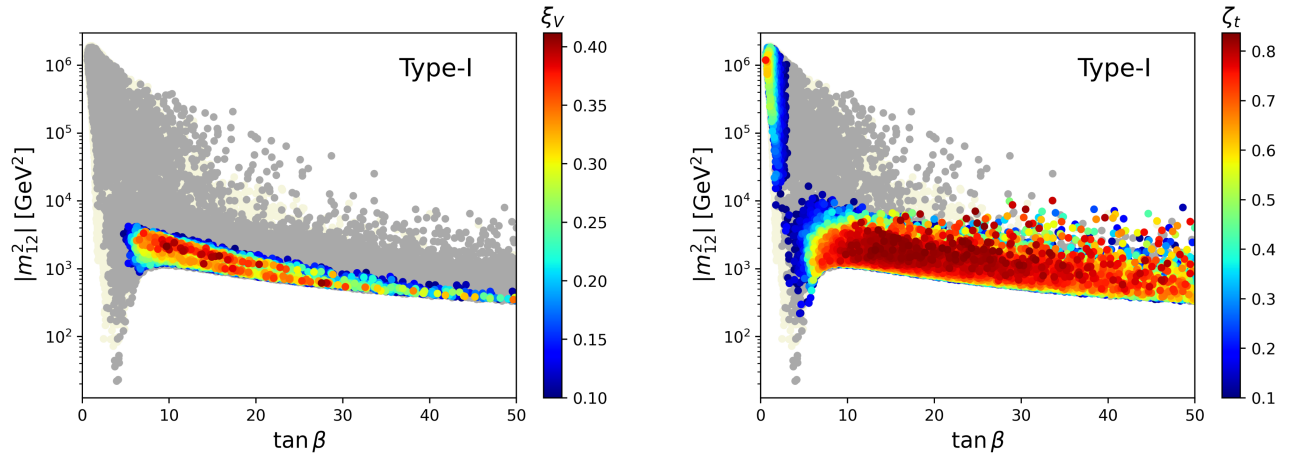


FIG. 5. Viable parameter points in the $(t_\beta, |m^2_{12}|)$ plane for the type-I C2HDM. The color bar indicates the value of the CPV measure ξ_V (left panel) and ζ_t (right panel). The beige and gray points follow the convention from Fig. 1, while the colored points represent the subset with $\xi_V > 0.1$ (left) and $\zeta_t > 0.1$ (right).

characterized by $|m_{12}^2| \sim \mathcal{O}(10^3)$ GeV². In contrast, the condition $\zeta_t > 0.1$ (right panel) permits two separate regions: a vertical band at low t_β with $|m_{12}^2| \sim \mathcal{O}(10^4\text{--}10^6)$ GeV² (corresponding to heavier H_2), and a horizontal band at larger t_β with $|m_{12}^2| \sim \mathcal{O}(10^3)$ GeV² (corresponding to lighter H_2). This distinction demonstrates that the preferred scale of $|m_{12}^2|$ depends critically on whether the CP violation originates in the gauge or Yukawa sector.

In conclusion, we summarize the parameter ranges identifying large CPV regions in type-I. For large gauge CPV ($\xi_V > 0.1$), the parameters are confined to

$$\begin{aligned} \text{type-I with } \xi_V > 0.1: & M_{H_2} \in [126.0, 148.7] \text{ GeV}, \\ & M_{H_3} \in [130.0, 194.7] \text{ GeV}, \\ & M_{H^\pm} \in [90.9, 229.4] \text{ GeV}, \\ & |m_{12}^2| \in [363.2, 4202.4] \text{ GeV}^2, \\ & t_\beta \in [5.1, 49.7]. \end{aligned} \quad (42)$$

For large Yukawa CPV ($\zeta_t > 0.1$), the allowed ranges bifurcate depending on the mass of H_2

$$\text{Type-I with } \zeta_t > 0.1 \quad (43)$$

Case 1 (Light H_2)

$$\begin{aligned} M_{H_2} &\in [126.0, 182.9] \text{ GeV}, \\ M_{H_3} &\in [130.0, 543.1] \text{ GeV}, \\ M_{H^\pm} &\in [90.9, 482.4] \text{ GeV}, \\ |m_{12}^2| &\in [219.0, 10991.1] \text{ GeV}^2, \\ t_\beta &\in [2.89, 50.0]. \end{aligned}$$

Case 2 (Heavy H_2)

$$\begin{aligned} M_{H_2}, M_{H_3} &\in [346.0, 1992.9] \text{ GeV}, \\ M_{H^\pm} &\in [285.4, 1991.9] \text{ GeV}, \\ |m_{12}^2| &\in [1.3 \times 10^4, 1.8 \times 10^6] \text{ GeV}^2, \\ t_\beta &\in [0.58, 3.07]. \end{aligned} \quad (44)$$

We emphasize that parameter points are not uniformly distributed across these ranges, as evident from the scatter plots presented in this subsection.

C. Characteristics of the large CP violation parameter space in type-II

We now turn our attention to the characteristics of the large CPV parameter space in type-II. In sharp contrast to type-I, the gauge-sector CPV measure ξ_V in type-II is extremely suppressed

$$\text{Type-II: } \xi_V < 3.28 \times 10^{-3}. \quad (45)$$

However, the Yukawa sector in type-II is richer, featuring two distinct CPV measures: ζ_t for the top quark and ζ_b for the bottom quark and τ lepton [see Eq. (28)]. Consequently, we focus our analysis on the viable parameter space satisfying $\zeta_t > 0.1$ and $\zeta_b > 0.1$.

Figure 6 displays the viable parameter points in the (M_{H_2}, M_{H_3}) plane, color-coded by ζ_t (left) and ζ_b (right). The beige and gray points adhere to the convention established in Fig. 1. Notably, unlike the suppressed gauge measure, the Yukawa CPV measures in type-II can be quite large, reaching maximum values of:

$$\text{Type-II: } \max(\zeta_t) \approx 0.747, \quad \max(\zeta_b) \approx 0.842. \quad (46)$$

A striking difference from type-I is immediately visible in the distribution of the beige points: Type-II requires both H_2 and H_3 to be significantly heavier than 125 GeV. This restriction stems primarily from the charged Higgs sector. Strong constraints from flavor physics, particularly the $B \rightarrow X_s \gamma$ measurement, impose a lower bound on the charged Higgs mass of $M_{H^\pm} \gtrsim 580$ GeV [28,29,72,88]. To satisfy electroweak precision constraints (specifically the oblique parameters S , T , U [70]), the neutral BSM scalars must effectively track the mass of this heavy charged Higgs, forcing them to be heavy as well. While the eEDM constraint (gray points) excludes a significant portion of the beige points, the excluded region is noticeably smaller than in type-I (compare with Fig. 1).

The requirement of large CP violation ($\zeta_t > 0.1$ or $\zeta_b > 0.1$) further constrains the mass spectrum. As shown by the colored points in Fig. 6, large Yukawa CP violation enforces a strong mass degeneracy between H_2 and H_3 , in contrast to the broader distribution of the gray points. However, the absolute mass bounds remain largely unchanged compared to the general viable parameter space: both M_{H_2} and M_{H_3} extend from approximately 580 GeV up to our 2 TeV scanning limit.

We next examine the parameter space in the $(M_{H^\pm}, t_\beta)$ plane, presented in Fig. 7. A comparison between the beige and gray points reveals a crucial feature: the eEDM measurements impose a stringent upper limit of $t_\beta \lesssim 4$. This constitutes a dramatic tightening of the parameter space relative to earlier constraints. For example, under the previous ACME I limit of $|d_e| < 9.3 \times 10^{-29} e \text{ cm}$ (95% CL) [89], t_β values as high as ~ 30 were permitted [48].

The impact of large CP violation on t_β depends on the specific flavor sector involved. The condition $\zeta_t > 0.1$ (left panel) moderately restricts the allowed range to $t_\beta \in [0.52, 2.18]$ and exhibits a clear correlation where larger ζ_t values favor smaller t_β . Crucially, this range ensures that the Yukawa couplings of the BSM Higgs bosons to top quarks are either enhanced or, at minimum, remain unsuppressed. This has significant implications for processes such as four-top-quark production [90–94].

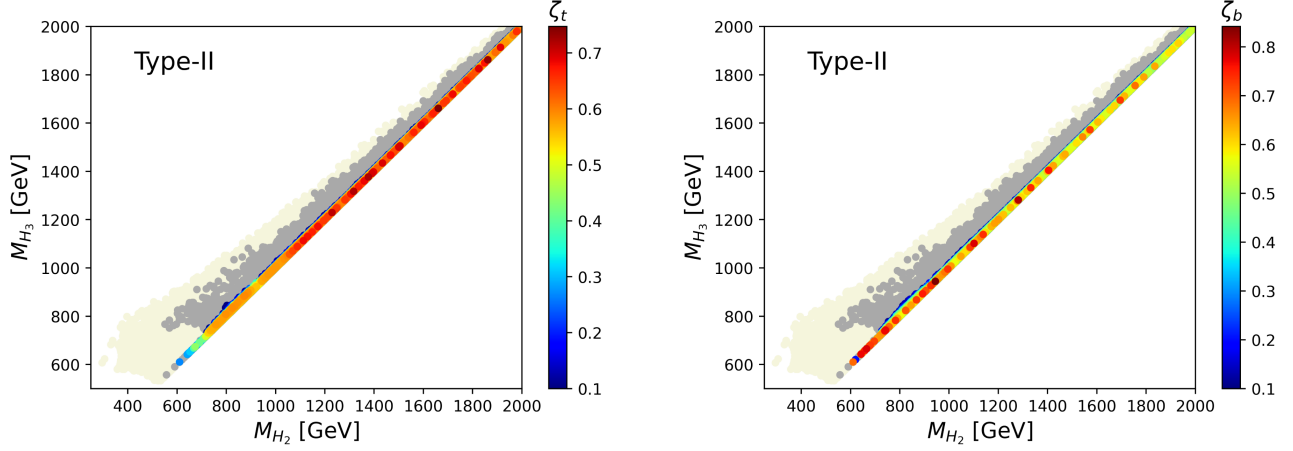


FIG. 6. Viable parameter points in the (M_{H_2}, M_{H_3}) plane for the type-II C2HDM. The color bar indicates the value of the CP -violating measure ζ_t (left panel) and ζ_b (right panel). The beige and gray points follow the convention from Fig. 1, while the colored points represent the subset with $\zeta_t > 0.1$ (left) and $\zeta_b > 0.1$ (right).

