Weak Decays of Excited B Mesons

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Weak decays of unstable b-mesons

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We investigate the decays of the excited $(b\bar{q})$ mesons as probes of the short-distance structure of the weak $\Delta B = 1$ transitions. These states are unstable under the electromagnetic or strong interactions although their widths are typically suppressed by phase space. As compared to the pseudoscalar B meson, the purely leptonic decays of the vector B^* are not chirally suppressed and are sensitive to different combinations of the underlying weak effective operators. An interesting example is $B_s^* \to \ell^+\ell^-$, which has a rate that can be accurately predicted in the standard model. The branching fraction is $\mathcal{B} \sim 10^{-11}$, irrespective of the lepton flavor and where the main uncertainty stems from the unmeasured and theoretically not-well known B_s^* width. We discuss the prospects for producing this decay mode at the LHC and explore the possibility of measuring the $B_s^* \to \ell\ell$ amplitude, instead, through scattering experiments at the B_s^* resonance peak.

Introduction: Heavy-light systems like the $(b\bar{q})$ mesons have a rich spectrum of excited states [1–4]. These mesons are unstable under electromagnetic or strong interactions, although they can have a narrow width because the mass-splittings in the spectrum are in general much smaller than the mass of the ground-state pseudoscalar B-meson they ultimately decay to. The corresponding lifetimes are of the order of 10^{-17} seconds or less and they typically do not live long enough to directly experience a weak disintegration induced by the b-quark flavor transition. However, with the high luminosities achieved at the e^+e^- colliders [5] and high production rates of $b\bar{b}$ pairs at the LHC, which already allow for sensitivities to branching fractions at the level of 10^{-10} [6], some of these modes may be detected and investigated.

Slightly heavier than the B is the vector B^*-meson and its decays have different sensitivities to the short-distance structure of the $\Delta B = 1$ transitions. Moreover, heavy-quark (HQ) symmetry relates matrix elements of these two mesons [2]. Thus, the interplay between B and B^* decays could prove useful in studies to test the standard model (SM) and search for new-physics (NP). This has immediate interest as various anomalies have been detected in charged- and neutral-current B decays [7–19].

In particular, the LHCb experiment has reported anomalies in $b \to s\ell\ell$ decays [20–24] including a signal of lepton-universality violation [25]; remarkably, they can be largely accommodated by a NP interaction selectively coupled to muons, described by operators of the type [26–32],

$$\mathcal{O}_{9(10)}^\mu = \frac{\alpha_{\rm em}}{4\pi}(\bar{s}\gamma^\mu P_L b)\,(\bar{\mu}\gamma^\mu(\gamma_5)\mu). \qquad (1)$$

Global fits to the $b \to s\ell\ell$ data point to scenarios where the NP contribution to their respective Wilson coefficients are $C_9^{\mu,\,\rm NP} \simeq -1$ or $C_9^{\mu,\,\rm NP} = -C_{10}^{\mu,\,\rm NP} \simeq -0.5$ [26–28, 33–36]. These anomalies suggest the presence of new particles with non-universal lepton couplings and with masses in the TeV range that could be searched at the LHC [27, 31, 37–52].

The interpretation of weak hadron decays is often obscured by the presence of long-distance QCD effects. Hadronic matrix elements of local operators can be parameterized in terms of form factors, whose description relies on the accuracy of different nonperturbative methods [29, 35, 53–61]. One also needs to take into account the "current-current" four-quark operators, $\mathcal{O}_1$ and $\mathcal{O}_2$ [62–64], which have large coefficients (C_1 and C_2), since they stem from the $b \to sc\bar{c}$ transition that is not suppressed by neither mixing angles nor loop factors in the SM. They contribute to $b \to s\ell\ell$ amplitudes through,

$$\mathcal{T}_i^\mu = i \int d^4x\, e^{i\,q\cdot x} T\{\mathcal{O}_i(0)\,, j_{\rm em}^\mu(x)\}, \qquad (2)$$

where the dilepton pair is produced by the off-shell photon from the electromagnetic current. Hadronic matrix elements of $\mathcal{T}_i^\mu$ receive dominant contributions from long-distance fluctuations of the charm-quark fields manifested as charmonium resonances.

At high q^2 one can analytically continue eq. (2) into the complex q^2-plane to perform an operator product expansion (OPE) which accurately describes it in terms of a series of matrix elements of local operators matched perturbatively to QCD [65–67]. Continuing the result back to real q^2 gives the physical rates. This is called "quark-hadron duality" and its validity is justified if q^2 is well above the resonant contributions. Violations to quark-hadron duality can be difficult to estimate and may be important for moderately high q^2. This is the case for the $b \to s\ell\ell$ exclusive decays for which $\sqrt{q^2} \leq m_B - m_K \lesssim 4790$ MeV while the heaviest charmonium state known is the $X(4660)$ [20, 68, 69].

In light of these difficulties, it is desirable to have alternative, theoretically cleaner processes probing the semileptonic operators in eq. (1) to confirm or to unambiguously characterize the putative NP effect. In this paper we investigate the purely leptonic decays of the B^*. In contrast to those of their pseudoscalar siblings, decays of the vector B^* are not chirally suppressed. This partly compensates for the shorter lifetime of the B^* and opens the possibility of probing the short-distance structure of the muonic and electronic decays, affording new tests of lepton-universality.

