Anomalous couplings of the third generation in rare *B* decays

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We study the potential effects of anomalous couplings of the third generation quarks to gauge bosons in rare *B* decays. We focus on the constraints from flavor changing neutral current processes such as $b\to s\gamma$ and $b\to sl^+l^-$. We consider both dimension-four and dimension-five operators and show that the latter can give large deviations from the standard model in the still unobserved dilepton modes, even after the bounds from $b\to s\gamma$ and precision electroweak observables are taken into account.

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I. INTRODUCTION

The continuing experimental success of the standard model (SM) suggests the possibility that additional particles and/or non-standard interactions may only be found at scales much larger than M_W. On the other hand, several questions remain unanswered within the SM framework that may require new dynamics in order to be addressed. Chief among these questions are the origin of electroweak symmetry breaking and of fermion masses. In principle, it could be argued that the energy scales of the new dynamics related to these questions may be so large as to be irrelevant to observables at the electroweak scale. However, it is known that the physics behind the Higgs sector, responsible for the breaking of the electroweak symmetry, cannot reside at scales much higher than few TeV. Furthermore, it is possible that the origin of the top quark mass might be related to electroweak symmetry breaking. Thus, at least in some cases, the dynamics associated with new physics may not reside at arbitrarily high energies and there might be some observable effects at lower energies.

The effects of integrating out the physics residing at some high energy scale $\Lambda \gg M_W$, can be organized in an effective field theory for the remaining degrees of freedom. Such a theory for the electroweak symmetry breaking (EWSB) sector of the SM involves the electroweak gauge bosons as well as the Nambu-Goldstone bosons (NGB) associated with the spontaneous breaking of $SU(2)_L\times U(1)_Y$ down to $U(1)_{\rm EM}$ [1,2]. The effective theory must be studied up to next-to-leading order for the possible departures from the SM to appear. This program resembles that of chiral perturbation theory for pions in low energy QCD where, for instance, the presence of the ρ resonance results in deviations from the low energy theorems. For the case of the electroweak interactions, a variety of electroweak precision measurements and flavor changing neutral current processes provide testing ground for possible deviations originating in the EWSB sector of the SM. The next-to-leading order terms in the effective theory will generally contribute to oblique corrections, triple and quartic anomalous gauge boson couplings, and corrections to the NGB propagators that result in four-fermion interactions [3].

In addition to the low energy description of the interactions of the EWSB sector (i.e. gauge bosons plus NGBs) one may consider the possibility that the new physics above Λ may also modify the effective interactions of the SM fermions to the electroweak gauge bosons. In principle, this also has a parallel in low energy QCD, as it is pointed out in Ref. [4], where symmetry alone is not enough to determine the axial coupling of nucleons to pions. In fact, the departure of this coupling from unity is a non-universal effect, only determined by the full theory of QCD. Thus, in Ref. [4] it is suggested that in addition to the effects in the EWSB sector of the theory, it is possible that the interactions of fermions with the NGBs are affected by the new dynamics above Λ, resulting in anomalous interactions with the electroweak gauge bosons. This is particularly interesting if fermion masses are dynamically generated, as is the case with the nucleon mass. Interestingly, the proximity of the top quark mass to the electroweak scale $v=246$ GeV, hints the possibility the top quark mass might be a dynamically generated ''constituent'' mass. Thus, it is of particular interest to study the couplings of third generation quarks to electroweak gauge bosons.

Processes involving FCNC transitions in *B* and *K* decays are a crucial complement to precision electroweak observables, when constraining the physics of the EWSB sector. The effects of anomalous triple gauge boson couplings [5], as well as of the corrections to NGB propagators [6] give in each case a distinct pattern of deviations from the SM expectations in rare *B* and *K* decays. On the other hand, the anomalous couplings of third generation quarks to the *W* and

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the Z can come from dimension-four and dimension-five operators. The indirect effects of the dimension-four operators have been considered in relation to electroweak observables in Refs. [4,7], as well as the $b\to s\gamma$ transitions [8]. The constraints on dimension-five operators from electroweak physics have been studied in Ref. [9]. In this paper, we consider the effects of all possible dimension-five operators in B flavor changing neutral current (FCNC) transitions such as $b\to s\gamma$ and $b\to s l^+ l^-$. For completeness, we also present the analysis of the dimension-four operators. We discuss the stringent limits found for one dimension-four parameter in Ref. [8] for $b\to s\gamma$ in light of underlying symmetries. Moreover, we will see that the effects of dimension-five operators are comparable and may even dominate over the supposedly leading lower dimension contributions.

In Sec. II we present a brief introduction to the effective theory approach and set our notation. We present the constraints from rare B decays on the coefficients of dimension-four operators in Sec. III. In Sec. IV we discuss the possible effects in rare B decays from all possible dimension-five operators involving the third generation quarks. Finally, we discuss the results and conclude in Sec. V.

II. THE EFFECTIVE THEORY

If the Higgs boson, responsible for the electroweak symmetry breaking, is very heavy, it can be effectively removed from the physical low-energy spectrum. In this case and for dynamical symmetry breaking scenarios relying on new strong interactions, one is led to consider the most general effective Lagrangian which employs a nonlinear representation of the spontaneously broken $SU(2)_L\times U(1)_Y$ gauge symmetry [10]. The resulting chiral Lagrangian is a nonrenormalizable nonlinear σ model coupled in a gauge-invariant way to the Yang-Mills theory. This model-independent approach incorporates by construction the low-energy theorems [11] that predict the general behavior of Goldstone boson amplitudes, irrespective of the details of the symmetry breaking mechanism. Unitarity requires that this low-energy effective theory should be valid up to some energy scale smaller than $4\pi v\simeq 3$ TeV, where new physics would come into play.

