Multiphoton signatures as a probe of *CP* violation in extended Higgs sectors

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We propose a novel signature with four-photon final states to probe CP-violating (CPV) extended Higgs sectors via $f\bar{f} \to Z^* \to H_1 H_2 \to 4\gamma$ processes with $H_{1,2}$ being additional neutral Higgs bosons. We focus on the nearly Higgs alignment scenario, in which the discovered Higgs boson almost corresponds to a neutral scalar state belonging to the isospin doublet field with the vacuum expectation value $v \simeq 246$ GeV. We show that the branching ratios of $H_{1,2} \to \gamma\gamma$ can simultaneously be sizable when CPV phases in the Higgs potential are of order one due to the enhancement of charged-Higgs boson loops. Such branching ratios can be especially significant when the fermiophobic scenario is taken into account. As a simple example, we consider the general two Higgs doublet model, and demonstrate that the cross section for the four-photon process can be 0.1 fb at LHC with the masses of $H_{1,2}$ to be a few 100 GeV in the Higgs alignment limit under the constraints from electric dipole moments (EDMs) and LHC Run-II data. We also illustrate that the searches for EDMs and diphoton resonances at high-luminosity LHC play complementary roles to explore CPV extended Higgs sectors.

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I. INTRODUCTION

CP-violation (CPV) is one of the necessary ingredients to explain the baryon asymmetry of the Universe [1]. Although nonzero CPV appears from the Kobayashi-Maskawa phase in the standard model (SM), its amount has been known to be too small to accommodate the observed value of the baryon asymmetry [2]. Therefore, new physics beyond the SM is required to provide additional sources of CPV.

A Higgs boson was discovered at LHC in 2012, and its properties, e.g., the mass, width, and couplings, have been measured from various production and decay channels. So far, the observed properties are consistent with those of the Higgs boson in the SM within the theoretical and experimental uncertainties [3,4]. This, however, does not necessarily mean that the Higgs sector is the minimal one assumed in the SM. In fact, it is indeed possible to realize nonminimal Higgs sectors with nearly Higgs alignment [5], in which couplings of the discovered Higgs boson (h) take almost the same values as those of the SM Higgs boson at tree level. Since the Higgs alignment can be compatible with CPV, e.g., in models with multi-Higgs doublets [6], it is now quite important to investigate CP-violating nonminimal Higgs sectors with the nearly Higgs alignment [7–9].

Searches for electric dipole moments (EDMs) can provide evidence for CPV. Currently, no clear signature for a nonzero EDM of elementary particles has been observed, and it has severely constrained a parameter space in nonminimal Higgs sectors with CPV. In particular, the magnitude of the electron EDM (eEDM) has been strongly constrained to be smaller than 4.1×10^{-30} e cm (90% CL) [10]. Future experiments for, e.g., the neutron EDM [11], would provide a clue of CPV. In addition to the EDMs, CP-violating effects can be tested at high energy collider experiments. For instance, the decay of neutral Higgs bosons into a tau-pair can be used to extract the CP-violating phase from the difference of the azimuthal angles defined by the tau decay plane [12–14], and the possibility of measuring the phase has been discussed at LHC [15–19] and at future electron-positron colliders [20,21]. The CP nature of the neutral Higgs boson can also be extracted via the top Yukawa coupling [22–24], the diboson decay [25] and also from Higgs to Higgs decays [26].

In this paper, we propose a novel approach to test nonminimal Higgs sectors with CPV at collider experiments. We focus on the four-photon final state driven by the electroweak (EW) pair production of additional neutral Higgs bosons H_1 and H_2 with their subsequent diphoton decays:

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$$f\bar{f} \to Z^* \to H_1 H_2 \to 4\gamma. \quad (1)$$

We show that the cross section for the above process can be significant in the presence of charged Higgs bosons when the CP-violating phase in the Higgs potential is sizable. We would like to emphasize that our approach can be applied to a plethora of extended Higgs sectors with CPV, and offers robust probe of CPV in the Higgs potential since the production part $f\bar{f} \to H_1 H_2$ is purely determined by the gauge coupling, by which the CP-violating nature can be extracted from the decays of $H_{1,2}$.

This paper is organized as follows. In Sec. II, we discuss the cross section and decay of H_1 and H_2 to realize the four-photon events in a model independent way. In Sec. III, we consider the general two Higgs doublet model (2HDM) as a representative example of the extended Higgs sector. Section IV is devoted to showing how large event numbers can be obtained for the four-photon process in the 2HDM in the region allowed by the existing experimental data. Finally, discussions and conclusions are given in Sec. V.

