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Constraints on Lepton Universality Violation from Rare B Decays

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The LHCb collaboration has very recently released a new study of $B^+ \to K^+\ell^+\ell^-$ and $B \to K^{*0}\ell^+\ell^-$ ($\ell = e, \mu$) decays, testing lepton universality with unprecedented accuracy using the whole Run 1 and 2 dataset. In addition, the CMS collaboration has recently reported an improved analysis of the branching ratios $B_{(d,s)} \to \mu^+\mu^-$. While these measurements offer, per se, a powerful probe of New Physics, global analyses of $b \to s\ell^+\ell^-$ transitions also rely on the assumptions about nonperturbative contributions to the decay matrix elements. In this work, we perform a global Bayesian analysis of New Physics in (semi)leptonic rare B decays, paying attention to the role of charming penguins which are difficult to evaluate from first principles. We find data to be consistent with the Standard Model once rescattering from intermediate hadronic states is included. Consequently, we derive stringent bounds on lepton universality violation in $|\Delta B| = |\Delta S| = 1$ (semi)leptonic processes.

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Since the first collisions in 2010, the Large Hadron Collider (LHC) allowed for a tremendous step forward in the electroweak (EW) sector of the Standard Model (SM) of Particle Physics – culminated with the discovery of the Higgs boson [1, 2] – while it has also excited the community with a few interesting hints of Physics Beyond the SM (BSM). In particular, the LHCb collaboration provided the first statistically relevant hint for Lepton Universality Violation (LUV) in flavor-changing neutral-current (FCNC) processes [3], measuring the ratio $R_K \equiv Br(B^+ \to K^+\mu^+\mu^-)/Br(B^+ \to K^+e^+e^-)$ in the dilepton invariant-mass range $q^2 \in [1, 6]$ GeV^2. These hints have been confirmed by subsequent measurements, always by the LHCb collaboration, namely R_K [4], R_{K^*} [5, 6], R_{K_S} and $R_{K^{*+}}$ [7].

Interestingly enough, these hints of LUV appeared in transitions where deviations from the SM were already claimed, see e.g. [8–11], on the basis of the measurements of angular distributions in $b \to s\mu^+\mu^-$ decays [12–23]. Claiming discrepancies from SM predictions in Branching Ratios (BRs) and angular distributions requires, however, full theoretical control on hadronic uncertainties in the matrix element calculation [24–26], and in particular on the so-called *charming penguins* [27], which might affect the vector coupling to the leptons even in regions of the dilepton invariant mass well below the charmonium threshold [28, 29] and bring the SM in agreement with experiment [30]. Combining angular distributions with LUV data strengthened the case for New Physics (NP), since a single NP contribution could reproduce the whole set of data [31–39]. On the other hand, charming penguins might affect the picture of NP behind LUV, since LUV ratios depend on the interplay of NP and hadronic contributions [38, 40–42]. While considerable progress has been made in estimating (at least part of) the charming-penguin amplitudes using light-cone sum rules [43, 44] and analyticity supplemented with perturbative QCD in the Euclidean q^2 region [45–48], calculating these hadronic contributions remains an open problem, as we discuss below.

Before presenting our results, we notice that very recently the experimental picture drawn so far has suddenly changed. Firstly, the CMS collaboration provided a new analysis of $\mathrm{BR}(B_{(d,s)} \to \mu^+\mu^-)$ with the full Run 2 dataset [49], bringing the HFLAV average

$$\mathrm{BR}(B_s \to \mu^+\mu^-) = (3.45 \pm 0.29) \cdot 10^{-9} \qquad (1)$$

into excellent agreement with the SM prediction $\mathrm{BR}(B_s \to \mu^+\mu^-) = (3.47 \pm 0.14) \cdot 10^{-9}$ [50, 51]. Being short-distance dominated, this FCNC process strongly constrains NP contributions involving, in particular, axial leptonic couplings [52, 53]. Furthermore, an updated LHCb analysis of R_K and R_{K^*} based on the full Run 1 and 2 dataset has been presented [54, 55]:

$$\begin{aligned} R_{K_{[0.1,1.1]}} &= 0.994\ ^{+0.090}_{-0.082}\ (\mathrm{stat})\ ^{+0.029}_{-0.027}\ (\mathrm{syst})\,, \qquad (2)\\ R_{K^*{}_{[0.1,1.1]}} &= 0.927\ ^{+0.093}_{-0.087}\ (\mathrm{stat})\ ^{+0.036}_{-0.035}\ (\mathrm{syst})\,,\\ R_{K_{[1.1,6]}} &= 0.949\ ^{+0.042}_{-0.041}\ (\mathrm{stat})\ ^{+0.022}_{-0.022}\ (\mathrm{syst})\,,\\ R_{K^*{}_{[1.1,6]}} &= 1.027\ ^{+0.072}_{-0.068}\ (\mathrm{stat})\ ^{+0.027}_{-0.026}\ (\mathrm{syst})\,, \end{aligned}$$

with correlations reported in Fig. 26 of ref. [55]. These new measurements dramatically change the scenario of possible LUV effects in FCNC B decays [56], questioning what in the last years served as fertile ground for model building, see for instance [57–80].