The $B_s^* \to \ell\ell$ decay rate: The B_s^* is a $(b\bar{s})$ mesons, with $J^{PC} = 1^{--}$, mass $m_{B_s^*} = 5415.4^{+2.4}_{-2.1}$ MeV [70] and an experimentally unknown width, estimated to be ~ 0.1 KeV

(see below). The SM amplitude for $B_s^* \to \ell^+\ell^-$ is:

$$
\mathcal{M}_{\ell\ell} = \frac{G_F}{2\sqrt{2}} \lambda_{ts} \frac{\alpha_{\rm em}}{\pi} \Big[\left(m_{B_s^*} f_{B_s^*} C_9 + 2 f_{B_s^*}^T m_b C_7 \right) \bar{\ell}\not{\varepsilon}\ell
$$
$$
+ m_{B_s^*} f_{B_s^*} C_{10} \bar{\ell}\not{\varepsilon}\gamma_5\ell
$$
$$
- 8\pi^2 \frac{1}{q^2} \sum_{i=1}^{6,8} C_i \langle 0|\mathcal{T}_i^\mu|B_s^*(p,\varepsilon)\rangle \, \bar{\ell}\gamma_\mu\ell \Big], \tag{3}
$$

where G_F is the Fermi constant, $\lambda_{ts} = V_{ts}^* V_{tb}$, $m_b(\mu)$ the running b-quark mass in the $\overline{MS}$ scheme and ε is the polarization vector of the B_s^*. Furthermore, $q^2 = m_{B_s^*}^2$. The information on the short distance structure of the $b \to s$ transition is carried by the (renormalization scale dependent) Wilson coefficients of the weak Hamiltonian for $\Delta B = 1$ processes [62–64]. In particular, $C_{9,10}$ are the ones related to the short-distance semileptonic operators, eq. (1), and C_7 is the coefficient of the "electromagnetic penguin operator" [71]. The operators in the third line of eq. (3), correspond to either the four-quark operators, including those of the current-current, $\mathcal{O}_{1,2}$, and the "QCD-penguins", $\mathcal{O}_{3,...,6}$, or the "chromo-magnetic penguin operator", $\mathcal{O}_8$ [118].

The nonperturbative contributions enter through two types of matrix elements. Those of the local operators $\mathcal{O}_{7,9,10}$ are described by two decay constants,

$$
\langle 0|\bar{s}\gamma^\mu b|B_s^*(p,\varepsilon)\rangle = m_{B_s^*} f_{B_s^*} \varepsilon^\mu,
$$
$$
\langle 0|\bar{s}\sigma^{\mu\nu} b|B_s^*(p,\varepsilon)\rangle = -i f_{B_s^*}^T (p^\mu \varepsilon^\nu - \varepsilon^\mu p^\nu), \tag{4}
$$

where $f_{B_s^*}^T(\mu)$ depends on the renormalization scale. In the HQ limit, these are related to the decay constant of the B_s [55, 72, 73], $\langle 0|\bar{s}\gamma^\mu\gamma_5 b|B_s(p)\rangle = -i f_{B_s} p^\mu$.

The second type of hadronic contribution enters, in the third line of eq. (3), through the matrix element of the operator in eq. (2), induced by all the four-quark and the chromomagnetic operators. At high $q^2 \sim m_b^2$, one can exploit the hierarchy of scales $\Lambda_{\rm QCD} \ll m_c \ll \sqrt{q^2} \sim m_b$ to expand this intrinsically nonlocal object into a series of local operators matched perturbatively to QCD [67]. The two leading operators of the resulting OPE are equivalent to $\mathcal{O}_7$ and $\mathcal{O}_9$ so that their matrix elements are described by the very same nonperturbative quantities $f_{B_s^*}$ and $f_{B_s^*}^T$. In other words, the leading effect in the OPE is implemented by the redefinitions $C_7(\mu) \to C_7^{\rm eff}(\mu, q^2)$ and $C_9(\mu) \to C_9^{\rm eff}(\mu, q^2)$, where the expressions of the matching are known up to next-to-leading order in α_s [67, 74, 75].

A remarkable feature of this OPE is that the subleading operators in the expansion are suppressed by either $\mathcal{O}(\alpha_s \times \Lambda_{\rm QCD}/m_b)$ or $\mathcal{O}(\Lambda_{\rm QCD}^2/m_b^2)$ [67, 68] and are numerically small [68]. Nevertheless, one needs to remember that the OPE is formally performed in the complex q^2 plane, away from the physical cuts and singularities [65–67]; there are the quark-hadron duality violations, not captured by the OPE to any order of α_s or $\Lambda_{\rm QCD}/m_b$ and known to appear in the analytic continuation to the physical region. These are not understood from first principles although it is believed they give

rise to the oscillations characteristic of the resonances and to decrease exponentially into the higher q^2 region [66]. For the kinematics of the B_s^* decay, $q^2 = m_{B_s^*}^2$ is well above the charmonium states (and far below the bottomonium states) where local quark-hadron duality is expected to apply.