II. GENERAL SETUP

Let us consider a general setup in the EW $SU(2)_I \times U(1)_Y$ gauge theory with extended Higgs sectors. Suppose that Φ and φ are respectively the isospin Higgs doublet with the hypercharge $1/2$ and a complex scalar multiplet with the hypercharge Y_φ containing a neutral component φ^0. We focus on the nearly Higgs alignment scenario as it is favored by the current LHC data [3,4], where the Fermi constant G_F is mainly given by the vacuum expectation value (VEV) v, i.e., $v \equiv \sqrt{2}\langle \Phi^0 \rangle \simeq (\sqrt{2} G_F)^{-1/2}$ and $h \simeq \sqrt{2}\Re\Phi^0 - v$. In this section, we consider the case with the exact Higgs alignment for simplicity. In the next section, we also discuss the consequence of a slight deviation from the alignment limit in a concrete model. The real part $\varphi_H \equiv \sqrt{2}\Re(\varphi^0)$ and the imaginary part $\varphi_A \equiv \sqrt{2}\Im(\varphi^0)$ can mix if the Higgs potential contains CP-violating phases. Their mass eigenstates are defined as

$$\begin{pmatrix} \varphi_H \\ \varphi_A \end{pmatrix} = R(\theta) \begin{pmatrix} H_1 \\ H_2 \end{pmatrix}, \qquad R(\theta) \equiv \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix}. \quad (2)$$

The $H_1 H_2 Z^\mu$ vertex is given by

$$|D_\mu \varphi|^2 \supset g_Z Y_\varphi (H_1 \overleftrightarrow{\partial}_\mu H_2) Z^\mu, \quad (3)$$

where $A\overleftrightarrow{\partial}_\mu B \equiv A(\partial_\mu B) - (\partial_\mu A)B$ and $g_Z \equiv g/\cos\theta_W$ with g and θ_W being the $SU(2)_I$ gauge coupling and the weak mixing angle, respectively. It is clear that $Y_\varphi \neq 0$ is required to obtain the nonvanishing interaction, so that φ should be an isospin nonsinglet field. The cross section is then expressed at leading order as

$$\hat{\sigma} = \frac{16\pi\alpha_{\rm em}^2 Y_\varphi^2}{3N_f^c s \sin^4 2\theta_W} \frac{v_f^2 + a_f^2}{\left(1 - \frac{m_Z^2}{s}\right)^2} \lambda^{3/2}\left(\frac{m_{H_1}^2}{s}, \frac{m_{H_2}^2}{s}\right), \quad (4)$$

where s is the squared center-of-mass energy, $N_f^c = 3(1)$ for f being quarks (leptons), and $v_f = I_f/2 - Q_f \sin^2\theta_W$ and $a_f = I_f/2$ with I_f being the third component of the isospin of a fermion f. The phase space function is given by $\lambda(x,y) = (1-x-y)^2 - 4xy$. The cross section for the four-photon process (1) is then estimated by $\hat{\sigma} \times \xi$ with

$$\xi \equiv {\rm BR}(H_1 \to \gamma\gamma) \times {\rm BR}(H_2 \to \gamma\gamma). \quad (5)$$

For f to be quarks, the cross section should be written as $\sigma \times \xi$ with σ being the hadronic production cross section for $pp \to H_1 H_2$.

Next, we discuss the decays of $H_{1,2}$. We introduce the following Yukawa and scalar interactions:

$$\mathcal{L}_{\rm int} = -\frac{\sqrt{2}m_f}{v}\hat{\zeta}_f \bar{f}_L f_R \varphi - v\sum_\alpha \lambda_{S_\alpha S_\alpha^* \varphi} S_\alpha S_\alpha^* \varphi + {\rm H.c.}, \quad (6)$$

where f_L (f_R) is a left-handed (right-handed) SM fermion, and S_α are charged scalars with the electric charge Q_α. We here do not specify the other properties of S_α such as the isospin. These interaction terms can be rewritten in the basis of H_a as

$$\mathcal{L}_{\rm int} = -\sum_{a=1,2}\left[\frac{m_f}{v}\bar{f}(\kappa_a^f + i\gamma_5\tilde{\kappa}_a^f)f + v\sum_\alpha \lambda_{S_\alpha S_\alpha^* H_a} S_\alpha S_\alpha^*\right] H_a, \quad (7)$$

with

$$\begin{pmatrix} \kappa_1^f \\ \tilde{\kappa}_1^f \end{pmatrix} = R(\theta)\begin{bmatrix} \Re(\hat{\zeta}_f) \\ \Im(\hat{\zeta}_f) \end{bmatrix}, \qquad \begin{pmatrix} \kappa_2^f \\ \tilde{\kappa}_2^f \end{pmatrix} = R(\theta)\begin{bmatrix} -\Im(\hat{\zeta}_f) \\ \Re(\hat{\zeta}_f) \end{bmatrix},$$
$$\begin{pmatrix} \lambda_{S_\alpha S_\alpha^* H_1} \\ \lambda_{S_\alpha S_\alpha^* H_2} \end{pmatrix} = \sqrt{2} R^T(\theta)\begin{bmatrix} \Re(\lambda_{S_\alpha S_\alpha^* \varphi}) \\ -\Im(\lambda_{S_\alpha S_\alpha^* \varphi}) \end{bmatrix}. \quad (8)$$

