SU(2)×U(1) Gauge Invariance and the Shape of New Physics in Rare B Decays

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Phys. Rev. Lett. **113**, 241802 — Published 9 December 2014

DOI: 10.1103/PhysRevLett.113.241802

$SU(2) \times U(1)$ gauge invariance and the shape of new physics in rare B decays

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New physics effects in B decays are routinely modeled through operators invariant under the strong and electromagnetic gauge symmetries. Assuming the scale for new physics is well above the electro-weak scale, we further require invariance under the full Standard-Model gauge symmetry group. Retaining up to dimension-six operators, we unveil new constraints between different new-physics operators that are assumed to be independent in the standard phenomenological analyses. We illustrate this approach by analyzing the constraints on new physics from rare B_q (semi-)leptonic decays.

Introduction. The exploration of the energy regime of electro-weak symmetry breaking (EWSB) at the Large Hadron Collider (LHC) has unveiled a scalar boson [1, 2] resembling the Standard Model (SM) singlet component of the Higgs doublet and no other particle. If one therefore assumes, as the experimental evidence suggests, that the scale of new physics (NP), Λ, is above the EWSB scale an effective field theory (EFT) built exclusively from SM fields can be used. In this widespread and fruitful scheme, higher-dimension operators suppressed by powers of the NP scale encode deviations of the SM in a generic and model-independent manner [3, 4]. The only requirements imposed on the operators are Lorentz and $SU(3)_c \times SU(2)_L \times U(1)_Y$ gauge symmetries.

These simple assumptions lead, as this letter is meant to show, to phenomenological consequences not only for physics at the EWSB scale but also for physics well below such scale. Furthermore, the consequences not only affect the "size" of the contribution of new physics to low energy processes, but also the "shape" or correlation among different operators. These extra constraints in the low energy Lagrangian are due simply to $SU(3)_c \times SU(2)_L \times U(1)_Y$ invariance. More specifically, there are three important ways in which the low energy EFT is further constrained:
(i) The operators must originate in those of an EFT with explicit electroweak symmetry;
(ii) The coefficients of operators are not all independent, as they may be related by their origin in the underlying spontaneously broken electroweak group; and
(iii) Some of the coefficients of the low energy EFT may be constrained by seemingly unrelated high energy processes.
The latter occurs, for example, when the low energy operator arises from integrating out a heavy field, like the Higgs, from an operator which itself produces effects observable in the decay of the heavy field.

To illustrate the aforementioned effects we consider rare, flavor changing-neutral (FCN) B-meson semi- or purely-leptonic decays, where, to our knowledge, such an analysis has not yet been carried out fully. The reduced set of observables that will be studied here allow us to focus on a subgroup of operators rather than the most general EFT consistent with electroweak symmetry, which is left for future work. We shall distinguish between three different scales: i) the NP scale Λ, ii) the EWSB scale, $\langle H^\dagger H\rangle = v^2/2$, and iii) the low scale, μ, in this case of the order of the bottom quark mass. We assume the following hierarchy of mass scales $\mu \ll v \ll \Lambda$.

Low energy: B-meson semi-leptonic Lagrangian. At energies around the bottom quark mass, the EFT Lagrangian is built from the light fields: the SM particle content except the W and Z bosons, the Higgs boson, and the top quark. In addition the EFT Lagrangian respects the gauge symmetries manifest at this scale, namely $SU(3)_c \times U(1)_{em}$. As we will show, not all of the possible operators constructed in this way are compatible with an effective Lagrangian invariant under $SU(2)_L \times U(1)_Y$.