The $B_s^* \to \ell\ell$ decay rate in the SM is then:

$$
\Gamma_{\ell\ell} = \frac{G_F^2 |\lambda_{ts}|^2 \alpha_{\rm em}^2}{96\pi^3} m_{B_s^*}^3 f_{B_s^*}^2 \times \tag{5}
$$
$$
\times \left(|C_9^{\rm eff}(m_{B_s^*}^2) + 2\frac{m_b f_{B_s^*}^T}{m_{B_s^*} f_{B_s^*}} C_7^{\rm eff}(m_{B_s^*}^2)|^2 + |C_{10}|^2 \right),
$$

where we have neglected $\mathcal{O}(m_\ell^2/m_{B_s^*}^2)$ contributions. For the implementation of the OPE in the present paper we follow [67] and consider $m_c \ll m_b$, so that an expansion up to $\mathcal{O}(m_c^2/m_b^2)$ is also implied. The relevant loop functions necessary for the matching at $\mathcal{O}(\alpha_s)$ are then obtained from refs. [74] and [76]. For the running Wilson coefficients C_{1-8} of the weak Hamiltonian we use the next-to-leading log results, while for $C_{9,10}$ we include the next-to-next-to-leading corrections calculated in [77]. The resulting renormalization scale dependence of the observables is very small, induced by either $C_9^{\rm eff}(\mu, q^2)$ at $\mathcal{O}(\alpha_s^2 \times C_{1,2}, \alpha_s \times C_{3-6})$ or by the combination $m_b(\mu) f_{B_s^*}^T(\mu) C_7^{\rm eff}(\mu, q^2)$ at $\mathcal{O}(\alpha_s^2)$ [67].

TABLE I: Values for input parameters employed in this work. All are obtained from the PDG averages [70], except for $|\lambda_{ts}|$ which is determined from $|V_{cb}|$ and $|V_{tb}^* V_{ts}|/|V_{cb}|$ following ref. [78], f_{B_s}, which is obtained from the $N_f = 2 + 1$ FLAG average [79] and $f_{B_s^*}/f_{B_s}$ which is taken from the HPQCD calculation in [80].

G_F		$1.1663787(6) \times 10^{-5}$ GeV^{-2}	
$m_b(m_b)$	4.18(3)	$\alpha_s(m_Z)$	0.1184(7)
$m_c(m_c)$	1.275(25)	$\alpha_{\rm em}(m_b)$	1/134
$|\lambda_{ts}|$	0.0416(9)	$f_{B_s^*}/f_{B_s}$	0.953(23)
$m_{B_s^*}$	$5415.4^{+2.4}_{-2.1}$ MeV	f_{B_s}	227.7(4.5) MeV

In Tab. I we show the values of the input parameters relevant for the numerical analysis. With these we obtain $C_9^{\rm eff}(m_b, m_{B_s^*}^2) = 4.560 + i\,0.612$ and $C_7^{\rm eff}(m_b, m_{B_s^*}^2) = -0.384 - i\,0.111$. In the HQ limit and up to $\mathcal{O}(\alpha_s)$ [55], $f_{B_s^*}/f_{B_s} = f_{B_s^*}^T(m_b)/f_{B_s} = 0.95$. Beyond the HQ limit, the $f_{B_s^*}/f_{B_s}$ ratios have been calculated using QCD sum rules [81–86] and, recently, on the lattice by the HPQCD collaboration, $f_{B_s^*}/f_{B_s} = 0.953(23)$ [80]. In this paper we will use this value as a benchmark for our predictions. We are not aware of any computation of the tensor decay constant, $f_{B_s^*}^T(m_b)$. For this, we use the result given in the HQ limit up to $\mathcal{O}(\alpha_s)$ [55] with an uncertainty $\mathcal{O}(\Lambda_{\rm QCD}/m_b) \sim 10\%$.

Our result for the decay rate then follows to be:

$$
\Gamma_{\ell\ell} = 1.12(5)(7) \times 10^{-18} \text{ GeV}, \tag{6}
$$

where the first error stems from the one in the combination of CKM parameters λ_{ts}, and the second from the decay constants added in quadratures. The error from the residual renormalization scale dependence is numerically very small, of the

order of 1% in the range from $\mu = m_b/2$ to $\mu = 2m_b$. Local quark-hadron duality violations to the OPE are estimated using the model of duality violation introduced in [66], fitted to the BES data on the $\sigma(e^+e^- \to \text{hadrons})$ across the charmonium region and adapted to the $b \to s\ell\ell$ [68]. Extrapolating the results of [68] to the region $q^2 \simeq m_{B_s^*}^2 \pm 2$ GeV^2 we observe that the duality-violating corrections to C_9^{eff} are estimated to be less than a 1.5% of its short-distance contribution.

The branching fraction and experimental prospects: The main difficulty for measuring this rare decay is that it has to compete with the dominant disintegration $B_s^* \to B_s\gamma$, which is an electromagnetic transition, although suppressed by a relatively small phase-space. The $B_s^* \to B_s\gamma$ rate is determined by a hadronic transition magnetic moment that can be estimated using HQ and chiral perturbation effective theories from the equivalent decays in the $D^{(*)}$ system [87–89]. A determination along these lines using current experimental data and recent lattice QCD results as input leads to $\Gamma = 0.10(5)$ KeV [119], which is consistent with the results of the earlier analyses and different quark-model calculations [90–92]. The error encompasses uncertainties from the truncation of the chiral and heavy-quak expansions. The conclusions of our study are then hindered by this large uncertainty in Γ; it is important to stress, though, that this concerns a single hadronic quantity that can be calculated in the lattice as recently demonstrated for the D_q^* system in refs. [93, 94]. Progress in this direction is essential for a conclusive assessment on the interest of this mode and for the experimental prospects for its detection and measurement.