When the Higgs alignment is exact, H_a do not decay into a weak boson pair, while they can decay into a fermion pair and/or a lighter additional Higgs boson associated with a (off-shell) weak boson at tree level. At one-loop level, H_a can decay into $\gamma\gamma$, $Z\gamma$, and gg. In order to discuss how the diphoton decay can be important, we give the decay rates into $f\bar{f}$, gg, and $\gamma\gamma$ as follows

$$\Gamma(H_a \to f\bar{f}) = \frac{N_f^c m_{H_a}^3}{32\pi v^2}\tau_a^f\left(|\kappa_a^f|^2 - \tau_a^f[\Re(\kappa_a^f)]^2\right)\sqrt{1 - \tau_a^f}, \quad (9)$$

$$\Gamma(H_a \to gg) = \frac{\alpha_s^2 m_{H_a}^3}{128\pi^3 v^2}\left\{\left|\sum_q \kappa_a^q I_{\frac{1}{2}}(\tau_a^q)\right|^2 + \left|\sum_q \tilde{\kappa}_a^q \tilde{I}_{\frac{1}{2}}(\tau_a^q)\right|^2\right\}, \tag{10}$$

$$\Gamma(H_a \to \gamma\gamma) = \frac{\alpha_{\rm em}^2 m_{H_a}^3}{256\pi^3 v^2}\left\{\left|\sum_f Q_f^2 N_f^c \kappa_a^f I_{\frac{1}{2}}(\tau_a^f) + \sum_\alpha Q_\alpha^2 \frac{v^2 \lambda_{S_\alpha S_\alpha^* H_a}}{m_{H_a}^2} I_0(\tau_a^{S_\alpha})\right|^2 + \left|\sum_f Q_f^2 N_f^c \tilde{\kappa}_a^f \tilde{I}_{\frac{1}{2}}(\tau_a^f)\right|^2\right\}, \tag{11}$$

where $\tau_a^X = 4m_X^2/m_{H_a}^2$. The loop functions are given by [27]

$$I_0(x) = 2[1 - xf(x)],$$
$$I_{\frac{1}{2}}(x) = 2x[(x-1)f(x) - 1], \qquad \tilde{I}_{\frac{1}{2}}(x) = 2xf(x), \tag{12}$$

with

$$f(x) = \begin{cases} \arcsin^2\sqrt{x^{-1}} & (x \geq 1), \\ -\frac{1}{4}\left[\ln\frac{1+\sqrt{1-x}}{1-\sqrt{1-x}} - i\pi\right]^2 & (x < 1). \end{cases} \tag{13}$$

In Eq. (11), the contribution from the W boson loop is neglected, because of the Higgs alignment condition. We note that the decay rates of $H_a \to Z\gamma$ can be comparable with those of $H_a \to \gamma\gamma$ as long as $m_{H_a} \gg m_Z$, which are included in our numerical analysis given below.

In the *CP*-conserving (CPC) limit, i.e., $\theta \to 0$ and $\Im(\hat{\zeta}_f) = \Im(\lambda_{S_\alpha^* S_\alpha \varphi}) = 0$, H_1 (H_2) behaves as a *CP*-even (*CP*-odd) scalar boson, and the $S_\alpha^\pm$ loop contribution to $H_2 \to \gamma\gamma$ vanishes. In this case, $\mathrm{BR}(H_2 \to \gamma\gamma)$ cannot be significant due to the dominant $H_2 \to f\bar{f}/gg$ modes. In fact, when we consider only the top-loop contribution to the $H_2 \to \gamma\gamma/gg$ modes, the ratio $\Gamma(H_2 \to \gamma\gamma)/\Gamma(H_2 \to gg)$ is given by $(\alpha_{\rm em} N_t^c Q_t^2/\sqrt{2}\alpha_s)^2 \simeq 4 \times 10^{-3}$. Thus, $H_2 \to \gamma\gamma$ cannot be the dominant mode. For $m_{H_2} \geq 2m_t$, $\mathrm{BR}(H_2 \to \gamma\gamma)$ is even more suppressed by the $H_2 \to t\bar{t}$ mode. On the other hand, for the case with CPV, the $S_\alpha^\pm$-loop contributes to the $H_2 \to \gamma\gamma$ mode, so that $\mathrm{BR}(H_2 \to \gamma\gamma)$ can be large. In particular, if both $\lambda_{S_\alpha^* S_\alpha H_1}$ and $\lambda_{S_\alpha^* S_\alpha H_2}$ are relatively larger than the $\hat{\zeta}_f$ parameters, *both* the branching ratios of $H_{1,2} \to \gamma\gamma$ can be sizable. Therefore, a larger value of ξ defined in Eq. (5) can be a telltale sign of CPV in the Higgs sector.

