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$\bar{B} \to X_s\gamma$ constraints on the top quark anomalous $t \to c\gamma$ coupling

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Observation of the top quark flavor changing neutral process $t \to c + \gamma$ at the LHC would be the signal of physics beyond the standard model. If anomalous $t \to c\gamma$ coupling exists, it will affect the precisely measured $\mathcal{B}(\bar{B} \to X_s\gamma)$. In this paper, we study the effects of a dimension 5 anomalous $tc\gamma$ operator in $\bar{B} \to X_s\gamma$ decay to derive constraints on its possible strength. It is found that, for real anomalous $t \to c\gamma$ coupling κ^{γ}_{tcR}, the constraints correspond to the upper bounds $\mathcal{B}(t \to c + \gamma) < 6.54 \times 10^{-5}$ (for $\kappa^{\gamma}_{tcR} > 0$) and $\mathcal{B}(t \to c + \gamma) < 8.52 \times 10^{-5}$ (for $\kappa^{\gamma}_{tcR} < 0$), respectively, which are about the same order as the 5σ discovery potential of ATLAS (9.4×10^{-5}) and slightly lower than that of CMS (4.1×10^{-4}) with 10 fb^{-1} integrated luminosity operating at $\sqrt{s} = 14$ TeV.

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I. INTRODUCTION

In the standard model (SM), a top quark lifetime is dominated by the $t \to bW^+$ process, and its flavor changing neutral current (FCNC) processes $t \to qV$ ($q = u, c$; $V = \gamma, Z, g$) are extremely suppressed by the Glashow-Iliopoulos-Maiani mechanism. It is known that the SM predicts a very tiny top FCNC branching ratio $\mathcal{B}(t \to qV)$, less than $\mathcal{O}(10^{-10})$ [1], which would be inaccessible at the CERN Large Hadron Collider (LHC). In the literature [2,3], however, a number of interesting questions have been intrigued by the large top quark mass which is close to the scale of electroweak symmetry breaking. For example, one may raise the question whether new physics (NP) beyond the SM could manifest itself in nonstandard couplings of the top quark which would show up as anomalies in the top quark productions and decays.

At present, the direct constraints on $\mathcal{B}(t \to qV)$ are still very weak. For its radiative decay, the available experimental bounds are $\mathcal{B}(t \to u\gamma) < 0.75\%$ from ZEUS [4] and $\mathcal{B}(t \to q\gamma) < 3.2\%$ from CDF [5] at 95% C.L., respectively. These constraints will be improved greatly by the large top quark sample to be available at the LHC, which is expected to produce 8×10^6 top quark pairs and another few million single top quarks per year at low luminosity (10 fb^{-1}/yr). Both ATLAS [6] and CMS [7] have analyses ready for hunting out top quark FCNC processes as powerful probes for NP. With 10 fb^{-1} data, it is expected that both ATLAS and CMS could observe $t \to q\gamma$ decays if their branching ratios are enhanced to $\mathcal{O}(10^{-4})$ by anomalous top quark couplings [6,7]. However, if the top quark anomalous couplings are present, they will affect some precisely measured qualities with virtual top quark contributions. Inversely, these qualities can also restrict the possible number of top quark FCNC decay signals at the LHC. The precisely measured inclusive decay $B \to X_s\gamma$ is one of the well-known sensitive probes for extensions of the SM, especially the NP which alter the strength of FCNCs [8]. Thus, when performing the study of the possible strength of $t \to c\gamma$ decays at the LHC, one should take into account the constraints from $B \to X_s\gamma$ [9,10].

In this paper, we will study the contribution of anomalous $t\gamma c$ operators to the $\bar{B} \to X_s\gamma$ branching ratio and derive constraints on its strength. In the next section, after a brief discussion of a set of model-independent dimension 5 effective operators relevant to $t \to c\gamma$ decay, we calculate the effects of operator $\bar{c}_L\sigma^{\mu\nu}t_RF_{\mu\nu}$ in $B \to X_s\gamma$ decay, which result in a modification to $C_{7\gamma}$. In Sec. III we present our numerical results of the constraints on its strength and the corresponding upper limits on the branching ratio of $t \to c\gamma$ decays. Finally, conclusions are made in Sec. IV. Calculation details are presented in Appendix A, and input parameters are collected in Appendix B.

II. TOP QUARK ANOMALOUS COUPLINGS AND THEIR EFFECTS IN $\bar{B} \to X_s\gamma$ DECAY

Without resorting to the detailed flavor structure of a specific NP model, the Lagrangian describing the top quark anomalous couplings can be written in a model-independent way with dimension 5 operators [11]

$$\begin{aligned}\mathcal{L}_5 = &-g_s \sum_{q=u,c,t} \frac{\kappa^g_{tqL}}{\Lambda}\bar{q}_R\sigma^{\mu\nu}T^a t_L G^a_{\mu\nu} \\ &- \frac{g}{\sqrt{2}}\sum_{q=d,s,b}\frac{\kappa^W_{tqL}}{\Lambda}\bar{q}_R\sigma^{\mu\nu}t_L W^-_{\mu\nu} \\ &- e\sum_{q=u,c,t}\frac{\kappa^{\gamma}_{tqL}}{\Lambda}\bar{q}_R\sigma^{\mu\nu}t_L F_{\mu\nu} \\ &- \frac{g}{2\cos\theta_W}\sum_{q=u,c,t}\frac{\kappa^Z_{tqL}}{\Lambda}\bar{q}_R\sigma^{\mu\nu}t_L Z_{\mu\nu} \\ &+ (R \leftrightarrow L) + \text{H.c.}, \end{aligned} \tag{1}$$

where κ is the complex coupling of its corresponding operator, θ_W is the weak angle, and T^a is the Gell-Mann matrix. Λ is the possible new physics scale, which is unknown but may be much larger than the electroweak scale. There are also Lagrangians describing the top quark anomalous interactions with dimension 4 and 6 operators, and the dimension 4 and 5 terms can be traced back to dimension 6 operators [12,13]. In fact, top quark anomalous interactions can be generally described by the gauge-invariant effective Lagrangian with dimension 6 operators in a form without redundant operators and parameters [10,14]. A recent full list of dimension 6 operators can be found in Ref. [15]. But for on-shell gauge bosons, the Lagrangian in Eq. (1) works and is commonly employed in high energy phenomenology analysis [3,6,16].

