Rare top decay $t \to c\gamma$ with flavor changing neutral scalar interactions in two Higgs doublet model

R. Gaitán* and J. H. Montes de Oca†

Departamento de Física, FES-Cuautitlán, UNAM, C.P. 54770, Cuautitlán Izacalli, Estado de México, México

E. A. Garcés‡

Departamento de Física, Centro de Investigación y de Estudios Avanzados del IPN, Apdo. Postal 14-740, 07000 Ciudad de México, Mexico

R. Martinez§

Departamento de Física, Universidad Nacional de Colombia, Ciudad Universitaria, K. 45 No. 26-85, Bogotá D.C., Colombia

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Models beyond the standard model with extra scalars have been highly motivated by the recent discovery of the Higgs boson. The two Higgs doublet model type III considers the most general case for the scalar potential, allowing mixing between neutral CP-even and CP-odd scalar fields. This work presents the results of the study on the $t \to c\gamma$ decay at one loop level if neutral flavor changing is generated by top-charm-Higgs coupling given by the Yukawa matrix. For instance, a value for the branching ratio $\mathrm{Br}(t \to c\gamma) \sim 10^{-6}$ for $\tan\beta = 2.5$ and general neutral Higgs mixing parameters, $1.16 \leq \alpha_1 \leq 1.5$, $-0.48 \leq \alpha_2 \leq -0.1$. The number of events for the $t \to c\gamma$ decay with an integrated luminosity of $300\ \mathrm{fb}^{-1}$ is estimated as $10 \lesssim N_{\mathrm{Eff}} \lesssim 100$ for the parameters of the model constrained by experimental data.

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I. INTRODUCTION

The observation of the scalarlike Higgs boson with a mass of 126 GeV at the Large Hadron Collider (LHC) [1,2] has motivated the study of extended models with multiple scalar multiplets. The mass hierarchy between the up-type and down-type quarks suggests the consideration of models with two complex $SU(2)_L$ doublet scalar fields, referred to as two Higgs doublet models (THDM). There are two versions of THDM, labeled as type I and type II, with invariance under a Z_2 discrete symmetry that ensures CP conservation in the scalar sector [3]. In the first case, all quarks acquire mass through one doublet [4,5] whereas in type-II models [6] one doublet gives mass to the up-type quarks while the other doublet gives mass to the down-type quarks.

In the so-called type-III models both doublets simultaneously give masses to all quark types, which we refer to as model III [7]. In any type of THDM five physical Higgs particles are predicted; three of them are neutral with CP-even or CP-odd states and a charged pair. An important feature in model III is the mixing between the CP-even and CP-odd states for neutral scalar fields given by the mixing parameters α_1, α_2, and α_3 [8–10]. Current measurements in the LHC imply that the 126 GeV scalar particle is in good agreement with the Higgs boson being CP even [11,12].

Model III without Z_2 discrete symmetry is a general version that generates flavor changing (FC) neutral scalar interactions (FCNSI) in Higgs-fermion Yukawa couplings and CP violation in the Higgs potential [13–17]. One motivation to look for new sources of CP violation beyond the SM is the matter-antimatter problem [18,19] as well as the fermion electric dipole moments [20–23]. On the side of the FCNSI, a motivation arises from the study of the flavor changing neutral current (FCNC) processes, which are extremely suppressed in the standard model (SM), for instance, $\mathrm{Br}(t \to q + x) \approx 10^{-17}$–$10^{-12}$ with $q = c$, u and $x = \gamma$, Z, g, H [24–32]. In particular, we are interested in the $t \to c\gamma$ rare decay. The LHC excludes the ranges of $\mathrm{Br}(t \to c\gamma) > 5.9 \times 10^{-3}$; meanwhile in future results it is expected to set an upper bound of order 10^{-5} [33]. Also for other rare top-quark decays with FCNC, current experimental limits are $\mathrm{Br}(t \to cg) < 1.6 \times 10^{-4}$ at 95% C.L. [34], $\mathrm{Br}(t \to cZ) < 2.1 \times 10^{-3}$ at 95% C.L. [33], $\mathrm{Br}(t \to ch) < 5.6 \times 10^{-3}$ [35].

A value for $\mathrm{Br}(t \to c\gamma) \sim 10^{-8}$ has been estimated with charged Higgs mass $m_{H^\pm} \sim 200$ GeV as well as small values of the β mixing parameter, $\tan\beta = 0.1$ [36]. A detailed study in the framework of model III with FCNC shows more feasible values for branching ratio in the range $10^{-12} < \mathrm{Br}(t \to c\gamma) < 10^{-7}$ with the masses of the scalars between 200 and 800 GeV [7,37–40]. For the different

*rgaitan@unam.mx
†josehalim@gmail.com
‡egarces@fisica.unam.mx
§remartinezm@unal.edu.co

THDM types, the $\mathrm{Br}(t \to c\gamma)$ is enhanced for specific regions of scalar masses and mixing parameters [30,41,42].

The rare top decay has been analyzed in extended models other than THDM, for instance, [43–47]. In a previous work [48,49], it was shown that $\mathrm{Br}(t \to c\gamma)$ is sensitive to $\tan\beta$ in the framework of model III, obtaining $\mathrm{Br}(t \to c\gamma) \sim 1 \times 10^{-6}$ for $8 \leq \tan\beta \leq 15$. The rare top-quark decays at one loop with FCNC coming from additional fermions and gauge bosons have been studied in several extensions of the SM such as the minimal supersymmetric standard model, left-right symmetry models, top color assisted technicolor, little Higgs and two Higgs doublets with four generations of quarks [7,32,36–40,50]. FCNC and CP violation between quarks and scalars can also contribute to interactions with the rare top decay [51,53]. A recent study on generic FCNC top decays can be found in [52].

The content of this paper is as follows. The next section introduces the model and the interactions between quarks and neutral Higgs bosons. In Sec. III, we calculate $\mathrm{Br}(t \to c\gamma)$ in the framework of model III with FCNSI including CP violation in the scalar sector; in Sec. IV we present the restrictions to the parameters involved in the rare top decay. We present the results and discussion of our analysis in Sec. V.