In order to specify the effective Lagrangian for the Goldstone bosons, we assume that the symmetry breaking pattern is $G=SU(2)_L\times U(1)_Y\to H=U(1)_{\rm em}$, leading to just three Goldstone bosons π^a $(a=1,2,3)$. With this choice, the building block of the chiral Lagrangian is the dimensionless unimodular matrix field Σ,

$$\Sigma=\exp\left(i\frac{\pi^a\tau^a}{v}\right), \tag{1}$$

where τ^a $(a=1,2,3)$ are the Pauli matrices. We implement the $SU(2)_C$ custodial symmetry by imposing a unique dimensionful parameter, v, for charged and neutral fields. Under the action of G the transformation of Σ is

$$\Sigma\to\Sigma'=L\Sigma R^\dagger,$$

where $L=\exp(i\alpha^a\tau^a/2)$ and $R=\exp(iy\tau^3/2)$, with α^a and y being the parameters of the transformation.

The gauge fields are represented by the matrices $\hat W_\mu=\tau^a W^a_\mu/(2i)$, $\hat B_\mu=\tau^3 B_\mu/(2i)$, while the associated field strengths are given by

$$\hat W_{\mu\nu}=\partial_\mu\hat W_\nu-\partial_\nu\hat W_\mu-g[\hat W_\mu,\hat W_\nu],$$

$$\hat B_{\mu\nu}=\partial_\mu\hat B_\nu-\partial_\nu\hat B_\mu.$$

In the nonlinear representation of the gauge group $SU(2)_L\times U(1)_Y$, the mass term for the vector bosons is given by the lowest order operator involving the matrix Σ. Therefore, the kinetic Lagrangian for the gauge bosons reads

$$\mathcal{L}_B=\frac{1}{2}{\rm Tr}(\hat W_{\mu\nu}\hat W^{\mu\nu}+\hat B_{\mu\nu}\hat B^{\mu\nu})+\frac{v^2}{4}{\rm Tr}(D_\mu\Sigma^\dagger D^\mu\Sigma), \tag{2}$$

where the covariant derivative of the field Σ is $D_\mu\Sigma=\partial_\mu\Sigma-g\hat W_\mu\Sigma+g'\Sigma\hat B_\mu$.

The effects of new dynamics on the couplings of fermions with the SM gauge bosons can be, in principle, also studied in an effective Lagrangian approach. For instance, if in analogy with the situation in QCD, fermion masses are dynamically generated in association with EWSB, residual interactions of fermions with Goldstone bosons could be important [4] if the $m_f\simeq f_\pi\simeq v$. Thus residual, non-universal interactions of the third generation quarks with gauge bosons could carry interesting information about both the origin of the top quark mass and EWSB.

In order to include fermions in this framework, we must define their transformation under G. Following Ref. [4], we postulate that matter fields feel directly only the electromagnetic interaction $f\to f'=\exp(iyQ_f)f$, where Q_f stands for the electric charge of fermion f. The usual left-handed fermion doublets are then defined with the following transformation under G:

$$\Psi_L=\Sigma\begin{pmatrix}f_1\\ f_2\end{pmatrix}_L\to\Psi'_L=L\ \exp(iyY/2)\Psi_L, \tag{3}$$

where $Q_{f_1}-Q_{f_2}=1$ and $Y=2Q_{f_1}-1$. Right-handed fermions are just the singlets f_R. This definition is useful since it permits the construction of linearly realized left-handed doublet fields in the same way that, when studying the breaking of $SU(2)_R\times SU(2)_L\to SU(2)_{R+L}$ in QCD, one introduces auxiliary fields for the nucleons which transform linearly under the broken axial group. In this framework, the lowest-order interactions between fermions and vector bosons that can be built are of dimension four, leading to anomalous vector and axial-vector couplings, which were analyzed in detail in Ref. [7].

In order to construct the most general Lagrangian describing these interactions, it is convenient to define the vector and tensor fields

$$\Sigma^a_\mu=-\frac{i}{2}{\rm Tr}(\tau^a V^R_\mu)=-\frac{i}{2}{\rm Tr}(\tau^a\Sigma^\dagger D_\mu\Sigma),$$

$$\Sigma^a_{\mu\nu}=-i\ {\rm Tr}[\tau^a\Sigma^\dagger[D_\mu,D_\nu]\Sigma]. \tag{4}$$

Under G transformations Σ^3_μ and $\Sigma^3_{\mu\nu}$ are invariant while

$$\Sigma^{\pm}_{\mu(\mu\nu)} \to \Sigma'^{\pm}_{\mu(\mu\nu)} = \exp(\pm iy)\Sigma^{\pm}_{\mu(\mu\nu)},$$

where $\Sigma^{\pm}_{\mu(\mu\nu)} = (\Sigma^1_{\mu(\mu\nu)} \mp i\Sigma^2_{\mu(\mu\nu)})/\sqrt{2}$.

The basic fermionic elements for the construction of neutral- and charged-current effective interactions are

$$\begin{aligned}
\Delta_X(q,q') &= \bar{q} P_X q', \\
\Delta^\mu_X(q,q') &= \bar{q}\gamma^\mu P_X q', \\
\tilde{\Delta}^\mu_X(q,q') &= \bar{q} P_X \tilde{D}^\mu q', \\
\Delta^{\mu\nu}_X(q,q') &= \bar{q}\sigma^{\mu\nu} P_X q',
\end{aligned} \tag{5}$$

where P_X ($X=0$, 5, L, and R) stands for I, γ^5, P_L, and P_R respectively, with I being the identity matrix and $P_{L(R)}$ the left (right) chiral projector. The fermionic field q (q') represents any quark flavor, and $\tilde{D}^\mu$ stands for the electromagnetic covariant derivative.

The most general dimension-four Lagrangian invariant under nonlinear transformations under G is

$$\begin{aligned}
\mathcal{L}_4 = {} & d^{NC}_L \Delta^\mu_L(t,t)\Sigma^3_\mu + d^{NC}_R \Delta^\mu_R(t,t)\Sigma^3_\mu + d^{CC}_L \Delta^\mu_L(t,b)\Sigma^+_\mu \\
& + d^{CC\dagger}_L \Delta^\mu_L(b,t)\Sigma^-_\mu + d^{CC}_R \Delta^\mu_R(t,b)\Sigma^+_\mu \\
& + d^{CC\dagger}_R \Delta^\mu_R(b,t)\Sigma^-_\mu .
\end{aligned} \tag{6}$$

In principle it is also possible to construct neutral current operators involving only the bottom quark. We will assume however, that these vertices are not modified by the dynamics of the symmetry breaking or, at most, that these modifications are suppressed as compared to those of the top quark.