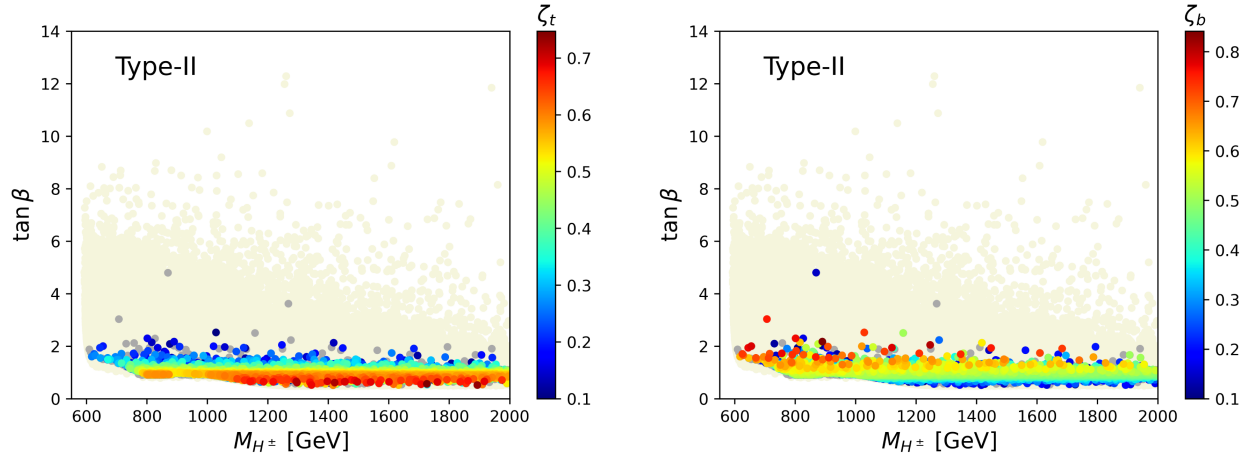


FIG. 7. Viable parameter points in the $(M_{H^\pm}, \tan \beta)$ plane for the type-II C2HDM. The color bar indicates the value of the CP -violating measure ζ_t (left panel) and ζ_b (right panel). The beige and gray points follow the convention from Fig. 1, while the colored points represent the subset with $\zeta_t > 0.1$ (left) and $\zeta_b > 0.1$ (right).

Conversely, the condition $\zeta_b > 0.1$ (right panel) allows for a t_β range similar to the general allowed (gray) points but exhibits an opposing correlation, where larger ζ_b values prefer larger t_β . Notably, since the parameter space satisfying either $\zeta_t > 0.1$ or $\zeta_b > 0.1$ excludes the high t_β regime, the heavy BSM Higgs bosons in type-II do not benefit from the t_β -enhancement of their Yukawa couplings to bottom quarks and τ leptons.

Regarding the charged Higgs mass, Fig. 7 confirms that demanding large CP violation does not significantly alter the lower bound, which increases only mildly to $M_{H^\pm} \approx 610$ GeV compared to the flavor-physics floor of ~ 580 GeV. Both the general allowed region and the large CPV subsets extend to the 2 TeV limit, confirming that the decoupling limit remains fully compatible with all constraints in type-II.

Finally, we investigate the soft Z_2 breaking scale in the type-II C2HDM by presenting the viable parameter points

in the $(t_\beta, |m_{12}^2|)$ plane, shown in Fig. 8. The color bars indicate the values of the CPV measures ζ_t (left panel) and ζ_b (right panel). The beige and gray points follow the convention established in Fig. 1, while the colored points represent the subsets satisfying $\zeta_t > 0.1$ (left) and $\zeta_b > 0.1$ (right).

Similar to the type-I case (see Fig. 5), the viable parameter points in type-II necessitate a nonzero soft breaking term m_{12}^2 , thereby prohibiting an exact Z_2 symmetry. However, we observe two distinct differences compared to type-I. First, the allowed values of $|m_{12}^2|$ are considerably larger in type-II. This enhancement is driven by the requirement for heavier BSM Higgs bosons in this scenario ($M_{H_2}, M_{H_3} \gtrsim 610$ GeV), which naturally scales up the soft breaking term. Second, the eEDM constraint (gray points) significantly raises the lower bound on $|m_{12}^2|$ compared to the theoretical and flavor constraints alone (beige points). This elevated lower bound remains robust

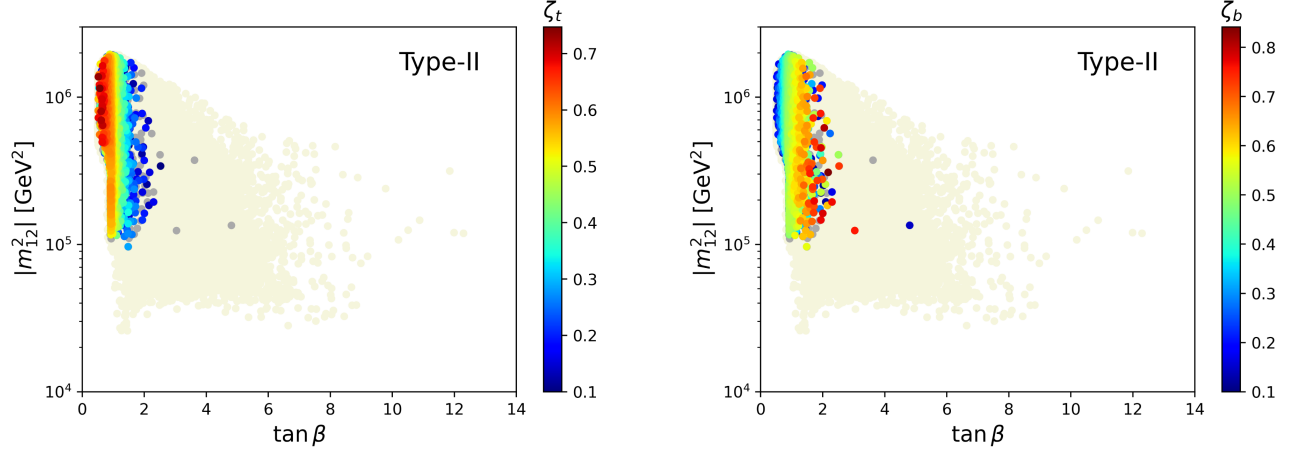


FIG. 8. Viable parameter points in the $(t_\beta, |m^2_{12}|)$ plane for the type-II C2HDM. The color bar indicates the value of the CP -violating measure ζ_t (left panel) and ζ_b (right panel). The beige and gray points follow the convention from Fig. 1, while the colored points represent the subset with $\zeta_t > 0.1$ (left) and $\zeta_b > 0.1$ (right).

even under the large CP violation conditions of $\zeta_t > 0.1$ or $\zeta_b > 0.1$. Furthermore, unlike in type-I, we observe no strong correlation between the magnitude of the CPV measures (ζ_t or ζ_b) and the scale of $|m^2_{12}|$.

In summary, the parameter regions supporting large Yukawa CP violation in type-II are defined as follows. For large top-quark Yukawa CPV ($\zeta_t > 0.1$)

$$\begin{aligned} \text{Type-II with } \zeta_t > 0.1: \quad & M_{H_2} \in [610.0, 2000] \text{ GeV}, \quad M_{H_3} \in [610.3, 2000] \text{ GeV}, \\ & M_{H^\pm} \in [617.3, 2000] \text{ GeV}, \\ & |m^2_{12}| \in [9.7 \times 10^4, 2.0 \times 10^6] \text{ GeV}^2, \quad t_\beta \in [0.52, 2.53]. \end{aligned} \quad (47)$$

For large bottom/tau Yukawa CPV ($\zeta_b > 0.1$)

$$\begin{aligned} \text{Type-II with } \zeta_b > 0.1: \quad & M_{H_2} \in [610.3, 2000] \text{ GeV}, \quad M_{H_3} \in [610.3, 2000] \text{ GeV}, \\ & M_{H^\pm} \in [617.3, 2000] \text{ GeV}, \\ & |m^2_{12}| \in [9.7 \times 10^4, 2.0 \times 10^6] \text{ GeV}^2, \quad t_\beta \in [0.52, 4.8]. \end{aligned} \quad (48)$$

As with type-I, the parameter distribution across these ranges is highly nonuniform, as evident from the corresponding figures.

IV. FUTURE PROSPECTS AND OBSERVABILITY OF LARGE CP

In the previous section, we delineated the intrinsic characteristics of the C2HDM parameter space capable of supporting large CP violation. We now turn to the phenomenological consequences and future prospects of these regions. We begin by evaluating the predicted eEDM values for both type-I and type-II, comparing them against the projected sensitivities of next-generation experiments. Having established these indirect constraints, we then examine the distinct collider signatures associated with the large CP -violation regimes in each model. For type-I, our analysis centers on the gauge sector, where large ξ_V necessitates a near-degenerate scenario between H_2 and the

125 GeV H_1 ; here, we demonstrate how the simultaneous contributions of these two states reproduce the observed Higgs signal strengths. Conversely, for type-II, we focus on the Yukawa sector, investigating the magnitude and observability of the CP -violating phases in the top-quark and τ -lepton couplings. Finally, we investigate the phenomenon of *hidden* CP violation near the Higgs alignment limit, where CP mixing persists in the heavy sector despite vanishing signatures in standard gauge interactions.