II. FLAVOR CHANGING NEUTRAL SCALAR INTERACTIONS

Given Φ_1 and Φ_2, two complex $SU(2)_L$ doublet scalar fields with hypercharge 1, the most general gauge invariant and renormalizable Higgs scalar potential is [54]

$$\begin{aligned} V = {} & m_{11}^2\Phi_1^+\Phi_1 + m_{22}^2\Phi_2^+\Phi_2 - [m_{12}^2\Phi_1^+\Phi_2 + \mathrm{H.c.}] + \frac{1}{2}\lambda_1(\Phi_1^+\Phi_1)^2 + \frac{1}{2}\lambda_2(\Phi_2^+\Phi_2)^2 \\ & + \lambda_3(\Phi_1^+\Phi_1)(\Phi_2^+\Phi_2) + \lambda_4(\Phi_1^+\Phi_2)(\Phi_2^+\Phi_1) \\ & + \left[\frac{1}{2}\lambda_5(\Phi_1^+\Phi_2)^2 + \lambda_6(\Phi_1^+\Phi_1)(\Phi_1^+\Phi_2) + \lambda_7(\Phi_2^+\Phi_2)(\Phi_1^+\Phi_2) + \mathrm{H.c.}\right], \end{aligned} \tag{1}$$

where m_{11}^2, m_{22}^2, and λ_1, λ_2, λ_3, λ_4 are real parameters and m_{12}^2, λ_5, λ_6, λ_7 can be complex parameters. The most general $U(1)_{\mathrm{EM}}$-conserving vacuum expectation values (VEV) are

$$\langle\Phi_1\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_1 \end{pmatrix}, \tag{2}$$

$$\langle\Phi_2\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_2 e^{i\xi} \end{pmatrix}, \tag{3}$$

where v_1 and v_2 are real and non-negative, $0 \leq |\xi| \leq \pi$, and $v^2 \equiv v_1^2 + v_2^2 = \frac{4M_W^2}{g^2} = (246\ \mathrm{GeV})^2$. Without loss of generality, the phase in Eq. (2) was eliminated through the $U(1)_Y$ global invariance, leaving the ξ phase in the VEV of Eq. (3). This ξ phase is a source of spontaneous CP violation that can be absorbed by redefining the free parameters [55].

The neutral components of the scalar Higgs doublets in the interaction basis are $\frac{1}{\sqrt{2}}(v_a + \eta_a + i\chi_a)$, where $a = 1, 2$. As a result of the explicit CP-symmetry breaking, a mixing matrix R relates the mass eigenstates h_i with the η_i as follows,

$$h_i = \sum_{j=1}^{3} R_{ij}\eta_j, \tag{4}$$

where the state orthogonal to the Goldstone boson associated to the Z boson is $\eta_3 = -\chi_1 \sin\beta + \chi_2 \cos\beta$ and R is parametrized as [56]

$$R = \begin{pmatrix} c_1c_2 & s_1c_2 & s_2 \\ -(c_1s_2s_3 + s_1c_3) & c_1c_3 - s_1s_2s_3 & c_2s_3 \\ -c_1s_2c_3 + s_1c_3 & -(c_1s_1 + s_1s_2c_3) & c_2c_3 \end{pmatrix}, \tag{5}$$

with $c_i = \cos\alpha_i$, $s_i = \sin\alpha_i$ for $-\frac{\pi}{2} \leq \alpha_{1,2} \leq \frac{\pi}{2}$, and $0 \leq \alpha_3 \leq \frac{\pi}{2}$. The neutral Higgs bosons h_i satisfy the mass relation $m_{h_1} \leq m_{h_2} \leq m_{h_3}$ [57–60]. In the CP conserving case η_1 and η_2 are CP even and mixed in a 2×2 matrix while η_3 is CP odd without mixing with η_1 and η_2. However, due to the CP-symmetry breaking in the general case, the neutral Higgs bosons $h_{1,2,3}$ do not have well-defined CP states. The most general structure for the Yukawa couplings among fermions and the scalar is

$$\begin{aligned} \mathcal{L}_{\mathrm{Yukawa}} = \sum_{i,j=1}^{3}\sum_{a=1}^{2} & (\bar{q}_{Li}^0 Y_{aij}^{0u}\tilde{\Phi}_a u_{Rj}^0 + \bar{q}_{Li}^0 Y_{aij}^{0d}\Phi_a d_{Rj}^0 \\ & + \bar{l}_{Li}^0 Y_{aij}^{0l}\Phi_a e_{Rj}^0 + \mathrm{H.c.}), \end{aligned} \tag{6}$$

where $Y_a^{u,d,l}$ are the 3×3 Yukawa matrices. q_L and l_L denote the left-handed fermion doublets under $SU(2)_L$, while u_R, d_R, l_R correspond to the right-handed singlets. The zero superscript in fermion fields stands for the interaction basis. After getting a correct spontaneous symmetry breaking by the VEV using Eqs. (2) and (3), the mass matrices become

$$M^{u,d,l} = \sum_{a=1}^{2} \frac{v_a}{\sqrt{2}} Y_a^{u,d,l}, \qquad (7)$$

where $Y_a^f = V_L^f Y_a^{0f} (V_R^f)^\dagger$, for $f = u,\ d,\ l$. The $V_{L,R}^f$ matrices are used to diagonalize the fermion mass matrices and relate the physical and interaction states. Note that in model III the diagonalization of mass matrices does not imply the diagonalization of the Yukawa matrices, as happens in type-I or II THDM. An important consequence of nondiagonal Yukawa matrices in physical states is the presence of FCNSI between neutral Higgs bosons and fermions.

The focus is on the up-type quark Yukawa interactions that contain the Feynman rules for the rare top decay. Replacing Eqs. (4) and (7) in the Yukawa Lagrangian of Eq. (6), the interactions between neutral Higgs bosons and fermions can be written as interactions of the THDM with CP conserving (type I or II) plus additional contributions, which arise from any of the $Y_{1,2}$ Yukawa matrices. The relation among the mass matrix M^F and the Yukawa matrices $Y_{1,2}^F$, for $F = u,\ d,\ l$, is used to write the Yukawa Lagrangian, Eq. (6), as a function only of one Yukawa matrix, Y_1^F or Y_2^F. We choose to write the interactions as a function of the Yukawa matrix Y_2, that is, $Y_1^F = \frac{\sqrt{2}}{v_1} M^F - \frac{v_2}{v_1} Y_2^F$ is replaced in Eq. (6). From now on, in order to simplify the notation, the subscript 2 in the Yukawa couplings is omitted. The interactions between quarks and Higgs bosons in the mass eigenstates are explicitly written as