The operators in Eq. (1) relevant to $t \to q\gamma$ decays read

$$\mathcal{L}_\gamma = -e \sum_{q=u,c} \frac{\kappa^\gamma_{tqL}}{\Lambda} \bar{q}_R \sigma^{\mu\nu} t_L F_{\mu\nu} - e \sum_{q=u,c} \frac{\kappa^\gamma_{tqR}}{\Lambda} \bar{q}_L \sigma^{\mu\nu} t_R F_{\mu\nu} + \text{H.c.} \quad (2)$$

It is understood that the Dirac matrix $\sigma_{\mu\nu}$ connects left-handed fields to right-handed fields. The $t \to c\gamma$ transition will involve two independent operators $m_q \bar{q}_R \sigma^{\mu\nu} t_L F_{\mu\nu}$ and $m_t \bar{q}_L \sigma^{\mu\nu} t_R F_{\mu\nu}$, where the mass factors must appear whenever a chirality flip $L \to R$ or $R \to L$ occurs. Because of the mass hierarchy $m_t \gg m_c$, the effect of $m_q \bar{q}_R \sigma^{\mu\nu} t_L F_{\mu\nu}$ can be neglected unless κ^γ_{tqL} is enhanced to be comparable to $\frac{m_t}{m_c} \kappa^\gamma_{tqR}$ by an unknown mechanism.

The anomalous $t\gamma q$ coupling affects $b \to s\gamma$ decays through the two Feynman diagrams depicted in Figs. 1(a) and 1(b). It is interesting to note that the Cabibbo-Kobayashi-Maskawa (CKM) factors in Figs. 1(a) and 1(b) are $V_{tb}V^*_{qs}$ and $V_{qb}V^*_{ts}$, respectively. Since $|V_{tb}V^*_{qs}| \gg |V_{qb}V^*_{ts}|$ for $q = u, c$, the contribution of Fig. 1(a) would be much stronger than that of Fig. 1(b). Furthermore, given the fact that the strengths of $t \to u\gamma$ and $t \to c\gamma$ are comparable, the contribution of Fig. 1(a) to $b \to s\gamma$ is still dominated by $t \to c\gamma$ because of $|V_{cs}| \gg |V_{us}|$. Hence we will only consider Fig. 1(a) with anomalous $tc\gamma$ coupling. From the Feynman diagram of Fig. 1(a), it is easy to observe that the large CKM factor $V_{tb}V_{cs} \approx 1$ makes $b \to s\gamma$ very sensitive to the strength of anomalous $tc\gamma$ coupling.

The calculation of Fig. 1(a) can be carried out straightforwardly. The calculation details are presented in Appendix A, and the final result reads

$$\begin{aligned} i\mathcal{M}(b \to s\gamma) &= \bar{s}[e\Gamma^\nu(k)]b\epsilon_\nu(k),\ e\Gamma^\nu(p,k) \\ &= ie\frac{G_F}{4\sqrt{2}\pi^2} V^*_{cs}V_{tb}[i\sigma^{\nu\mu}k_\mu(m_s f_L(x)L \\ &\quad + m_b f_R(x)R)]. \end{aligned} \quad (3)$$

Usually the m_s term can be neglected, and the function $f_R(x)$ is calculated to be

$$f_R(x) = \frac{\kappa^\gamma_{tcR}}{\Lambda} 2m_t \Big[-\frac{1}{(x_c-1)(x_t-1)} - \frac{x_c^2}{(x_c-1)^2(x_c-x_t)} \ln x_c + \frac{x_t^2}{(x_t-1)^2(x_c-x_t)} \ln x_t \Big], \quad (4)$$

with $x_q = m_q^2/m_W^2$. Now we are ready to incorporate the NP contribution into its SM counterpart for $\bar{B} \to X_s\gamma$ decay.

In the SM, it is known that $\bar{B} \to X_s\gamma$ decay is governed by the effective Hamiltonian at scale $\mu = \mathcal{O}(m_b)$ [17]

$$\mathcal{H}_{\text{eff}}(b \to s\gamma) = -\frac{4G_F}{\sqrt{2}} V^*_{ts}V_{tb} \Big[\sum_{i=1}^{6} C_i(\mu)Q_i(\mu) + C_{7\gamma}(\mu)O_{7\gamma}(\mu) + C_{8g}(\mu)O_{8g}(\mu) \Big], \quad (5)$$

where $C_i(\mu)$ are the Wilsion coefficients, $O_{i=1-6}$ are the effective four quark operators, and

$$\begin{aligned} O_{7\gamma} &= \frac{e}{16\pi^2} m_b(\bar{s}_L \sigma^{\mu\nu} b_R) F_{\mu\nu}, \\ O_{8g} &= \frac{g}{16\pi^2} m_b(\bar{s}_L \sigma^{\mu\nu} T^a b_R) G^a_{\mu\nu}. \end{aligned} \quad (6)$$

For calculating $\mathcal{B}(\bar{B} \to X_s\gamma)$, instead of the original Wilson coefficients C_i, it is convenient to use the so-called "effective coefficients" [18]