In a very general parametrization, the dimension-four anomalous couplings of third generation quarks can be written in terms of the usual physical fields as

$$\begin{aligned}
\mathcal{L}_4 = {} & -\frac{g}{\sqrt{2}}[C_L(\bar{t}_L\gamma_\mu b_L) + C_R(\bar{t}_R\gamma_\mu b_R)]W^{+\mu} \\
& - \frac{g}{2c_W}[N^t_L(\bar{t}_L\gamma_\mu t_L) + N^t_R(\bar{t}_R\gamma_\mu t_R)]Z^\mu + \text{H.c.},
\end{aligned} \tag{7}$$

where s_W (c_W) is the sine (cosine) of the weak mixing angle, θ_W. The parameters $C_{L,R}$, $N^t_{L,R}$ can be written in terms of the constants $d^{NC,CC}_{L,R}$ of Eq. (6) and contain the residual, non-universal effects associated with the new dynamics, perhaps responsible for the large top quark mass. Then, if we assume that the new couplings are CP conserving [12], there are four new parameters. They are constrained at low energies by a variety of experimental information, mostly from electroweak precision measurements and the rate of $b \to s\gamma$.

In the case of dimension-five operators, the most general neutral-current interactions, which are invariant under nonlinear transformations under G, are [13]

$$\begin{aligned}
\mathcal{L}^{NC}_5 = {} & a^{NC}_1 \Delta_0(t,t)\Sigma^+_\mu \Sigma^{-\mu} + a^{NC}_2 \Delta_0(t,t)\Sigma^3_\mu \Sigma^{3\mu} + i a^{NC}_3 \Delta_5(t,t)\partial^\mu \Sigma^3_\mu + i b^{NC}_1 \Delta^{\mu\nu}_0(t,t)\mathrm{Tr}[T\hat{W}_{\mu\nu}] + b^{NC}_2 \Delta^{\mu\nu}_0(t,t)B_{\mu\nu} \\
& + i b^{NC}_3 \Delta^{\mu\nu}_0(t,t)(\Sigma^+_\mu \Sigma^-_\nu - \Sigma^+_\nu \Sigma^-_\mu) + i c^{NC}_1 [\tilde{\Delta}^\mu_0(t,t) - \overline{\tilde{\Delta}^\mu_0(t,t)}]\Sigma^{3\mu},
\end{aligned} \tag{8}$$

and the charged-current interactions are

$$\begin{aligned}
\mathcal{L}^{CC}_5 = {} & a^{CC}_{1L} \Delta_L(t,b)\Sigma^+_\mu \Sigma^{3\mu} + a^{CC}_{1R} \Delta_R(t,b)\Sigma^+_\mu \Sigma^{3\mu} + i a^{CC}_{2L} \Delta_L(t,b)\tilde{D}^\mu \Sigma^+_\mu + i a^{CC}_{2R} \Delta_R(t,b)\tilde{D}^\mu \Sigma^+_\mu + b^{CC}_{1L} \Delta^{\mu\nu}_L(t,b)\Sigma^+_{\mu\nu} \\
& + b^{CC}_{1R} \Delta^{\mu\nu}_R(t,b)\Sigma^+_{\mu\nu} + b^{CC}_{2L} \Delta^{\mu\nu}_L(t,b)(\Sigma^+_\mu \Sigma^3_\nu - \Sigma^+_\nu \Sigma^3_\mu) + b^{CC}_{2R} \Delta^{\mu\nu}_R(t,b)(\Sigma^+_\mu \Sigma^3_\nu - \Sigma^+_\nu \Sigma^3_\mu) + i c^{CC}_{1L} \tilde{\Delta}^\mu_L(t,b)\Sigma^+_\mu \\
& + i c^{CC}_{1R} \tilde{\Delta}^\mu_R(t,b)\Sigma^+_\mu + \text{H.c.}
\end{aligned} \tag{9}$$

In general, since chiral Lagrangians are related to strongly interacting theories, it is hard to make firm statements about the expected magnitude of the couplings. Requiring the loop corrections to the effective operators to be of the same order of the operators themselves suggests that these coefficients are of $\mathcal{O}(1)$ [14]. However, this is not necessarily the case when the operators result from small symmetry breaking effects, as we will see below.

In the unitary gauge, we can rewrite the interactions (8) and (9) as a scalar, a vector, and a tensorial Lagrangian involving the physical fields. For the Lagrangian involving scalar currents we have

$$\begin{aligned}
\mathcal{L}_S = {} & \frac{g^2}{2\Lambda}\left[\bar{t}t\left(\alpha^{NC}_1 W^+_\mu W^{-\mu} + \frac{\alpha^{NC}_2}{2c^2_W} Z^\mu Z_\mu\right)\right] + i\frac{g}{2c_W\Lambda}\alpha^{NC}_3 \bar{t}\gamma^5 t\partial^\mu Z_\mu + \frac{g^2}{2\sqrt{2}\Lambda c_W}\{\bar{t}[\alpha^{CC}_{1L}(1-\gamma^5) + \alpha^{CC}_{1R}(1+\gamma^5)]bW^+_\mu Z^\mu \\
& + \bar{b}[\alpha^{CC}_{1L}(1+\gamma^5) + \alpha^{CC}_{1R}(1-\gamma^5)]tW^-_\mu Z^\mu\} + i\frac{g}{2\sqrt{2}\Lambda}\{\bar{t}[\alpha^{CC}_{2L}(1-\gamma^5) + \alpha^{CC}_{2R}(1+\gamma^5)]b(\partial^\mu W^+_\mu + ieA^\mu W^+_\mu) \\
& - \bar{b}[\alpha^{CC}_{2L}(1+\gamma^5) + \alpha^{CC}_{2R}(1-\gamma^5)]t(\partial^\mu W^-_\mu - ieA^\mu W^-_\mu)\}.
\end{aligned} \tag{10}$$