$$\begin{aligned}
\mathcal{L} = {} & \frac{1}{v\cos\beta}\sum_{ijk}\bar{u}_i M_{ij}^u (A_k P_L + A_k^* P_R) u_j h_k + \frac{1}{v\cos\beta}\sum_{ijk}\bar{d}_j M_{ij}^d (A_k^* P_L + A_k P_R) d_j h_k \\
& + \frac{1}{\cos\beta}\sum_{ijk}\bar{u}_i Y_{ij}^u (B_k P_L + B_k^* P_R) u_j h_k + \frac{1}{\cos\beta}\sum_{ijk}\bar{d}_i Y_{ij}^d (B_k^* P_L + B_k P_R) d_j h_k \\
& + \Bigg[\frac{\sqrt{2}}{\cos\beta}\sum_{ij}\bar{u}_i ((KY^d)_{ij} P_R - (Y^u K)_{ij} P_L) d_j H^+ \\
& + \frac{\sqrt{2}}{v}\tan\beta\sum_{ij}\bar{u}_i (-(KM^d)_{ij} P_R + (M^u K)_{ij} P_L) d_j H^+ \\
& + \frac{\sqrt{2}}{v}\bar{u}_i\sum_{ij}((M^d K)_{ij} P_R - (M^u K)_{ij} P_L) d_j G_W^+ + \text{H.c.}\Bigg], \qquad (8)
\end{aligned}$$

where we define

$$\begin{aligned}
A_k &= R_{k1} - iR_{k3}\sin\beta, \\
B_k &= R_{k2}\cos\beta - R_{k1}\sin\beta + iR_{k3}. \qquad (9)
\end{aligned}$$

The fermion spinors are denoted as $(u_1, u_2, u_3) = (u, c, t)$, where the indexes $i,\ j = 1,\ 2,\ 3$ denote the family generations in Eq. (8), while $k = 1,\ 2,\ 3$ is used for the neutral Higgs bosons and $P_{R,L} = \frac{1}{2}(1 \pm \gamma_5)$. Note that a CP conserving case is obtained only if two neutral Higgs bosons are mixed with well-defined CP states, for instance, $\alpha_2 = \alpha_3 = 0$ is the usual limit. The CP conserving case has been studied by Kim *et al.* in [61] in the alignment limit, where $\sin(\alpha_1 - \beta) = 1$ holds.

III. RARE TOP DECAY $t \to c\gamma$

The expression for the $t \to c\gamma$ decay amplitude is a magnetic transition written as

$$\mathcal{M} = \bar{u}(p')[F_1\sigma_{\mu\nu} + F_2\sigma_{\mu\nu}\gamma_5] q^\nu u(p)\epsilon^\mu(q), \qquad (10)$$

where $p' = p - q$, $\epsilon^\mu(q)$ is the photon polarization; when the photon is on shell, $q^2 = 0$, and $\epsilon^\mu(q) q_\mu = 0$. The invariant amplitudes $F_{1,2}$ are obtained in terms of the model parameters as shown in Eq. (11). Equation (10) corresponds to a five-dimensional operator, and then the on-shell $t \to c\gamma$ amplitude must be represented by a set of loop diagrams. Fig. 1 shows the dominant contributions for the rare top decay $t \to c\gamma$ at one loop coming from neutral and charged Higgs bosons. The charged contributions, see Figs. 1(b) and 1(c), are suppressed by the bottom-quark mass compared to the top-quark mass in the neutral Higgs contribution. In order to study the effects of FCNSI we analyze only the dominant contribution; see Fig. 1(a). In order to obtain the partial width of the $t \to c\gamma$ decay in model III we apply the method previously used in [32]. Integrating over the internal momentum, the partial width is

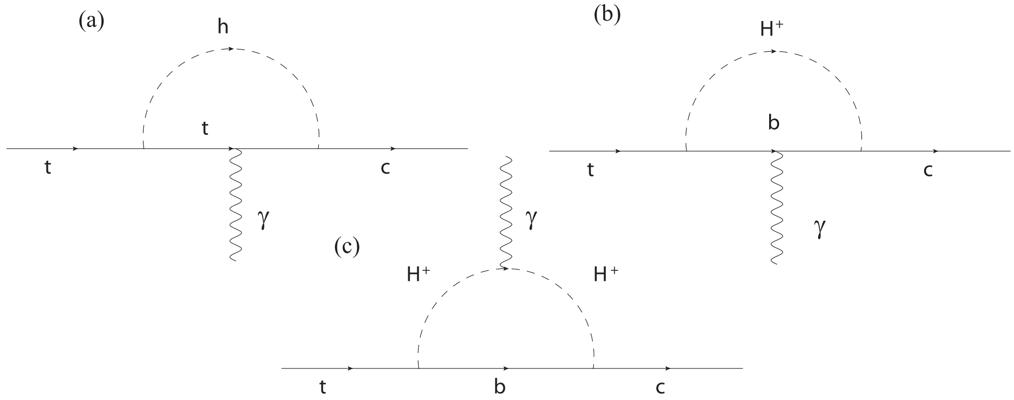

FIG. 1. One loop Feynman diagram with a Higgs boson in the internal line, (a) flavor changing neutral scalar contribution, (b) and (c) charged contributions.

$$\Gamma(t \to c\gamma) = \frac{\alpha G_F m_t^3}{192\pi^4 \cos^4\beta} |Y_{ct}^u|^2 \times \sum_k |f_1(\hat{m}_k) A_k^* B_k + f_2(\hat{m}_k) A_k B_k^*|^2, \quad (11)$$

where $G_F^{-1} = \sqrt{2}v^2$, $v = 246$ GeV, $\alpha \approx 1/128$ at electroweak scale and the functions $f_{1,2}$ are defined as

$$f_1(\hat{m}_k) = \int_0^1 dx \int_0^{1-x} dy \frac{x(x+y-1)}{x^2 + xy - (2-\hat{m}_k^2)x + 1}, \quad (12)$$

$$f_2(\hat{m}_k) = \int_0^1 dx \int_0^{1-x} dy \frac{(x-1)}{x^2 + xy - (2-\hat{m}_k^2)x + 1}, \quad (13)$$

with $\hat{m}_i = \frac{m_{h_i}}{m_t}$ for $i = 1, 2, 3$. The branching ratio can be approximated as

$$\mathrm{Br}(t \to c\gamma) \approx \frac{\Gamma(t \to c\gamma)}{\Gamma_{\rm top}}, \quad (14)$$

where $\Gamma_{\rm top}$ at next-to-leading order is given by [33]

$$\Gamma_{\rm top} = \frac{G_f m_t^3}{8\pi\sqrt{2}} \left(1 - \frac{M_W^2}{m_t^2}\right)^2 \left(1 - 2\frac{M_W^2}{m_t^2}\right) \times \left[1 - \frac{2\alpha_s}{3\pi}\left(\frac{2\pi^2}{3} - \frac{5}{2}\right)\right]. \quad (15)$$