To leading order in $G_F = 1/(\sqrt{2}v^2)$ the effective Lagrangian for $\Delta B = 1$ processes is [5–7]

$$\mathcal{L}_{\rm eff} = -\frac{4G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_{ps} \left(C_1 \mathcal{O}_1^p + C_2 \mathcal{O}_2^p + \sum_{i=3}^{10} C_i \mathcal{O}_i \right), \tag{1}$$

with $\lambda_{ps} = V_{pb} V_{ps}^*$. The "current-current" operators, $\mathcal{O}_{1,2}^p$, "QCD-penguins," $\mathcal{O}_{3,\ldots,6}$, and "chromo-magnetic operator," $\mathcal{O}_8$, do not contribute to $B_s \to l\bar{l}$ and their contribution to $B \to K^{(*)} l\bar{l}$ requires an electromagnetic interaction (they contribute to the "non-factorizable" corrections, in the language of QCD factorization [8]). We therefore focus on the electromagnetic penguin, $\mathcal{O}_7$ and the semileptonic operators, $\mathcal{O}_{9,10}$, defined as

$$\mathcal{O}_7 = \frac{e}{(4\pi)^2} \overline{m}_b [\bar{s}\sigma^{\mu\nu} P_R\, b] F_{\mu\nu}, \tag{2}$$

$$\mathcal{O}_9 = \frac{e^2}{(4\pi)^2} [\bar{s}\gamma_\mu P_L b][\bar{l}\gamma^\mu l],\ \mathcal{O}_{10} = \frac{e^2}{(4\pi)^2} [\bar{s}\gamma_\mu P_L b][\bar{l}\gamma^\mu \gamma_5 l],$$

where b, s, l stand for the bottom and strange quarks and a charged lepton, respectively, $F_{\mu\nu}$ is the photon field strength, and $P_{R,L} = (1 \pm \gamma_5)/2$.

In addition, beyond the SM (BSM) physics can generate chirally-flipped ($b_{L(R)} \to b_{R(L)}$) versions of these operators,

$\mathcal{O}'_{7,\dots,10}$, and also four scalar and two tensor operators [9],

$$\mathcal{O}_S^{(\prime)} = \frac{e^2}{(4\pi)^2}[\bar{s}P_{R(L)}b][\bar{l}l],\ \mathcal{O}_P^{(\prime)} = \frac{e^2}{(4\pi)^2}[\bar{s}P_{R(L)}b][\bar{l}\gamma_5 l], \tag{3}$$

$$\mathcal{O}_T = \frac{e^2}{(4\pi)^2}[\bar{s}\sigma_{\mu\nu}b][\bar{l}\sigma^{\mu\nu}l],\ \mathcal{O}_{T5} = \frac{e^2}{(4\pi)^2}[\bar{s}\sigma_{\mu\nu}b][\bar{l}\sigma^{\mu\nu}\gamma_5 l]. \tag{4}$$

Note that there are only two possible non-vanishing tensor operators [36]. These, together with those in Eq. (2) and their chirally-flipped counterparts, constitute the most general basis for the Lagrangian describing B_s (semi-)leptonic rare decays. In this construction, the 12 coefficients in the EFT Lagrangian of these 12 distinct operators are *a priori* independent. However, as discussed in the next section, if the NP lies above the EW scale the number of free coefficients is reduced to 8.

The same effective Lagrangian (and following discussions) can be applied to $b \to d$ decays by replacing a d-quark for the s-quark in Eq. (1), and accounting for the difference in CKM elements by replacing λ_{pd} for λ_{ps} throughout. The Wilson coefficients are not necessarily the same as in the $b \to d$ transitions, and comparing the two sets of coefficients would give information about flavor violation of presumed NP. Similarly, Wilson coefficients could also depend on the family of the leptons, which would result in lepton universality violation.

High energy: New Physics above the EWSB scale. If the operators appearing in the effective Lagrangian are generated by physics at a scale Λ above the EWSB scale, $v \ll \Lambda$ they must originate from operators manifestly $SU(3)_c \times SU(2)_L \times U(1)_Y$ invariant. The fields at our disposal for the construction of such Lagrangian are the chiral fermions $q_L = (u_L, d_L)^T$, $\ell_L = (\nu_L, l_L)^T$, u_R, d_R, e_R, the Higgs doublet H and covariant derivatives containing gluons, weak-isospin and hypercharge vector bosons. We will work in the basis in which the down-type Yukawa matrix is diagonal and write the quark doublets as $q_d = (u_{jL}V^*_{jd}, d_L)$, $q_s = (u_{jL}V^*_{js}, s_L)$ and $q_b = (u_{jL}V^*_{jb}, b_L)$.