The Lagrangian containing vectorial current is given by

$$\mathcal{L}_V=\frac{1}{2\Lambda}\left\{ i\frac{g}{c_W}\gamma^{NC}\overline{t}(\tilde{D}_\mu t)Z^\mu - i\frac{g}{c_W}\gamma^{NC}\overline{(\tilde{D}_\mu t)}tZ^\mu + i\frac{g}{\sqrt{2}}\overline{t}[\gamma_L^{CC}(1-\gamma^5)+\gamma_R^{CC}(1+\gamma^5)](\tilde{D}_\mu b)W^{+\mu} - i\frac{g}{2c_W}\overline{(\tilde{D}_\mu b)} \times[\gamma_L^{CC}(1+\gamma^5)+\gamma_R^{CC}(1-\gamma^5)]tW^{-\mu}\right\}. \tag{11}$$

Finally, the piece involving a tensorial structure is

$$\mathcal{L}_T=\frac{1}{4\Lambda}\left[\overline{t}\sigma^{\mu\nu}t\left(\beta_1^{NC}eF_{\mu\nu}+\beta_2^{NC}\frac{g}{c_W}Z_{\mu\nu}+4ig^2\beta_3^{NC}W_\mu^+W_\nu^-\right)\right]+\frac{g}{2\sqrt{2}\Lambda}\left\{\overline{t}\sigma^{\mu\nu}[\beta_{1L}^{CC}(1-\gamma^5)+\beta_{1R}^{CC}(1+\gamma^5)]b[W_{\mu\nu}^+ + ie \times(A_\mu W_\nu^+ - A_\nu W_\mu^+)]+\overline{b}\sigma^{\mu\nu}[\beta_{1L}^{CC}(1+\gamma^5)+\beta_{1R}^{CC}(1-\gamma^5)]t[W_{\mu\nu}^- - ie(A_\mu W_\nu^- - A_\nu W_\mu^-)]+i\frac{g}{c_W}\overline{t}\sigma^{\mu\nu}[\beta_{2L}^{CC}(1-\gamma^5) + \beta_{2R}^{CC}(1+\gamma^5)]b(Z_\mu W_\nu^+ - Z_\nu W_\mu^+) - i\frac{g}{c_W}\overline{b}\sigma^{\mu\nu}[\beta_{2L}^{CC}(1+\gamma^5)+\beta_{2R}^{CC}(1-\gamma^5)]t(Z_\mu W_\nu^- - Z_\nu W_\mu^-)\right\}. \tag{12}$$

The couplings constants α's, β's and γ's are linear combinations of the a's, b's and c's in Eqs. (8) and (9). In writing the interactions (10) and (12), the coupling constants were defined in such a way that we have a factor $g/(2c_W)$ per Z boson, $g/\sqrt{2}$ per $W^\pm$, and a factor e per photon. Similar interactions were obtained in Ref. [13] and for a linearly realized symmetry group, in Ref. [15].

As an example of the above anomalous couplings, we show their couplings for the SM with a heavy Higgs boson integrated out. In this case, we can perform the matching between the full theory and the effective Lagrangian [16]. For instance, if we concentrate on the non-decoupling effects, the leading contributions come at one-loop order [17]. Setting $m_b=0$ and keeping only the leading terms of the order $m_t\log(M_H^2)$, we find that only the first two effective operators of Eq. (10) are generated with coefficients,

$$\alpha_1^{NC}=\alpha_2^{NC}=\frac{g^2m_t\Lambda}{16\pi^2M_W^2}\log\frac{M_H^2}{m_t^2}.$$

III. RESULTS FOR THE $b\to s\gamma$ AND $b\to sl^+l^-$ TRANSITIONS

For the $b\to s\gamma$ and $b\to sl^+l^-$ transitions it is useful to cast the contributions of the dimension-four and dimension-five anomalous couplings as shifts in the matching conditions at M_W for the Wilson coefficient functions in the weak effective Hamiltonian

$$H_{\rm eff}=-\frac{4G_F}{\sqrt{2}}V_{tb}^*V_{ts}\sum_{i=1}^{10}C_i(\mu)O_i(\mu), \tag{13}$$

with the operator basis defined in Ref. [18]. Of interest in our analysis are the electromagnetic penguin operator

$$\mathcal{O}_7=\frac{e}{16\pi^2}m_b(\overline{s_L}\sigma_{\mu\nu}b_R)F^{\mu\nu}, \tag{14}$$

and the four-fermion operators corresponding to the vector and axial-vector couplings to leptons,

$$\mathcal{O}_9=\frac{e^2}{16\pi^2}(\overline{s_L}\gamma_\mu b_L)(\overline{l}\gamma^\mu l), \tag{15}$$

and

$$\mathcal{O}_{10}=\frac{e^2}{16\pi^2}(\overline{s_L}\gamma_\mu b_L)(\overline{l}\gamma^\mu\gamma_5 l). \tag{16}$$

The operators above are already present in the SM. In addition, the dimension-five anomalous couplings generate the operators

$$\mathcal{O}_{11}=\frac{e^2}{16\pi^2}\frac{m_b}{M_Z^2}[\overline{s_L}\sigma_{\mu\nu}(iQ^\nu)b_R](\overline{l}\gamma^\mu\gamma_5 l), \tag{17}$$

$$\mathcal{O}_{12}=\frac{e^2}{16\pi^2}\frac{m_b}{M_Z^2}[\overline{s_L}\sigma_{\mu\nu}(iQ^\nu)b_R](\overline{l}\gamma^\mu l). \tag{18}$$

However, these new operators will not lead to important effects as will see below, due to the fact that they are further suppressed by the weak scale.

The anomalous couplings of Eq. (7), (10), (11) and (12) will induce shifts in the Wilson coefficient functions $C_i(\mu)$ at the matching scale, which we take to be $\mu=M_W$. We make use of the next-to-leading order calculation of the Wilson coefficients as described in Ref. [19].