IV. CONSTRAINTS ON RARE TOP DECAY PARAMETERS

We note that Eq. (11) contains free parameters of the THDM, such as the masses of the neutral Higgs bosons, the mixing angles α_i, β, and Yukawa couplings. In order to set allowed values for free parameters we first review the possible constraints that the $b \to s\gamma$ decay can impose on the Y_{tc} coupling. Following Refs. [62–66], the branching ratio of the $b \to s\gamma$ decay is a function of the Wilson coefficients and it can be written as

$$\mathrm{Br}(B \to X_s\gamma) \approx a + a_{77}\delta C_7^2 + a_{88}\delta C_8^2 + \mathrm{Re}(a_7 \delta C_7) + \mathrm{Re}(a_8 \delta C_8) + \mathrm{Re}(a_{78} \delta C_7 \delta C_8^*), \quad (16)$$

with $a \approx 3.0 \times 10^{-4}$, $a_{77} \approx 4.7 \times 10^{-4}$, $a_{88} \approx 0.8 \times 10^{-4}$, $a_7 \approx (-7.2 + 0.6i) \times 10^{-4}$, $a_8 \approx (-2.2 - 0.6i) \times 10^{-4}$, and $a_{78} \approx (2.5\text{–}0.9i) \times 10^{-4}$. The main contributions due to Wilson coefficients, beyond the W-boson contribution, are given by charged Higgs and FC Yukawa couplings, $\delta C_{7,8} = C^{H^\pm_{7,8}} + C^{H,FC}_{7,8}$. The charged-Higgs contribution is

$$C^{H^\pm_{7,8}} = \frac{1}{3\tan^2\beta} f_{7,8}^{(1)}(y_t) + f_{7,8}^{(2)}(y_t), \quad (17)$$

while the FC contribution is

$$C_{7,8}^{H,FC} = \frac{2M_W}{g m_t K_{ts} \cos\beta} (Y^u K)_{ts} f_{7,8}^{(2)}(y_t) + \frac{2M_W}{g m_b K_{tb} \cos\beta} (K Y^d)_{tb} f_{7,8}^{(2)}(y_t) \quad (18)$$

with $y_t = m_t^2/M_H^2$ and the explicit relations $f_{7,8}^{(1),(2)}(x)$ can be found in Refs. [62–66]. Using the hierarchy of the Kobayashi-Maskawa matrix (K) we have the following approximations: $(Y^u K)_{ts} \approx Y_{tc} K_{cs}$ and $(K Y^d)_{tb} \approx K_{tb} Y_{bb}$. In order to have a bound to the Y_{tc} FC Yukawa coefficient, it was considered that $(K Y^d)_{tb}$ gives the most important contribution. The limits on the $B \to X_s\gamma$ decay come from *BABAR*, Belle, and CLEO [67–71]. The current world average for $E > 1.6$ GeV, given by HFAG [72], is

$$\mathrm{Br}(B \to X_s\gamma) = (3.43 \pm 0.21 \pm 0.07) \times 10^{-4}. \quad (19)$$

This result provides an important constraint on the (Y_{tc}, Y_{bb}) space, Fig. 2, with $m_{H^\pm} = 500$ GeV and $0 < \tan\beta < 20$.

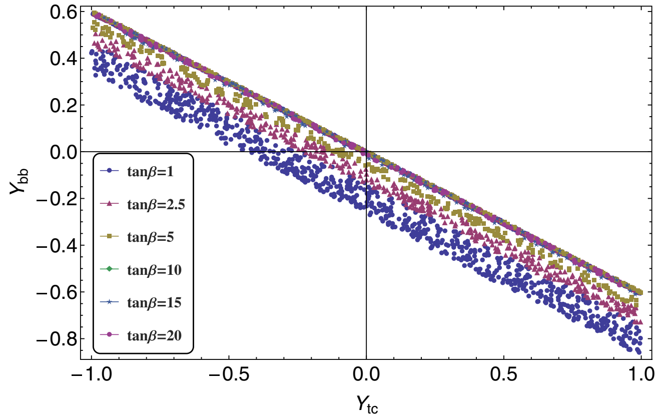

FIG. 2. Allowed values for the Yukawa couplings in a scatter plot with points compatible with the experimental value of the $BR(B \to X_s\gamma)$ and $300 \text{ GeV} \le m_H^{\pm} \le 600 \text{ GeV}$ and $\tan\beta = 1$, 2.5, 5, 10, and 15.

The second constraint considered is based on the branching ratio of the SM Higgs boson decay to bottom-quark pairs, which has a reported value of $Br(H \to b\bar{b}) = 5.77 \times 10^{-1}{}^{+3.2\%}_{-3.3\%}$ [33]. The width decay in the THDM for $h_1 \to b\bar{b}$ is given by

$$\Gamma_{h_1 \to b\bar{b}} = \frac{N_c m_{h1}}{8\pi}\left(1 - 4\frac{m_b^2}{m_{h1}^2}\right)^{\frac{1}{2}}\left[C^2\left(1 - 4\frac{m_b^2}{m_{h1}^2}\right) + D^2\right], \tag{20}$$

where

$$C = \frac{m_b}{v\cos\beta}R_{11} + \frac{Y_{bb}}{\cos\beta}(R_{12}\cos\beta - R_{11}\sin\beta) \tag{21}$$

and

$$D = -\frac{m_b}{v\cot\beta}R_{13} + \frac{Y_{bb}}{\cos\beta}R_{13}. \tag{22}$$

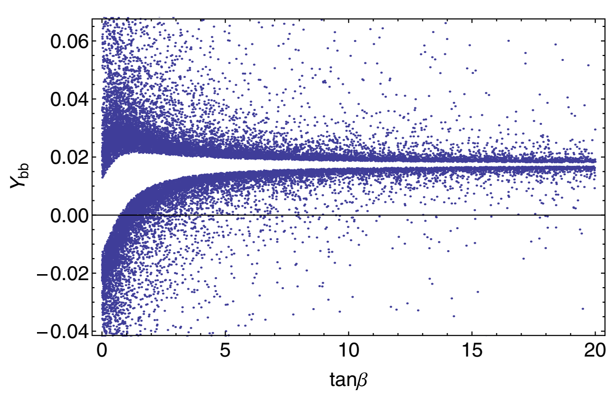

FIG. 3. Allowed values for the Yukawa couplings Y_{bb} in a scatter plot with compatible points with the experimental value of the $BR(H \to b\bar{b})$ and $\pi/2 \le \alpha_{1,2} \le \pi/2$.