In this *Letter* we provide a reassessment of NP effects in $b \to s\,\mu^+\mu^-$ transitions in view of the experi-

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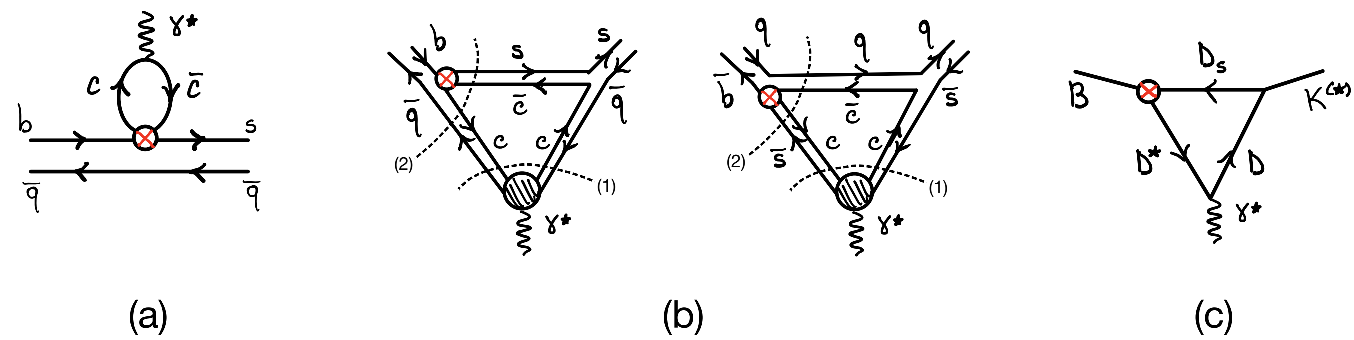

FIG. 1. *Example of charming-penguin diagrams contributing to the $B \to K^{(*)}\ell^+\ell^-$ amplitude. Diagram **(a)** represents the class of charming-penguin amplitudes related to $c-\bar{c}$ state that subsequently goes into a virtual photon, see refs. [43, 45–48]. Diagram **(b)** and **(c)** represent the kind of contributions from rescattering of intermediate hadronic states, at the quark and meson level respectively. The phenomenological relevance of rescattering for the SM prediction of the $B \to K^{(*)}\ell^+\ell^-$ decays has been recently considered in ref. [38].*

mental novelties discussed above. Adopting the model-independent language of the Standard Model Effective Theory (SMEFT) [81, 82], we present an updated analysis of $|\Delta B| = |\Delta S| = 1$ (semi)leptonic processes and show that current data no longer provide strong hints for NP. Indeed, updating the list of observables considered in our previous global analysis [38] with the results in eqs. (1) and (2), the only remaining measurements deviating from SM expectations and not affected by hadronic uncertainties are the LUV ratios R_{K_S} and $R_{K^{*+}}$ [7], for which a re-analysis by the LHCb collaboration is mandatory in view of what discussed in [54, 55].

The anatomy of the $B \to K^{(*)}\ell^+\ell^-$ decay can be characterized in terms of helicity amplitudes [24, 83], that in the SM at a scale close to the bottom quark mass m_b can be written as:

$$H_V^\lambda \propto \left\{ C_9^{\rm SM}\widetilde{V}_{L\lambda} + \frac{m_B^2}{q^2}\left[\frac{2m_b}{m_B}C_7^{\rm SM}\widetilde{T}_{L\lambda} - 16\pi^2 h_\lambda\right]\right\} ,$$

$$H_A^\lambda \propto C_{10}^{\rm SM}\widetilde{V}_{L\lambda} \ , \ H_P \propto \frac{m_\ell\, m_b}{q^2}\, C_{10}^{\rm SM}\left(\widetilde{S}_L - \frac{m_s}{m_b}\widetilde{S}_R\right) ,$$

with $\lambda = 0, \pm$ and $C_{7,9,10}^{\rm SM}$ the SM Wilson coefficients of the semileptonic operators of the $|\Delta B| = |\Delta S| = 1$ weak effective Hamiltonian [84–86], normalized as in ref. [41]. The naively factorizable contributions to the above amplitudes can be expressed in terms of seven q^2-dependent form factors, $\widetilde{V}_{0,\pm}$, $\widetilde{T}_{0,\pm}$ and $\widetilde{S}$ [87, 88]. At the loop level, non-local effects parametrically not suppressed (neither by small Wilson coefficients nor by small CKM factors) arise from the insertion of the following four-quark operator:

$$Q_2^c = (\bar{s}_L\gamma_\mu c_L)(\bar{c}_L\gamma^\mu b_L)\,, \tag{3}$$

that yields non-factorizable power corrections in H_V^λ via the hadronic correlator $h_\lambda(q^2)$ [26, 30, 89], receiving the main contribution from the time-ordered product:

$$\frac{\epsilon_\mu^*(\lambda)}{m_B^2}\int d^4x\; e^{iqx}\langle\bar{K}^*|\mathcal{T}\{j^\mu_{\rm em}(x)Q_2^c(0)\}|\bar{B}\rangle\,, \tag{4}$$

with $j^\mu_{\rm em}(x)$ the electromagnetic (quark) current.

This correlator receives two kinds of contributions. The first corresponds to diagrams of the form of diagram **(a)** in Fig. 1, where the initial B meson decays to the $K^{(*)}$ plus a $c\bar{c}$ state that subsequently goes into a virtual photon. This contribution has been studied in detail in the context of light-cone sum rules in the regime $q^2 \ll 4m_c^2$ in [43]; in the same reference, dispersion relations were used to extend the result to larger values of the dilepton invariant mass. While the operator product expansion performed in ref. [43] was criticized in ref. [29], and multiple soft-gluon emission may represent an obstacle for the correct evaluation of this class of hadronic contributions [30, 40, 90, 91], refs. [45, 46] have exploited analyticity in a more refined way than [43]. In those works the negative q^2 region – where perturbative QCD is supposed to be valid – has been used to further constrain the amplitude. Building on these works, together with unitarity bounds [47], ref. [48] found a very small effect in the large-recoil region.