III. CONCRETE MODELS

Let us discuss the four-photon process (1) in the general 2HDM without imposing any additional symmetries as a prototype of an extended Higgs sector. The multiplet φ is then identified with another doublet field Φ' with $Y_{\Phi'} = 1/2$, and the charged scalars $S_\alpha^\pm$ are identified with the charged component of Φ', i.e., $\Phi'^\pm (\equiv H^\pm)$. We can take $\langle\Phi'\rangle = 0$ without loss of generality, because Φ' can be regarded as the field defined in the Higgs basis [5].

The most general Higgs potential is written as

$$\begin{aligned} V = {} & m^2|\Phi|^2 + M^2|\Phi'|^2 - (\mu^2\Phi^\dagger\Phi' + \text{H.c.}) + \frac{\lambda_1}{2}|\Phi|^4 \\ & + \frac{\lambda_2}{2}|\Phi'|^4 + \lambda_3|\Phi|^2|\Phi'|^2 + \lambda_4|\Phi^\dagger\Phi'| \\ & + \left(\frac{\lambda_5}{2}\Phi^\dagger\Phi' + \lambda_6|\Phi|^2 + \lambda_7|\Phi'|^2\right)(\Phi^\dagger\Phi') + \text{H.c.}, \end{aligned} \tag{14}$$

where μ^2 and $\lambda_{5,6,7}$ are generally complex parameters. Imposing the stationary conditions, we can eliminate the parameters m^2 and μ^2. The mass matrix for the neutral Higgs bosons are then given in the basis of $(\sqrt{2}\Re\Phi^0, \sqrt{2}\Re\Phi'^0, \sqrt{2}\Im\Phi'^0)$ as

$$v^2\begin{pmatrix} \lambda_1 & \Re\lambda_6 & -\Im\lambda_6 \\ \Re\lambda_6 & \frac{M^2}{v^2} + \frac{\lambda_3+\lambda_4+\Re\lambda_5}{2} & -\frac{\Im\lambda_5}{2} \\ -\Im\lambda_6 & -\frac{\Im\lambda_5}{2} & \frac{M^2}{v^2} + \frac{\lambda_3+\lambda_4-\Re\lambda_5}{2} \end{pmatrix}. \tag{15}$$

We can remove the phase of λ_5 by the field redefinition without loss of generality. The Higgs alignment condition is given by

$$\lambda_6 = 0, \tag{16}$$

in which the mass matrix takes the diagonal form, i.e., $\theta = 0$ in Eq. (2), and we can identify $(\sqrt{2}\Re\Phi^0, \sqrt{2}\Re\Phi'^0, \sqrt{2}\Im\Phi'^0) = (h, H_1, H_2)$ among which h can be regarded as the discovered Higgs boson with the mass of 125 GeV. The basis invariant form of CPV in the 2HDM has been found in Ref. [28] as follows:

$$J_1 \propto \Im[\lambda_5^*\lambda_6^2], \qquad J_2 \propto \Im[\lambda_5^*\lambda_7^2], \qquad J_3 \propto \Im[\lambda_6^*\lambda_7], \tag{17}$$

where *CP*-symmetry is broken if at least one of the three invariants is nonzero. Therefore, our scenario $\lambda_6 = \Im\lambda_5 = 0$ with $\Im\lambda_7 \neq 0$ gives $J_2 \neq 0$, and we definitely have CPV in the potential. Under $\lambda_6 = \Im\lambda_5 = 0$, J_2 can also be written as

$$J_2 \propto (m_{H_1}^2 - m_{H_2}^2)\Im[\lambda_7^2]. \tag{18}$$

In this scenario, the scalar couplings $\lambda_{S_\alpha S_\alpha^* H_a}$ defined in Eq. (7) are expressed as $\lambda_{H^+H^-H_1} = -v\Re[\lambda_7]$ and