With these caveats, we proceed to combine eq. (6) with our estimate of Γ and obtain a branching fraction that in the SM is in the range:

$$\mathcal{B}^{\text{SM}}(B_s^* \to \ell\ell) = 1.12(9) \times 10^{-11} \left(\frac{0.10(5) \text{ KeV}}{\Gamma} \right), \quad (7)$$

which yields $\mathcal{B}^{\text{SM}}(B_s^* \to \ell\ell) = (0.7 - 2.2) \times 10^{-11}$ using $\Gamma = 0.10(5)$ KeV and irrespective of the lepton flavor. This is a very small branching fraction, lying an order of magnitude below $\mathcal{B}^{\text{SM}}(B_d \to \mu\mu)$ [78] and the rarest decay ever detected in an experiment, $K \to \pi\nu\bar{\nu}$ [95].

Given the large production rates of $b\bar{b}$ pairs in high-energy pp collisions, this mode could be searched for at the LHC. Taking into account that 100 $B_s \to \mu\mu$ events have been detected by the combined analysis of the LHCb (3 fb^{-1}) and CMS (25 fb^{-1}) [96] (see e.g. [97] for ATLAS), the increased $b\bar{b}$ integrated production luminosity at $\sqrt{s} = 14$ TeV [98, 99] we estimate that of the order of 10 (100) $B_s^* \to \ell\ell$ decays could be produced by the end of the run III (High-Luminosity phase) of the LHC [120]. Whether or not this could be measured by the LHC experiments will depend on a careful assessment of the backgrounds, but in general, we would expect the signal to manifest as a separate peak to the right of the B_s distribution in the invariant dilepton mass of the $B_{d,s} \to \mu\mu$ measurements. The estimate for $B_s^* \to ee$ differs from the previous one because of the different detector efficiencies for

muons and electrons. Interestingly, the electronic mode has no background from the $B_s \to ee$ decay since this mode is very suppressed.

TABLE II: Branching fractions of purely leptonic $B_i^{(*)}$ decays ($i = u$, c and $\ell = e$, μ) in the SM. B_i^* uncertainties are dominated by their total widths while for B_i they are a few percent of central value.

	$B_i^* \to \ell\bar{\nu}$	$B_i \to e\bar{\nu}$	$B_i \to \mu\bar{\nu}$
$i = u$	$0.6^{+0.3}_{-0.2} \times 10^{-9}$	1.2×10^{-12}	4.9×10^{-7}
$i = c$	$1.3^{+0.4}_{-0.2} \times 10^{-5}$	2.6×10^{-9}	1.6×10^{-4}

The idea of studying the weak disintegrations of the unstable heavy-light systems can be straightforwardly applied to the charged-current leptonic decays of the excited $B_i^{*\pm}$ states, where $i = u$, c. The complementarity between the decays of the purely leptonic decays of the B_i and B_i^* can be explored by modifying the characteristic charged-current $V - A$ interaction of the SM as,

$$\mathcal{L}_{\text{c.c.}} = -\frac{4G_F}{\sqrt{2}} V_{ib} \left((1 + \epsilon_L^{i\ell})(\bar{u}_i\gamma^\mu P_L b)\left(\bar{\ell}\gamma_\mu P_L \nu\right) \right.$$
$$\left. + \epsilon_R^{i\ell}(\bar{u}_i\gamma^\mu P_R b)\left(\bar{\ell}\gamma_\mu P_L \nu\right) \right), \quad (8)$$

where $\epsilon_{L,R}^{\ell i}$ are NP Wilson coefficients encoding left- or right-handed lepton-dependent currents. Contributions of this type are among the possible explanations for the different anomalies found in the $b \to u\ell\nu$ and $b \to c\ell\nu$ transitions [12–19]. The $B_i^{(*)} \to \ell\nu$ decay rates are:

$$\Gamma_{\nu\ell} = \frac{G_F^2}{8\pi}|V_{ib}|^2(1 + \epsilon_L^{i\ell} - \epsilon_R^{i\ell})^2 m_{B_i} f_{B_i}^2 m_\ell^2, \quad (9)$$

$$\Gamma_{\nu\ell}^* = \frac{G_F^2}{12\pi}|V_{ib}|^2(1 + \epsilon_L^{i\ell} + \epsilon_R^{i\ell})^2 m_{B_i^*}^3 f_{B_i^*}^2, \quad (10)$$

where we neglect subleading $\mathcal{O}(m_\ell^2/m_{B_i^*}^2)$ corrections. Extending the effective field theory calculation above to the $B_i^{*\pm}$, we estimate $\Gamma_u = 0.50(25)$ KeV and $\Gamma_c = 0.030(7)$ KeV. In Tab. II we show the SM predictions for the $B_i^{(*)} \to \ell\nu$ branching fractions [121].

Resonant B_s^* production in $\ell^+\ell^-$ scattering: We investigate here a completely different experiment to measure the $B_s^* \to \ell\ell$ rate and we briefly study its feasibility. It consists of producing a B_s^* through resonant $\ell^+\ell^-$ scattering, where the ℓ could either be an electron or a muon. The idea is that the loop- and CKM-suppression of the amplitude is largely compensated by the resonant enhancement in the cross-section from the small width of the B_s^*. Moreover, we expect that the production of a single b or $\bar{b}$ quark at $\sqrt{s} \sim 5.5$ GeV from a $\ell^+\ell^-$ collision would give such a distinct experimental signature that it could be easily disentangled from other electromagnetically produced $\ell^+\ell^- \to$ hadron events.