We restrict attention to BSM operators of dimension 6, which yield the leading corrections given the hierarchy of scales assumed. The effective Lagrangian takes the form $\mathcal{L}_{BSM} = \frac{1}{\Lambda^2}\sum_i C_i Q_i$. The relevant operators for the study of rare (semi-)leptonic decays in the B_q system are either dipole-like,

$$\begin{aligned} &Q_{dW} = g_2(\bar{q}_s\sigma^{\mu\nu}b_R)\tau^I H\, W^I_{\mu\nu},\ Q_{dB} = g_1(\bar{q}_s\sigma^{\mu\nu}b_R)HB_{\mu\nu}, \\ &Q'_{dW} = g_2 H^\dagger\tau^I(\bar{s}_R\sigma^{\mu\nu}q_b)W^I_{\mu\nu},\ Q'_{dB} = g_1 H^\dagger(\bar{s}_R\sigma^{\mu\nu}q_b)B_{\mu\nu}, \end{aligned} \tag{5}$$

Higgs-current times fermion-current,

$$\begin{aligned} Q^{(1)}_{Hq} &= \left(H^\dagger\, i\overleftrightarrow{D}_\mu H\right)(\bar{q}_s\gamma^\mu q_b) \\ Q^{(3)}_{Hq} &= H^\dagger\, i(\tau^I\overrightarrow{D}_\mu - \overleftarrow{D}_\mu\tau^I)H(\bar{q}_s\tau^I\gamma^\mu q_b) \\ Q_{Hd} &= \left(H^\dagger\, i\overleftrightarrow{D}_\mu H\right)(\bar{s}_R\gamma^\mu b_R) \end{aligned} \tag{6}$$

or four-fermion,

$$\begin{aligned} Q^{(1)}_{\ell q} &= (\bar{\ell}\gamma_\mu\ell)(\bar{q}_s\gamma^\mu q_b), & Q^{(3)}_{\ell q} &= (\bar{\ell}\gamma_\mu\tau^I\ell)(\bar{q}_s\gamma^\mu\tau^I q_b), \\ Q_{ed} &= (\bar{l}_R\gamma_\mu l_R)(\bar{s}\gamma^\mu b_R), & Q_{\ell d} &= (\bar{\ell}\gamma_\mu\ell)(\bar{s}\gamma^\mu b_R), \\ Q_{qe} &= (\bar{q}_s\gamma_\mu q_b)(\bar{l}\gamma^\mu l_R), & Q_{\ell edq} &= (\bar{q}_s b_R)(\bar{l}_R\ell), \\ Q'_{\ell edq} &= (\bar{\ell}\, l_R)(\bar{s}_R q_b), \end{aligned} \tag{7}$$

where color and weak-isospin indices are omitted and τ^I stand for the Pauli matrices in $SU(2)$-space. Primed operators correspond to a different flavor entry of the hermitian conjugate of the unprimed operator.

This Lagrangian cannot be compared still with that of Eq. (1); one has to integrate out the heavy degrees of freedom, *i.e.*, Z, W, t and H, and run it down to μ_b. The first step yields four-fermion and dipole operators as in Eqs. (2, 3). Remarkably, no new tensor-like operators (4) appear after integration of W and Z bosons at leading order. By contrast, new contributions to the coefficients of $\mathcal{O}^{(\prime)}_{9,10}$ are indeed generated by the operators in Eq. (6).