A. Prospect of the eEDM in the large CP violation parameter space

The search for the eEDM has reached an unprecedented level of sensitivity, providing a powerful indirect probe of new physics. Currently, the most stringent upper limit on the eEDM is set by the JILA experiment using trapped HfF^+ ions, $|d_e| < 4.1 \times 10^{-30} e \text{ cm}$ at 90% confidence level [41]. This bound is notably more restrictive than

the previous benchmark set by the ACME II collaboration ($|d_e| < 1.1 \times 10^{-29} e\text{ cm}$ [42]), placing severe pressure on the parameter space of CP -violating extensions of the Standard Model like the C2HDM.

Looking ahead, the experimental frontier is poised to improve sensitivity by another one to two orders of magnitude. The ACME III collaboration targets the $10^{-31} e\text{ cm}$ regime using a cryogenic beam of thorium monoxide (ThO) [42], while the next-generation JILA experiment aims for a similar precision of a few times $10^{-31} e\text{ cm}$ [41]. Complementing these efforts, the NL-eEDM experiment utilizes barium fluoride (BaF) to provide a systematic cross-check with a projected sensitivity between 10^{-31} and $10^{-32} e\text{ cm}$ [95]. Achieving these goals will extend the reach of eEDM measurements far beyond the direct search capabilities of the LHC [96–99].

Given the current experimental constraints and future prospects, the following questions arise: (i) what underlying mechanisms allow a sizable parameter space to persist despite the stringent JILA eEDM constraint; (ii) what values of $|d_e|$ are predicted across the viable C2HDM parameter space, particularly in regions of maximal CP violation; and (iii) how does $|d_e|$ correlate with the sensitivity of future Higgs precision measurements?

The resilience of the C2HDM parameter space against eEDM bounds is rooted in the internal cancellation structure of the two-loop Barr–Zee contributions [82]. In Fig. 9, we decompose these contributions for type-I and type-II scenarios. The cancellation mechanism differs markedly between the two Yukawa types.

In type-I, individual contributions to $|d_e|$ for viable parameter points (gray) are typically of $\mathcal{O}(10^{-30}) e\text{ cm}$, comparable to the current JILA bound. Nevertheless, the opposite sign of the H^+H^- loop contribution relative to the other contributions leads to a partial cancellation in the total d_e .

In type-II, individual contributions can exceed the JILA bound by up to three orders of magnitude. Satisfying the experimental constraints therefore requires significant mutual cancellations, as reflected in the strong anticorrelation between the W^+W^- contribution and the remaining terms. Intriguingly, in regions with maximally allowed Yukawa CP violation ($\zeta_{t,b} \gtrsim 0.7$), the individual contributions themselves become suppressed, with each remaining near or below the JILA bound. This behavior is consistent with the findings of Ref. [44], where large CP violation was shown to be compatible with experimental limits without requiring excessive fine-tuning.

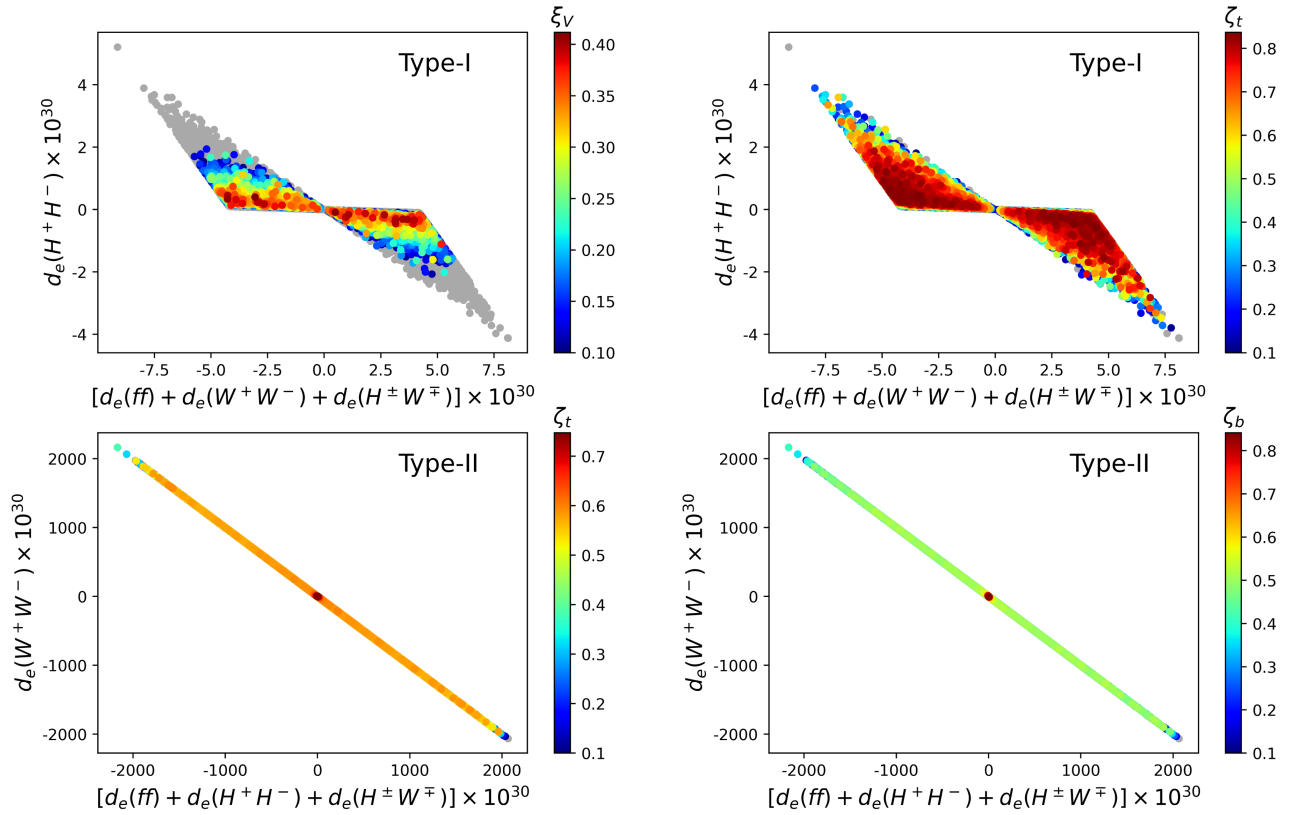


FIG. 9. Anatomy of individual two-loop Barr–Zee contributions to the electron EDM for type-I (upper) and type-II (lower) parameter space under the JILA bound. Here, $d_e(X_i X_j)$ denotes the contribution from the $X_i X_j$ loop class to the total eEDM. Gray and colored points satisfy the same theoretical and experimental constraints as in Fig. 1. Color bars indicate the CP -violating measures ξ_V , ξ_t , and ξ_b .

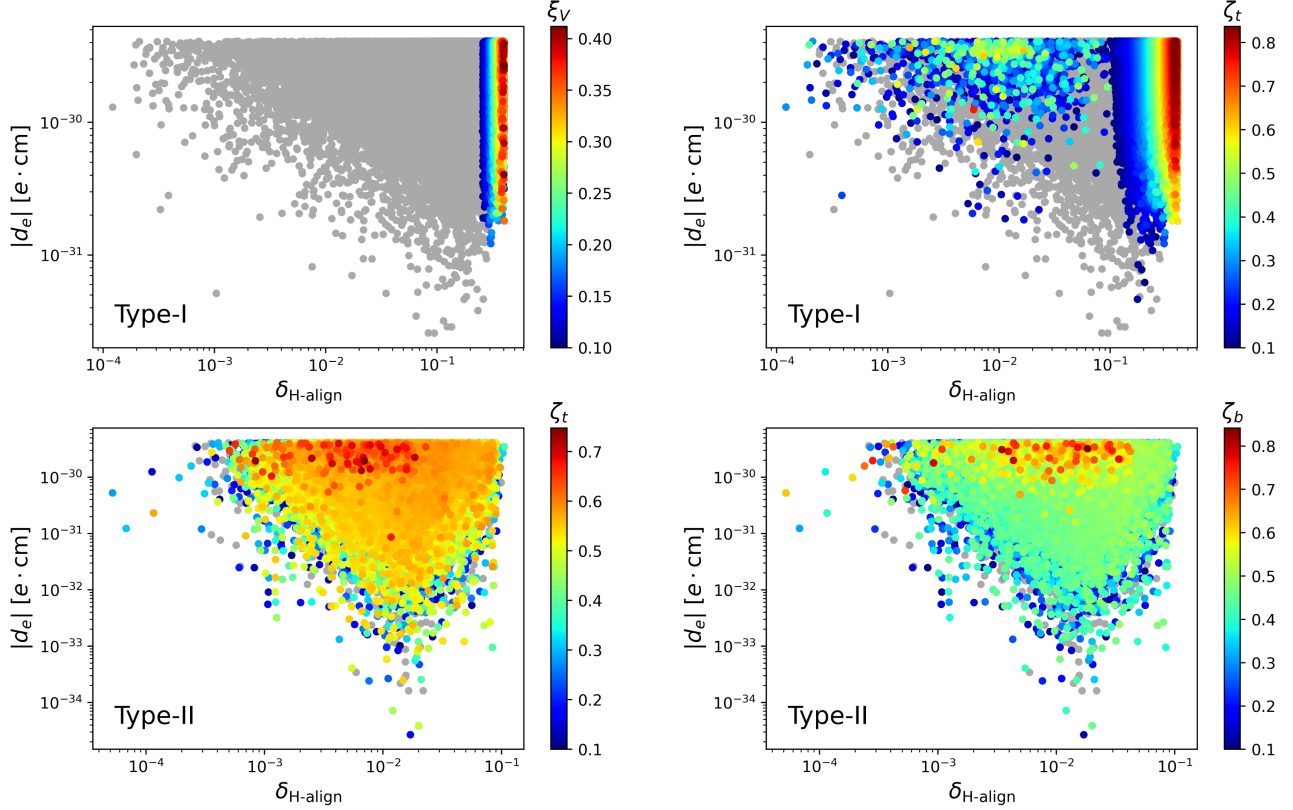


FIG. 10. Predicted $|d_e|$ as a function of the alignment measure $\delta_{H\text{-align}}$ for type-I (upper) and type-II (lower). Beige, gray, and colored points satisfy the same conditions as in Fig. 1. The color scales represent the CP -violating measures ξ_V (upper left), ζ_t (upper right), ζ_t (lower left), and ζ_b (lower right).