A. Effects of the dimension-four operators

The dimension-four operators defined in Eq. (7) induce new contributions to the $b\to s\gamma$ and $b\to sZ$ loops as well as the box diagram. They appear in the effective Hamiltonian formulation as shifts of the Wilson coefficients $C_7(M_W)$, $C_9(M_W)$ and $C_{10}(M_W)$. The contributions from the $b\to s\gamma$ loops to $C_7(M_W)$ are

$$\delta C_7 = -\frac{m_t}{m_b} C_R \left\{ \frac{1}{12(x-1)^2}(5x^2-31x+20) + \frac{1}{2(x-1)^3}x(3x-2)\log(x) \right\} + C_L \left\{ \frac{1}{4(x-1)^4}x^2 \times(3x-2)\log(x) - \frac{1}{24(x-1)^3}x(8x^2+5x-7) \right\}, \tag{19}$$

where we have defined the dimensionless quantity $x = m_t^2/M_W^2$. We should notice that the above result is finite, i.e. independent of Λ, and agrees with the previous result in the literature [8]. On the other hand, the result for all other operators is not finite and in this case we have kept only the leading non-analytic, i.e. logarithmic, dependence on the new physics scale Λ. In this way, for $C_9(M_W)$ we have

$$\delta C_9^\gamma = -\frac{1}{12} C_L (3x-16)\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{20}$$

The corrections arising from $b\to sZ$ loops to $C_9(M_W)$ and $C_{10}(M_W)$ are

$$\delta C_{10}^Z = \frac{-1}{1-4s_W^2}\,\delta C_9^Z = \frac{1}{16 s_W^2}(4N_L^t - N_R^t + C_L)x\log\left(\frac{\Lambda^2}{M_W^2}\right), \tag{21}$$

while box loops contributions can be written as

$$\delta C_9^{\rm box} = -\,\delta C_{10}^{\rm box} = \frac{1}{16 s_W^2} C_L (x-16)\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{22}$$

The measured $b\to s\gamma$ branching ratio imposes a stringent bound on C_R as its contribution to Eq. (19) is enhanced by the factor m_t/m_b. This has been discussed in the literature [8], where the obtained bounds on C_R: $-0.05 < C_R < 0.01$. Although this is very small, it is possible to generate such value for C_R in a large variety of generic strongly coupled theories. For instance, the pseudo Nambu-Goldstone bosons (PNGBs) of extended technicolor (ETC) that result from the breaking of the various fermion chiral symmetries [20], generate at one loop a small C_R proportional to m_b:

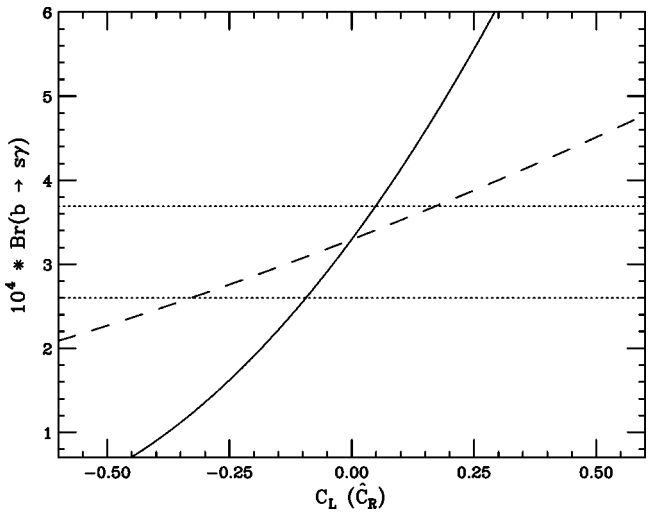

FIG. 1. The $b\to s\gamma$ branching ratio vs $\hat{C}_R$ (solid) and C_L (dashed). The horizontal band corresponds to the 1σ interval from the latest CLEO result [23].

$$C_R \simeq \frac{1}{4\pi}\frac{m_b m_t}{f_\pi^2}\log\left(\frac{m_\pi^2}{m_t^2}\right). \tag{23}$$

This is well within experimental bounds in all ETC incarnations, and it is even smaller in modern ETC theories such as top-color-assisted technicolor [21], where the top quark mass entering in Eq. (23) is only a few GeV [22]. Thus, C_R is a consequence of the explicit ETC-breaking of chiral symmetry responsible for m_b. Another hint of this, is the fact that in general C_R contributes to the renormalization of the b-quark line with a term which does not vanish with m_b:

$$\Sigma(m_b) = \frac{g^2}{32\pi^2} C_R\, m_t\,(x-4)\,\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{24}$$

Thus if we take into account the potential role of chiral symmetry in suppressing C_R and we rescale this coefficient by defining $\hat{C}_R$ as

$$C_R = \frac{m_b}{\sqrt{2}v}\hat{C}_R \tag{25}$$

(where $v = 246$ GeV), the rescaled bounds on $\hat{C}_R$ are $\mathcal{O}(1)$, which is a more natural value of this coefficient once the underlying symmetry is taken into account.

In Fig. 1 we plot the $b\to s\gamma$ branching fraction as a function of $\hat{C}_R$. We also include the effect of C_L, which is now comparable for similar values of the coefficients. The horizontal band corresponds to the latest CLEO result [23] $Br(b\to s\gamma) = (3.15 \pm 0.35 \pm 0.32 \pm 0.26)\times 10^{-4}$, where we take a 1σ interval after adding the statistical, systematic and model-dependent errors in quadrature.

On the other hand, the effect in $b\to s l^+ l^-$ is dominated by the coefficients C_L, N_L^t and N_R^t in Eqs. (20), (21), and

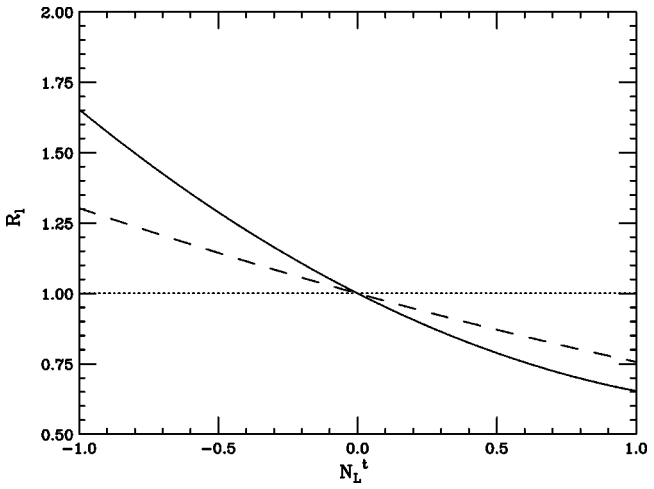

FIG. 2. The $b\to sl^+l^-$ branching ratio (normalized to the SM) vs $C_L=N_L^t$ (solid) and $N_L^t=N_R^t$ (dashed).