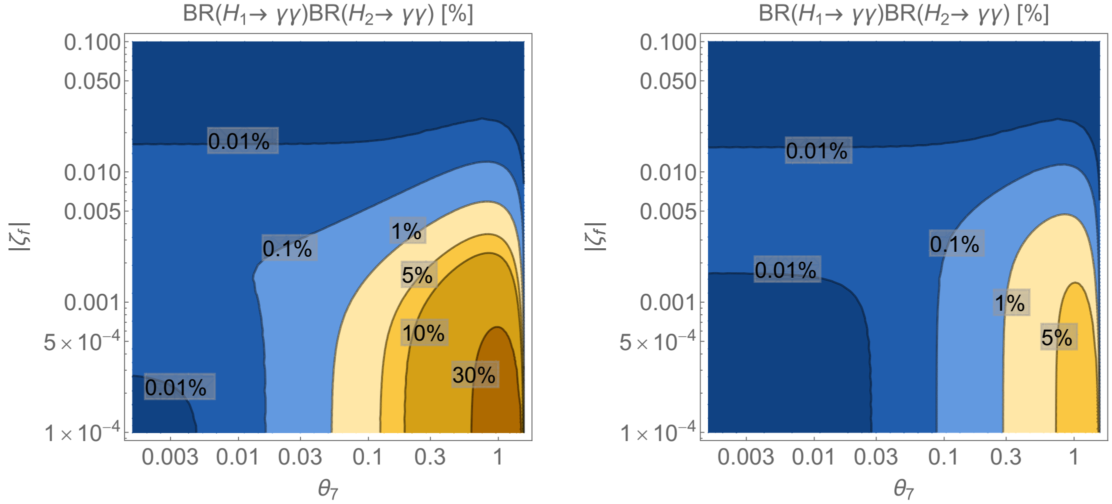

FIG. 1. Contour plots of $\xi = \mathrm{BR}(H_1 \to \gamma\gamma) \times \mathrm{BR}(H_2 \to \gamma\gamma)$ on the $\theta_7 - |\zeta_f|$ plane in the 2HDM with the exact Higgs alignment, i.e., $\lambda_6 = 0$. The left and right panels show the case with $m_{H_2} - m_{H_1}$ to be 5 and 10 GeV, respectively. For all the plots, we take $\zeta_f = \zeta_u = \zeta_d = \zeta_e$, $\arg(\zeta_f) = 0$, $m_{H_1} = m_{H^\pm} = 250$ GeV, and $|\lambda_7| = 1$.

$\lambda_{H^+H^-H_2} = v\Im[\lambda_7]$. Using these expressions, the invariant J_2 is also written as

$$J_2 \propto (m_{H_1}^2 - m_{H_2}^2)\lambda_{H^+H^-H_1}\lambda_{H^+H^-H_2}. \tag{19}$$

This suggests that J_2 becomes significant when the product of the scalar couplings $\lambda_{H^+H^-H_1}\lambda_{H^+H^-H_2}$ and/or the mass difference between m_{H_1} and m_{H_2} are taken to be larger values. In particular, as we can see in Eq. (11) a larger value of $\lambda_{H^+H^-H_1}\lambda_{H^+H^-H_2}$ can lead to larger branching ratios of $H_{1,2} \to \gamma\gamma$, i.e., larger ξ defined in Eq. (5). It is also worth to mention here that the phase of λ_7 turns out to be unphysical when $m_{H_1} = m_{H_2}$ as it is seen in Eq. (18). Therefore, a large number of the four-photon events via the H_1H_2 production can be a probe of CPV in the 2HDM if two different masses are reconstructed from the events.

The Yukawa interactions are generally given in the mass basis for fermions as

$$\begin{aligned}\mathcal{L}_Y = -\frac{\sqrt{2}}{v}[&\bar{Q}_L^u M_u i\tau_2(\Phi^* + \rho_u\Phi'^*)u_R \\ &+ \bar{Q}_L^d M_d(\Phi + \rho_d\Phi')d_R + \bar{L}_L M_e(\Phi + \rho_e\Phi')e_R] \\ &+ \mathrm{H.c.},\end{aligned} \tag{20}$$

where $Q_L^d = (V^\dagger u_L, d_L)^T$ and $Q_L^u = (u_L, V d_L)^T$ with V being the Cabibbo-Kobayashi-Maskawa matrix. In the above expression, M_f $(f = u,\ d,\ e)$ are the diagonalized mass matrix, and ρ_f are general complex 3×3 matrices. In order to avoid flavor-changing neutral currents via Higgs boson mediations at tree level, we impose the Yukawa alignment [29], i.e.,

$$\rho_f = \zeta_f I_{3\times 3} \quad (f = u, d, e), \tag{21}$$

where ζ_f are complex parameters and $I_{3\times 3}$ is the 3×3 unit matrix. Comparing Eq. (6), we can identify $\hat{\zeta}_f = \zeta_f$.

In Fig. 1, we show the contour of ξ on the $\theta_7(\equiv \arg[\lambda_7]) - |\zeta_f|$ plane in the 2HDM. We take the mass difference $\Delta m \equiv m_{H_2} - m_{H_1}$ to be 5 GeV (left) and 10 GeV (right). As expected, in the CPC limit $\theta_7 \to 0$, ξ is given to be of order 0.1% or smaller, because $H_2 \to \gamma\gamma$ cannot be significant. On the other hand, ξ takes larger values for larger θ_7 and/or smaller $|\zeta_f|$. It is also seen that a larger value of ξ is realized for smaller Δm, because the decay mode $H_2 \to H_1 Z^*$ is phase space suppressed.