Building on this understanding of the eEDM constraints, Fig. 10 illustrates the resulting distribution of $|d_e|$ as a function of the Higgs alignment measure $\delta_{H\text{-align}}$ for type-I (upper) and type-II (lower). We first note that the apparent absence of both CP -conserving configurations and exact Higgs alignment is a natural consequence of the numerical random scan over the multidimensional parameter space. While neither the CP -conserving limit ($\alpha_2 = 0$, $\alpha_3 \in \{0, \pm\pi/2\}$) nor the exact Higgs alignment limit ($\beta = \alpha_1$, $\alpha_2 = 0$, $\alpha_3 \in \{0, \pm\pi/2\}$) is theoretically excluded, these limits correspond to measure-zero regions that are statistically inaccessible in a stochastic sampling.

Since our primary objective is to investigate the phenomenology of sizable CP violation—and noting that CP is conserved in the exact alignment limit—we do not fine-tune our scan to populate these regions. As a result, the ensemble of viable points naturally clusters away from these limits, leading to an apparent lower threshold for $|d_e|$, which we interpret as the characteristic scale of CP -violating signals within the considered parameter space.

It is also important to comment on the role of the chosen mass range ($M_{H_{2,3}}, M_{H^\pm} < 2$ TeV). In principle, pushing the model toward the decoupling limit with much heavier Higgs masses would suppress $|d_e|$. However, this limit is not phenomenologically viable in the present framework,

as the quartic couplings $|\lambda_{1,\dots,5}|$ scale with the square of the Higgs masses [58], leading to violations of perturbative unitarity at large mass scales.

To demonstrate this, we performed additional scans with fixed heavy scalar masses $M = 3, 4, 5$ TeV, with m_{H_b} and $M_{H^\pm}$ fixed to M . We find that no parameter points survive theoretical constraints for $M = 4$ and 5 TeV, while only a negligible number of points remain viable at $M = 3$ TeV, and these are incompatible with current eEDM bounds. Therefore, the distributions shown in Fig. 10 represent the physically consistent range of $|d_e|$ predictions within the model at realistic mass scales.

Beyond these constraints, Fig. 10 reveals distinct phenomenological features between the two Yukawa types. In type-I, we observe a pronounced correlation between the eEDM magnitude and the alignment measure $\delta_{H\text{-align}}$: as the model approaches the alignment limit (smaller $\delta_{H\text{-align}}$), the minimum values of $|d_e|$ realized in our scan—interpreted as the characteristic eEDM scale within the considered parameter space—shift to larger values. Notably, no comparable correlation is observed between $|d_e|$ and the gauge-sector CPV measure ξ_V .

In contrast, type-II predictions correlate strongly with the magnitudes of the Yukawa CPV measures ζ_t and ζ_b , where increased CP violation directly drives larger EDM values.

This carries a critical implication: should future eEDM experiments continue to yield null results, they will strictly constrain the possibility of large CP violation within the type-II Yukawa sector.

In summary, eEDM measurements offer a powerful, type-dependent test of the C2HDM. In type-I, the viable points are concentrated in a range of $|d_e|$ well-matched to the reach of next-generation experiments. While the type-II model can accommodate much smaller eEDM values, it exhibits strong correlations with both the CPV measures and the Higgs alignment parameter. Consequently, future EDM bounds will serve as a potent indirect probe of the type-II C2HDM, tightly constraining both the allowed CP violation and deviations from the Higgs alignment limit.

B. Higgs signal strengths of the near degenerate scenario in type-I

A key finding of our comprehensive scan in the type-I C2HDM is that large CP violation in the gauge sector strongly favors a configuration where H_2 is nearly degenerate with the 125 GeV state H_1 . As discussed in Sec. III A, our HIGSSIGNALS analysis confirms that such near-degeneracy is fully compatible with current Higgs precision data. To understand the underlying mechanism, we must examine how the individual contributions from H_1 and H_2 combine to reproduce the observed Higgs signals. For the purpose of this analysis, we designate the subset of parameter points satisfying $M_{H_2} \in [126, 128]$ GeV and $\xi_V > 0.1$ as the near-degenerate CPV scenario.

The first distinctive feature of this scenario is the suppression of the total width of the 125 GeV state H_1 . We find

$$\Gamma_{\text{tot}}(H_1) \in [2.65, 3.49] \text{ MeV}, \quad (49)$$

which corresponds to a reduction of approximately 15% – 35% relative to the SM prediction $\Gamma_{\text{tot}}^{\text{SM}}(h_{125}) \simeq 4.1$ MeV. The corresponding branching ratios are

$$\begin{aligned} \text{Br}(H_1 \rightarrow b\bar{b}) &\in [57.25, 59.05]\%, \\ \text{Br}(H_1 \rightarrow W^+W^-) &\in [20.83, 22.93]\%, \\ \text{Br}(H_1 \rightarrow gg) &\in [7.53, 7.77]\%, \\ \text{Br}(H_1 \rightarrow \tau^+\tau^-) &\in [6.16, 6.35]\%, \\ \text{Br}(H_1 \rightarrow c\bar{c}) &\in [2.80, 2.89]\%, \\ \text{Br}(H_1 \rightarrow ZZ) &\in [2.60, 2.86]\%, \\ \text{Br}(H_1 \rightarrow \gamma\gamma) &\in [0.215, 0.323]\%, \\ \text{Br}(H_1 \rightarrow Z\gamma) &\in [0.152, 0.181]\%. \end{aligned} \quad (50)$$

These branching ratios are obtained using C2HDM_HDECAY [48], which evaluates Higgs decay rates via the ANYHDECAY interface based on the HDECAY package [100,101]. This setup ensures state-of-the-art accuracy, including NLO/NNLO QCD and electroweak corrections where available, as well as off-shell effects.

A second notable feature concerns the heavier state H_2 . In this scenario, its total width is even more suppressed

$$\Gamma_{\text{tot}}(H_2) \in [0.163, 1.37] \text{ MeV}, \quad (51)$$

which is significantly smaller than the SM prediction for a 125 GeV Higgs boson. The corresponding branching ratios for $M_{H_2} \in [126, 128]$ GeV are shown in Fig. 11. They remain nearly constant across this narrow mass window and closely resemble those of the SM Higgs boson. The only qualitative difference is the opening of the decay mode $H_2 \rightarrow H^\pm W^\mp$ when it is kinematically accessible. As indicated by the color bar, a lighter $M_{H^\pm}$ enhances this

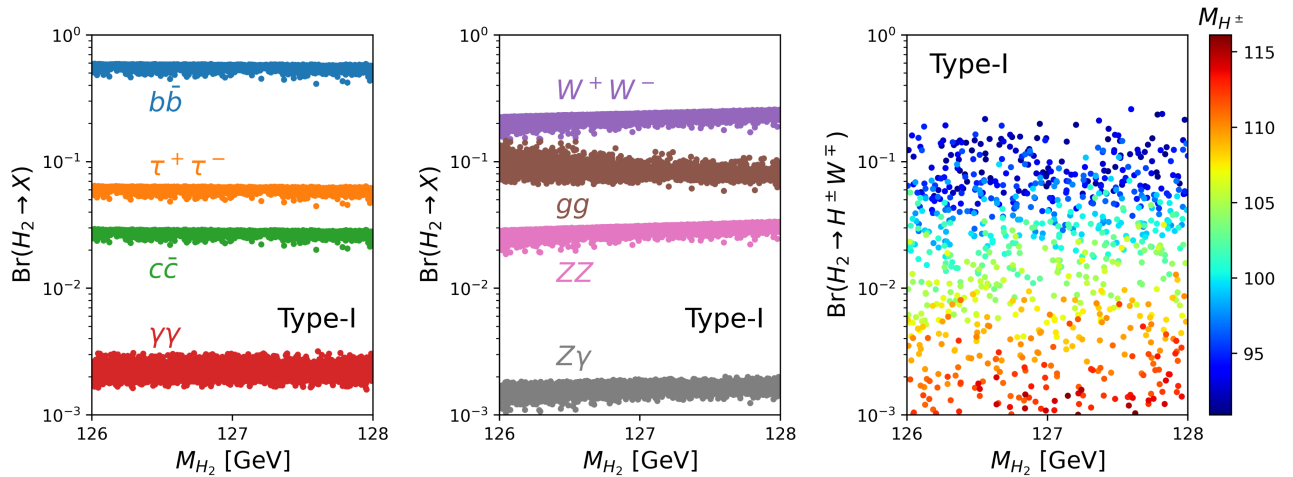


FIG. 11. Branching ratios of H_2 in the near-degenerate CPV scenario where $M_{H_2} \in [126, 128]$ GeV and $\xi_V > 0.1$ in type-I C2HDM. The decay modes into $b\bar{b}$, $\tau^+\tau^-$, $c\bar{c}$, $\gamma\gamma$ are shown in the left panel; W^+W^- , gg , ZZ , $Z\gamma$ in the middle; and $H^\pm W^\mp$ in the right panel. In the $H^\pm W^\mp$ panel, the color bar indicates $M_{H^\pm}$.

mode, although the branching fraction stays below the percent level for $M_{H^\pm} \gtrsim 100$ GeV.

These characteristics of H_2 —a suppressed total width combined with SM-like branching ratios—have two critical consequences for Higgs precision observables. First, they imply an approximately universal rescaling of the partial widths of H_2 into SM final states f_{SM}

$$\Gamma(H_2 \rightarrow f_{\text{SM}}) \simeq \mu(H_2) \Gamma^{\text{SM}}(h_{125} \rightarrow f_{\text{SM}}) \quad \text{for } \mu(H_2) < 1. \quad (52)$$

Because the non-SM decay mode $H_2 \rightarrow H^\pm W^\mp$ remains subdominant, the total width is well-approximated by the relation $\Gamma_{\text{tot}}(H_2) \simeq \mu(H_2) \Gamma_{\text{tot}}(h_{125})$, which explains the significant suppression of the total width observed in Eq. (51). Crucially, this factor naturally suppresses the gluon-fusion production rate of H_2 , thereby ensuring that this additional state remains compatible with current Higgs precision data.