(22). In principle, these coefficients are constrained by electroweak precision measurements, most notably $\epsilon_1=\Delta\rho=\alpha T$ and ϵ_b [7]:

$$\epsilon_1=\frac{G_F}{2\sqrt{2}\pi^2}3m_t^2(-N_L^t+N_R^t+C_L)\log\left(\frac{\Lambda^2}{M_W^2}\right)$$

$$\epsilon_b=\frac{G_F}{2\sqrt{2}\pi^2}3m_t^2\left(-\frac{1}{4}N_R^t+N_L^t\right)\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{26}$$

In general, the bounds obtained on a particular coupling from electroweak observables strongly depend on assumptions about the other couplings. For instance, enforcing custodial isospin symmetry in order to avoid the strong constraints from T will imply that $N_L^t=C_L$ and $N_R^t=0$. On the other hand if $C_L=0$, then the combination $(N_L^t-N_R^t)$ is strongly constrained since it breaks custodial isospin symmetry and contributes to T. Imposing $C_L=N_L^t$, then $N_R^t<0.02$ [4,24] since it is the only linear source of isospin breaking.

We study here three cases in which the stringent constraints from electroweak observables can be evaded.

(i) $C_L\approx N_L^t$. In this case the contributions of C_L and N_L^t to the T parameter cancel, leaving N_R^t as the only seriously constrained quantity. However, R_b still gives the bound $-0.03<N_L^t<0.15$ (for $\Lambda=1$ TeV). In Fig. 2 we plot the $b\to sl^+l^-$ branching ratio, normalized to the SM expectation, as a function of $C_L=N_L^t$ (solid line). From this plot it can be seen that, when incorporating the R_b constraint, the effect in $b\to sl^+l^-$ is bound to be smaller than roughly a 10% deviation.

(ii) $N_L^t\approx N_R^t$. In this scenario the measurement of the T parameter greatly constrains C_L, which prompts us to take this coefficient as equal to zero in this portion of the analysis. The dashed line in Fig. 2 corresponds to the effect of $N_L^t=N_R^t$ in the dilepton branching fraction. As we can see, the effect in this decay is rather small.

(iii) Finally and for completeness, we consider the case $N_L^t\approx N_R^t/4$. With this approximate relation these two coefficients cancel in R_b leaving no sizeable effect from the dimension-four Lagrangian (7) in this quantity. However, this relation leads to a potentially large contribution to the T parameter proportional to $(C_L+3N_R^t/4)$. When this bound is incorporated, the effect in $b\to sl^+l^-$ branching ratio is constrained to be below 15%.

In sum, we have seen that the leading effects of the dimension−four operators in rare B decays are given by $\hat{C}_R$ in $b\to s\gamma$, and the effects in $b\to sl^+l^-$ due to C_L, N_L^t and N_R^t are below 15% deviations once the constraints from electroweak precision measurements are considered. This distinction comes from the fact that Z-pole quantities are not significantly sensitive to $\hat{C}_R$. The effects of $\hat{C}_R$ in $b\to sl^+l^-$ can be significant, but $b\to s\gamma$ is considerably more sensitive to this parameter.

B. Effects of the dimension-five operators

Although in principle dimension-five operators are considered sub-leading with respect to the operators in Eq. (7) due to the additional suppression by the high energy scale Λ, they can still induce large deviations in both electroweak observables and FCNC processes. In Ref. [9] bounds on the coefficients of dimension-five operators were derived from data at the Z pole. Here we consider the effect of these operators in $b\to s\gamma$ and $b\to sl^+l^-$. They induce new contributions to the $b\to s\gamma$ and $b\to sZ$ loops as well as the box diagram. They appear in the effective Hamiltonian formulation as shifts of the Wilson coefficients $C_7(M_W)$, $C_9(M_W)$ and $C_{10}(M_W)$, $C_{11}(M_W)$, and $C_{12}(M_W)$.

The contribution from the $b\to s\gamma$ loops to these coefficients are

$$\delta C_7=-\frac{1}{12}\frac{M_W^2}{m_b\Lambda}[4\alpha_{2R}^{CC}x-4\beta_{1R}^{CC}(3x-7)+\gamma_R^{CC}(x+2)]\log\left(\frac{\Lambda^2}{M_W^2}\right), \tag{27}$$

and

$$\delta C_{12}^{\gamma}=\frac{1}{24}\frac{M_Z^2}{m_b\Lambda}\{[3(2\alpha_{2R}^{CC}-\gamma_R^{CC})x+4(3\alpha_{2R}^{CC}-6\beta_{1R}^{CC}+\gamma_R^{CC})]\}\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{28}$$

The corrections arising from $b\to sZ$ loops to the different coefficients are

$$\delta C_{10}^{Z}=\frac{-1}{1-4s_W^2}\,\delta C_9^{Z}=-\frac{1}{96s_W^2}\frac{m_t}{\Lambda}[3(2\alpha_{1L}^{CC}+12\alpha_{2L}^{CC}+12\beta_{2L}^{CC}-\gamma_L^{CC}+8\gamma^{NC})x-18c_W^2(\alpha_{2L}^{CC}+2\beta_{1L}^{CC})x$$
$$+2s_W^2(\gamma_L^{CC}-12\alpha_{2L}^{CC})x-9(2\alpha_{1L}^{CC}-4\beta_{1L}^{CC}+12\beta_{2L}^{CC}-\gamma_L^{CC}-4\gamma^{NC})-18c_W^2(\alpha_{2L}^{CC}-2\beta_{1L}^{CC})$$
$$-6s_W^2(\gamma_L^{CC}+12\beta_{1L}^{CC})]\log\left(\frac{\Lambda^2}{M_W^2}\right), \tag{29}$$