The second kind of contribution to the correlator in eq. (4) originates from the triangle diagrams depicted in Fig. 1 **(b)**, in which the photon can be attached both to the quark and antiquark lines and we have not drawn explicitly the gluons exchanged between quark-antiquark pairs. An example of an explicit hadronic contribution of this kind is depicted in Fig. 1 **(c)**.[1] The D_sD^* pair is produced by the weak decay of the initial B meson with low momentum, so that no color transparency argument holds and rescattering can easily take place. Furthermore, the recent observation of tetraquark states in $e^+e^- \to K(D_sD^* + D_s^*D)$ by the BESIII collaboration [93] confirms the presence of nontrivial nonperturbative dynamics of the intermediate state.

One could think of applying dispersive methods also

[1] See ref. [92] for a very recent estimate of similar diagrams with up quarks, rather than charm quarks, in the internal loop.

to this kind of contributions, but the analytic structure of triangle diagrams is quite involved, depending on the values of external momenta and internal masses. A dispersion relation in q^2 of the kind used in refs. [43, 45–48], based on the cut denoted by (1) in Fig. 1 **(b)**, could be written if the B invariant mass were below the threshold for the production of charmed intermediate states. However, when the B invariant mass raises above the threshold for cut (2), an additional singularity moves into the q^2 integration domain, requiring a nontrivial deformation of the path (see for example the detailed discussion in ref. [94]). Another possibility would be to get an order-of-magnitude estimate of contributions as the one in Fig. 1 **(c)** using an approach similar to ref. [92].

To be conservative, and in the absence of a first-principle calculation of the diagrams in Fig. 1, we adopt a data-driven approach based on the following parameterization of the hadronic contributions, inspired by the expansion of the correlator of eq. (4) as originally done in ref. [24], and worked out in detail in ref. [91]:

$$
\begin{aligned}
H_V^- \propto & \frac{m_B^2}{q^2}\left[\frac{2m_b}{m_B}\left(C_7^{\rm SM}+h_-^{(0)}\right)\widetilde{T}_{L-} - 16\pi^2 h_-^{(2)}\, q^4\right] \\
&+ \left(C_9^{\rm SM}+h_-^{(1)}\right)\widetilde{V}_{L-}\,, \\
H_V^+ \propto & \frac{m_B^2}{q^2}\left[\frac{2m_b}{m_B}\left(C_7^{\rm SM}+h_-^{(0)}\right)\widetilde{T}_{L+} - 16\pi^2 \Big(h_+^{(0)} \right. \\
&\left. + h_+^{(1)}\, q^2 + h_+^{(2)}\, q^4\Big)\right] + \left(C_9^{\rm SM}+h_-^{(1)}\right)\widetilde{V}_{L+}\,, \\
H_V^0 \propto & \frac{m_B^2}{q^2}\left[\frac{2m_b}{m_B}\left(C_7^{\rm SM}+h_-^{(0)}\right)\widetilde{T}_{L0} - 16\pi^2 \sqrt{q^2}\Big(h_0^{(0)} \right. \\
&\left. + h_0^{(1)}\, q^2\Big)\right] + \left(C_9^{\rm SM}+h_-^{(1)}\right)\widetilde{V}_{L0}\,. \qquad (5)
\end{aligned}
$$

This parameterization – while merely rooted on a phenomenological basis – has the advantage of making transparent the interplay between hadronic and possible NP contributions. Indeed, the coefficients $h_-^{(0)}$ and $h_-^{(1)}$ have the same effect of a lepton universal shift due to NP in the real part of the Wilson coefficients C_7 and C_9, respectively. Consequently, the theoretical assumptions on the size of these hadronic parameters crucially affect the extraction of NP contributions to $C_{7,9}$ from global fits. Within the SM, the new measurements in eqs. (1)-(2) do not affect the knowledge of the h_λ coefficients; the most up-to-date data-driven extraction of the hadronic parameters introduced in eq. (5) can be found in Table 1 of ref. [38].

Moving to the analysis of NP, current constraints from direct searches at the LHC reasonably suggest in this context that BSM physics would arise at energies much larger than the electroweak scale. Then, a suitable framework to describe such contributions is given by the SMEFT, in particular by adding to the SM the following dimension-six operators:[2]

$$
\begin{aligned}
O_{2223}^{LQ^{(1)}} &= (\bar{L}_2\gamma_\mu L_2)(\bar{Q}_2\gamma^\mu Q_3)\,, \\
O_{2223}^{LQ^{(3)}} &= (\bar{L}_2\gamma_\mu \tau^A L_2)(\bar{Q}_2\gamma^\mu \tau^A Q_3)\,, \\
O_{2322}^{Qe} &= (\bar{Q}_2\gamma_\mu Q_3)(\bar{e}_2\gamma^\mu e_2)\,, \\
O_{2223}^{Ld} &= (\bar{L}_2\gamma_\mu L_2)(\bar{d}_2\gamma^\mu d_3)\,, \\
O_{2223}^{ed} &= (\bar{e}_2\gamma_\mu e_2)(\bar{d}_2\gamma^\mu d_3)\,, \qquad (6)
\end{aligned}
$$

where in the above $\tau^{A=1,2,3}$ are Pauli matrices, a sum over A is understood, L_i and Q_i are $SU(2)_L$ doublets, e_i and d_i singlets, and flavor indices are defined in the basis where the down-quark Yukawa matrix is diagonal. For concreteness, we normalize SMEFT Wilson coefficients to a NP scale $\Lambda_{\rm NP} = 30$ TeV and we only consider NP contributions to muons.[3] The matching between the weak effective Hamiltonian and the SMEFT operators implies the following contributions to the SM operators and to the chirality-flipped ones denoted by primes [97]:

$$
\begin{aligned}
C_9^{\rm NP} &= \mathcal{N}_\Lambda\left(C_{2223}^{LQ^{(1)}} + C_{2223}^{LQ^{(3)}} + C_{2322}^{Qe}\right)\,, \\
C_{10}^{\rm NP} &= \mathcal{N}_\Lambda\left(C_{2322}^{Qe} - C_{2223}^{LQ^{(1)}} - C_{2223}^{LQ^{(3)}}\right)\,, \\
C_9^{\prime,\rm NP} &= \mathcal{N}_\Lambda\left(C_{2223}^{ed} + C_{2223}^{Ld}\right)\,, \\
C_{10}^{\prime,\rm NP} &= \mathcal{N}_\Lambda\left(C_{2223}^{ed} - C_{2223}^{Ld}\right)\,, \qquad (7)
\end{aligned}
$$

with $\mathcal{N}_\Lambda = (\pi v^2)/(\alpha_e V_{ts} V_{tb}^* \Lambda_{\rm NP}^2)$. As evident from the above equation, operators $O_{2223}^{LQ^{(1,3)}}$ always enter as a sum. Hence we denote their Wilson coefficient as C_{2223}^{LQ}.

We perform a Bayesian fit to the data in refs. [13, 17–23, 49, 54, 55, 98–104] employing the `HEPfit` code [105, 106]. For the form factors and input parameters, we follow the same approach used in our previous refs. [30, 38, 40–42, 90, 91]. In particular, we use the same inputs as in ref. [38], with the only exception of CKM parameters, which have been updated according to the results of ref. [51]. We compute $B \to K^{(*)}\ell^+\ell^-$ and $B_s \to \phi\ell^+\ell^-$ decays using QCD factorization [107].

As already mentioned discussing Fig. 1, a global analysis of $b \to s\ell^+\ell^-$ transitions can be sensitive to hadronic contributions that are difficult to compute from first principles and that can yield important phenomenological effects. Therefore, in what we denote below as *data driven* scenario, we assume a flat prior in a sufficiently large range for the $h_\pm^{(0,1,2)}$ and $h_0^{(0,1)}$ parameters, which are then determined from data simultaneously with the NP coefficients.[4] To clarify the phenomenological relevance of charming penguins, we compare the results of the *data*

[2] Note that these operators may be generated via renormalization group effects, see, e.g., refs. [95, 96].

[3] This choice is mainly motivated by the fact that $B_s \to \mu^+\mu^-$ is one of the key observables of the present study.

[4] As in ref. [38], we assume exact SU(3) flavor for the h parameters and add additional ones for $B \to K$.

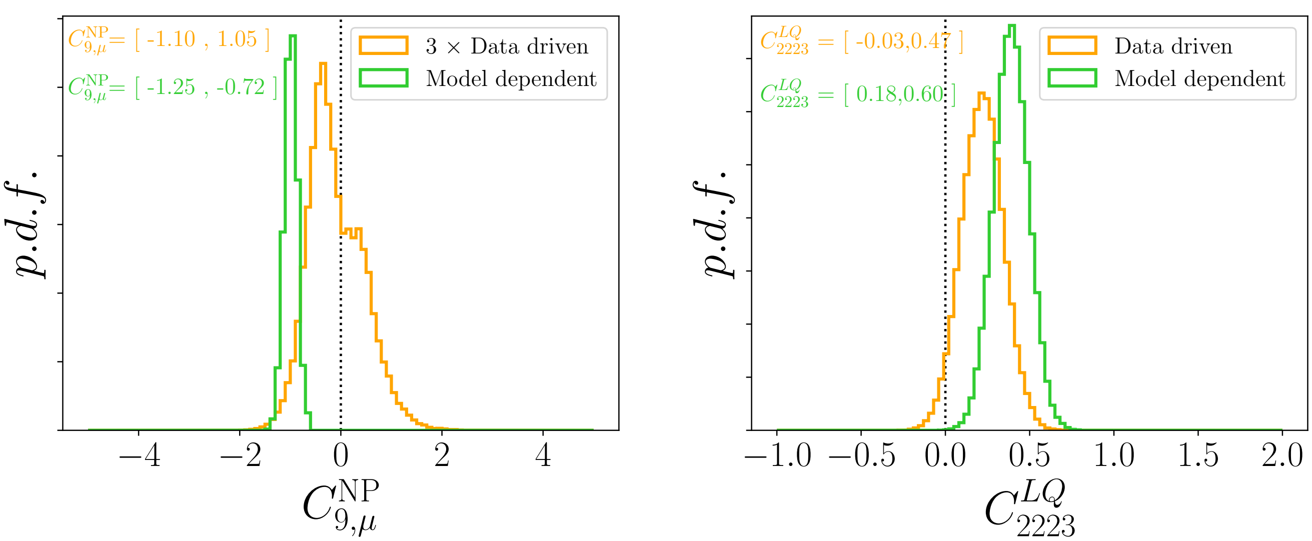

FIG. 2. *Left panel: Posterior p.d.f. for the NP coefficient $C_{9,\mu}^{\rm NP}$. Right panel: Posterior p.d.f. for the SMEFT Wilson coefficient C_{2223}^{LQ}. For both panels, we show the p.d.f. in green and orange on the basis of the hadronic approach adopted in the global analysis (see the text for more details).*