Explicitly, the connection with the Lagrangian of Eq. (1), at the scale M_W is [37], for scalar and tensor type operators:

$$\begin{aligned} C^l_S &= -C^l_P = \frac{4\pi^2}{e^2\lambda_{ts}}\frac{v^2}{\Lambda^2}C_{\ell edq}, \\ C^{l\prime}_S &= C^{l\prime}_P = \frac{4\pi^2}{e^2\lambda_{ts}}\frac{v^2}{\Lambda^2}C'_{\ell edq}, \\ C_T &= C_{T5} = 0, \end{aligned} \tag{8}$$

for dipole operators:

$$C_7^{(\prime)} = \frac{8\pi^2}{y_b\lambda_{ts}}\frac{v^2}{\Lambda^2}\left(C^{(\prime)}_{dB} - C^{(\prime)}_{dW}\right),$$

and for the current-current type of leptonic operators:

$$\begin{aligned} C_9 =& \frac{4\pi^2}{e^2\lambda_{ts}}\frac{v^2}{\Lambda^2}\left(C_{qe} + C^{(1)}_{\ell q} + C^{(3)}_{\ell q} - (1-4s^2_W)(C^{(1)}_{Hq} + C^{(3)}_{Hq})\right), \\ C_{10} =& \frac{4\pi^2}{e^2\lambda_{ts}}\frac{v^2}{\Lambda^2}\left(C_{qe} - C^{(1)}_{\ell q} - C^{(3)}_{\ell q} + (C^{(1)}_{Hq} + C^{(3)}_{Hq})\right), \\ C'_9 =& \frac{4\pi^2}{e^2\lambda_{ts}}\frac{v^2}{\Lambda^2}\left(C_{ed} + C_{\ell d} - (1-4s^2_W)C_{Hd}\right), \\ C'_{10} =& \frac{4\pi^2}{e^2\lambda_{ts}}\frac{v^2}{\Lambda^2}\left(C_{ed} - C_{\ell d} + C_{Hd}\right). \end{aligned}$$

Equation (8) shows explicitly what has been advertised in the introduction:
(i) Some operators cannot be generated in the EFT ($C_T = C_{T5} = 0$).
(ii) There are correlations between nonvanishing coefficients ($C_S = -C_P$ and $C'_S = C'_P$).
(iii) The contributions to some EFT coefficients may be subject to constraints arising purely from high energies (*e.g.*, $Q^{(\prime)}_{dW}$, $Q^{(\prime)}_{dB}$, Q_{Hq} and Q_{Hd} contribute to flavor-violating Z and H decays).

The reduction in the number of structures occurs only for scalar, pseudo-scalar and tensor operators. The reason for

this reduction is invariance under hypercharge: the tensor-like operators simply cannot be promoted to be $U(1)_Y$ invariant, and for scalar and pseudo-scalar $U(1)_Y$ requires the leptons to have definite chirality dependent on the b quark chirality. For the remaining operators the coefficients are independent linear combinations. However, note that there are additional correlations between the neutral current and the charged current version of the operators that arise from operators involving doublets. While these play no role directly in FCN leptonic decays of B mesons, they may give rise to additional constraints on the effects of NP.

Violations to the relations of Eq. (8) of order v^2/Λ^2 arise from dimension-8 operators like $\bar{q}Hb_R\,\bar{\ell}Hl_R$ and possibly of order $g_{EW}^2/16\pi^2$ from 1-loop matching.

Consequences in $B_q^0 \to l^+l^-$. A powerful probe of NP is the decay $B_q^0 \to l^+l^-$. In the SM it is first induced at 1-loop level and is chirally suppressed. Moreover, the hadronic matrix element is determined fully by B_q^0 decay constants F_{B_q}, which are calculated in lattice QCD [14].

The SM predictions for the branching fractions, $\overline{\mathcal{B}}$, have been worked out to high accuracy. For the muonic and electronic modes they currently are [15]:

$$\begin{aligned}\overline{\mathcal{B}}_{s\mu} &= 3.65(23)\times 10^{-9}, \quad \overline{\mathcal{B}}_{d\mu} = 1.06(9)\times 10^{-10},\\ \overline{\mathcal{B}}_{se} &= 8.54(55)\times 10^{-14}, \quad \overline{\mathcal{B}}_{de} = 2.48(21)\times 10^{-15},\end{aligned} \tag{9}$$

where the overline indicates untagged, time-integrated rates (as required by the sizable width difference in the $\bar{B}_s - B_s$ system, although not for B_d [16]).