Let us discuss how the branching ratio is modified if we consider the case with a slight deviation from the Higgs alignment limit, i.e., $\lambda_6 \neq 0$. In Fig. 2, we show the branching ratios of H_1 (left) and H_2 (right) as a function of λ_6. It is clear that both the branching ratios of H_1 and H_2 into diphoton are highly suppressed as λ_6 increases, while the $H_1 \to WW/ZZ$ and $H_2 \to hZ$ modes dominate. In Fig. 3, we show the contour plot for the value of ξ on the λ_6 and Δm plane. The result in the 2HDM is shown in the left panel, in which we see that ξ can be $\mathcal{O}(10)\%$ for λ_6 and Δm to be smaller than about 10^{-4} and 10 GeV, respectively. For fixed values of λ_6 and Δm, a larger value of ξ can be obtained by introducing additional charged scalars in the 2HDM. For instance, if we introduce an additional charged singlet scalar[1] with the electric charge Q_S, the same mass and $\lambda_{S_a S_a^* H_a}$ as those of $H^\pm$, then values of ξ are enhanced as shown in the center (the case with $Q_S = 1$) and the right (the case with $Q_S = 2$) panels of Fig. 3.

[1]Charged singlet scalars are introduced in models with radiative generation of neutrino masses, e.g., the model proposed by Babu and Zee [30,31].

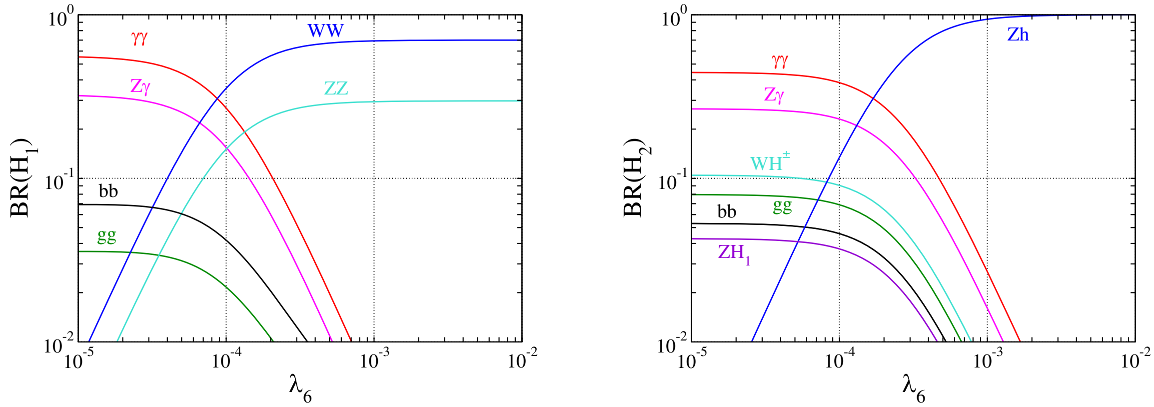

FIG. 2. Branching ratios of H_1 (left) and H_2 (right) as a function of λ_6 in the 2HDM. We take $m_{H_1} = m_{H^\pm} = M = 250$ GeV, $m_{H_3} = 255$ GeV, $\zeta_f = \zeta_u = \zeta_d = \zeta_e = 10^{-3}$, $\arg(\zeta_f) = 0$, and $|\lambda_7| = 1$.

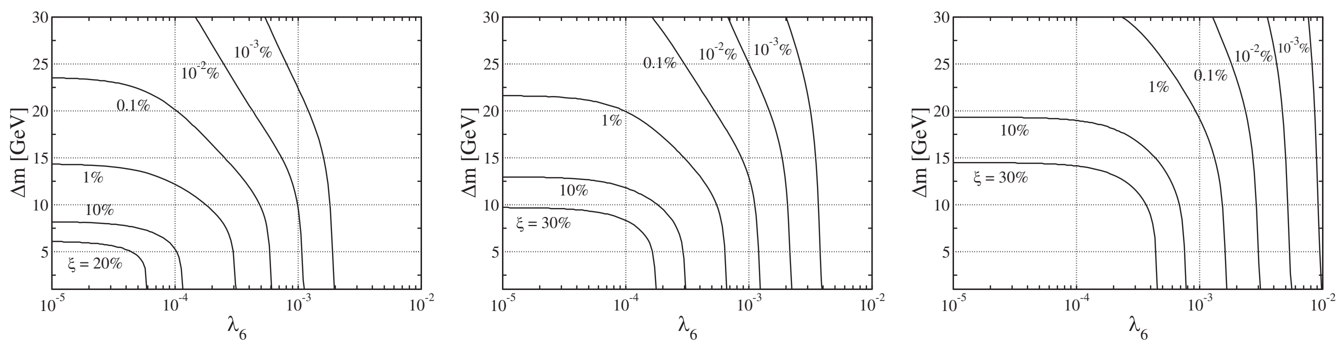

FIG. 3. Contour plots of $\xi = \mathrm{BR}(H_1 \to \gamma\gamma) \times \mathrm{BR}(H_2 \to \gamma\gamma)$ on the $\lambda_6 - \Delta m(\equiv m_{H_2} - m_{H_1})$ plane in the 2HDM (left panel) and that with an additional charged-singlet scalar having the electric charge Q_S (center panel for $Q_S = 1$ and right panel for $Q_S = 2$). We take $\zeta_f = \zeta_u = \zeta_d = \zeta_e = 10^{-3}$, $\arg(\zeta_f) = 0$, $m_{H_1} = m_{H^\pm} = M = 250$ GeV, $|\lambda_7| = 1$, and $\theta_7 = \pi/4$.