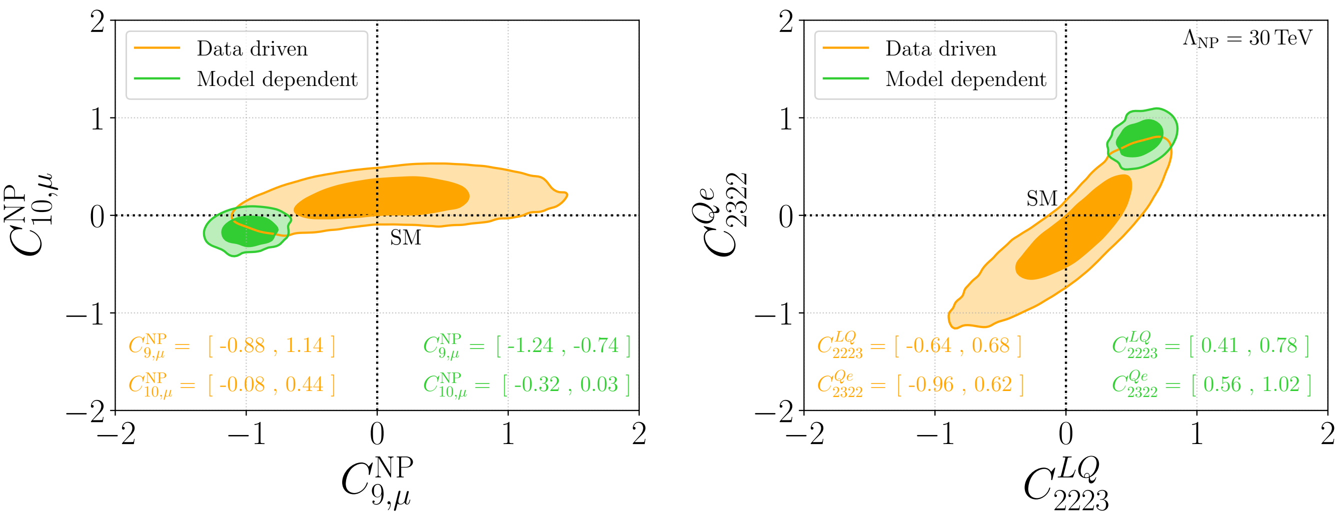

FIG. 3. *Left panel: Joint posterior p.d.f. for $C_{9,\mu}^{\rm NP}$ and $C_{10,\mu}^{\rm NP}$. Right panel: Joint posterior p.d.f. for the SMEFT Wilson coefficients C_{2223}^{LQ} and C_{2322}^{Qe}. For both panels, we show 68% and 95% probability regions in green and orange on the basis of the hadronic approach adopted in the global analysis (see the text for more details).*

driven approach against what we denote instead as *model dependent* treatment of hadronic uncertainties, in which we assume that the contributions generated by the diagrams in Fig. 1 **(b)** (or **(c)**) are negligible and that the correlator in eq. (4) is well described by the approach of refs. [43–48], yielding a subleading effect to the hadronic effects computable in QCD factorization.

In both approaches to QCD long-distance effects, we obtain a sample of the posterior joint probability density function (p.d.f.) of SM parameters, including form factors, and, in the *data driven* scenario, h_λ parameters, together with NP Wilson coefficients. From each posterior p.d.f. we compute the highest probability density intervals (HPDIs), which represent our best knowledge of the model parameters after the new measurements. We also perform model comparison using the information criterion [108], defined as:

$$IC \equiv -2\overline{\log \mathcal{L}} + 4\sigma^2_{\log \mathcal{L}} \,, \tag{8}$$

where the first and second terms are the mean and variance of the log-likelihood posterior distribution. The first term measures the quality of the fit, while the second one is related to the effective degrees of freedom involved, penalizing more complicated models. Models with smaller

	95% HPDI	ΔIC
$C_{9,\mu}^{\rm NP}$	[-1.10 , 1.05]	-1.1
	[-1.25 , -0.72]	65
$\{C_{9,\mu}^{\rm NP}, C_{10,\mu}^{\rm NP}\}$	{[-0.88 , 1.14], [-0.08 , 0.44]}	0.3
	{[-1.24 , -0.74], [-0.32 , 0.03]}	59
$\{C_{9,\mu}^{\rm NP}, C_{9,\mu}^{\prime,\rm NP}\}$	$\{[-1.22, 1.41], [-2.77, 1.46],$	-2.3
	$\{[-1.34, -0.80], [-0.04, 0.82],$	64
$\{C_{9,\mu}^{\rm NP}, C_{10,\mu}^{\prime,\rm NP}\}$	$\{[-1.12, 1.34], [-0.28, 0.22],$	-2.2
	$\{[-1.38, -0.79], [-0.36, 0.06],$	62
$\{C_{9,\mu}^{\rm NP}, C_{10,\mu}^{\rm NP}, C_{9,\mu}^{\prime,\rm NP}, C_{10,\mu}^{\prime,\rm NP}\}$	$\{[-1.10, 1.40], [-0.18, 0.60], [-2.66, 1.32], [-0.33, 0.47]\}$	-1.5
	$\{[-1.39, -0.81], [-0.40, 0.05], [-0.51, 0.77], [-0.43, 0.19]\}$	62

TABLE I. *HPDI for the Wilson coefficients of the low-energy weak Hamiltonian in all the considered NP scenarios along with the corresponding ΔIC. White rows correspond to results obtained in the data driven scenario, while model dependent scenario results are shaded in gray. See the text for the definition of the two scenarios.*