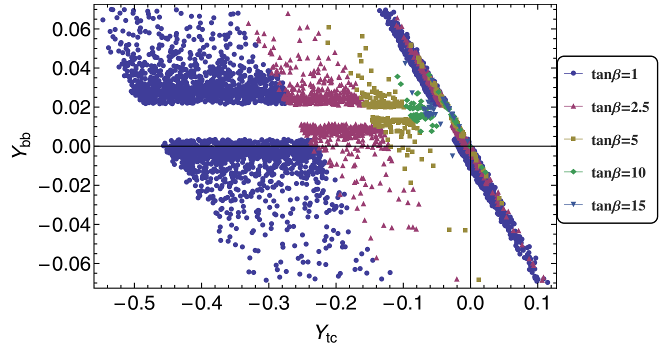

FIG. 4. Allowed values for the Yukawa couplings in a scatter plot with compatible points with the experimental values for $Br(H \to b\bar{b})$ and $Br(B \to X_s\gamma)$.

Note that the matrix elements R_{11}, R_{12}, and R_{13} are independent of the mixing parameter α_3. Figure 3 shows the behavior of Y_{bb} as a function of $\tan\beta$ for random values of $\alpha_{1,2}$. After that previous constraints are imposed; the allowed values for Yukawa couplings are $-0.02 \le Y_{bb} \le 0.06$ and $-0.5 \le Y_{tc} \le 0.02$ for $1 \le \tan\beta \le 15$; see Fig. 4. The nondiagonal elements of the Yukawa matrix responsible for the FCNSI, shown in Eq. (8), must be suppressed [73].

V. RESULTS AND DISCUSSION

Focusing on the rest of the parameters, note that the masses of the h_i neutral Higgs bosons are set so that the mass of the lightest Higgs boson h_1 is equal to the mass value of the observed scalar reported by ATLAS and CMS, $m_{h_1} \approx 126$ GeV [1,2]. Contributions to Eq. (14) from h_2 and h_3 are negligible for masses $m_{h_2}, m_{h_3} > 600$ GeV.

Also, note that the contribution from h_1 is independent of the mixing parameter α_3; see the first row in matrix Eq. (5). Therefore, the set of free parameters considered in the partial width Eq. (11) is reduced only to the mixing angles $\{\alpha_1, \alpha_2, \beta\}$. Figures 5 and 6 show the allowed values for mixing parameters α_1 and α_2 when the current limit for

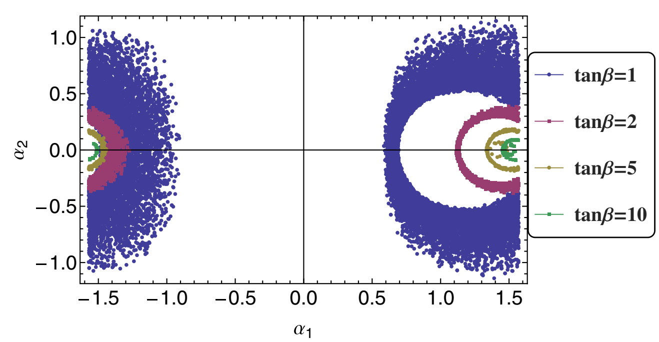

FIG. 5. Allowed values for the mixing parameter $\alpha_{1,2}$ in a scatter plot with points compatible with the experimental values for $Br(H \to b\bar{b})$, $Br(B \to X_s\gamma)$, and $Br(t \to c\gamma) < 5.9 \times 10^{-3}$, for fixed values of $\tan\beta = 1$, 2.5, 5, 10, and $Y_{tc} = 0.01$.

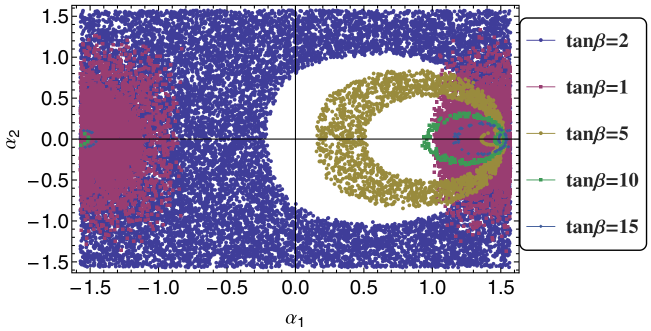

FIG. 6. Allowed values for the mixing parameters $\alpha_{1,2}$ in a scatter plot with points compatible with the experimental values for $\mathrm{Br}(H \to b\bar{b})$, $\mathrm{Br}(B \to X_s\gamma)$, and $\mathrm{Br}(t \to c\gamma) < 5.9 \times 10^{-3}$, for fixed values of $\tan\beta = 1, 2.5, 5, 10$, and $Y_{tc} = -0.04$.

the $\mathrm{Br}(t \to c\gamma) < 5.9 \times 10^{-3}$ is considered [33]. Based in Fig. 4 the Yukawa coupling Y_{tc} was fixed with the two representative values $Y_{tc} = -0.04, 0.01$.

In order to analyze the $\mathrm{Br}(t \to c\gamma)$ we consider the allowed regions for the mixing parameters α_1 and α_2 previously fixed in [74,75]. The following regions can be obtained for α_1 and α_2 from $0.5 \le R_{\gamma\gamma} \le 2$ with $m_{H^\pm} = 300$ GeV and $\tan\beta = 2.5$ [57–60]:

$$R_1 = \{-1.39 \le \alpha_1 \le -1.2 \quad \text{and} \quad -0.13 \le \alpha_2 \le 0\}, \tag{23}$$

and

$$R_2 = \{1.16 \le \alpha_1 \le 1.5 \quad \text{and} \quad -0.48 \le \alpha_2 \le -0.1\}. \tag{24}$$

