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Journal
Physical Review Letters, 116(14)

ISSN
0031-9007

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Publication Date
2016-04-08

DOI
10.1103/physrevlett.116.141801

Peer reviewed

Weak decays of unstable b-mesons

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Abstract

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. Finally we also discuss the charged-current $B_{u,c}^* \to \ell\nu$ decay which can provide complementary information on the $b \to u\ell\nu$ and $b \to c\ell\nu$ transitions.

I. 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 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 $\sim 10^{-10}$ [6], some of these modes could become accessible to detection and investigation.

Of particular interest is the B^* which is the partner of the B in the heavy-meson doublet of the $(b\bar{q})$ system [2]. The B^* are vectors and their $\Delta B = 1$ decays have different sensitivities to the short-distance structure of the transition as compared to those of the pseudoscalar B mesons. Moreover, the hadronic matrix elements of these two mesons, which give the long-distance contributions to their decays, are related by heavy-quark symmetry. 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 actual and immediate interest as various anomalies have been detected in different charged- and neutral-current B decays. For instance, there is a long-standing discrepancy between the inclusive and exclusive determinations of the Cabibbo-Kobayashi-Maskawa (CKM) matrix element V_{ub} [7] or tensions between the SM predictions and the measured $B \to D^{(*)}\tau\nu$ decay rates [8–11]. These could be explained by NP altering the $V - A$ structure of the charged-current interaction characteristic of the SM [12–19].

Moreover, in the course of the analyses done over run I, the LHCb experiment has reported a series of anomalies in various (neutral-current) rare $b \to s\ell\ell$ decays [20–24] including a signal of lepton-universality violation [25]; remarkably, they can be largely accommodated by a NP contribution to low-energy "semileptonic" operators selectively coupled to the muons, of the type [26–32],

$$\mathcal{O}_9^\mu = \frac{\alpha_{\rm em}}{4\pi}(\bar{s}\gamma^\mu P_L b)\,(\bar{\mu}\gamma^\mu\mu), \qquad \mathcal{O}_{10}^\mu = \frac{\alpha_{\rm em}}{4\pi}(\bar{s}\gamma^\mu P_L b)\,(\bar{\mu}\gamma^\mu\gamma_5\mu). \tag{1}$$

More specifically, 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 degrees of freedom with non-universal lepton couplings and with masses in the TeV range that could be accessible to the direct searches at the LHC [27, 31, 37–52].

Unfortunately, the interpretation of weak hadron decays is often obscured by the presence of long-distance QCD effects whose impact in the analyses needs to be carefully assessed. This is specially true for some of the $b \to s\ell\ell$ anomalies which are found in observables constructed from the decay rates of the semileptonic processes $B \to K^{(*)}\mu\mu$ [20–23, 25] or $B_s \to \phi\mu\mu$ [24]. On one hand there are the hadronic matrix elements of local operators that can be parameterized in terms of functions of the invariant squared dilepton mass q^2 or form factors and whose description relies on the accuracy of different nonperturbative methods [29, 35, 53–61]. On the other, one needs to take into account the "current-current" four-quark operators, $\mathcal{O}_1$ and $\mathcal{O}_2$ [62–64], which in the SM stem from the tree-level decay $b \to sc\bar{c}$. Hence, they come accompanied by large Wilson coefficients, C_1 and C_2, and are not suppressed by either mixing angles or loop factors with respect to the contributions of the semileptonic operators. They contribute to the neutral-current semileptonic decay amplitudes through an operator of the type,

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

produced by the contraction of the $c\bar{c}$ pair with the electromagnetic current and the off-shell photon eventually decaying into the dilepton pair. The hadronic matrix element of these nonlocal operators receives dominant contributions from long-distance fluctuations of the charm-quark fields manifested as charmonium resonances above the $c\bar{c}$ threshold.