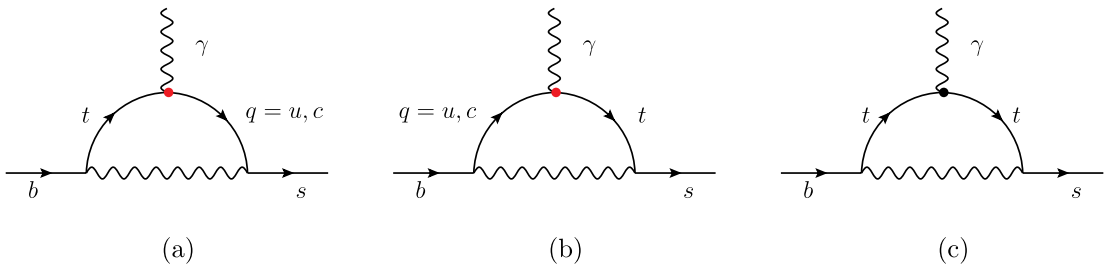

FIG. 1 (color online). Feynman diagrams for $b \to s\gamma$. (a) and (b) are the penguin diagrams with the anomalous $tq\gamma$ couplings. (c) Sample LO penguin diagram in the SM.

$$C_{7\gamma}^{(0)\mathrm{eff}}(m_b) = \eta^{16/23} C_{7\gamma}^{(0)\mathrm{SM}}(M_W) + \frac{8}{3}(\eta^{14/23} - \eta^{16/23}) C_{8g}^{(0)\mathrm{SM}}(M_W) + C_2^{(0)\mathrm{SM}}(M_W) \sum_{i=1}^{8} h_i \eta^{a_i}, \quad (7)$$

where $\eta = \alpha_s(\mu_W)/\alpha_s(\mu_b)$, and

$$h_i = \left(\frac{626\,126}{272\,277} \quad -\frac{56\,281}{51\,730} \quad -\frac{3}{7} \quad -\frac{1}{14} \quad -0.6494 \quad -0.0380 \quad -0.0185 \quad -0.0057\right), \quad (8)$$

$$a_i = \left(\frac{14}{23} \quad \frac{16}{23} \quad \frac{6}{23} \quad -\frac{12}{23} \quad 0.4086 \quad -0.4230 \quad -0.8994 \quad 0.1456\right). \quad (9)$$

To the leading order approximation, the $\mathcal{B}(\bar{B} \to X_s\gamma)$ is proportional to $|C_{7\gamma}^{(0)\mathrm{eff}}(m_b)|^2$ [19].

In terms of the operator basis in Eq. (5), the contribution of the anomalous $t \to c\gamma$ couplings in Eq. (3) would result in the deviation of

$$C_{7\gamma}(M_W) \to C'_{7\gamma}(M_W) = C_{7\gamma}(M_W) + C_{7\gamma}^{\mathrm{NP}}(M_W) \quad (10)$$

and $C_{7\gamma}^{\mathrm{NP}}(M_W)$ can be read from Eq. (3) as

$$C_{7\gamma}^{\mathrm{NP}}(M_W) = \frac{\kappa_{tcR}^{\gamma}}{\Lambda} \frac{V_{cs}^*}{V_{ts}^*} m_t \left[\frac{1}{(x_c-1)(x_t-1)} + \frac{x_c^2}{(x_c-1)^2(x_c-x_t)} \log x_c - \frac{x_t^2}{(x_t-1)^2(x_c-x_t)} \ln x_t\right]. \quad (11)$$

From this equation, one can see that the NP contribution is suppressed by a factor of m_t/Λ but enhanced by V_{cs}/V_{ts}.

Since the NP contribution does not bring about any new operator, the renormalization group evolution of $C_{7\gamma}^{\mathrm{eff}}$ from M_W to the m_b scale is just the same as the SM one in Eq. (7). For $m_t = 172$ GeV, $m_b = 4.67$ GeV, $\alpha_s(M_Z) = 0.118$, and $\Lambda = 1$ TeV, we have

$$\begin{aligned} C_{7\gamma}'^{\mathrm{eff}}(m_b) &= \eta^{16/23}[C_{7\gamma}^{(0)\mathrm{SM}}(M_W) + C_{7\gamma}^{(0)\mathrm{NP}}(M_W)] + \frac{8}{3}(\eta^{14/23} - \eta^{16/23}) C_{8g}^{(0)\mathrm{SM}}(M_W) \\ &\quad + C_2^{(0)\mathrm{SM}}(M_W) \sum_{i=1}^{8} h_i \eta^{a_i} = 0.665[C_{7\gamma}^{(0)\mathrm{SM}}(M_W) + C_{7\gamma}^{(0)\mathrm{NP}}(M_W)] \\ &\quad + 0.093 C_{8g}^{(0)\mathrm{SM}}(M_W) - 0.158 C_2^{(0)\mathrm{SM}}(M_W) = 0.665[-0.189 + \kappa_{tcR}^{\gamma}(-1.092)] \\ &\quad + 0.093(-0.095) - 0.158. \end{aligned} \quad (12)$$