Second, the SM-like nature of the observed H_1 (corresponding to a small deviation from the exact Higgs alignment limit, $\alpha_1 = \beta$ and $\alpha_2 = 0$) leads to an approximate sum rule: $c^2(H_1 VV) + c^2(H_2 VV) \simeq 1$. This relation is derived by introducing a small misalignment parameter $\epsilon = \beta - \alpha_1 = \alpha_2 \ll 1$, which yields $c^2(H_1 VV) + c^2(H_2 VV) \simeq 1 - \epsilon^2(1 + s_{2\alpha_3})$. This sum rule ensures that

the two neutral states, H_1 and H_2 , collectively saturate the SM-like coupling strength to vector bosons. Consequently, for any channel driven by the Higgs-gauge coupling—whether it is electroweak production via vector-boson fusion (VBF) and associated production (VH), or decay into weak bosons (WW, ZZ)—the combined signal strength reproduces the SM expectation

$$\mu_{\text{tot}}^{(V)} = \mu^{(V)}(H_1) + \mu^{(V)}(H_2) \simeq 1, \quad (53)$$

where the superscript (V) collectively denotes any process involving vector bosons in either the production or decay.

We now present the individual signal strengths $\mu_{\text{production}}^{\text{decay}}$ for H_1 and H_2 , defined as

$$\begin{aligned} \mu_{\text{production}}^{\text{decay}}(H_i) &\equiv \frac{[\sigma(pp \rightarrow H_i)_{\text{production}} \times \text{BR}(H_i \rightarrow \text{decay})]_{\text{C2HDM}}}{[\sigma(pp \rightarrow H)_{\text{production}} \times \text{BR}(H \rightarrow \text{decay})]_{\text{SM}}}. \end{aligned} \quad (54)$$

The branching ratios are computed as described above. The production cross sections are obtained using state-of-the-art tools: gluon fusion (ggF) at NNLO QCD via SUSHI [102,103], VH production at NNLO QCD via VH@NNLO [104,105], and VBF via HIGGSBOUNDS [80] (which rescales

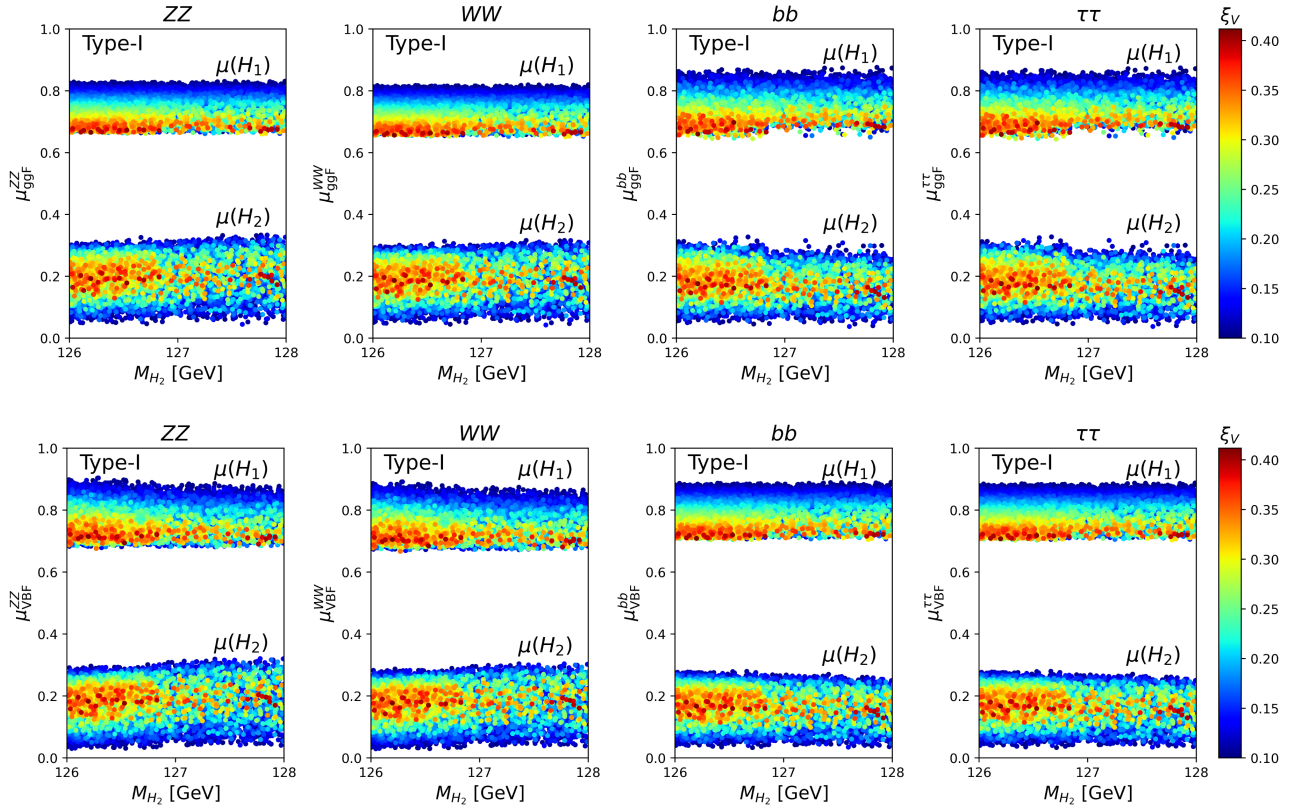


FIG. 12. Signal strengths $\mu_{\text{production}}^{\text{decay}}$ of H_1 and H_2 in the near-degenerate CPV scenario: $M_{H_2} \in [126, 128]$ GeV and $\xi_V > 0.1$ in type-I C2HDM. Shown are decay modes into ZZ , W^+W^- , $b\bar{b}$, and $\tau^+\tau^-$ for gluon fusion (top) and VBF (bottom) productions. The color code denotes ξ_V .

the SM cross sections from the LHCHSWG to NNLO QCD + NLO EW accuracy).

In Fig. 12, we display the resulting signal strengths of H_1 and H_2 for the main decay modes into vector bosons and fermions (ZZ , W^+W^- , $b\bar{b}$, and $\tau^+\tau^-$), comparing gluon-fusion (upper panels) and VBF (lower panels) production. It is evident that H_1 dominates the total signal, with $\mu(H_1)$ ranging roughly from 0.67 to 0.85, while H_2 contributes the remaining fraction, with $\mu(H_2)$ falling between 0.1 and 0.3. This distribution directly reflects the coupling sum rule $c^2(H_1VV) + c^2(H_2VV) \simeq 1$. Hence, although each state individually yields a sub-SM rate, their combined signal remains close to unity (within $\sim 10\%$), effectively satisfying the current Higgs precision data.

A remarkable feature of Fig. 12 is the strong correlation between the signal strengths and the CPV measure ξ_V . Larger values of ξ_V confine $\mu(H_1)$ and $\mu(H_2)$ to specific ranges; for instance, demanding $\xi_V \gtrsim 0.4$ implies $\mu(H_1) \simeq 0.67$ and $\mu(H_2) \sim 0.2$. In this high-CPV regime, the total signal strength—comprising the sum of the H_1 and H_2 contributions—falls distinctly below the SM expectation. Consequently, improved precision in future Higgs measurements will effectively probe the CPV measure ξ_V within this scenario.

We now turn to the loop-induced final states $H_i \rightarrow \gamma\gamma$ and $H_i \rightarrow Z\gamma$. As these decays are radiatively generated, they offer a unique window into BSM physics [106,107], being susceptible to both modifications of the SM-like couplings and new loop contributions from the charged Higgs boson. Given the light $M_{H^\pm}$ preferred in the large- ξ_V regime (see Fig. 2), these additional loops can induce observable deviations. Moreover, the interference between W and top loops differs between the $Z\gamma$ and $\gamma\gamma$ modes, making them complementary probes.

Figure 13 presents the predicted $\mu^{\gamma\gamma}$ and $\mu^{Z\gamma}$ for both ggF (upper panels) and VBF (lower panels) production. Since the individual values for $\mu^{\gamma\gamma}(H_1)$ and $\mu^{\gamma\gamma}(H_2)$ exhibit a broad spread, one might question whether the total rate is consistent with the current Higgs precision data. To address this, the second column of Fig. 13 explicitly displays the combined rate $\mu^{\gamma\gamma}(H_1) + \mu^{\gamma\gamma}(H_2)$. This demonstrates that the total diphoton signal remains well within the current ATLAS best-fit value of $\mu^{\gamma\gamma} = 1.05 \pm 0.10$ [108]. In contrast, for the $Z\gamma$ channel, both the individual and combined signal strengths show a much narrower spread, with the combined rate consistently suppressed below ~ 0.75 . While this suppression is compatible with the current experimental value of $\mu^{Z\gamma} = 1.3^{+0.6}_{-0.5}$ [107], such a reduced rate is a robust