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 the real q^2 gives the physical rates. This is called "quark-hadron duality" and its validity is justified if q^2 is large and above the resonant contributions. More care is required when using the OPE in a region with resonances where the violations to quark-hadron duality can be difficult to estimate. This is the case for the $b \to s\ell\ell$ exclusive decays since they are restricted to a region $q^2 \leq (m_B - m_K)^2 \lesssim 22$ GeV^2 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. A paradigmatic example is the $B_{d,s} \to \ell\ell$ decay, which depends on only one hadronic quantity, a decay constant, that has been accurately determined using lattice simulations [70]. The contribution to the amplitude of C_9^ℓ, together with those of C_1 and C_2, vanish due to the conservation of the vector current and the decay rate becomes sensitive only to C_{10}^ℓ.

In this paper we investigate the purely leptonic decays of the B^* which are a particularly interesting class of decays. In contrast to those of their pseudoscalar siblings, with the B^* being vector they are not chirally suppressed. This partly compensates for the shorter lifetime of the B^* and makes them interesting to probe the short-distance structure of the muonic and electronic decays, specially in search for lepton-universality violation effects. They only depend on decay constants, which are calculable functions of the one of the B in the heavy-quark limit and can be accurately computed on the lattice. Interestingly, the neutral-current decay $B_s^* \to \ell\ell$ becomes sensitive to C_9^ℓ while the kinematics of the decay are such that $q^2 \simeq 28$ GeV^2, which is well above the region of the charmonium resonances and the quark-hadron duality-violation to the contributions from C_1 and C_2 is expected to be much less of a concern. We will discuss the prospects for producing this decay mode at the LHC and will also explore the possibility of measuring the $B_s^* \to \ell\ell$ amplitude, instead, through scattering experiments at the B_s^* resonance peak. We finish discussing also the charged-current $B_{u,c}^* \to \ell\nu$ decays which can provide complementary information on the $b \to u\ell\nu$ and $b \to c\ell\nu$ transitions.

II. THE $B_s^* \to \ell\ell$ DECAY

A. Anatomy of the decay amplitude

The state we are interested in is the partner of the B_s in the ground-state spin doublet of $(b\bar{s})$ mesons. Its quantum numbers are $J^{PC} = 1^{--}$, with a mass $m_{B_s^*} = 5415.4^{+2.4}_{-2.1}$ MeV [71] and a width that is experimentally unknown although estimated to be of the order of 0.1 KeV (see Appendix). In the SM, and neglecting electromagnetic corrections, the amplitude of the decay of a B_s^* into a dilepton pair is:

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

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" [72]. The operators in the second 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", $\mathcal{O}_{3,\ldots,6}$, or the"chromo-magnetic penguin operator", $\mathcal{O}_8$. [1]

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,

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

where $f^T_{B_s^*}(\mu)$ depends on the renormalization scale. In the heavy-quark limit, these are related to the decay constant of the B_s [55],

$$f_{B_s^*} = f_{B_s}\left(1 - \frac{2\alpha_s}{3\pi}\right), \qquad f^T_{B_s^*} = f_{B_s}\left[1 + \frac{2\alpha_s}{3\pi}\left(\log\left(\frac{m_b}{\mu}\right) - 1\right)\right], \tag{5}$$

where $\langle 0|\bar{s}\gamma^\mu\gamma_5 b|B_s(q)\rangle = -i f_{B_s}\, q^\mu$ and we have neglected $\mathcal{O}(\alpha_s^2)$ corrections.

The second type of hadronic contribution enters, in the second 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 $f^T_{B_s^*}$. 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, 73, 74].

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 $\mathcal{O}(\Lambda^2_{\rm QCD}/m_b^2)$ [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

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

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.