In principle, $C_{7\gamma}'^{\mathrm{eff}}(m_b)$ will receive corrections from anomalous $t \to cg$ couplings in Eq. (1) which will cause a deviation to $C_{8g}^{(0)\mathrm{SM}}(M_W)$. However, as shown by Eq. (12), the coefficient $\eta^{16/23}$ of $C_{7\gamma}^{(0)}(M_W)$ is about 1 order larger than $\frac{8}{3}(\eta^{14/23} - \eta^{16/23})$ of $C_{8g}^{(0)\mathrm{NP}}(M_W)$. Given the relative strength of $C_{8g}^{(0)\mathrm{NP}}(M_W)$ to $C_{8g}^{(0)\mathrm{SM}}(M_W)$ at 10% level, $C_{7\gamma}'^{\mathrm{eff}}(m_b)$ will be shifted by only a few percentage. For simplifying the numerical analysis, we neglect the contribution of the anomalous $t \to cg$ couplings. We also find that the operator $\bar{q}_R \sigma^{\mu\nu} t_L F_{\mu\nu}$ contributes to $\bar{B} \to X_s\gamma$ only through the term $m_s \bar{s} \sigma_{\mu\nu}(1-\gamma_5) b$ as shown by Eqs. (3) and (A7). Combined with the previous remarks on this operator, the effects of $\bar{q}_R \sigma^{\mu\nu} t_L F_{\mu\nu}$ could be safely neglected.

III. NUMERICAL RESULTS AND DISCUSSIONS

The current average of experimental results of $\mathcal{B}(\bar{B} \to X_s\gamma)$ by the Heavy Flavor Average Group is [20]

$$\mathcal{B}^{\mathrm{exp}}(\bar{B} \to X_s\gamma) = (3.55 \pm 0.24 \pm 0.09) \times 10^{-4}. \quad (13)$$

On the theoretical side, the next-to-leading order (NLO) calculation has been completed [19,21] and gives

$$\mathcal{B}(\bar{B} \to X_s\gamma) = (3.57 \pm 0.30) \times 10^{-4}. \quad (14)$$

The recent estimation at next-to-next-to-leading order (NNLO) [22] gives $\mathcal{B}(\bar{B} \to X_s\gamma) = (3.15 \pm 0.23) \times 10^{-4}$, which is about 1σ lower than the experimental average in Eq. (13). Thus the experimental measurement of $\mathcal{B}(\bar{B} \to X_s\gamma)$ is in good agreement with the SM predictions with roughly 10% errors on each side. The agreement would provide strong constraints on the top quark anomalous interactions beyond the SM [9,10].

The decay amplitude of $t \to c\gamma$ has been calculated up to NLO [16]. For a consistent treatment of the constraints from $t \to c\gamma$ and $b \to s\gamma$ decays, we use the NLO formulas in Ref. [21] to calculate $\mathcal{B}(\bar{B} \to X_s\gamma)$. The experimental inputs and main formulas are collected in Appendix B. For numerical analysis, we will use the notation $\kappa_{tcR}^{\gamma} = |\kappa_{tcR}^{\gamma}| e^{i\theta_{tcR}^{\gamma}}$ and set $\Lambda = 1$ TeV.

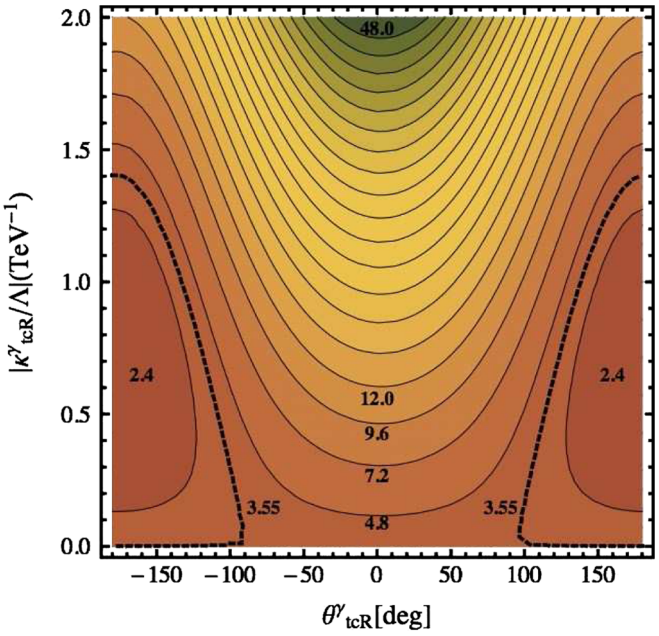

FIG. 2 (color online). The contour plot describes the dependence of $\mathcal{B}(\bar{B} \to X_s\gamma)(\times 10^{-4})$ on $|\kappa^{\gamma}_{tcR}/\Lambda|$ and θ^{γ}_{tcR}. The dashed lines correspond to the experimental center value of $\mathcal{B}(\bar{B} \to X_s\gamma)$.

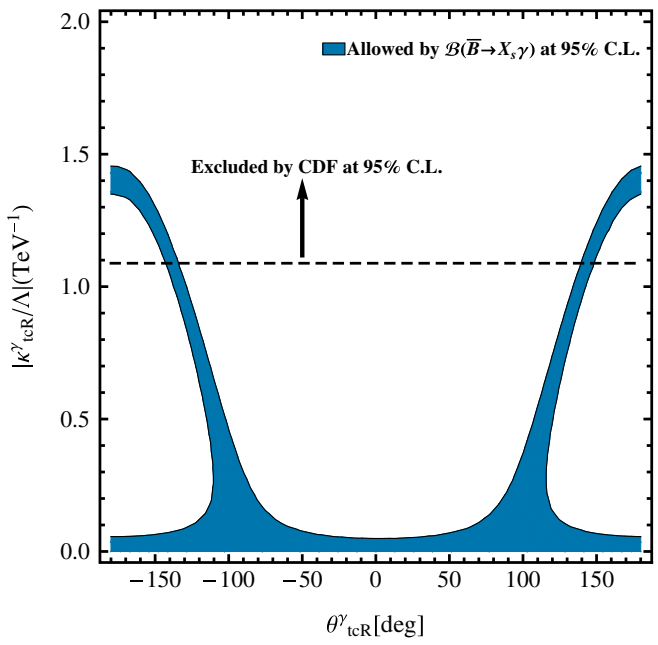

FIG. 3 (color online). The 95% C.L. upper bounds on anomalous coupling $|\kappa^{\gamma}_{tcR}/\Lambda|$ as a function of θ^{γ}_{tcR}. The shadowed region is allowed by $\mathcal{B}^{\mathrm{exp}}(\bar{B} \to X_s\gamma)$ and the dashed line is the CDF [5] upper limit.