A calculation of the cross section of $\ell^+\ell^- \to B_s^* \to B_s\gamma$ and its charged-conjugate (we omit CP-violation effects) gives:

$$\sigma(s) = \frac{24\pi m_{B_s^*}^2 (s - m_{B_s}^2)^3}{s(m_{B_s^*}^2 - m_{B_s}^2)^3} \frac{\Gamma_{\ell\ell}\Gamma}{(s - m_{B_s^*}^2)^2 + m_{B_s^*}^2\Gamma^2}, \quad (11)$$

where we have assumed $s \simeq m_{B_s^*}^2$ so that the rates $\Gamma_{\ell\ell}$ and Γ are evaluated for the B_s^* on-shell and have neglected lepton mass effects and non-resonant contributions to the process. It follows that:

$$\sigma_0 = \sigma(m_{B_s^*}^2) = \frac{24\pi}{m_{B_s^*}^2} \mathcal{B}(B_s^* \to \ell\ell) \qquad (12)$$

and using the results in eq. (7), we obtain that $\sigma_0 = (7-22)$ fb, where the large error originates again from Γ. This is a small cross section, characteristic of other weak processes like neutrino-nucleon scattering which occurs at $\sigma_{\nu N} \sim 1-10$ fb.

In order to assess if this process is accessible to experimental study, we need to consider the fact that the energy of the particles in the beams distribute over certain range whose size is quantified by the "energy spread" of the accelerator, ΔE. For current e^+e^- colliders, and for the center-of-mass energies under consideration, $\Delta E_e \sim 1$ MeV, which is much larger than Γ so that only a small fraction of the collisions would occur where the cross-section is maximal. A better control over the energy spread could be achieved at a $\mu^+\mu^-$ collider, although the minimum that has been projected for such hypothetical facility is $\Delta E_\mu \sim 100$ KeV for the energies of interest [100], which is also much larger than Γ.

For the sake of simplicity let us assume that the energy of the particles in the colliding beams spreads uniformly within the interval $m_{B_s^*}/2 \pm \Delta E$, and that $\Gamma \ll \Delta E$. In this case, the average cross-section $\bar{\sigma}$ is:

$$\bar{\sigma} = \frac{1}{(2\Delta E)^2} \int dE_1\, dE_2\, \sigma(s) = \frac{\pi}{4}\frac{\Gamma}{\Delta E}\sigma_0, \qquad (13)$$

(where the integration limits are $m_{B_s^*}/2 \pm \Delta E$) and using the σ_0 and the ΔE discussed above, $\bar{\sigma} \sim 1$ ab and $\bar{\sigma} \sim 10$ ab for the e^+e^- and $\mu^+\mu^-$ colliders, respectively. Producing these processes experimentally might be at reach in the future as, for example, SuperKEKB expects to produce more than 10 ab^{-1}/yr of e^+e^- collisions within the next decade [5]. Finally, we want to note that P-wave orbital excitations of the $b\bar{s}$ system [101–103] could also be produced in resonant $\ell^+\ell^-$ scattering. Their widths are in the range $\Gamma \sim 0.01 - 1$ MeV [1–4, 104–110] and the variety in their quantum numbers can lead to different sensitivities to the short-distance structure of the $b \to s\ell\ell$ weak transition.

Conclusions: The vector B^* states are very narrow resonances because of the phase-space suppression suffered by their dominant electromagnetic decays. The fact that the purely leptonic decays of the B^* are not chirally suppressed compensates for their short lifetimes and the resulting branching fractions are not much smaller (for muons) or are even larger (for electrons) than those of the leptonic decays of the pseudoscalar B mesons.

The $B_s^* \to \ell\ell$ decay is especially interesting since it can provide complementary information on the semileptonic $b \to s\ell\ell$ operators, especially $\mathcal{O}_9$. This decay is theoretically very clean because *(i)* the amplitude only depends on decay constants which are determined accurately in the lattice; and *(ii)*

the invariant mass of the process is well above the charmonium resonances and the application of an operator-product expansion for the nonlocal contributions of eq. (2) via quark-hadron duality (which always accompany the contributions of $\mathcal{O}_9$) is well justified.

The $B_s^* \to \ell\ell$ decay rate can be accurately predicted in the standard model. Using some estimates for the unmeasured width of the B_s^*, we obtained that the branching fraction for this process is $\sim 10^{-11}$ which could be within reach in the next series of experiments at the LHC. More accurate determinations of the width, for example using lattice techniques, are important since this remains the major obstacle for an accurate calculation of the branching fraction of the decay.

The same amplitudes can be measured using a different strategy based on resonant $\ell^+\ell^- \to B_s^* \to B_s\gamma$ scattering. The idea is that the strong suppression of the amplitude is compensated by a large enhancement from the small width of the resonance. Taking into account the energy spread of the beams we estimated that the effective cross-section would be of the order of $1-10$ ab for the current or projected accelerators. Other orbitally excited $(b\bar{s})$ states are also interesting as they have broader widths and can present different sensitivities to the same underlying effective operators.

[1] S. Godfrey and R. Kokoski, Phys. Rev. **D43**, 1679 (1991).

[2] N. Isgur and M. B. Wise, Phys. Rev. Lett. **66**, 1130 (1991).