B. Numerical analysis

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

$$\Gamma_{\ell\ell} = \frac{G_F^2\,|\lambda_{ts}|^2\alpha_{\rm em}^2}{96\pi^3} m_{B_s^*}^3 f_{B_s^*}^2 \left(|C_9^{\rm eff}(m_{B_s^*}^2) + 2\frac{m_b\, f_{B_s^*}^T}{m_{B_s^*}\, f_{B_s^*}}\, C_7^{\rm eff}(m_{B_s^*}^2)|^2 + |C_{10}|^2 \right), \tag{6}$$

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. [73] and [75]. 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 [76]. 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 the relevant input parameters employed in the calculations of this work. All are obtained from the PDG averages [71], except for $|\lambda_{ts}|$ which is determined from $|V_{cb}|$ and $|V_{tb}^* V_{ts}|/|V_{cb}|$ following ref. [77], f_{B_s}, which is obtained from the $N_f = 2+1$ FLAG average [70] and $f_{B_s^*}/f_{B_s}$ which is taken from the HPQCD calculation in [78].

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
$\lvert\lambda_{ts}\rvert$	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 of this work. With these we obtain $C_9^{\rm eff}(m_b, m_{B_s^*}^2) = 4.560 + i\,0.612$ and $C_7^{\rm eff}(m_b, m_{B_s^*}^2) = -0.384 - i\,0.111$. For the decay constants, one can relate them to f_{B_s} using eqs. (5), which have been calculated accurately by different lattice collaborations [70]. One obtains $f_{B_s^*}/f_{B_s} = f_{B_s^*}^T(m_b)/f_{B_s} =$

0.95. Beyond the heavy-quark limit, the $f_{B_s^*}/f_{B_s}$ ratios have been calculated using QCD sum rules [79–84] and, recently, on the lattice by the HPQCD collaboration [78]. Interestingly, most of the QCD sum-rule calculations obtain $f_{B_s^*}/f_{B_s} \simeq 1.00 - 1.15 > 1$ [79–81], while the latest sum-rule study [83, 84] and the lattice computation obtain a value that is consistent with the one in the heavy-quark limit, *viz.* $f_{B_s^*}/f_{B_s} = 0.953(23)$ [78]. In this paper we will use this value as a benchmark for our predictions. It would be important, for further improvements of our analysis, to have independent calculations of this ratio. In addition, we are not aware of any computation of the tensor decay constant, $f^T_{B_s^*}(m_b)$. For this, we use the result given in the heavy-quark limit by the second equation in (5) with an uncertainty $\mathcal{O}(\Lambda_{\rm QCD}/m_b) \sim 10\%$.

Our result for the decay rate then follows to be:

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

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 difficult to estimate, especially because the kinematics of the decay fall in a region where the data of $\sigma(e^+e^- \to \text{hadrons})$ is scarce. In any case, we observe that the few data points in the $\sqrt{s} \in [5, 6]$ GeV region of this process are consistent with the result from QCD [71]. A better estimate could be obtained 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 $b \to s\ell\ell$, as done in [68]. Extrapolating the results of this reference to the region $q^2 \simeq m^2_{B_s^*} \pm 2$ GeV^2 we observe that the duality-violating corrections to $C_9^{\rm eff}$ are estimated to be less than a 1.5% of its short-distance contribution.

C. The branching fraction and prospects for experimental production

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. The latter is suppressed by a relatively small phase space [2] while the vector nature of the B_s^* makes the former not to be chirally suppressed, as it is the case of $B_s \to \mu\mu$. In order to calculate the branching fraction and

[2] The three-momentum of the recoiling particles is $|\vec{k}| = (m^2_{B_s^*} - m^2_{B_s})/(2m_{B_s^*}) = 0.048$ GeV to be compared with the one of the leptonic rare decay, $|\vec{k}| \simeq m_{B_s^*}/2 \simeq 2.71$ GeV. Strong decays of the B_s^* are forbidden by phase space.

study the feasibility of measuring this decay mode we thus need the B_s^* width produced by the electromagnetic transition that is not measured and theoretically not very well known.