At first, we analyze the dependence of $\mathcal{B}^{\mathrm{SM+NP}}(\bar{B} \to X_s\gamma)$ on the new physics parameters $|\kappa^{\gamma}_{tcR}/\Lambda|$ and θ^{γ}_{tcR}, which is shown in Fig. 2. From the figure, one can find that the contribution of anomalous $t \to c\gamma$ coupling is constructive to the SM one for $\theta^{\gamma}_{tcR} \in [-50°, 50°]$; thus $\mathcal{B}(\bar{B} \to X_s\gamma)$ is very sensitive to $|\kappa^{\gamma}_{tcR}|$. However, when $|\theta^{\gamma}_{tcR}| \in [80°, 130°]$, the sensitivity of $\mathcal{B}(\bar{B} \to X_s\gamma)$ to $|\kappa^{\gamma}_{tcR}|$ becomes weak. For $|\theta^{\gamma}_{tcR}| \sim 180°$, the contribution of anomalous $t \to c\gamma$ coupling is destructive to the SM one and there are two separated possible strengths for $|\kappa^{\gamma}_{tcR}/\Lambda|$.

The allowed region for the parameters $|\kappa^{\gamma}_{tcR}/\Lambda|$ and θ^{γ}_{tcR} under the constraints from $\mathcal{B}(\bar{B} \to X_s\gamma)$ at 95% C.L. is shown in Fig. 3. The corresponding 95% C.L. upper bound on $\mathcal{B}(t \to c\gamma)$ is shown in Fig. 4.

Now we discuss our numerical results. From Eq. (12), the explicit relation between the SM and the $t \to c\gamma$ coupling contributions is

$$C^{\mathrm{eff}}_{7\gamma}(m_b) = -0.293 - 0.726\kappa^{\gamma}_{tcR}. \qquad (15)$$

Obviously, when $\mathrm{Re}\kappa^{\gamma}_{tcR} > 0$, the interference between them is constructive, and it turns out to be destructive when $\theta^{\gamma}_{tcR} > 90°$. Thus the features of these constraints shown in Figs. 3 and 4 for different θ^{γ}_{tcR} are as follows:

(i) The bound on $|\kappa^{\gamma}_{tcR}/\Lambda|$ is very strong for $\theta^{\gamma}_{tcR} \in [-50°, 50°]$. For $\theta^{\gamma}_{tcR} \approx 0°$, as shown in Fig. 3, we obtain the most restrictive upper bound $|\kappa^{\gamma}_{tcR}/\Lambda| < 4.9 \times 10^{-5}\ \mathrm{GeV}^{-1}$, which implies $\mathcal{B}(t \to c\gamma) < 6.54 \times 10^{-5}$.

(ii) The bound on $|\kappa^{\gamma}_{tcR}/\Lambda|$ is rather weak for θ^{γ}_{tcR} around 110°. For such a case, $\mathrm{Re}\kappa^{\gamma}_{tcR}$ is destructive to the SM contribution as shown by Eq. (15), so, the allowed strength for the anomalous coupling is

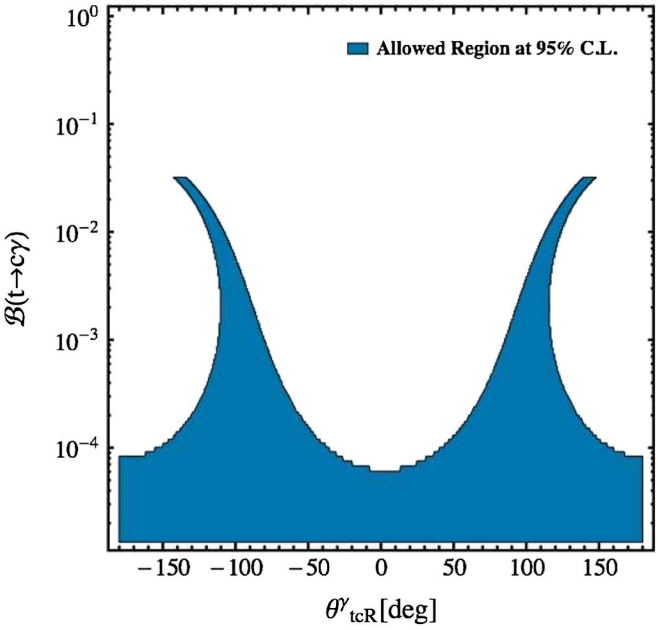

FIG. 4 (color online). $\mathcal{B}(t \to c\gamma)$ as a function of θ^{γ}_{tcR}. The shadowed region is allowed by the combined constraints of $\mathcal{B}(\bar{B} \to X_s\gamma)$ and CDF searching at 95% C.L.

TABLE I. The 95% C.L. constraints on the anomalous $t \to c\gamma$ coupling by $\mathcal{B}(\bar{B} \to X_s\gamma)$ and $\mathcal{B}(t \to c\gamma)$ for some specific θ^{γ}_{tcR} values.