	95% HPDI	ΔIC
C_{2223}^{LQ}	[-0.03 , 0.47]	1.2
	[0.18 , 0.60]	29
$\{C_{2223}^{LQ}, C_{2322}^{Qe}\}$	{[-0.64 , 0.68], [-0.96 , 0.62]}	0.3
	{[0.41 , 0.78] , [0.56 , 1.02]}	59
$\{C_{2223}^{LQ}, C_{2223}^{ed}\}$	{[-0.03 , 0.50], [-0.44 , 0.29]}	1.0
	{[0.35 , 0.67], [-1.02 , -0.12]}	41
$\{C_{2223}^{LQ}, C_{2223}^{Ld}\}$	{[-0.05 , 0.67], [-0.23 , 0.53]}	-0.5
	{[0.28 , 0.58], [-0.07 , 0.30]}	36
$\{C_{2223}^{LQ}, C_{2322}^{Qe}, C_{2223}^{Ld}, C_{2223}^{ed}\}$	$\{[-0.85, 0.90], [-1.22, 0.72], [-0.88, 1.82], -1.03, 1.92]\}$	-1.5
	$\{[0.40, 0.85], [0.63, 1.19], [-0.48, 0.13], [-0.64, 0.59]\}$	62

TABLE II. *Same as Tab. I for SMEFT Wilson coefficients.*

IC should then be preferred [109]. While the posterior distributions for SM parameters are unaffected by LUV measurements, the SM IC of course depends on the latter: indeed, the SM in the *data driven* scenario provides an excellent description of current data, leading to very small or even negative values of $\Delta\text{IC} \equiv \text{IC}_{\rm SM} - \text{IC}_{\rm NP}$. Conversely, the agreement of the SM with angular observables remains poor in the *model dependent* approach, implying for this case large values of ΔIC, signaling a statistically significant preference for NP.

We now discuss several NP configurations, in order of increasing complexity. We start by allowing a single nonvanishing NP Wilson coefficient, either $C_{9,\mu}^{\rm NP}$, defined in the low-energy weak Hamiltonian, or the Wilson coefficient C_{2223}^{LQ}, belonging to the SMEFT. The p.d.f.s for the two NP Wilson coefficients are reported in Fig. 2, while the corresponding numerical results for the 95% HPDIs are reported in the first row of Tables I and II. As anticipated above, no significant preference for NP is seen in the *data driven* scenario, while NP contributions are definitely needed in the *model dependent* scenario, with a clear preference for $C_{9,\mu}^{\rm NP} \neq 0$.

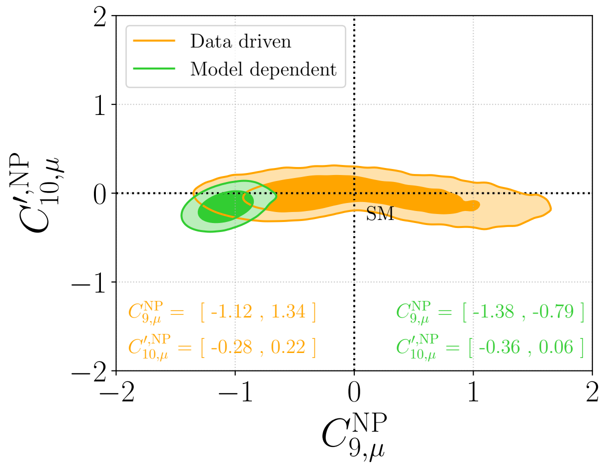

FIG. 4. *Joint posterior p.d.f. for $C_{9,\mu}^{\rm NP}$ and $C_{10,\mu}^{\prime\rm NP}$. We show 68% and 95% probability regions in green and orange on the basis of the hadronic approach adopted in the global analysis (see the text for more details).*

Figure 3 displays the allowed regions in the $C_{9,\mu}^{\rm NP} - C_{10,\mu}^{\rm NP}$ and $C_{2223}^{LQ} - C_{2322}^{Qe}$ planes, while the corresponding HPDIs are reported in the second row of Tables I and II respectively. Again, no evidence for NP is seen in the *data driven* case, while clear evidence for a nonvanishing $C_{9,\mu}^{\rm NP}$ appears in the *model dependent* approach. Deviations from zero of $C_{10,\mu}^{\rm NP}$ are strongly constrained by BR$(B_s \to \mu^+\mu^-)$, corresponding to the strong correlation $C_{2223}^{LQ} \sim C_{2322}^{Qe}$ seen in the right panel of Fig. 3.

Next, we consider NP models in which right-handed $b \to s$ transitions arise. In the weak effective Hamiltonian, we allow for nonvanishing $C_{9,\mu}^{\rm NP}$ and $C_{9,\mu}^{\prime,\rm NP}$ or $C_{10,\mu}^{\prime,\rm NP}$. In particular, in Fig.4 we present the results of the fit in the $C_{9,\mu}^{\rm NP} - C_{10,\mu}^{\prime,\rm NP}$ case, which we considered in ref. [41] as the best fit one in view of the deviation from one of the ratio R_K/R_{K^*} [110]. With the current experimental situation, this is not the case anymore, and $C_{10,\mu}^{\prime,\rm NP}$ is again strongly constrained by BR$(B_s \to \mu^+\mu^-)$. In the SMEFT, we consider nonvanishing C_{2223}^{LQ} and C_{2223}^{ed} or C_{2223}^{Ld}. The numerical results for the NP coefficients can be found in the third and fourth rows of Tables I and II.