The ratio $R_{\gamma\gamma}$ given by

$$R_{\gamma\gamma} = \frac{\sigma(gg \to h_1)\mathrm{Br}(h_1 \to \gamma\gamma)}{\sigma(gg \to h_{\mathrm{SM}})\mathrm{Br}(h_{\mathrm{SM}} \to \gamma\gamma)} \tag{25}$$

allows us to compare the prediction of the THDM with the SM prediction for the Higgs boson diphoton decay. Figure 7 shows $\mathrm{Br}(t \to c\gamma)$ as a function of α_1 and α_2 in the allowed regions R_1 and R_2 with $\tan\beta = 2.5$. The $\mathrm{Br}(t \to c\gamma)$ can be enhanced up to 10^{-6} in the regions $R_{1,2}$. The limits obtained in model III are less restrictive than those obtained in type-I and type-II THDM, which are of the order of 10^{-8} [27,29]. In 2021, the LHC is expected to reach an integrated luminosity of the order of 300 fb^{-1} [76]. Experiments in LHC Run 3 with this amount of data could find evidence of new physics beyond the SM, in particular, processes with the FCNC. The expected number of events can be naively estimated with the following approximation,

$$N \approx \sigma(p\bar{p} \to t\bar{t})\mathrm{Br}(\bar{t} \to \bar{b}W)\mathrm{Br}(t \to c\gamma)\mathcal{L}_{\mathrm{int}}, \tag{26}$$

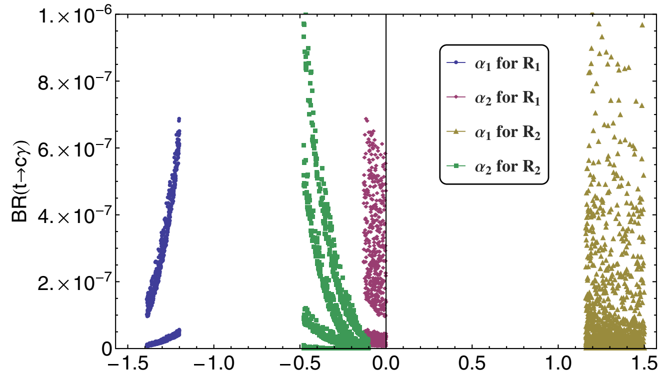

FIG. 7. The model III branching ratio for $t \to c\gamma$ as a function of α_1–α_2 in regions R_1 and R_2.

where $\sigma(p\bar{p} \to t\bar{t}) \approx 176$ pb [33], $\mathcal{L}_{\mathrm{int}}$ is the integrated luminosity ∼300 fb^{-1}, $\mathrm{Br}(\bar{t} \to \bar{b}W) \approx 1$, and $\mathrm{Br}(t \to c\gamma)$ is the obtained result in model III, Eq. (14). Because of trigger and selection cuts only a fraction of the produced events is detected by the experiments. An efficiency of 2.4% is achieved by CMS from simulation of $tc\gamma$ signal events taking into account all selection criteria [77]. Therefore, a more realistic estimation of the effective number of events has to be written as $N_{\mathrm{Eff}} \approx 0.2 \times N$.

The limit that is expected to be reached in future experiments is $\mathrm{Br}(t \to c\gamma) \sim 10^{-5}$ [76]. If we consider this expected limit as $\mathrm{Br}(t \to c\gamma) \sim (1\text{–}10)^{-5}$ with $N_{\mathrm{Eff}} \ge 1$ and impose the restrictions discussed in the previous section, then the N_{Eff} can be estimated for fixed values of $\tan\beta$. Figure 8 shows N_{Eff} as a function of Y_{tc}. The mixing parameters $\alpha_{1,2}$ are also bounded by the same constraints and the allowed values of the $\alpha_{1,2}$ are shown in Fig. 9 for fixed $\tan\beta$. The numerical values for $\tan\beta$ are fixed by the representative values $\tan\beta = 1.56, 2.5, 5, 10, 15$; however, the N_{Eff} as a function of $\tan\beta$ with the above restrictions is shown in Fig. 10. We find that the effective number of events is greater than 1, $N_{\mathrm{Eff}} \ge 1$, from $\tan\beta \ge 1.56$. The

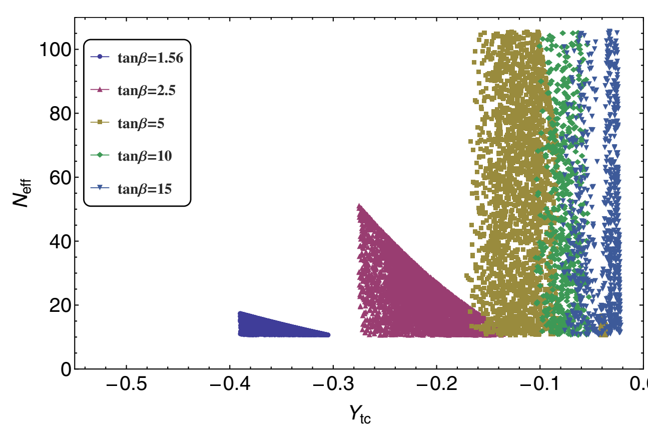

FIG. 8. Effective number of events for $t \to c\gamma$ as a function of Y_{tc} for $\tan\beta = 1.56, 2.5, 5, 10, 15$ expected in LHC Run 3.

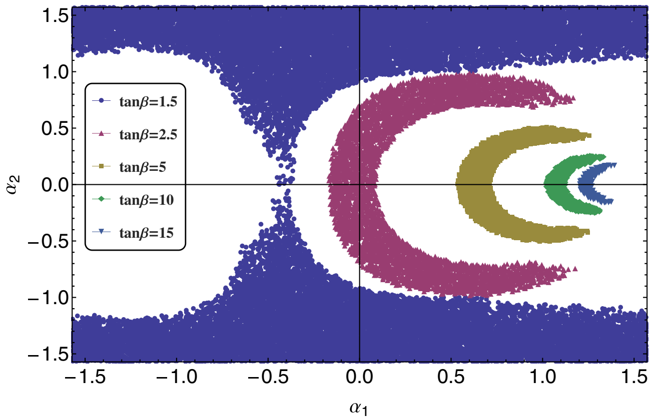

FIG. 9. Allowed regions for α_1 and α_2 when $\mathrm{Br}(t \to c\gamma) \sim 10^{-5}$ is assumed with $-0.385 \leq Y_{tc} \leq -0.307$ for $\tan\beta = 1.56$, $-0.267 \leq Y_{tc} \leq -0.135$ for $\tan\beta = 2.5$, $-0.173 \leq Y_{tc} \leq -0.035$ for $\tan\beta = 5$, $-0.105 \leq Y_t c \leq -0.02$ for $\tan\beta = 10$ and $-0.08 \leq Y_{tc} \leq -0.02$ for $\tan\beta = 15$. In LHC Run 3.