[3] E. J. Eichten, C. T. Hill, and C. Quigg, Phys. Rev. Lett. **71**, 4116 (1993), hep-ph/9308337.

[4] D. Ebert, V. O. Galkin, and R. N. Faustov, Phys. Rev. **D57**, 5663 (1998), [Erratum: Phys. Rev.D59,019902(1999)], hep-ph/9712318.

[5] A. J. Bevan et al. (Belle, BaBar), Eur. Phys. J. **C74**, 3026 (2014), 1406.6311.

[6] V. Khachatryan et al. (LHCb, CMS), Nature **522**, 68 (2015), 1411.4413.

[7] A. Matyja et al. (Belle Collaboration), Phys.Rev.Lett. **99**, 191807 (2007), 0706.4429.

[8] J. Lees et al. (BaBar Collaboration), Phys.Rev.Lett. **109**, 101802 (2012), 1205.5442.

[9] J. Lees et al. (The BaBar Collaboration), Phys.Rev. **D88**, 072012 (2013), 1303.0571.

[10] Y. Amhis et al. (Heavy Flavor Averaging Group (HFAG)) (2014), 1412.7515.

[11] R. Aaij et al. (LHCb) (2015), 1506.08614.

[12] A. Crivellin, Phys. Rev. **D81**, 031301 (2010), 0907.2461.

[13] A. J. Buras, K. Gemmler, and G. Isidori, Nucl. Phys. **B843**, 107 (2011), 1007.1993.

[14] A. Crivellin and S. Pokorski, Phys. Rev. Lett. **114**, 011802 (2015), 1407.1320.

[15] R. Aaij et al. (LHCb), Nature Phys. **11** (2015), 1504.01568.

[16] S. Fajfer, J. F. Kamenik, and I. Nisandzic, Phys.Rev. **D85**, 094025 (2012), 1203.2654.

[17] D. Becirevic, N. Kosnik, and A. Tayduganov, Phys.Lett. **B716**, 208 (2012), 1206.4977.

[18] M. Tanaka and R. Watanabe (2012), 1212.1878.

[19] M. Freytsis, Z. Ligeti, and J. T. Ruderman (2015), 1506.08896.

[20] R. Aaij et al. (LHCb collaboration), Phys.Rev.Lett. **111**, 112003 (2013), 1307.7595.

[21] R. Aaij et al. (LHCb collaboration), Phys.Rev.Lett. **111**, 191801 (2013), 1308.1707.

[22] R. Aaij et al. (LHCb), JHEP **1307**, 084 (2013), 1305.2168.

[23] R. Aaij et al. (LHCb), JHEP **1406**, 133 (2014), 1403.8044.

[24] R. Aaij et al. (LHCb) (2015), 1506.08777.

[25] R. Aaij et al. (LHCb collaboration), Phys.Rev.Lett. **113**, 151601 (2014), 1406.6482.

[26] S. Descotes-Genon, J. Matias, and J. Virto, Phys.Rev. **D88**, 074002 (2013), 1307.5683.

[27] W. Altmannshofer and D. M. Straub, Eur.Phys.J. **C73**, 2646 (2013), 1308.1501.

[28] F. Beaujean, C. Bobeth, and D. van Dyk, Eur.Phys.J. **C74**, 2897 (2014), 1310.2478.

[29] R. R. Horgan, Z. Liu, S. Meinel, and M. Wingate, Phys.Rev.Lett. **112**, 212003 (2014), 1310.3887.

[30] R. Alonso, B. Grinstein, and J. Martin Camalich, Phys.Rev.Lett. **113**, 241802 (2014), 1407.7044.

[31] G. Hiller and M. Schmaltz, Phys.Rev. **D90**, 054014 (2014), 1408.1627.

[32] D. Ghosh, M. Nardecchia, and S. Renner, JHEP **1412**, 131 (2014), 1408.4097.

[33] T. Hurth, F. Mahmoudi, and S. Neshatpour, JHEP **1412**, 053 (2014), 1410.4545.

[34] W. Altmannshofer and D. M. Straub (2014), 1411.3161.

[35] S. Jäger and J. Martin Camalich (2014), 1412.3183.

[36] W. Altmannshofer and D. M. Straub (2015), 1503.06199.

[37] W. Altmannshofer, S. Gori, M. Pospelov, and I. Yavin, Phys. Rev. **D89**, 095033 (2014), 1403.1269.

[38] B. Gripaios, M. Nardecchia, and S. Renner (2014), 1412.1791.

[39] S. Sahoo and R. Mohanta (2015), 1501.05193.

[40] A. Crivellin, G. D'Ambrosio, and J. Heeck (2015), 1501.00993.

[41] A. Crivellin, G. D'Ambrosio, and J. Heeck, Phys.Rev. **D91**, 075006 (2015), 1503.03477.

[42] I. de Medeiros Varzielas and G. Hiller (2015), 1503.01084.

[43] S. D. Aristizabal, F. Staub, and A. Vicente (2015), 1503.06077.

[44] A. Celis, J. Fuentes-Martin, M. Jung, and H. Serodio (2015), 1505.03079.

[45] S. M. Boucenna, J. W. F. Valle, and A. Vicente (2015), 1503.07099.

[46] A. Crivellin, L. Hofer, J. Matias, U. Nierste, S. Pokorski, et al. (2015), 1504.07928.