The muonic modes have been recently measured by LHCb [17, 18] and CMS [19], and an average of the results leads to [20]:

$$\overline{\mathcal{B}}_{s\mu}^{\rm expt} = 2.9(7)\times 10^{-9}, \quad \overline{\mathcal{B}}_{d\mu}^{\rm expt} = 3.6^{+1.6}_{-1.4}\times 10^{-10}, \tag{10}$$

where the $\overline{\mathcal{B}}_{d\mu}$ mode is not statistically significant yet ($< 3\sigma$). For the electronic modes we currently have only upper bounds at 95% C.L. [21]:

$$\overline{\mathcal{B}}_{se}^{\rm expt} < 2.8\times 10^{-7}, \qquad \overline{\mathcal{B}}_{de}^{\rm expt} < 8.3\times 10^{-8}. \tag{11}$$

Useful quantities to compare the theory to are the ratios [16]:

$$\overline{R}_{ql} = \frac{\overline{\mathcal{B}}_{ql}}{\left(\overline{\mathcal{B}}_{ql}\right)_{\rm SM}} = \frac{1+\mathcal{A}_{\Delta\Gamma}^{ll}\,y_q}{1+y_q}\left(|S|^2+|P|^2\right), \tag{12}$$

where $y_q = \tau_{B_q}\Delta\Gamma_q/2$, $\mathcal{A}_{\Delta\Gamma}^{ll}$ is the mass eigenstate rate asymmetry [16] and:

$$S = \sqrt{1-\frac{4m_l^2}{m_{B_q}^2}}\,\frac{C_S-C_S'}{r_{ql}}, \quad P = \frac{C_{10}-C_{10}'}{C_{10}^{\rm SM}} + \frac{C_P-C_P'}{r_{ql}}, \quad \text{where} \quad r_{ql} = \frac{2m_l\,(m_b+m_q)C_{10}^{SM}}{m_{B_q}^2}. \tag{13}$$

The contributions of $C_S^{(\prime)}$ and $C_P^{(\prime)}$ are enhanced by the factor m_B/m_l, so below we will neglect the NP in $C_{10}^{(\prime)}$ for simplicity. The decay rate is only sensitive to the differences $(C_P - C_P')$ and $(C_S - C_S')$ so the sums, $(C_P + C_P')$ and $(C_S + C_S')$, need to be constrained through other means.

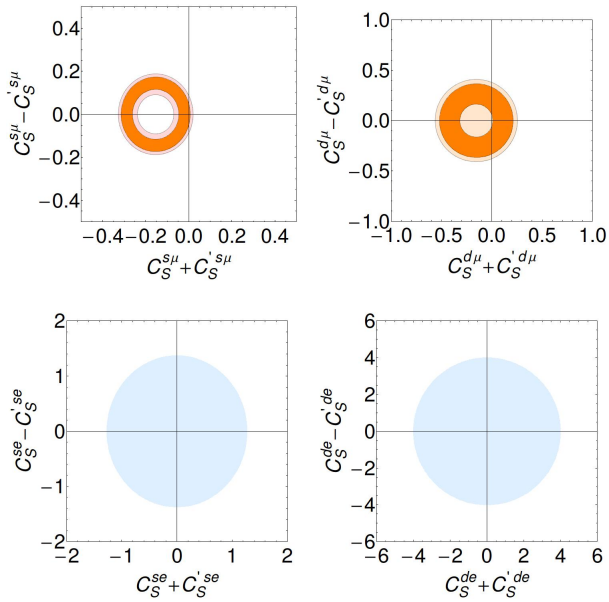

FIG. 1: In the upper panel we show the limits at 68% C.L. and 95% C.L. on the scalar Wilson coefficients that are induced by the experimental $\overline{\mathcal{B}}_{q\mu}$ in Eq. (10), where the corrections by mixing have been taken into account. For the electronic modes in the lower panel, we only show the 95% C.L. allowed regions (11). In both cases the Wilson coefficients are understood to be renormalized at $\mu = m_b$.