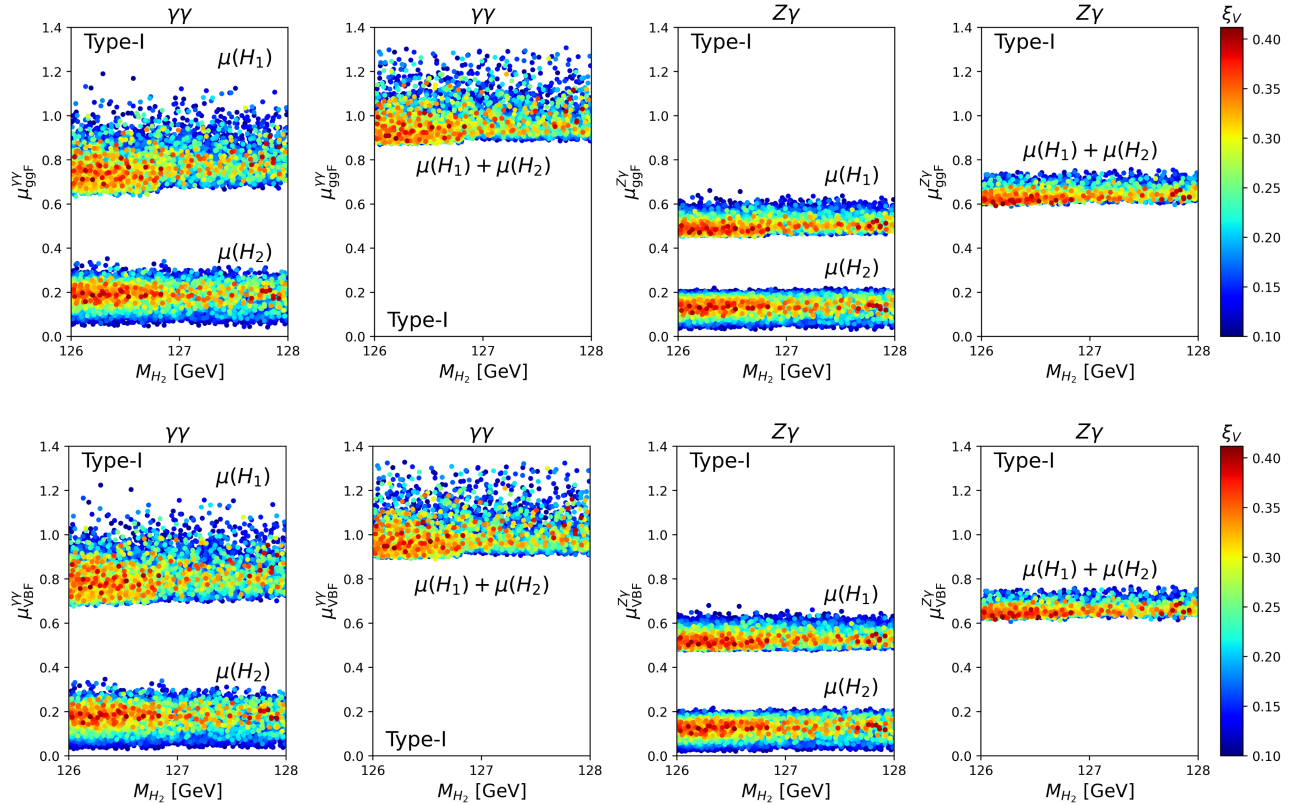


FIG. 13. Signal strengths of H_1 and H_2 for the decay modes into $\gamma\gamma$ and $Z\gamma$ in the near-degenerate CPV scenario ($M_{H_2} \in [126, 128]$ GeV, $\xi_V > 0.1$). Results for gluon fusion are shown in the top panels, and for VBF in the bottom. The color code denotes ξ_V .

prediction of the near-degenerate CPV scenario. Consequently, a future measurement of $\mu^{Z\gamma}$ significantly below unity would corroborate the robust suppression predicted by this near-degenerate CPV scenario.

Experimental outlook.—The near-degenerate CPV scenario predicts that the observed 125 GeV Higgs peak arises from a superposition of two states, H_1 and H_2 . While their combined signal strengths mimic SM-like behavior, the individual states possess reduced couplings and narrower widths. Given the current LHC invariant-mass resolution of $\Delta m \sim 1\text{--}2$ GeV in the $H \rightarrow \gamma\gamma$ and $H \rightarrow ZZ \rightarrow 4\ell$ channels [106,109], our predicted mass splitting of $\Delta M = M_{H_2} - M_{H_1} \lesssim 3$ GeV typically results in the two states appearing as a single, broadened resonance. Furthermore, mass-insensitive channels spanning the 122–130 GeV window naturally accommodate this scenario [110–112]. As demonstrated in our analysis, the H_2 contribution remains subdominant ($\sim 10\%$ of the H_1 signal), preserving consistency with the observed resonance structure.

The theoretical foundation for analyzing such overlapping resonances was established by Pilaftsis [113], who systematically studied the lineshape effects of nearly degenerate, CP -violating neutral Higgs bosons using a propagator-level resonant-mixing formalism. This approach was further developed for the LHC in the context of the CP -violating MSSM [114], employing a full propagator-matrix formalism to consistently treat the mixing, production, and decay of nearly degenerate states. Based on this framework, a high-statistics *lineshape analysis* at the HL-LHC is expected to serve as a particularly powerful probe. Specifically, a splitting of $\Delta M \sim 1$ GeV would produce a subtle non-Breit-Wigner distortion, while $\Delta M \gtrsim 2$ GeV could manifest as a visible shoulder or double-peak structure. Both ATLAS and CMS are pursuing such analyses, and the increased sensitivity of the HL-LHC will be decisive.

Ultimately, future e^+e^- colliders (ILC, FCC-ee, CEPC) offer the definitive resolution to this ambiguity. Through precise threshold scans, these machines could not only disentangle the split resonance structure but also provide a model-independent measurement of the total width Γ_{tot} [115–118]. This would allow for a critical test of the scenario’s primary prediction: a reduced H_1 width in the range $\Gamma_{\text{tot}}(H_1) \in [2.65, 3.49]$ MeV. Confirming such a deviation from the SM expectation of ~ 4.1 MeV would serve as a smoking-gun signature of the extended Higgs sector, even if the mass splitting remains near the resolution limit.

C. The CP structure of top and tau couplings in type-II

While the near-degenerate CPV scenario in type-I highlights the possibility of large CP violation in the gauge sector (ξ_V), the type-II model exhibits a fundamentally different behavior. As observed in our global scan, type-II is severely constrained to the regime of negligible gauge-sector CP violation, effectively forcing $\xi_V \lesssim 3 \times 10^{-3}$.

Consequently, the primary manifestation of CP violation in type-II shifts to the Yukawa sector. Here, the interactions of the top quark and the τ lepton serve as the definitive probes of the Higgs CP parity.

To quantify the CP structure in this sector, we parametrize the Yukawa couplings between a neutral Higgs boson H_i and fermions f in terms of a coupling strength modifier $\kappa_f^{H_i}$ and a CP -mixing angle $\alpha_f^{H_i}$

$$\mathcal{L}_{\text{Yukawa}} \supset - \sum_{i=1,2,3} \sum_{f=t,\tau} \frac{m_f}{v} \kappa_f^{H_i} [\cos \alpha_f^{H_i} + i\gamma_5 \sin \alpha_f^{H_i}] \bar{f} f H_i. \quad (55)$$

In this notation, $\alpha_f = 0$ corresponds to a purely CP -even interaction, $\alpha_f = 90^\circ$ represents a purely CP -odd state, while intermediate values indicate a CP -mixed Higgs boson.

We begin with the CP properties of the observed 125 GeV Higgs boson, H_1 . Current measurements indicate that H_1 is dominantly CP -even [81,119–123], though a small CP -odd admixture remains a viable possibility. A powerful avenue for probing this possibility involves its direct Yukawa interactions. At the LHC, the CP structure of the top–Higgs coupling is primarily accessed through the $t\bar{t}h_{125}$ production mode, which offers two key advantages: (i) strong suppression of QCD backgrounds, and (ii) sensitivity to CP -odd effects through spin-correlation and angular observables involving the top quarks and the Higgs boson [17,124–127]. Recent ATLAS and CMS analyses on $t\bar{t}h_{125}$ production and decay measurements find no significant deviation from the SM CP -even hypothesis, constraining the CP -mixing angle to $|\alpha_t^{H_1}| \lesssim 40^\circ$ (95% CL) [128,129].

Similarly, the CP -mixing angle $\alpha_\tau^{H_1}$ can be probed through the transverse spin correlations of τ -leptons in $H_1 \rightarrow \tau\tau$ decays. These correlations are imprinted on the directions of the decay products, specifically in the acoplanarity angle Φ^* between the two τ decay planes. The differential decay rate is modulated by $\cos(\Phi^* - 2\alpha_\tau^{H_1})$ [130–133]. Since reconstructing Φ^* is experimentally challenging due to the missing energy from neutrinos, the constraint on $\alpha_\tau^{H_1}$ is weaker than on $\alpha_t^{H_1}$. Standard analyses by ATLAS and CMS, which utilize template-based fits to τ -pair kinematics [81,123], currently yield loose constraints: $\alpha_\tau^{H_1} = 9^\circ \pm 34^\circ$ at 95% CL.

Given these experimental constraints, a key question is how large the CP violation in the top– H_1 and τ – H_1 couplings can be within the type-II C2HDM. In Fig. 14, we present the viable parameter space, showing the $(\alpha_t^{H_1}, \kappa_t^{H_1})$ plane in the left panel and the $(\alpha_\tau^{H_1}, \kappa_\tau^{H_1})$ plane in the right panel. The gray points denote the general viable parameter space, while the colored points represent the subset satisfying the large CP violation conditions ($\zeta_t > 0.1$ for the left panel and $\zeta_b > 0.1$ for the right). As evident from the overlap, the accessible