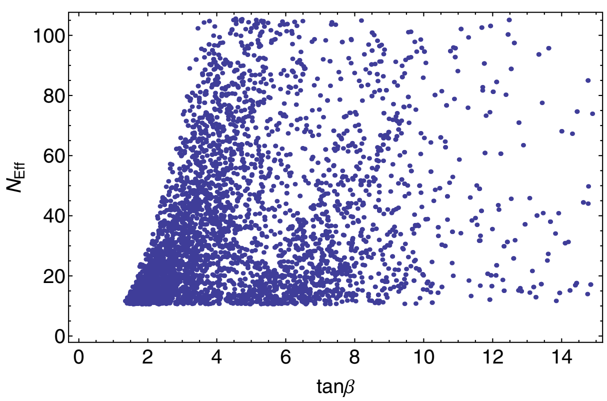

FIG. 10. Effective number of events for $t \to c\gamma$ as a function of $\tan\beta$ expected in LHC Run 3.

expression for the rare top decay $t \to c\gamma$ was calculated at one loop due to the FCNSI in an extended model with two scalar doublets. The SM predicted value for the $\mathrm{Br}(t \to c\gamma)$ is extremely suppressed from LHC sensitivity, while in the considered type-III THDM with mixing in the neutral scalars the same branching ratio has been increased, making it possible to test rare decays in future experiments. In this work we have studied a theoretical framework where $\mathrm{Br}(t \to c\gamma) \sim 10^{-5}$ can be viable for specific values of mixing parameters. If the $t \to c\gamma$ decay is observed in the LHC, it will provide important evidence of physics beyond the SM. With the allowed regions for the α_1, α_2 and

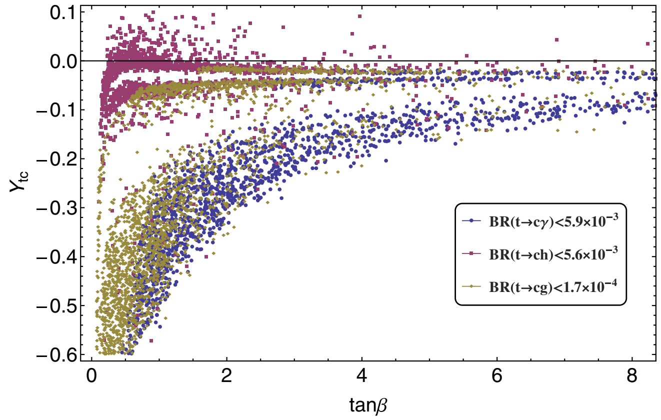

FIG. 11. Y_{tc} coupling as a function of $\tan\beta$ for different FCNC top quark decays. Current experimental limits are taken into account.

$\tan\beta \simeq 2.5$, model III predicts $\mathrm{Br}(t \to c\gamma) \sim 10^{-6}$. Model III, with an integrated luminosity of 300 fb^{-1}, predicts up to $N_{\mathrm{Eff}} \approx 100$ events for the $t \to c\gamma$ decay with α_1, α_2 and $\tan\beta$ given in the previous section. There are also other relevant decays, for instance, it is straightforward to obtain $\mathrm{Br}(t \to cg)$ in the framework of model III; in this case $\mathrm{Br}(t \to cg)$ can be found from Eq. (11) by replacing the factor α by $3\alpha_s$, where α_s is the strong coupling constant, and within the framework of our analysis the parameter space can be phenomenologically restricted from the experimental limits of this decay; we find the approximate region $-0.6 < Y_{tc} < 0$ for $0 < \tan\beta < 8$. On the other hand one can also calculate the $\mathrm{Br}(t \to ch)$ at tree level in model III. From the current experimental limit on $\mathrm{Br}(t \to ch)$ we obtain that the flavor changing Yukawa coupling Y_{tc} can be restricted to $-0.2 < Y_{tc} < 0.1$ for $0 < \tan\beta < 8$. The phenomenological restrictions for Y_{tc} as a function of $\tan\beta$ for the three FCNC top-quark decays are shown in Fig 11.