and

$$\delta C_{11}^{Z}=\frac{-1}{1-4s_W^2}\,\delta C_{12}^{Z}=-\frac{1}{48s_W^2}\frac{M_Z^2}{m_b\Lambda}[3(\alpha_{1R}^{CC}+\alpha_{2R}^{CC}-2\beta_{2R}^{CC})x+3c_W^2(\alpha_{2R}^{CC}-6\beta_{1R}^{CC}+\gamma_R^{CC})x-2s_W^2(4\alpha_{2R}^{CC}-\gamma_R^{CC})x$$
$$+3(2\alpha_{1R}^{CC}+8\beta_{1R}^{CC}-4\beta_{2R}^{CC}-\gamma_R^{CC})+6c_W^2(\alpha_{2R}^{CC}+6\beta_{1R}^{CC}+\gamma_R^{CC})-4s_W^2(8\beta_{1R}^{CC}-\gamma_R^{CC})]\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{30}$$

The box loops contributions can be written as

$$\delta C_9^{\rm box}=-\delta C_{10}^{\rm box}=-\frac{1}{96s_W^2}\frac{m_t}{\Lambda}[4\alpha_{2L}^{CC}(2x-5)-36\beta_{1L}^{CC}-\gamma_L^{CC}(x-1)]\log\left(\frac{\Lambda^2}{M_W^2}\right), \tag{31}$$

and

$$\delta C_{12}^{\rm box}=-\delta C_{11}^{\rm box}=-\frac{1}{16s_W^2}\frac{M_Z^2}{m_b\Lambda}[2\alpha_{2R}^{CC}x-4\beta_{1R}^{CC}-\gamma_R^{CC}(x-2)]\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{32}$$

In order to quantify the effect of the operators of Eqs. (10), (11) and (12) we classify them into two groups: the right-handed and left-handed couplings. The RH couplings that are of interest due to their potential effects are α_{2R}^{CC}, β_{1R}^{CC} and γ_R^{CC}. Just as for the dimension-four coefficient C_R, the effects of these coefficients can be very large in operators such as $\mathcal{O}_7$ generating important deviations in $b\to s\gamma$, as seen in Eq. (27). This is particularly so since their contributions appear to be unsuppressed by m_b. As we noted for the dimension-four coefficient C_R, the coefficients α_{2R}^{CC} and β_{1R}^{CC} are chirally unsuppressed in Eq. (10), which results in an renormalization of the b-quark mass. Although, unlike C_R, the dimension-five coefficients are not tightly bound by current experiments, we could also build strongly coupled theories where these are small due to chiral symmetry. If this suppression is present, it would make their effect on the operator O_7 negligible. On the other hand, the coefficient γ_R^{CC} has a chiral suppression already present in the accompanying operator in Eq. (11). Because of this its contribution to the renormalization of the b-quark line vanishes with m_b:

$$\Sigma(m_b)=\frac{-3g^2}{128\pi^2\Lambda}\gamma_R^{CC}m_b^2(x+1)\log\left(\frac{\Lambda^2}{M_W^2}\right). \tag{33}$$

Thus it is possible that some of the RH coefficients are naturally suppressed while others may not be. We will illustrate the effects of the RH couplings using γ_R^{CC} as an example. The effects of α_{2R}^{CC} and β_{1R}^{CC} would be of similar size if they are not chirally suppressed. The contribution of γ_R^{CC} to the penguin operator gives rise to a deviation of the $b\to s\gamma$ branching ratio from its SM expectation. In Fig. 3 we plot this branching fraction as a function of this coefficient. This measurement is the most constraining bound on these type of operators. It can be seen that even for rather small values of γ_R^{CC} there could be considerable deviations from the SM expectations. On the other hand, the effect is less dramatic in $b\to s l^+ l^-$, as shown in Fig. 4, where an observable devia-

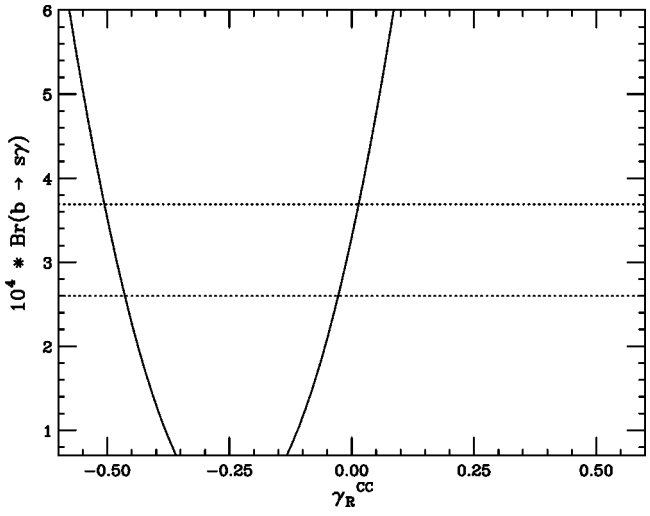

FIG. 3. The $b\to s\gamma$ branching ratio vs γ_R^{CC}.

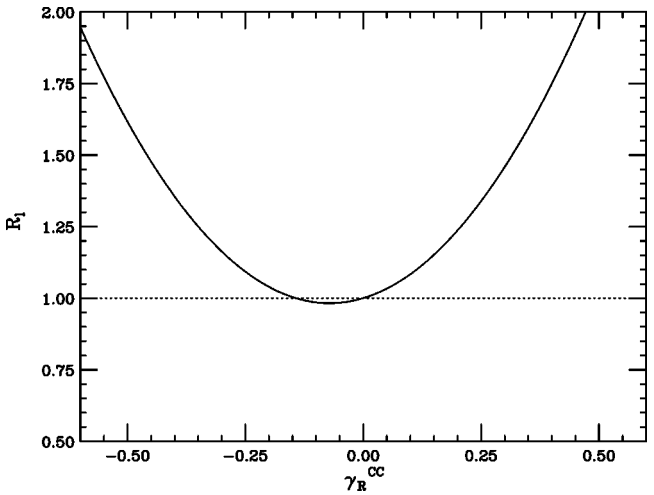

FIG. 4. The $b\to sl^+l^-$ branching ratio (normalized to the SM prediction) vs γ_R^{CC}.