Finally, we present the results of a combined fit in which all the four NP Wilson coefficients considered above are allowed to float simultaneously, namely C_{2223}^{LQ}, C_{2322}^{Qe}, C_{2223}^{Ld} and C_{2223}^{ed}, or equivalently, in the language of the weak effective Hamiltonian, $C_{9,\mu}^{\rm NP}$, $C_{10,\mu}^{\rm NP}$ and the cor-

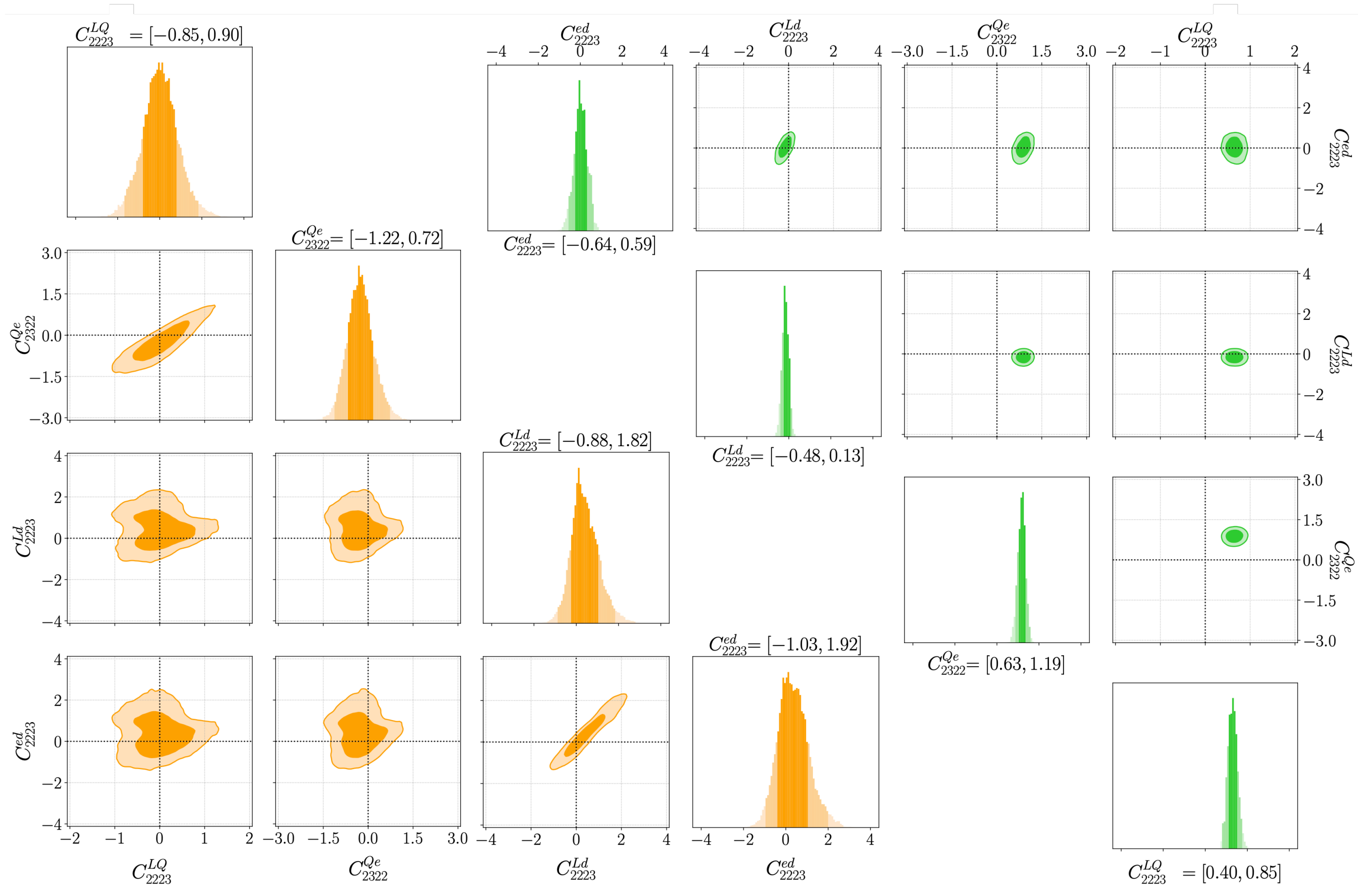

FIG. 5. *Two- and one-dimensional marginalized joint p.d.f. for the set of SMEFT Wilson coefficients* C^{LQ}_{2223}, C^{Qe}_{2322}, C^{ed}_{2223} *and* C^{Ld}_{2223}. *For both panels, we show the* 68% *and* 95% *probability regions in green and orange on the basis of the hadronic approach adopted in the global analysis (see the text for more details).*

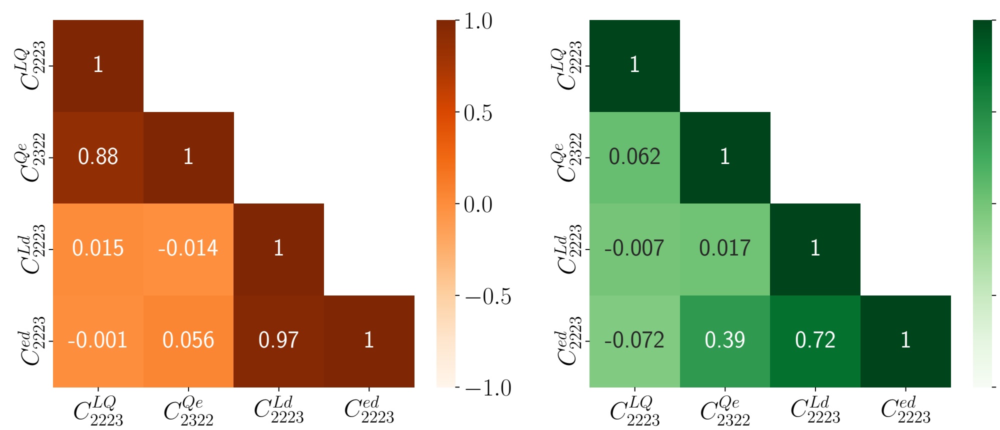

FIG. 6. *Correlation matrix of the Wilson coefficients of the SMEFT operators studied in this work under the "data driven" (left panel, orange) and the "model dependent" (right panel, green) approaches to hadronic uncertainties in our global analysis.*