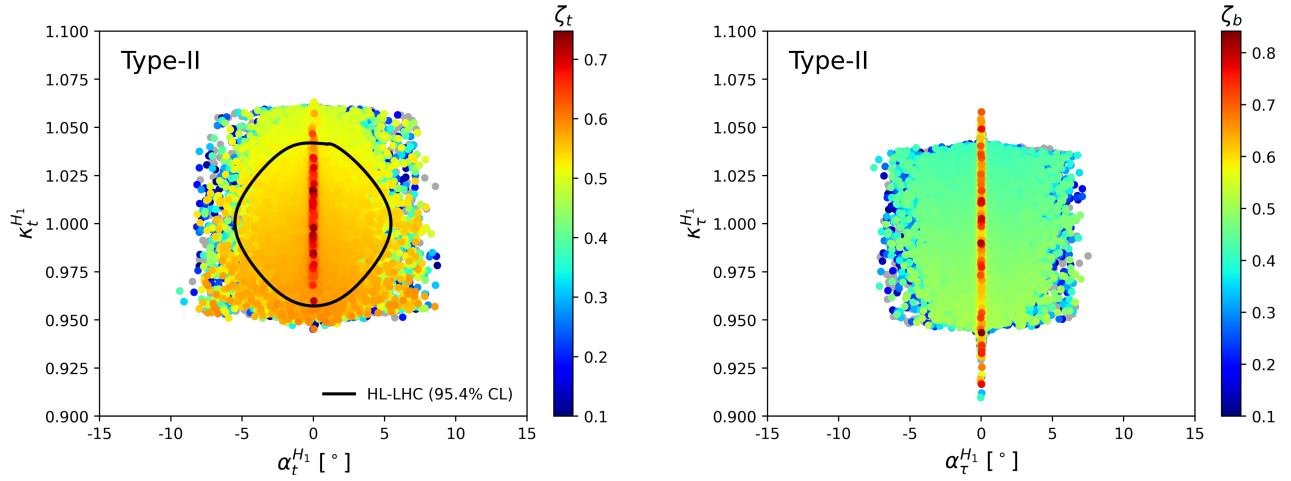


FIG. 14. Viable parameter points in the $(\alpha_t^{H_1}, \kappa_t^{H_1})$ plane (left) and $(\alpha_\tau^{H_1}, \kappa_\tau^{H_1})$ plane (right) for the type-II C2HDM. Gray points represent the full set of viable parameter points, while colored points indicate the subset satisfying the large CP violation condition: $\zeta_t > 0.1$ (left) and $\zeta_b > 0.1$ (right). The color code denotes the magnitude of ζ_t and ζ_b , respectively. The solid black curve in the left panel displays the projected 95% CL exclusion contour for the HL-LHC (14 TeV, 3 ab^{-1}), obtained using the GNN-based method.

ranges for the Yukawa coupling parameters are nearly identical for both the general and large CP violation parameter sets.

The left panel of Fig. 14 reveals that the CP violation in the top- H_1 coupling is quite limited. The magnitude $\kappa_t^{H_1}$ deviates from the SM prediction by at most $\sim 6\%$, and the mixing angle $\alpha_t^{H_1}$ remains within approximately $\pm 8^\circ$. Remarkably, we observe an inverse correlation: higher values of ζ_t tend to correspond to smaller values of $\alpha_t^{H_1}$. This indicates that the top- H_1 coupling is not the most efficient probe for the global CPV measure ζ_t .

To assess the HL-LHC's ability to probe this restricted parameter space, we show in the left panel of Fig. 14 the projected 95% CL exclusion contour (solid black line) for $\sqrt{s} = 14 \text{ TeV}$ and $\mathcal{L}_{\text{tot}} = 3 \text{ ab}^{-1}$. This projection is obtained by extending the machine learning (ML) framework of Ref. [127] to the general $\kappa_t^{H_1}$ case using the $t\bar{t}h_{125}$ production channel. The conditional graph neural network (GNN) approach of Ref. [127] significantly improves sensitivity by directly learning event topology and angular correlations from reconstructed objects. This offers a distinct advantage over traditional cut-and-count analyses, which suffer from suboptimal signal regions and limited sensitivity to CP -odd observables [121,122,134–137].

However, even with these advanced techniques, the HL-LHC will lack the statistical power to conclusively test this scenario. For instance, in the optimistic case where $\alpha_t^{H_1} \approx 8^\circ$, the HL-LHC would detect a nonzero CP phase with a significance of only $\sim 3.2\sigma$. Conversely, as indicated by the 95% CL contour, a null result would not exclude the high ζ_t regime. Crucially, because the parameter points exhibiting the largest global CP violation (e.g., $\zeta_t \gtrsim 0.5$) cluster near $\alpha_t^{H_1} \approx 0$, they mimic the purely CP -even H_1 state predicted by the SM. Thus, merely increasing the

statistical precision of the $t\bar{t}H_1$ measurement at future colliders will not suffice to uncover this Yukawa-sector CP violation.

The situation is equally challenging for the τ - H_1 coupling. The right panel of Fig. 14 shows the viable parameter points in the $(\alpha_\tau^{H_1}, \kappa_\tau^{H_1})$ plane. Deviations from the SM-like limit are again small: $\kappa_\tau^{H_1}$ is within $\sim 5\%$ of the SM value, and $\alpha_\tau^{H_1}$ is confined to $\pm 7^\circ$. These values are well within current experimental limits. While advanced deep learning techniques, such as the graph transformer network proposed in Ref. [138], can significantly improve sensitivity (reaching $|\alpha_\tau| \sim 20^\circ$ at 95% CL with 100 fb^{-1}), the predicted mixing angles in the type-II C2HDM are substantially smaller than even these optimized projections. Thus, probing $\alpha_\tau^{H_1}$ at the HL-LHC remains a formidable challenge.

Given that the CP -violating effects in the top and τ couplings of the 125 GeV Higgs boson are effectively beyond the realistic reach of the HL-LHC, we turn to the heavier neutral Higgs states as the complementary probes. Figure 15 displays the viable points in the (α_t, κ_t) and $(\alpha_\tau, \kappa_\tau)$ planes for H_2 (upper panels) and H_3 (lower panels). Gray points show the full viable parameter space, while the colored points highlight regions with large global CPV measures, $\zeta_t > 0.1$ (left panels) or $\zeta_b > 0.1$ (right panels).

The most striking feature to emerge from our analysis is that strong global CP violation forces the heavy-Higgs sector into a state of maximal mixing. As ζ_t and ζ_b increase, the mixing angles $\alpha_t^{H_{2,3}}$ and $\alpha_\tau^{H_{2,3}}$ simultaneously converge to $\pi/4$, signifying that the heavy Higgs bosons manifest as maximally mixed CP -even/ CP -odd states. This distinct phenomenological pattern identifies the heavy neutral sector as the primary reservoir of CP -violating effects, offering a far more robust signature than the SM-like state. We further observe that H_2 and H_3 exhibit very similar

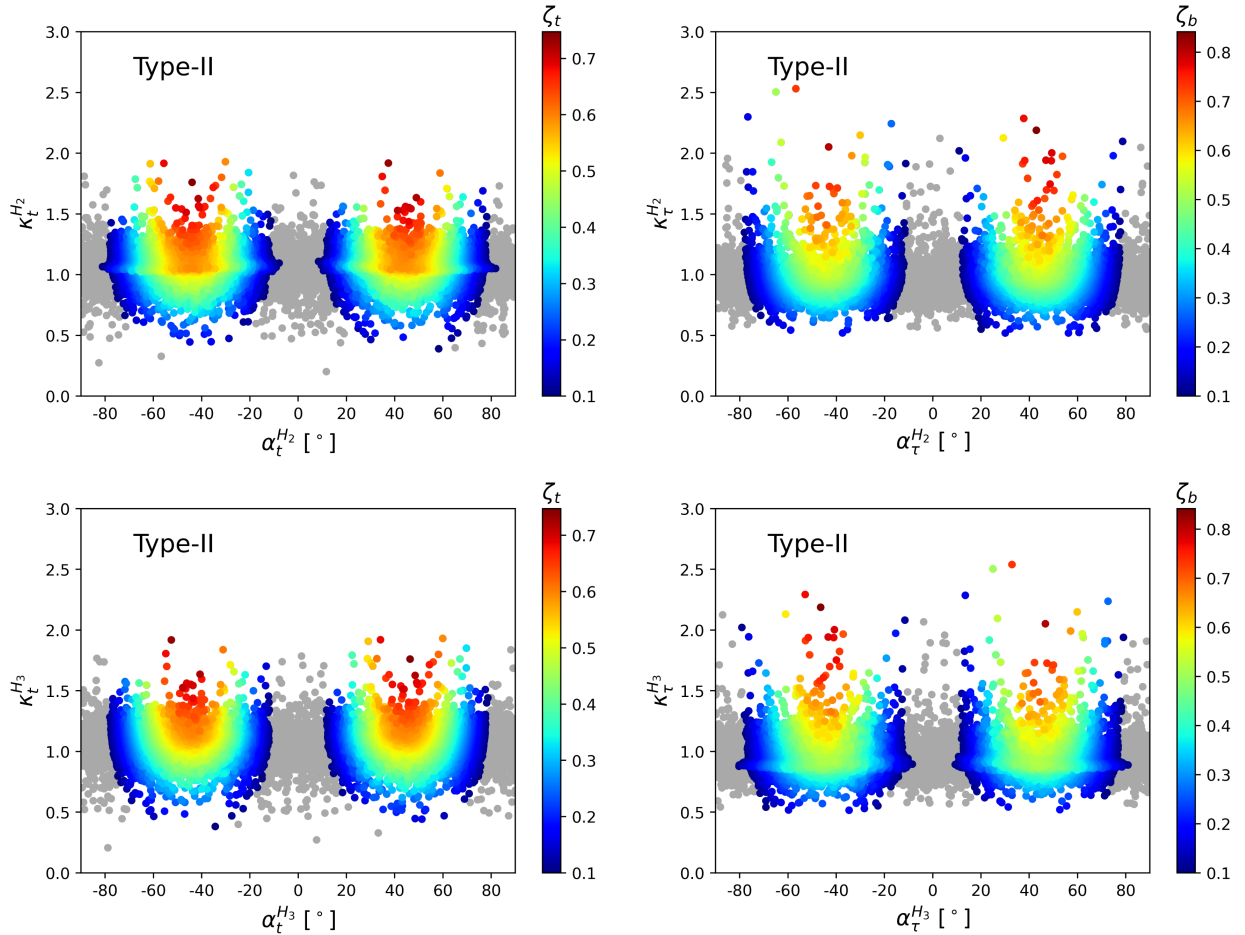


FIG. 15. Viable parameter points in the (α_t, κ_t) (left panels) and $(\alpha_\tau, \kappa_\tau)$ (right panels) planes for H_2 (upper panels) and H_3 (lower panels) in the type-II C2HDM. Gray points represent the full viable space; colored points satisfy the large CP violation conditions ($\zeta_t > 0.1$ left, $\zeta_b > 0.1$ right). The color code indicates the magnitude of ζ_t and ζ_b .

behavior in this regard, while their coupling magnitudes $\kappa_t^{H_{2,3}}$ and $\kappa_\tau^{H_{2,3}}$ remain broadly SM-like (typically within 1 ± 0.5 , though values up to ~ 2.5 are possible in rare high-CPV cases). Consequently, extending CP -violation searches to these BSM Higgs bosons at the HL-LHC is essential, as they offer significantly greater sensitivity to the underlying CP structure than measurements restricted to the SM-like Higgs boson.