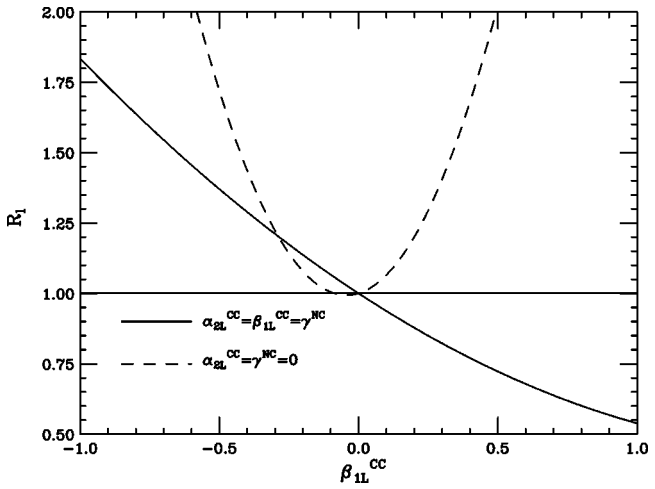

FIG. 5. The $b\to sl^+l^-$ branching ratio (normalized to the SM prediction) vs β_{1L}^{CC}, for $\alpha_{2L}^{CC}=\beta_{1L}^{CC}=\gamma^{NC}$ (solid line) and $\alpha_{2L}^{CC}=\gamma^{NC}=0$ (dashed line).

tion from the SM will result only if γ_R^{CC} is large enough to dominate the $b\to s\gamma$ branching ratio.

The effects of the new operators $\mathcal{O}_{11}$ and $\mathcal{O}_{12}$ are negligible. Although the presence of m_b in the denominators in Eqs. (28), (30) and (32) suggests the possibility of an enhancement, this is not enough. The effect is suppressed by an effective scale given by $m_b\Lambda/M_Z\simeq 55$ GeV, which should be compared with the typical momentum transfers in B decays.

The remaining group of coefficients we dubbed left-handed includes α_{1L}^{CC}, α_{2L}^{CC}, β_{1L}^{CC}, β_{2L}^{CC}, γ_L^{CC} plus γ^{NC} which actually is the coefficient of a vector operator, but since is not chirally suppressed is included with the LH in this part of the analysis. These operators affect mainly the $b\to sl^+l^-$ rates. Then we could imagine that the underlying new physics preserves chiral symmetry in such a way that does not generate sizeable contributions to the RH operators. *Thus it is possible a scenario where no deviations in* $b\to s\gamma$ *occur; but where the new interactions give rise to large effects in the dilepton modes*. Although these have not been observed yet, the current experimental sensitivity is very close to the SM predictions and it will reach them in the near future. The leading effects in $b\to sl^+l^-$ come from the coefficients α_{2L}^{CC}, β_{1L}^{CC} and γ^{NC}. For simplicity we only consider these and plot in Fig. 5 the branching ratio, normalized to the SM one, for two cases: $\alpha_{2L}^{CC}=\beta_{1L}^{CC}=\gamma^{NC}$ (solid line) and $\alpha_{2L}^{CC}=\gamma^{NC}=0$ (dashed line). From Fig. 5 it is apparent that cancellations occur when the three coefficients are similar. The effect of considering only β_{1L}^{CC} shows than even larger effects are possible. In any case, sizeable deviations in $b\to sl^+l^-$ are possible even in the absence of effects in $b\to s\gamma$.

IV. DISCUSSION

Processes involving FCNC transitions in B and K decays are a crucial complement to precision electroweak observables, when constraining the physics of the EWSB sector. In this paper, we have considered the effects of anomalous couplings of third generation quarks to the W and Z gauge bosons. We computed the effects of all possible dimension-five operators in B FCNC transitions such as $b\to s\gamma$ and $b\to sl^+l^-$. For completeness, we have also also presented the analysis of the dimension-four operators.

We have shown that with the natural assumption of chiral symmetry, in fact enforcing vanishing fermion mass renormalization in the chiral limit, the effects of the dimension-four operators with coefficient C_R for $b\to s\gamma$, are not as dramatic as found in Ref. [8], and somehow smaller than those of C_L, which can produce important deviations in the branching ratio that could be resolved in the next round of experiments at B factories.

The effects in $b\to sl^+l^+$ due to C_L, N_L^t and N_R^t are below 15% deviations once the constraints from electroweak precision measurements are taken into account.

On the other hand, we have found that the dimension-five operators with RH coefficients can give rise to observable deviations of the $b\to s\gamma$ branching ratio from its SM expectation. We illustrated these effects for the coefficient γ_R^{CC}, for which no additional chiral suppression is expected. Left handed operators, on the other hand, affect mainly the $b\to sl^+l^-$ rates and we have illustrated that in several scenarios sizeable deviations in $b\to sl^+l^-$ are possible *even in the absence of effects in* $b\to s\gamma$.

The dimension-five operators, just as in the case of the more studied dimension-four operators, can be generated at high energies scales by the presence of new particles and/or interactions. For instance, as a simple example, a heavy scalar sector with both charged and neutral states, would give contributions to many of the coefficients of the Lagrangian in Eqs. (10). Richer dynamics at the TeV scale might generate also some of the vector and/or tensor couplings of Eqs. (11) and (12).

The e^+e^- B factories at Cornell, KEK and SLAC are expected to reach better measurements of the $b\to s\gamma$ branching ratio, which will largely constrain the dimension-five coefficient γ_R^{CC} and to a lesser extent the dimension-four coefficients C_L and $\hat{C}_R$. Furthermore, these experiments as well

as those at the Fermilab Tevatron, will reach the SM sensitivity for the $b \to s l^+ l^-$ branching fraction. The present analysis, together with previous ones addressing the effects of other anomalous higher dimensional operators in these decay modes, will enable us to interpret a possible pattern of deviations from the SM and perhaps point to its dynamical origin.

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