Introducing the hypothesis of this work, we impose (8) in Eq. (12) and (13) so that now

$$\overline{R}_{ql} \simeq \frac{|C_S - C_S'|^2}{r_{ql}^2} + \left|1 - \frac{C_S + C_S'}{r_{ql}}\right|^2, \tag{14}$$

where we have neglected $y_s = 0.075(12)\%$ [22] and the phase space factor for clarity. In addition to the reduction of free parameters from 4 to 2 in the scalar and pseudo-scalar sector, now these two parameters enter the decay rate in two orthogonal linear combinations. As a result the $B_q \to l^+l^-$ branching fraction alone bounds *all directions* in our two parameter space. In particular, for real Wilson coefficients, the bound of Eq. (14) defines a circle in parameter space centered at $(C_S + C_S', C_S - C_S') = (r_{ql}, 0)$ with radius $|r_{ql}|\sqrt{\overline{R}_{ql}^{\rm expt}}$.

The contour plots in Fig. 1 show these circular bounds with the radius in the muonic cases determined by $|r_{q\mu}| \simeq 0.16$. This shape is in contrast with the bands, experimentally unconstrained in one direction, that would be obtained in the standard analysis. Note that improving the experimental accuracy in these modes will only reduce the width of the ring and that breaking the degeneracy will require other observables. One attractive possibility is the observable $\mathcal{A}_{\Delta\Gamma}^{\mu\mu}$, which may be obtained by measuring the effective $B_s \to \mu^+\mu^-$ lifetime [16].

For the electronic modes, $|r_{qe}| \sim 10^{-3}$ and the strength of the limits in the parameter space is governed by the size of

$\overline{R}_{qe}^{\rm expt}$. Hence, improved experimental bounds on the branching fractions have the potential to probe extremely high energies through the scalar operators.

To quantify this, the bounds on the scale Λ for the different decay rates in Eq. (9) can be computed making use of Eq. (8) and assuming naturalness, namely $C^{(\prime)}_{ledq}(\Lambda) \simeq 1$. To compute these bounds the running is taken into account in the two intervals: *i)* from m_b to $\sim M_W$ where QCD at 1 loop dominates [9] and strengthens the bound by ~ 0.74, *ii)* from M_W to Λ using the recently-computed anomalous dimension [10–13], which gives an extra factor ~ 0.66. The bounds from $\overline{\mathcal{B}}_{s\mu}$, $\overline{\mathcal{B}}_{d\mu}$, $\overline{\mathcal{B}}_{se}$, and $\overline{\mathcal{B}}_{de}$, are respectively 78 TeV, 130 TeV, 36 TeV, and 49 TeV. Note that B_d mesons decays do better at constraining new physics due to the CKM suppressed SM background. In particular, $\overline{\mathcal{B}}_{de}$ can supersede the present bounds for a precision of 10^{-10}.

Consequences in $B \to K^{(*)} l^+ l^-$. The exclusive semileptonic B decays are also powerful flavor laboratories. As 3- and 4-body decays, their angular distributions lead to a rich and non-trivial phenomenology which could potentially unveil NP surfacing through various operators (see e.g. Refs. [9, 23–26] and references therein).

The effects of scalar or tensor operators in $B \to K^{(*)} l^+ l^-$ have been considered by several authors [9, 23, 24]. The first immediate consequence of our analysis in these decays is that tensor operators can be ignored altogether (up to $\mathcal{O}(v^2/\Lambda^2)$ corrections). This observation should lead to a considerable simplification in the theoretical analyses of the angular observables [9, 24].

The scalar operators contribute to the total decay rates $B \to K^{(*)} l^+ l^-$, providing another experimental input to resolve degeneracies. In practice, however, any sensitivity is blurred by the SM contribution which depends on quite uncertain hadronic form factors [25]. As an example, the coefficient $I_6^c(q^2)$ in the angular distribution in the K^* mode is directly proportional to the combination $|C_S - C'_S|^2$, and it is a null test of the SM [23, 25] but the contribution is suppressed by m_l so that the observable is not competitive with purely leptonic decays.