Let us comment on the other possibilities for the multiplet φ than the isospin doublet. As mentioned in Sec. II, φ cannot be an isospin singlet because of $Y_\varphi = 0$. For an isospin triplet, we can consider the one with $Y_\varphi = 1$, but this model does not contain a physical *CP*-violating phase in the potential [32]. Thus, ξ cannot be large. The same thing holds for models with φ whose isospin is larger than triplet except for the case with φ to be quadruplet with $Y_\varphi = 1/2$. For the latter, the potential contains two terms $(\Phi\varphi^*)^2$ and $(\Phi\Phi^*\Phi^*\varphi)$, and one of the phases for these couplings can be physical, so that a large ξ value can be realized. For models with more than one extra scalar fields, e.g., a model with two triplets [32,33], physical *CP*-violating phases can appear in the potential, and a larger value of ξ can be realized.

IV. FOUR-PHOTON PROCESS AT LHC

We discuss how large cross section of the four-photon process (1) can be obtained at LHC in the 2HDM with the Higgs alignment.

We take into account the constraints from the eEDM experiments, $|d_e| < 4.1 \times 10^{-30}$ e cm (90% CL) [10]. Detailed discussions for EDMs have been performed in the Yukawa aligned 2HDM in Ref. [6], and we apply the expressions given in that paper to survey the parameter region allowed by the data. As we will see below, the eEDM data excludes the case with ζ_f to be larger than 10^{-2} with θ_7 being $\mathcal{O}(1)$, while most of the region of interest, i.e., ζ_f to be smaller than 10^{-3}, is allowed. We note that the constraints from the other EDMs such as the neutron EDM [34] do not further exclude the parameter space allowed by the eEDM. In addition, we impose the following two constraints from LHC: (A) searches for a diphoton resonance [35] and (B) those for multiphoton ($\geq 3\gamma$) final states [36]. For (A), we consider the gluon fusion (ggF) $gg \to H_a$ [37] and the EW $q\bar{q}' \to H^\pm H_a$ [38] production processes. We estimate the production cross section for ggF using `SusHi` [39,40] at NNLO in QCD. Since the cross section for ggF is proportional to $|\zeta_u|^2$, the limit from ggF is negligible for $|\zeta_f| \ll 1$. However, the EW production remains crucial and deliver a stringent limit on the

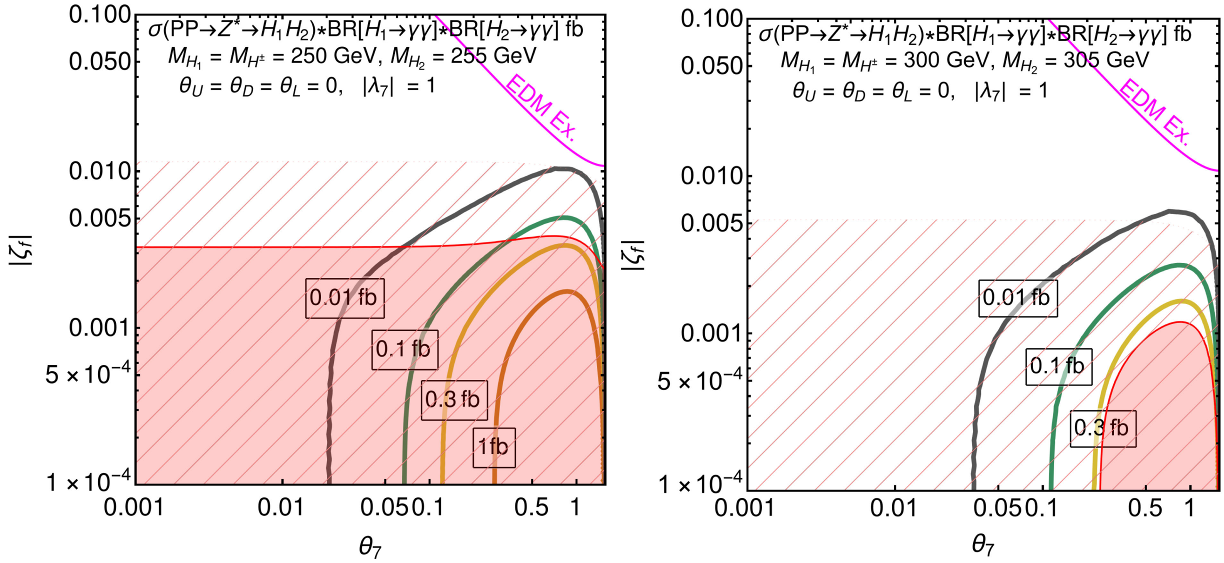

FIG. 4. Contour plots of the cross section for the four-photon process given in (1) at LHC on the $\theta_7 - |\zeta_f|$ plane in the 2HDM. The left and right panels show the case with $m_{H_1} = m_{H^\pm} = 250$ and 300 GeV, respectively. For all the plots, we take $m_{H_2} - m_{H_1} = 5$ GeV, $|\lambda_7| = 1$ and $\theta_f = 0$. The red shaded region is excluded by the constraint from the diphoton searches, while the region above the magenta solid curve is excluded by the constraint from the eEDM. The hashed region is expected to be explored with more than 2σ level from diphoton searches at the HL-LHC with 3 ab^{-1}.