	$\theta^{\gamma}_{tcR} = 0°$	$\theta^{\gamma}_{tcR} = \pm 180°$ S1	$\theta^{\gamma}_{tcR} = \pm 180°$ S2	$\theta^{\gamma}_{tcR} = \pm 110°$
$\mathcal{B}(\bar{B} \to X_s\gamma)$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 0.049$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 0.056$	$1.35 < \lvert\kappa^{\gamma}_{tcR}\rvert < 1.45$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 0.55$
$\mathcal{B}(t \to c\gamma)$ CDF bounds [5]	$\lvert\kappa^{\gamma}_{tcR}\rvert < 1.09$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 1.09$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 1.09$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 1.09$
Combined bounds	$\lvert\kappa^{\gamma}_{tcR}\rvert < 0.049$	$\lvert\kappa^{\gamma}_{tcR}\rvert < 0.056$	...	$\lvert\kappa^{\gamma}_{tcR}\rvert < 0.55$
$\mathcal{B}(t \to c\gamma)$	$< 6.54 \times 10^{-5}$	$< 8.52 \times 10^{-5}$	...	$< 8.17 \times 10^{-3}$

much larger than the one for real κ^{γ}_{tcR}. When $|\theta^{\gamma}_{tcR}| \approx 135°$ and $|\kappa^{\gamma}_{tcR}| \approx 0.571$, $C'^{\rm eff}_{7\gamma}(m_b)$ is almost imaginary since $\mathrm{Re}C'^{\rm eff}_{7\gamma}(m_b) \approx 0$. Then the restriction on $|\kappa^{\gamma}_{tcR}/\Lambda|$ is provided by the CDF search for $\mathcal{B}(t \to c\gamma)$ [5].

(iii) As shown in Fig. 3, when $\theta^{\gamma}_{tcR} \sim \pm 180°$, there are two solutions for $|\kappa^{\gamma}_{tcR}/\Lambda|$. The larger one $|\kappa^{\gamma}_{tcR}/\Lambda| \sim 1.4 \times 10^{-3}$ GeV^{-1} (S2 column in Table I) corresponds to the situation in which the sign of $C^{\rm eff}_{7\gamma}$ is flipped. However, it has been excluded by the CDF upper bound of $\mathcal{B}(t \to c\gamma) < 0.032$ [5]. The other solution (S1 column in Table I) $|\kappa^{\gamma}_{tcR}/\Lambda| < 5.6 \times 10^{-5}$ GeV^{-1} will result in the upper limit $\mathcal{B}(t \to c\gamma) < 8.52 \times 10^{-5}$.

Taking $\theta^{\gamma}_{tcR} = 0°, \pm 180°$ and $\pm 110°$ as benchmarks, we summarize our numerical constraints on κ^{γ}_{tcR} and their corresponding upper limits on $\mathcal{B}(t \to c\gamma)$ in Table I. From the table, we can find that our indirect bound on real κ^{γ}_{tcR} is much stronger than the CDF direct bound. The corresponding upper limits on $\mathcal{B}(t \to c\gamma)$ are about the same order as the ATLAS sensitivity $\mathcal{B}(t \to c\gamma) > 9.4 \times 10^{-5}$ [6] and CMS sensitivity $\mathcal{B}(t \to c\gamma) > 4.1 \times 10^{-4}$ [7] with an integrated luminosity of 10 fb^{-1} of the LHC operating at $\sqrt{s} = 14$ TeV [6].

IV. CONCLUSIONS

In this paper, starting with model-independent dimension 5 anomalous $tc\gamma$ operators, we have studied their contributions to $\mathcal{B}(\bar{B} \to X_s\gamma)$. It is noted that the $t \to c\gamma$ transition will involve two independent operators $\kappa^{\gamma}_{tcR}\bar{c}_L\sigma^{\mu\nu}t_R F_{\mu\nu}$ and $\kappa^{\gamma}_{tcL}\bar{c}_R\sigma^{\mu\nu}t_L F_{\mu\nu}$. The first operator will produce a left-handed photon in $t \to c\gamma$ decay, while the second one will produce a right-handed photon. It is found that $\bar{B} \to X_s\gamma$ is sensitive to the first operator, but not to the second one.

For real κ^{γ}_{tcR}, the constraint on the presence of $\kappa^{\gamma}_{tcR}\bar{c}_L\sigma^{\mu\nu}t_R F_{\mu\nu}$ is very strong, which corresponds to the indirect upper limits $\mathcal{B}(t \to c\gamma) < 6.54 \times 10^{-5}$ (for positive κ^{γ}_{tcR}) and $\mathcal{B}(t \to c\gamma) < 8.52 \times 10^{-5}$ (for negative κ^{γ}_{tcR}), respectively. These upper limits for $\mathcal{B}(t \to c\gamma)$ are close to the 5σ discovery sensitivities of ATLAS [6] and slightly lower than that of CMS [7] with 10 fb^{-1} integrated luminosity operating at $\sqrt{s} = 14$ TeV. For nearly imaginary κ^{γ}_{tcR}, the constraints are rather weak since $C_{7\gamma}$ in the SM is a real number. If $\mathcal{B}(t \to c\gamma)$ were found to be of the order of $\mathcal{O}(10^{-3})$ at the LHC in the future, it would imply the weak phase of κ^{γ}_{tcR} to be around $\pm 100°$. However, such a coupling might be ruled out by the other observable in B meson decays [23].

In summary, we have studied the interesting interplay between the precise measurement of $b \to s\gamma$ decay at B factories and the possible $t \to c\gamma$ decay at the LHC. For real anomalous coupling, it is shown that $\mathcal{B}(t \to c\gamma)$ has been restricted to be below 10^{-4} at 95% C.L. by $\bar{B} \to X_s\gamma$ decay, which is already 2 orders lower than the direct upper bound from CDF [5]. The result also implies that one may need a data sample much larger than 10 fb^{-1} to hunt out $t \to c\gamma$ signals at the LHC.