In the case of the K mode the two angular observables A_{FB} and F_H [9] are also null tests of the SM and receive contributions from $(C_S + C'_S)$ and $(C_P + C'_P)$. In the standard analysis these observables provide sensitivity to the orthogonal directions scanned in $B_q \to l^+ l^-$ and in our case could lift the degeneracies in Fig. 1. However, at low q^2 the scalars appear suppressed by either m_l in A_{FB} or by a kinematical factor in F_H [9] [38].

As a final example let us comment on the impact of our analysis in lepton universality violation in $B^+ \to K^+ l^+ l^-$ decays [27]. Recently the LHCb collaboration [28] has reported a deficit in muonic decays with respect to electronic ones in the $[1, 6]$ GeV^2 bin with a significance of 2.6σ:

$$R_K \equiv \frac{\text{Br}\,(B^+ \to K^+ \mu\mu)}{\text{Br}\,(B^+ \to K^+ ee)} = 0.745^{+0.090}_{-0.074}(\text{stat}) \pm 0.036(\text{syst}). \tag{15}$$

In the SM, R_K is given very accurately, $R_K = 1.0003(1)$ [9], since the hadronic contributions cancel in the ratio to very good approximation. In Ref. [9] possible scenarios with sizable scalar operators were shown to produce large effects in R_K. Our analysis shows that the bounds from the fully leptonic decay suffice to exclude the possibility of scalar operators accounting for (15), since at 95% C.L. we have:

$$R_K \in [0.982, 1.007]. \tag{16}$$

In light of this and the absence of tensors, we conclude that a large lepton universality violation in R_K could be only produced by the operators $\mathcal{O}_9^{(\prime)}$ and $\mathcal{O}_{10}^{(\prime)}$. Unfortunately these are not very well bounded, especially for the electronic case, so different scenarios of NP could currently explain (15). For example one could entertain the possibility of a sizable and negative effect in C_9 affecting only the muonic mode, $\delta C_9^\mu = -1$. In this scenario one obtains $R_K \simeq 0.79$. As a side remark, it is worth emphasizing that such a negative NP contribution to $\mathcal{O}_9$ has been argued to be necessary to understand the current $b \to s\mu\mu$ data set [30–33].

Conclusions. We have discussed a novel approach to the study of new-physics effects in the B_q (semi-)leptonic decays. This relies on the assumption that the new dynamics enter at a scale $\Lambda \gg v$, and it is based on the (tree-level) matching of the effective weak Lagrangian customarily used in the phenomenological analyses, $\mathcal{L}_W$, to the most general 6-dimensional Lagrangian invariant under the SM gauge group (as done, customarily, in the analysis of other weak hadronic processes like nuclear and neutron β-decays [34, 35]).

As a direct consequence of $SU(2)_L \times U(1)_Y$ invariance, new constraints correlate the operators in $\mathcal{L}_W$. For example, in rare B_q (semi-)leptonic decays the coefficients of the a priori four possible scalar operators are reduced to two and the tensor operators are forbidden. The phenomenology of this reduced set of operators in $B_q \to l^+ l^-$ decays was studied.

The present approach could be extended to other low energy processes but also combined with EW scale physics to narrow down possible NP operators. Finally let us remark that, with the growing experimental data, the type of correlations discussed here is likely to play an important role in the determination of the nature of the new physics to appear.

Acknowledgments We are happy to thank S. Jäger for helpful discussions. This work was supported in part by DOE grants DE-SC000991 and DE-SC0009919. JMC has received funding from the People Programme (Marie Curie Actions) of the European Union's Seventh Framework Programme (FP7/2007-2013) under REA grant agreement n PIOF-GA-2012-330458 and acknowledges the Spanish Ministerio de Economía y Competitividad and european FEDER funds under the contract FIS2011-28853-C02-01 for support.

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[37] For a complete connection, the coefficients should be previously run from Λ to M_W ([10–13]), this is taken into account for the numerical bounds on Λ.

[38] The kinematical suppression in F_H is lifted in the low-recoil region [24].