parameter space. For (B), the EW production H_1H_2 with their $\gamma\gamma$ and/or $Z\gamma$ decays can give rise to the multiphoton signal. In addition to the constraints (A) and (B), there are the other constraints from direct searches for additional Higgs bosons at LHC. In the scenario of our interest, $H_{1,2}$ can decay into a fermion pair such as $H_{1,2} \to b\bar{b}$, see Fig. 2 with $\lambda_6 \ll 1$, and the charged Higgs bosons mainly decay into $H^\pm \to tb$. Constraints on such decay modes from current LHC data have been studied e.g., in Ref. [41], in which a case with $\tan\beta \lesssim 2$ (the ratio of two VEVs) and a few hundred GeV of the masses for additional Higgs bosons has been excluded in the type-I 2HDM. Our choice $|\zeta_f| \ll 1$ can correspond to the case with $\tan\beta \gg 1$ in the type-I 2HDM, so that we can safely avoid these constraints. In what follows, we consider the case with $m_{H_1} = m_{H^\pm}$ and $m_{H_2} > m_{H_1}$, so that the decay $H_2 \to H_1 Z^*$ provides $H_1H_1Z^*$ in the intermediate state, and it can also contribute to the multiphoton signal.

Figure 4 shows one of the main results of our analysis in the 2HDM for $m_{H_1} = 250$ GeV (left panel) and 300 GeV (right panel) with $m_{H_2} - m_{H_1} = 5$ GeV. Each contour shows the cross section for the four-photon final state under all the experimental constraints explained above. We see that the severe bound exists from the bound (A) indicated by the red shaded region. On the other hand, the limit from the multiphoton searches (B) do not appear in these figures, because only the Run-I data with at 8 TeV and 20 fb^{-1} are available, which are rather weak to exclude the parameter region shown here. The constraint from the eEDM excludes the region with larger $|\zeta_f|$ and/or larger θ_7. We also show the region expected to be explored with more than 2σ level at the high-luminosity LHC (HL-LHC) by the hashed region, which is obtained by extrapolating the current result of the diphoton search [35]. We see that the searches for EDM and the diphoton resonance at LHC play complementary roles to each other. We would like to emphasize that the diphoton signal, although far-reaching, does not specify the CP nature of the scalar potential. Searching for the proposed four-photon signal is essentially important to probe CPV in the extended Higgs sector at LHC.

V. DISCUSSIONS AND CONCLUSIONS

We comment on four-photon final states realized in the other scenarios. In the CPC type-I 2HDM, one can consider the process $gg \to h \to H_1H_1 \to 4\gamma$ [42]. There are two crucial differences between the above process and ours, i.e., (i) the invariant mass distribution for the diphoton system shows only one peak at m_{H_1} in the above but two peaks at m_{H_1} and m_{H_2} in our process and (ii) a deviation from the Higgs alignment is required to obtain the $h \to H_1H_1$ decay in the above process. We also note that the 2HDMs with a softly-broken Z_2 symmetry, including the type-I 2HDM, can provide a nonzero CP-violating phase in the potential, while this phase introduces a mixing among three neutral Higgs bosons. Therefore, such 2HDMs may be able to give a larger value of ξ, but they also introduce a larger deviation in the couplings of h from the SM values.

To conclude, it is worthwhile to systematically investigate the multiphoton process at LHC in addition to diphoton processes. As we have shown in Fig. 4, the diphoton search at the HL-LHC can explore most regions of interest as a discovery mode of extra Higgs bosons, in which the cross section of the four-photon events can maximally be a few fb level under the constraints from current LHC data. The important thing is that the diphoton events can be significant even in the CPC case, while the four-photon process can be sizable only when the CP-violation is realized in the Higgs sector. We thus would like to emphasize that the four-photon search would play an important role to test the CPV nature of extended Higgs sectors after the extra Higgs bosons are discovered via the diphoton search. We have also shown that to obtain substantial $\mathcal{O}(1)\%$ four-photon branching ratio, the magnitudes of the Higgs alignment parameter λ_6, the mass difference Δm and the Yukawa couplings ζ_f have to typically be smaller than of order $10^{-3}-10^{-4}$, 10 GeV, and 10^{-2}, respectively.

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