ACKNOWLEDGMENTS

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APPENDIX A: THE CALCULATION OF $C^{\rm NP}_{7\gamma}(\mu_W)$

Using the Feynman rules in Fig. 5(a), the amplitude of the penguin diagram in Fig. 5(b) can be written as

$$i\mathcal{M} = \bar{u}_s(p')[e\Gamma^{\nu}(p,k)]u_b(p)\epsilon_{\nu}(k), \qquad \text{(A1)}$$

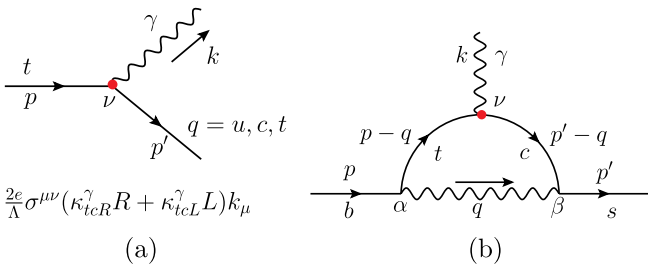

FIG. 5 (color online). (a) The Feynman rules of $t\gamma c$ interactions in the Lagrangian of Eq. (1). (b) Penguin diagram contribution to $b \to s\gamma$ with top quark anomalous interactions.

$$\Gamma^{\nu}(p,k) = -\frac{ig^2}{\Lambda} V_{cs}^* V_{tb} \int \frac{d^4q}{(2\pi)^4} \frac{N}{[(p'-q)^2 - m_c^2 + i\epsilon][(p-q)^2 - m_t^2 + i\epsilon][q^2 - m_W^2 + i\epsilon]}, \tag{A2}$$

$$\begin{aligned} N = {} & \gamma_\alpha L(p\!\!\!/' - q\!\!\!/ + m_q)\sigma^{\mu\nu}(\kappa_{tcR}^{\gamma} R + \kappa_{tcL}^{\gamma} L) \\ & \times (p\!\!\!/ - q\!\!\!/ + m_t)\gamma_\beta L g^{\alpha\beta} k_\mu, \end{aligned} \tag{A3}$$

with $R = (1+\gamma^5)/2$ and $L = (1-\gamma^5)/2$. By Dirac algebra

$$\gamma_\alpha L q\!\!\!/ \sigma^{\mu\nu}(\kappa_{tcR}^{\gamma} R + \kappa_{tcL}^{\gamma} L) q\!\!\!/ \gamma_\beta L = 0. \tag{A4}$$

The terms with q^2 in N vanish and N becomes

$$\begin{aligned} N = {} & m_c \kappa_{tcL}^{\gamma}[2(p\!\!\!/ - q\!\!\!/)\sigma^{\mu\nu} + (4-D)\sigma^{\mu\nu}(p\!\!\!/ - q\!\!\!/)] \\ & \times L k_\mu + m_t \kappa_{tcR}^{\gamma}[2\sigma^{\mu\nu}(p\!\!\!/' - q\!\!\!/) \\ & + (4-D)(p\!\!\!/' - q\!\!\!/)\sigma^{\mu\nu}] R k_\mu. \end{aligned} \tag{A5}$$

Thus, there is no divergence in $\Gamma^{\nu}(p,k)$. After integrating out q in the $\Gamma^{\nu}(p,k)$ and using the on-shell condition, $\Gamma^{\nu}(p,k)$ can be written in the following form:

$$\begin{aligned} e\Gamma^{\nu}(p,k) = {} & ie \frac{G_F}{4\sqrt{2}\pi^2} V_{cs}^* V_{tb}[i\sigma^{\nu\mu} k_\mu (m_s f_L(x) L \\ & + m_b f_R(x) R)], \end{aligned} \tag{A6}$$

where

$$\begin{aligned} f_L(x) = {} & \frac{\kappa_{tcL}^{\gamma}}{\Lambda} 2m_c \Big[-\frac{1}{(x_c-1)(x_t-1)} \\ & - \frac{x_c^2}{(x_c-1)^2(x_c-x_t)} \ln x_c \\ & + \frac{x_t^2}{(x_t-1)^2(x_c-x_t)} \ln x_t \Big], \end{aligned} \tag{A7}$$

$$\begin{aligned} f_R(x) = {} & \frac{\kappa_{tcR}^{\gamma}}{\Lambda} 2m_t \Big[-\frac{1}{(x_c-1)(x_t-1)} \\ & - \frac{x_c^2}{(x_c-1)^2(x_c-x_t)} \ln x_c \\ & + \frac{x_t^2}{(x_t-1)^2(x_c-x_t)} \ln x_t \Big]. \end{aligned} \tag{A8}$$

Using the convention of Ref. [19], we have

$$\begin{aligned} C_{7\gamma}^{(0)\mathrm{NP}}(M_W) = {} & -\frac{1}{2}\frac{V_{cs}^*}{V_{ts}^*} f_R(x) \\ = {} & \frac{V_{cs}^*}{V_{ts}^*} m_t \frac{\kappa_{tcR}^{\gamma}}{\Lambda} \Big[\frac{1}{(x_c-1)(x_t-1)} \\ & + \frac{x_c^2}{(x_c-1)^2(x_c-x_t)} \ln x_c \\ & - \frac{x_t^2}{(x_t-1)^2(x_c-x_t)} \ln x_t \Big]. \end{aligned} \tag{A9}$$