responding operators with right-handed quark currents $C'^{,\mathrm{NP}}_{9,\mu}$, $C'^{,\mathrm{NP}}_{10,\mu}$. Several interesting features emerge in this fit. First, the updated experimental value of $\mathrm{BR}(B_s \to \mu^+\mu^-)$ forces $C^{\mathrm{NP}}_{10,\mu}$ and $C'^{,\mathrm{NP}}_{10,\mu}$ to be small, corresponding to the correlations visible in the two-dimensional projections on the C^{LQ}_{2223} vs C^{Qe}_{2322} and C^{Ld}_{2223} vs C^{ed}_{2223} planes and reported in Fig. 6. Second, the SM point is well inside the 68% probability regions in the *data driven* approach, while in the *model dependent* scenario there is evidence of a nonvanishing $C^{\mathrm{NP}}_{9,\mu}$, or equivalently of a nonvanishing

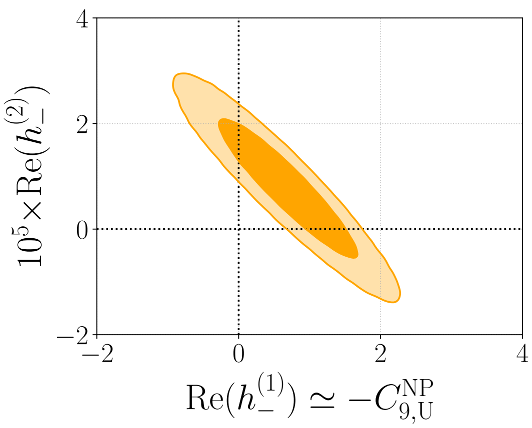

FIG. 7. *Joint posterior p.d.f. for* $\mathrm{Re}(h_-^{(1)})$ *and* $\mathrm{Re}(h_-^{(2)})$ *in a SM fit in the "data driven" scenario. Darker (lighter) regions correspond to* 68% (95%) *probability. Notice that according to our hadronic parameterization given in eq.* (4)*,* $\mathrm{Re}(h_-^{(1)})$ *can be reinterpreted as a flavor universal NP contribution,* $C_{9,\mathrm{U}}^{\mathrm{NP}}$*.*

$C_{2223}^{LQ} \sim C_{2322}^{Qe}$, stemming from BRs and angular distributions of $b \to s\mu^+\mu^-$ transitions. In the *data driven* scenario the latter are reproduced thanks to the charming penguin contributions. Eventually, notice that the allowed ranges for NP coefficients are much larger in the *data driven* scenario since the uncertainties on charming penguins leak into the determination of NP Wilson coefficients.

Before concluding, we comment briefly on the possibility of a lepton universal NP contribution to C_9, that we denote here $C_{9,\mathrm{U}}^{\mathrm{NP}}$, affecting only absolute BRs and angular distributions of $b \to s\ell^+\ell^-$ decays, but leaving LUV ratios as in the SM. This possibility was already discussed in detail in ref. [38], and the experimental situation has not changed since then. Therefore, we just summarize here the main findings of ref. [38] for the reader's convenience. Performing a fit to experimental data within the SM in the *data driven* scenario, one finds that several h_λ parameters are determined to be different from zero at 95% probability, supporting the picture of sizable rescattering in charming penguin amplitudes (see Table 1 in ref. [38]). In particular, there is an interesting correlation between $\mathrm{Re}(h_-^{(1)}) \simeq -C_{9,\mathrm{U}}^{\mathrm{NP}}$ and $\mathrm{Re}(h_-^{(2)})$, as is evident from Fig. 7. Data definitely require a nonvanishing combination of the two parameters; if charming penguins are treated à la [43–48], $\mathrm{Re}(h_-^{(2)})$ is put to zero and $\mathrm{Re}(h_-^{(1)})$ is identified with a lepton universal contribution $C_{9,\mathrm{U}}^{\mathrm{NP}}$, leading to an evidence of NP inextricably linked to the assumptions on charming-penguin amplitudes.

Summarizing, we performed a Bayesian analysis of possible LUV NP contributions to $b \to s\ell^+\ell^-$ transitions in view of the very recent updates on $\mathrm{BR}(B_{(d,s)} \to \mu^+\mu^-)$ by the CMS collaboration [49] and on R_K and R_{K^*} by the LHCb collaboration [54, 55]. As pointed out in refs. [24, 26, 30, 38, 40–42, 90, 91], the NP sensitivity of these transitions is spoilt by possible long-distance effects, see Fig. 1. Thus, in the *data driven* scenario we determined simultaneously hadronic contributions, parameterized according to eq. (4), and NP Wilson coefficients, finding no evidence for LUV NP. Conversely, evidence for NP contributions is found if charming penguins are assumed to be well described by the approach of refs. [43–48], as reported in Tables I and II.

Finally, we considered the case of a lepton universal NP contribution to C_9, which is phenomenologically equivalent to the effect of $h_-^{(1)}$ in our *data driven* analysis, confirming our previous findings in ref. [38]: in the context of the *data driven* approach, we found several hints of nonvanishing h_λ^i parameters, but no evidence of a nonvanishing $\mathrm{Re}(h_-^{(1)}) \simeq -C_{9,\mathrm{U}}^{\mathrm{NP}}$; evidence for $C_{9,\mathrm{U}}^{\mathrm{NP}}$ only arises in the *model dependent* scenario in which all genuine hadronic contributions are phenomenologically negligible. Future improvements in theoretical calculations and in experimental data will hopefully allow clarifying this last point.

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