[47] D. Bečirević, S. Fajfer, and N. Košnik (2015), 1503.09024.

[48] C.-J. Lee and J. Tandean (2015), 1505.04692.

[49] R. Alonso, B. Grinstein, and J. Martin Camalich (2015), 1505.05164.

[50] A. Greljo, G. Isidori, and D. Marzocca (2015), 1506.01705.

[51] L. Calibbi, A. Crivellin, and T. Ota (2015), 1506.02661.

[52] W. Altmannshofer and I. Yavin (2015), 1508.07009.

[53] N. Isgur and M. B. Wise, Phys. Rev. **D42**, 2388 (1990).

[54] G. Burdman and J. F. Donoghue, Phys. Lett. **B270**, 55 (1991).

[55] A. V. Manohar and M. B. Wise, Camb. Monogr. Part. Phys. Nucl. Phys. Cosmol. **10**, 1 (2000).

[56] M. Beneke and T. Feldmann, Nucl. Phys. **B592**, 3 (2001), hep-ph/0008255.

[57] B. Grinstein and D. Pirjol, Phys. Lett. **B533**, 8 (2002), hep-ph/0201298.

[58] S. Jäger and J. Martin Camalich, JHEP **1305**, 043 (2013), 1212.2263.

[59] R. R. Horgan, Z. Liu, S. Meinel, and M. Wingate, Phys.Rev. **D89**, 094501 (2014), 1310.3722.

[60] S. Descotes-Genon, L. Hofer, J. Matias, and J. Virto, JHEP **1412**, 125 (2014), 1407.8526.

[61] A. Bharucha, D. M. Straub, and R. Zwicky (2015), 1503.05534.

[62] B. Grinstein, M. J. Savage, and M. B. Wise, Nucl.Phys. **B319**, 271 (1989).

[63] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev.Mod.Phys. **68**, 1125 (1996), hep-ph/9512380.

[64] K. G. Chetyrkin, M. Misiak, and M. Munz, Phys.Lett. **B400**, 206 (1997), hep-ph/9612313.

[65] J. Chay, H. Georgi, and B. Grinstein, Phys. Lett. **B247**, 399 (1990).

[66] B. Chibisov, R. D. Dikeman, M. A. Shifman, and N. Uraltsev, Int. J. Mod. Phys. **A12**, 2075 (1997), hep-ph/9605465.

[67] B. Grinstein and D. Pirjol, Phys. Rev. **D70**, 114005 (2004), hep-ph/0404250.

[68] M. Beylich, G. Buchalla, and T. Feldmann, Eur. Phys. J. **C71**, 1635 (2011), 1101.5118.

[69] J. Lyon and R. Zwicky (2014), 1406.0566.

[70] K. Olive et al. (Particle Data Group), Chin.Phys. **C38**, 090001 (2014).

[71] B. Grinstein, R. P. Springer, and M. B. Wise, Phys. Lett. **B202**, 138 (1988).

[72] D. J. Broadhurst and A. G. Grozin, Phys. Rev. **D52**, 4082 (1995), hep-ph/9410240.

[73] S. Bekavac, A. G. Grozin, P. Marquard, J. H. Piclum, D. Seidel, and M. Steinhauser, Nucl. Phys. **B833**, 46 (2010), 0911.3356.

[74] D. Seidel, Phys. Rev. **D70**, 094038 (2004), hep-ph/0403185.

[75] C. Greub, V. Pilipp, and C. Schupbach, JHEP **12**, 040 (2008), 0810.4077.

[76] M. Beneke, T. Feldmann, and D. Seidel, Nucl.Phys. **B612**, 25 (2001), hep-ph/0106067.

[77] P. Gambino, M. Gorbahn, and U. Haisch, Nucl. Phys. **B673**, 238 (2003), hep-ph/0306079.

[78] C. Bobeth, M. Gorbahn, T. Hermann, M. Misiak, E. Stamou, and M. Steinhauser, Phys. Rev. Lett. **112**, 101801 (2014), 1311.0903.

[79] S. Aoki, Y. Aoki, C. Bernard, T. Blum, G. Colangelo, et al., Eur.Phys.J. **C74**, 2890 (2014), 1310.8555.

[80] B. Colquhoun, C. T. H. Davies, R. J. Dowdall, J. Kettle, J. Koponen, G. P. Lepage, and A. T. Lytle (HPQCD), Phys. Rev. **D91**, 114509 (2015), 1503.05762.

[81] M. Neubert, Phys. Rev. **D46**, 1076 (1992).

[82] T. Huang and C.-W. Luo, Phys. Rev. **D53**, 5042 (1996).

[83] P. Gelhausen, A. Khodjamirian, A. A. Pivovarov, and D. Rosenthal, Phys.Rev. **D88**, 014015 (2013), 1305.5432.

[84] S. Narison (2014), 1404.6642.

[85] W. Lucha, D. Melikhov, and S. Simula, in *8th International Workshop on the CKM Unitarity Triangle (CKM2014) Vienna, Austria, September 8-12, 2014* (2014), 1411.3890, URL http://inspirehep.net/record/1328087/files/arXiv:1411.3890.pdf.

[86] W. Lucha, D. Melikhov, and S. Simula, Phys. Rev. **D91**,

116009 (2015), 1504.03017.

[87] P. L. Cho and H. Georgi, Phys. Lett. **B296**, 408 (1992), [Erratum: Phys. Lett.B300,410(1993)], hep-ph/9209239.