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[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] J. F. Gunion, H. E. Haber, G. L. Kane, and S. Dawson, *The Higgs Hunters Guide* (Westview Press, Boulder, 2000).
[4] H. E. Haber, G. L. Kane, and T. Sterling, Nucl. Phys. **B161**, 493 (1979).
[5] L. J. Hall and M. B. Wise, Nucl. Phys. **B187**, 397 (1981).
[6] J. F. Donoghue and L. F. Li, Phys. Rev. D **19**, 945 (1979).
[7] D. Atwood, L. Reina, and A. Soni, Phys. Rev. D **55**, 3156 (1997).
[8] A. Barroso, P. M. Ferreira, R. Santos, and J. P. Silva, Phys. Rev. D **86**, 015022 (2012).
[9] L. Lavoura and J. P. Silva, Phys. Rev. D **50**, 4619 (1994).
[10] D. Fontes, J. C. Romo, R. Santos, and J. P. Silva, J. High Energy Phys. 06 (2015) 060.
[11] ATLAS Collaboration, Report No. ATLAS-CONF-2013-013.
[12] V. Khachatryan *et al.* (CMS Collaboration), Phys. Rev. D **92**, 012004 (2015).
[13] T. P. Cheng and M. Sher, Phys. Rev. D **35**, 3484 (1987).
[14] A. Crivellin, A. Kokulu, and C. Greub, Phys. Rev. D **87**, 094031 (2013).
[15] H. Fritzsch, Phys. Lett. B **70B**, 436 (1977).
[16] H. Fritzsch, Phys. Lett. B **73B**, 317 (1978).
[17] H. Fritzsch, Nucl. Phys. **B155**, 189 (1979).
[18] J. Shu and Y. Zhang, Phys. Rev. Lett. **111**, 091801 (2013).
[19] D. E. Morrissey and M. J. Ramsey-Musolf, New J. Phys. **14**, 125003 (2012).
[20] S. Inoue, M. J. Ramsey-Musolf, and Y. Zhang, Phys. Rev. D **89**, 115023 (2014).
[21] A. E. Carcamo Hernandez, R. Martinez, and J. A. Rodriguez, Eur. Phys. J. C **50**, 935 (2007).
[22] G. Cvetic, S. S. Hwang, and C. S. Kim, Int. J. Mod. Phys. A **14**, 769 (1999).
[23] G. Cvetic, C. S. Kim, and S. S. Hwang, Phys. Rev. D **58**, 116003 (1998).
[24] J. A. Aguilar-Saavedra, Acta Phys. Pol. B **35**, 2695 (2004).
[25] B. Grzadkowski, J. F. Gunion, and P. Krawczyk, Phys. Lett. B **268**, 106 (1991).
[26] B. Mele, arXiv:hep-ph/0003064.
[27] B. Mele, S. Petrarca, and A. Soddu, Phys. Lett. B **435**, 401 (1998).
[28] E. Gabrielli and B. Mele, Phys. Rev. D **83**, 073009 (2011).
[29] G. Eilam, J. L. Hewett, and A. Soni, Phys. Rev. D **44**, 1473 (1991); **59**, 039901(E) (1998).
[30] J. A. Aguilar-Saavedra and B. M. Nobre, Phys. Lett. B **553**, 251 (2003).
[31] F. Larios, R. Martinez, and M. A. Perez, Int. J. Mod. Phys. A **21**, 3473 (2006).
[32] J. L. Diaz-Cruz, R. Martinez, M. A. Perez, and A. Rosado, Phys. Rev. D **41**, 891 (1990).
[33] K. A. Olive *et al.* (Particle Data Group Collaboration), Chin. Phys. C **38**, 090001 (2014).
[34] The ATLAS Collaboration, Report No. ATLAS-CONF-2013-063.
[35] CMS Collaboration, Report No. CMS-PAS-HIG-13-034.
[36] M. E. Luke and M. J. Savage, Phys. Lett. B **307**, 387 (1993).
[37] D. Atwood, L. Reina, and A. Soni, Phys. Rev. D **53**, 1199 (1996).
[38] D. Atwood, L. Reina, and A. Soni, Phys. Rev. D **54**, 3296 (1996).
[39] D. Atwood, L. Reina, and A. Soni, Phys. Rev. Lett. **75**, 3800 (1995).
[40] A. Arhrib, Phys. Rev. D **72**, 075016 (2005).
[41] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher, and J. P. Silva, Phys. Rep. **516**, 1 (2012).
[42] J. A. Aguilar-Saavedra, Phys. Rev. D **67**, 035003 (2003); **69**, 099901(E) (2004).
[43] X. F. Han, L. Wang, and J. M. Yang, Phys. Rev. D **80**, 015018 (2009).
[44] H. Hong-Sheng, Phys. Rev. D **75**, 094010 (2007).
[45] G. A. Gonzalez-Sprinberg, R. Martinez, and J. A. Rodriguez, Eur. Phys. J. C **51**, 919 (2007).
[46] J. J. Cao, G. Eilam, M. Frank, K. Hikasa, G. L. Liu, I. Turan, and J. M. Yang, Phys. Rev. D **75**, 075021 (2007).
[47] T. Han, K. i. Hikasa, J. M. Yang, and X. m. Zhang, Phys. Rev. D **70**, 055001 (2004).
[48] R. A. Diaz, R. Martinez, and J. Alexis Rodriguez, arXiv:hep-ph/0103307.
[49] R. Gaitan-Lozano, R. Martinez, and J. H. M. de Oca, arXiv:1407.3318.
[50] A. Dedes, M. Paraskevas, J. Rosiek, K. Suxho, and K. Tamvakis, J. High Energy Phys. 11 (2014) 137.
[51] E. Accomando *et al.*, arXiv:hep-ph/0608079.
[52] D. Bardhan, G. Bhattacharyya, D. Ghosh, M. Patra, and S. Raychaudhuri, Phys. Rev. D **94**, 015026 (2016).
[53] P. Niezurawski, A. F. Zarnecki, and M. Krawczyk, J. High Energy Phys. 02 (2005) 041.
[54] H. E. Haber and R. Hempfling, Phys. Rev. D **48**, 4280 (1993).
[55] I. F. Ginzburg and M. Krawczyk, Phys. Rev. D **72**, 115013 (2005).
[56] A. W. El Kaffas, W. Khater, O. M. Ogreid, and P. Osland, Nucl. Phys. **B775**, 45 (2007).
[57] L. Basso, A. Lipniacka, F. Mahmoudi, S. Moretti, P. Osland, G. M. Pruna, and M. Purmohammadi, J. High Energy Phys. 11 (2012) 011.
[58] A. Arhrib, E. Christova, H. Eberl, and E. Ginina, J. High Energy Phys. 04 (2011) 089.
[59] M. Krawczyk, D. Sokolowska, P. Swaczyna, and B. Swiezewska, J. High Energy Phys. 09 (2013) 055.
[60] C. Y. Chen, S. Dawson, and Y. Zhang, J. High Energy Phys. 06 (2015) 056.
[61] C. S. Kim, Y. W. Yoon, and X. B. Yuan, J. High Energy Phys. 12 (2015) 038.
[62] G. Degrassi, P. Gambino, and G. F. Giudice, J. High Energy Phys. 12 (2000) 009.
[63] M. Misiak *et al.*, Phys. Rev. Lett. **98**, 022002 (2007).
[64] E. Lunghi and J. Matias, J. High Energy Phys. 04 (2007) 058.
[65] M. E. Gomez, T. Ibrahim, P. Nath, and S. Skadhauge, Phys. Rev. D **74**, 015015 (2006).
[66] G. Barenboim, C. Bosch, M. L. Lpez-Ibaez, and O. Vives, J. High Energy Phys. 11 (2013) 051.

[67] S. Chen *et al.* (CLEO Collaboration), Phys. Rev. Lett. **87**, 251807 (2001).
[68] K. Abe *et al.* (Belle Collaboration), Phys. Lett. B **511**, 151 (2001).
[69] J. P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **86**, 052012 (2012).
[70] J. P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **86**, 112008 (2012).
[71] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **77**, 051103 (2008).
[72] Y. Amhis *et al.* (Heavy Flavor Averaging Group Collaboration), arXiv:1412.7515.
[73] S. L. Glashow and S. Weinberg, Phys. Rev. D **15**, 1958 (1977).
[74] R. Gaitán, R. Martinez, J. H. Montes de Oca, and S. R. Romo, Eur. Phys. J. C **74**, 2788 (2014).
[75] R. Gaitan, E. A. Garces, J. H. M. de Oca, and R. Martinez, Phys. Rev. D **92**, 094025 (2015).
[76] ATLAS Collaboration, arXiv:1307.7292.
[77] V. Khachatryan *et al.* (CMS Collaboration), J. High Energy Phys. 04 (2016) 035.