D. Hidden CP violation in the near-alignment limit

We conclude our analysis by examining the CP violation arising from the mixing of the heavy neutral states H_2 and H_3 —a particularly compelling phenomenon that emerges in the near-alignment regime. We define this regime as

$$\text{Near-alignment limit: } \delta_{\text{H-align}} < 0.1. \quad (56)$$

In this limit, the 125 GeV Higgs boson H_1 is approximately CP -even and largely decoupled from the heavy sector. Nevertheless, sizable CP violation can still be realized

within the heavy sector through the mixing angle α_3 . Since this CP violation does not significantly manifest in the couplings of the observed 125 GeV state, we refer to it as *hidden CP violation*.

Probing hidden CP violation via the gauge sector is experimentally challenging because the near-alignment condition inherently suppresses the couplings of the heavy neutral scalars to a weak gauge boson pair VV ($V = W^\pm, Z$)

$$\begin{aligned} c(H_2 VV) &\simeq -c(H_3 VV) \\ &\simeq (c_{\alpha_3} - s_{\alpha_3})\delta_{\text{H-align}} + \mathcal{O}(\delta_{\text{H-align}}^3). \end{aligned} \quad (57)$$

Furthermore, the off-diagonal neutral Higgs couplings to the Z boson involving the SM-like state—specifically the $[H_1, H_2, Z^\mu]$ and $[H_1, H_3, Z^\mu]$ vertices—are similarly suppressed by factors of $\delta_{\text{H-align}}$.⁴ In striking contrast, the off-diagonal H_2 - H_3 - Z vertex remains robust [48]

⁴The diagonal neutral Higgs couplings to the Z boson, $[H_i, H_i, Z^\mu]$, vanish identically due to Bose symmetry.

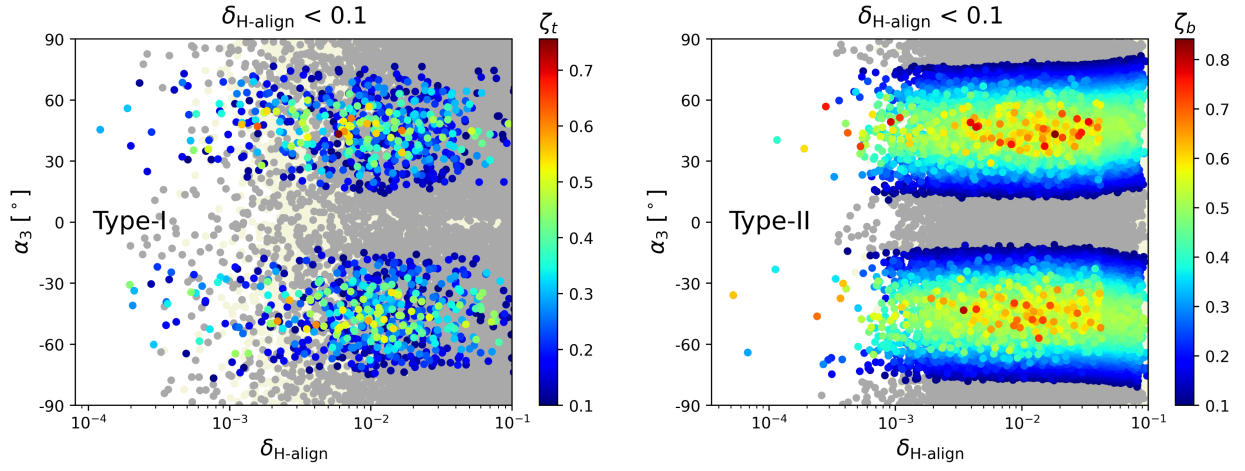


FIG. 16. Viable parameter points in the $(\delta_{\text{H-align}}, \alpha_3)$ plane for type-I (left panel) and type-II (right panel) C2HDM within the near-alignment regime ($\delta_{\text{H-align}} < 0.1$). Beige, gray, and colored points satisfy the same conditions as in Fig. 1.

$$[H_2, H_3, Z^\mu] \simeq \frac{g}{2c_W} (p_{H_2} - p_{H_3})^\mu + \mathcal{O}(\delta_{\text{H-align}}^2). \quad (58)$$

This indicates that the heavy states maintain a robust scalar-to-scalar transition through the Z boson. Crucially, this vertex is determined by the gauge structure of the model and remains independent of the CP -violating mixing. Such a robust vertex suggests a potential production channel at the LHC via $pp \rightarrow Z^* \rightarrow H_2 H_3$. To directly probe the CP -violating structure of the heavy sector, we propose examining the Yukawa couplings, which retain a strong dependence on α_3 even in the near alignment limit [see Eq. (35)]. This CP -violating structure remains fundamentally inaccessible through H_1 precision measurements alone.

To quantify the scope of this hidden CP violation, Fig. 16 presents the distribution of viable parameter points in the $(\delta_{\text{H-align}}, \alpha_3)$ plane. Here, we focus specifically on the near-alignment regime by restricting the visualization to points satisfying $\delta_{\text{H-align}} < 0.1$. This selective plotting highlights the behavior of the heavy-sector mixing in the limit where H_1 is most SM-like. The color coding of the points follows the conventions established in Fig. 1, allowing for a consistent interpretation of the underlying physics across the different parameter scans.

Two significant features emerge from the analysis of Fig. 16. First, our results demonstrate that sizable CP violation persists in the heavy neutral Higgs sector even as the 125 GeV state approaches the exact alignment limit. Specifically, for parameter points characterized by large CP violation (colored points), the mixing angle α_3 remains significantly displaced from the CP -conserving values of $\{0, \pm\pi/2\}$. This persistence reflects the nondecoupling behavior identified by Pilaftsis and Wagner [139], whereby CP violation decouples from the SM-like Higgs sector at large mass scales but remains operative within the heavy neutral Higgs system. In this regime, the strong mixing between the heavy states can facilitate a resonant enhancement of CP -violating effects

through scalar-pseudoscalar transitions, particularly when H_2 and H_3 are nearly degenerate. Our global scan confirms that the C2HDM provides a robust manifestation of this behavior, supporting sizable CP -violating observables while remaining consistent with the full suite of current theoretical and experimental constraints.

Second, the mixing angle α_3 exhibits a strong correlation with the Yukawa-sector CPV measures. In the near-alignment regime ($\delta_{\text{H-align}} \ll 1$), the leading contributions to these measures simplify to

$$\zeta_{23} \approx s_{2\alpha_3}^2 c_\beta^2, \quad \zeta_{13} \approx s_{2\alpha_3}^2 s_\beta^2, \quad (59)$$

where their specific relationships to ζ_t and ζ_b in type-I and type-II are detailed in Eq. (28). From these expressions, it is evident that both ζ_t and ζ_b are maximized at $\alpha_3 = \pm\pi/4$, corresponding to maximal mixing between the heavy neutral states.

The near-alignment regime is therefore the physically relevant region for heavy-sector CP violation. In this regime, α_3 retains physical meaning as the mixing angle of the heavy neutral mass eigenstates. Although the neutral-gauge couplings of $[H_{2,3}, V, V]$ and $[H_1, H_{2,3}, Z]$ are strongly suppressed near alignment, interactions involving $[H_2, H_3, Z]$ can still provide complementary information on the heavy-sector mixing structure.

V. CONCLUSIONS

In this work, we have performed a comprehensive global analysis of the complex two-Higgs-doublet model (C2HDM) with softly broken Z_2 symmetry, identifying the intrinsic features of the parameter space that support large CP violation while satisfying all current theoretical, collider, flavor, and EDM constraints. Our analysis reveals that type-I and type-II models follow sharply distinct pathways toward sizable CP violation and exhibit unique structural behaviors near the Higgs alignment limit.

In type-I, both the gauge- and Yukawa-sector CPV measures can reach significant magnitudes, with ξ_V rising to ~ 0.4 and ζ_I reaching values as high as ~ 0.8 , but only within a narrow region where the second neutral Higgs boson is nearly degenerate with the 125 GeV state H_1 . For the ensemble of parameter points satisfying the full set of theoretical and experimental constraints within our motivated mass range, the predicted electron EDM exhibits a characteristic scale of $|d_e| \gtrsim 10^{-31} e \text{ cm}$. This suggests that the physically viable type-I parameter space largely lies within the projected sensitivity of next-generation eEDM experiments, irrespective of the size of the CP -violating measures.

By contrast, type-II strongly suppresses gauge-sector CP violation while still allowing large Yukawa-sector CP violation. In addition, destructive interference among different contributions to the eEDM can suppress $|d_e|$ down to values of order $10^{-35} e \text{ cm}$. No phenomenologically relevant lower bound on $|d_e|$ emerges within the presently viable parameter space.

Finally, we have identified a novel phenomenon in the near-alignment limit: hidden CP violation. Although the heavy neutral scalars lose their couplings to gauge-boson pairs in this limit, their mutual CP mixing—parametrized by the angle α_3 —remains unconstrained and can become nearly maximal as the Yukawa-sector CPV measures increase. We have shown that this hidden CP violation can nevertheless be probed through CP -violating Yukawa couplings, providing a novel and necessary pathway to experimentally expose this otherwise elusive source of CP violation at future colliders.

Taken together, our analysis demonstrates that large CP violation in the C2HDM is highly structured and tightly constrained. The results presented here provide a systematic map of the viable parameter space and outline clear strategies for testing these scenarios through future EDM measurements, Higgs precision studies, and dedicated searches for the heavy Higgs sector.

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DATA AVAILABILITY

The datasets (parameter points) investigated in this work were generated numerically using the publicly available tool ScannerS. The specific model configurations, scanning boundaries, and constraints applied are fully detailed within the manuscript, allowing for the complete reproduction of the data.

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