APPENDIX B: MAIN FORMULAS AND INPUTS

Following the notation in Ref. [21], the branching ratio of $\bar{B} \to X_s\gamma$ can be expressed as

$$\begin{aligned} & \mathcal{B}[\bar{B} \to X_s\gamma]_{E_\gamma > E_0} \\ & = \mathcal{B}^{\exp}[\bar{B} \to X_c e\bar{\nu}] \left| \frac{V_{ts}^* V_{tb}}{V_{cb}} \right|^2 \frac{6\alpha_{\mathrm{em}}}{\pi C}[P(E_0) + N(E_0)], \end{aligned} \tag{B1}$$

where $P(E_0)$ is the perturbative ratio

$$\frac{\Gamma[\bar{B} \to X_s\gamma]_{E_\gamma > E_0}}{|V_{cb}/V_{ub}|^2 \Gamma[\bar{B} \to X_u e\bar{\nu}]} = \left| \frac{V_{ts}^* V_{tb}}{V_{cb}} \right|^2 \frac{6\alpha_{\mathrm{em}}}{\pi} P(E_0), \tag{B2}$$

which includes the Wilson coefficients of Eq. (7). $N(E_0)$ denotes the nonperturbative corrections. The semileptonic phase space factor

$$C = \left| \frac{V_{ub}}{V_{cb}} \right|^2 \frac{\Gamma[\bar{B} \to X_c e\bar{\nu}]}{\Gamma[\bar{B} \to X_u e\bar{\nu}]} \tag{B3}$$

can be obtained from a fit of the experimental spectrum of the $\bar{B} \to X_c l\bar{\nu}$ [24].

For calculating $\mathcal{B}(t \to c\gamma)$, we use the NLO formulas in Refs. [16,25]. Because $t \to bW$ is the dominant top quark decay mode, the branching ratio of $t \to c\gamma$ is defined as

$$\mathcal{B}(t \to c\gamma) = \frac{\Gamma(t \to c\gamma)}{\Gamma(t \to bW)}. \tag{B4}$$

The partial width $\Gamma(t \to c\gamma)$ at the NLO can be found in Ref. [16], namely,

$$\begin{aligned} & \Gamma_{\mathrm{NLO}}(t \to c\gamma) \\ & = \frac{2\alpha_s}{9\pi} \Gamma_0(t \to c\gamma) \Big[-3\log\Big(\frac{\mu^2}{m_t^2}\Big) - 2\pi^2 + 8 \Big], \end{aligned} \tag{B5}$$

where $\Gamma_0(t \to c\gamma) = \alpha m_t^3 (\kappa_{tcR}^{\gamma}/\Lambda)^2$ is the LO partial decay width.

The partial width of $t \to bW$ has been calculated in Ref. [25] at the NLO, which reads

TABLE II. Experimental inputs for calculating the branching ratios of $\bar{B} \to X_s\gamma$ and $t \to c\gamma$.

Experimental inputs	
$\alpha_{\rm em} = 1/137.036$ [26]	$M_Z = 91.1876 \pm 0.0021$ GeV [26]
$\alpha_s(M_Z) = 0.1184 \pm 0.0007$ [26]	$M_W = 80.399 \pm 0.023$ GeV [26]
$G_F = 1.16637 \times 10^{-5}$ GeV^{-2} [26]	$m_b^{1S} = 4.67^{+0.18}_{-0.06}$ GeV [26]
$A = 0.812^{+0.013}_{-0.027}$ [27]	$m_c(m_c) = (1.224 \pm 0.017 \pm 0.054)$ GeV [28]
$\lambda = 0.225\,43 \pm 0.000\,77$ [27]	$m_{t,\rm pole} = 172.0 \pm 0.9 \pm 1.3$GeV [26]
$\bar{\rho} = 0.144 \pm 0.025$ [27]	$\mathcal{B}^{\rm exp}[\bar{B} \to X_c e \bar{\nu}] = (10.64 \pm 0.17 \pm 0.06)\%$ [29]
$\bar{\eta} = 0.342^{+0.016}_{-0.015}$ [27]	$C = 0.580 \pm 0.016$ [24]
$\lvert V_{ts}^* V_{tb}/V_{cb}\rvert^2 = 0.9625$	$\epsilon_{\rm ew} = 0.0071$ [21,30]
$(V_{us}^* V_{ub})/(V_{ts}^* V_{tb}) = -0.007 + 0.018i$	$N(E_0) = 0.0036 \pm 0.0006$ [21]
$V_{cs}^*/V_{ts}^* = -24.023 - 0.432i$	$E_0 = 1.6$ GeV

$$\Gamma_{\rm NLO}(t \to bW) = \Gamma_0(t \to bW)\Big\{1 + \frac{2\alpha_s}{3\pi}\Big[2\Big(\frac{(1-\beta_W^2)(2\beta_W^2-1)(\beta_W^2-2)}{\beta_W^4(3-2\beta_W^2)}\Big)\ln(1-\beta_W^2) - \frac{9-4\beta_W^2}{3-2\beta_W^2} \times \ln\beta_W^2 + 2{\rm Li}_2(\beta_W^2) - 2{\rm Li}_2(1-\beta_W^2) - \frac{6\beta_W^4 - 3\beta_W^2 - 8}{2\beta_W^2(3-2\beta_W^2)} - \pi^2\Big]\Big\} \tag{B6}$$

with $\Gamma_0(t \to bW) = (G_F m_t^3/8\sqrt{2}\pi)|V_{tb}|^2\beta_W^4(3-2\beta_W^2)$ and $\beta_W \equiv (1 - m_W^2/m_t^2)^{1/2}$.

The experimental inputs are collected in Table II, in which the CKM factors are derived from the Wolfenstein parameters A, λ, $\bar{\rho}$, and $\bar{\eta}$.

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