[88] J. F. Amundson, C. G. Boyd, E. E. Jenkins, M. E. Luke, A. V. Manohar, J. L. Rosner, M. J. Savage, and M. B. Wise, Phys. Lett. **B296**, 415 (1992), hep-ph/9209241.

[89] I. W. Stewart, Nucl. Phys. **B529**, 62 (1998), hep-ph/9803227.

[90] J. L. Goity and W. Roberts, Phys. Rev. **D64**, 094007 (2001), hep-ph/0012314.

[91] D. Ebert, R. N. Faustov, and V. O. Galkin, Phys. Lett. **B537**, 241 (2002), hep-ph/0204089.

[92] H.-M. Choi, Phys. Rev. **D75**, 073016 (2007), hep-ph/0701263.

[93] D. Becirevic and B. Haas, Eur. Phys. J. **C71**, 1734 (2011), 0903.2407.

[94] G. C. Donald, C. T. H. Davies, J. Koponen, and G. P. Lepage, Phys. Rev. Lett. **112**, 212002 (2014), 1312.5264.

[95] A. V. Artamonov et al. (E949), Phys. Rev. Lett. **101**, 191802 (2008), 0808.2459.

[96] V. Khachatryan et al. (CMS Collaboration, LHCb Collaboration) (2014), 1411.4413.

[97] G. Aad et al. (ATLAS), Phys. Lett. **B713**, 387 (2012), 1204.0735.

[98] L. collaboration (LHCb collaboration), Tech. Rep. CERN-LHCC-2012-007. LHCb-TDR-12, CERN, Geneva (2012), URL http://cds.cern.ch/record/1443882.

[99] C. Collaboration (CMS) (2013).

[100] J.-P. Delahaye et al., in *Community Summer Study 2013: Snowmass on the Mississippi (CSS2013) Minneapolis, MN, USA, July 29-August 6, 2013* (2013), 1308.0494, URL http://inspirehep.net/record/1246163/files/arXiv:1308.0494.pdf.

[101] V. M. Abazov et al. (D0), Phys. Rev. Lett. **100**, 082002 (2008), 0711.0319.

[102] T. Aaltonen et al. (CDF), Phys. Rev. Lett. **100**, 082001 (2008), 0710.4199.

[103] R. Aaij et al. (LHCb), Phys. Rev. Lett. **110**, 151803 (2013), 1211.5994.

[104] A. F. Falk and T. Mehen, Phys. Rev. **D53**, 231 (1996), hep-ph/9507311.

[105] M. Di Pierro and E. Eichten, Phys. Rev. **D64**, 114004 (2001), hep-ph/0104208.

[106] W. A. Bardeen, E. J. Eichten, and C. T. Hill, Phys. Rev. **D68**, 054024 (2003), hep-ph/0305049.

[107] J. Vijande, A. Valcarce, and F. Fernandez, Phys. Rev. **D77**, 017501 (2008), 0711.2359.

[108] E. B. Gregory et al., Phys. Rev. **D83**, 014506 (2011), 1010.3848.

[109] P. Colangelo, F. De Fazio, F. Giannuzzi, and S. Nicotri, Phys. Rev. **D86**, 054024 (2012), 1207.6940.

[110] F. Bernardoni, B. Blossier, J. Bulava, M. Della Morte, P. Fritzsch, N. Garron, A. Grardin, J. Heitger, G. von Hippel, and H. Simma (2015), 1505.03360.

[111] J. E. Bartelt et al. (CLEO), Phys. Rev. Lett. **80**, 3919 (1998), hep-ex/9711011.

[112] J. Lees et al. (BaBar), Phys.Rev.Lett. **111**, 111801 (2013), 1304.5657.

[113] D. Becirevic, B. Blossier, E. Chang, and B. Haas, Phys. Lett. **B679**, 231 (2009), 0905.3355.

[114] W. Detmold, C. J. D. Lin, and S. Meinel, Phys. Rev. **D85**, 114508 (2012), 1203.3378.

[115] F. Bernardoni, J. Bulava, M. Donnellan, and R. Sommer (ALPHA), Phys. Lett. **B740**, 278 (2015), 1404.6951.

[116] T.-W. Chiu, T.-H. Hsieh, C.-H. Huang, and K. Ogawa (TWQCD), Phys. Lett. **B651**, 171 (2007), 0705.2797.

[117] P. Gambino and C. Schwanda, Phys. Rev. **D89**, 014022 (2014), 1307.4551.

[118] For the definitions of these operators in this paper we follow the notations and basis introduced in ref. [64].

[119] We use the one-loop chiral result in [88] and the data on $D^{*\pm} \to D^{\pm}\gamma$ as input [111, 112]. For the the effective pion-B coupling in the loops we use recent lattice calculations, $g_1 \simeq 0.50$ [113–115].

[120] We also need to know the fraction of B_s^* produced by the hadronization of the b-quark as compared to the one for B_s meson. In the HQ limit this can be derived by simple helicity arguments that suggest that the B_s^* are produced 3 times more often than the B_s [55].

[121] For the masses we take the PDG averages [70] and we use lattice calculations for the rest of the hadronic inputs [79, 80, 116]. Finally, we use $|V_{ub}| = 4.13(49) \times 10^{-3}$ [70] and $|V_{cb}| = 0.0424(9)$ [117].