The $B_s^* \to B_s\gamma$ rate is determined by a hadronic transition magnetic moment, μ_{bs}, that can be estimated using heavy-quark and chiral perturbation effective theories from the equivalent decays in the $D^{(*)}$ system [85–87]. In the Appendix we define this quantity and show a determination along these lines using current experimental data and recent lattice QCD results as input. We obtain $\Gamma = 0.10(5)$ KeV, which is consistent with the results of the earlier analyses and different quark-model calculations [88–90]. 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 $D^{(*)}$ system in ref. [91]. 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. (7) with our estimate of Γ and obtain a branching fraction that in the SM is in the range:

$$\mathcal{B}^{\rm SM}(B_s^* \to \ell\ell) = 1.12(9)\times 10^{-11}\left(\frac{0.10(5)\,\text{KeV}}{\Gamma}\right) = (0.7-2.2)\times 10^{-11}, \tag{8}$$

irrespective of the lepton flavor, and where we have added in quadratures the uncertainties in the $\Gamma_{\ell\ell}$ rate. This is a very small branching fraction, lying an order of magnitude below $\mathcal{B}^{\rm SM}(B_d \to \mu\mu)$ [77] and the rarest decay ever detected in an experiment, $K \to \pi\nu\bar{\nu}$ [92].

Measuring $B_s^* \to \ell\ell$ is far from the reach of the Super B-factories, as for example, Belle II expects to collect no more than 5×10^8 of B_s^* after 5 ab^{-1} at the $\Upsilon(5S)$ [5]. On the other hand, and given the large production rates of $b\bar{b}$ pairs in high-energy pp collisions, it could be searched for at the LHC. To make an estimate, let us start with the 100 $B_s \to \mu\mu$ events expected after run I by the combined analysis of the LHCb (3 fb^{-1}) and CMS (25 fb^{-1}) [93], while the full results from the ATLAS experiment have not been reported yet (see e.g. [94]). In the course of runs II and III the experiments at CERN will collect ~ 10 times more data [95, 96] and the $b\bar{b}$ production will be boosted further by a factor ~ 2 by the higher cross section at $\sqrt{s} = 14$ TeV; after the high-luminosity (HL-LHC) upgrade, a factor ~ 10 more of data is expected [95, 96]. Rescaling up naively the current $B_s \to \mu\mu$ events, we estimate $\sim 3\times 10^3$ events by the end of run III and $\sim 3\times 10^4$ after the full run in the HL-LHC phase.

In order to use this to estimate the number of $B_s^* \to \ell\ell$ we 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 heavy-quark 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]. This has been confirmed for the B^* system in measurements at the Z^0 peak by LEP [97]. For the B_s^* system this factor is even larger in the production at the $\Upsilon(5S)$ [5, 71]. This means that most of the B_s^0 mesons detected at the LHC should have been produced through a $B_s^*\gamma$ decay. Taking this into account and that the branching fraction of the $B_s \to \mu\mu$ is ~ 300 times larger than (8) we finally estimate that of the order of 10 (100) $B_s^* \to \ell\ell$ decays could be produced by the end of the run III (HL-LHC). 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.

D. Resonant B_s^* production in $\ell^+\ell^-$ scattering

We speculate here about 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:

$$\sigma(s) = \frac{24\pi\, m_{B_s^*}^2}{s}\left(\frac{s - m_{B_s}^2}{m_{B_s^*}^2 - m_{B_s}^2}\right)^3 \frac{\Gamma_{\ell\ell}\Gamma}{(s - m_{B_s^*}^2)^2 + m_{B_s^*}^2\Gamma^2}, \tag{9}$$

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:

$$\sigma_0 = \sigma(m_{B_s^*}^2) = \frac{24\pi}{m_{B_s^*}^2}\mathcal{B}(B_s^* \to \ell\ell) \tag{10}$$

and using the results in eq. (8), we obtain:

$$\sigma_0 = (7 - 22)\,\text{fb}, \tag{11}$$

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 [98], 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:

$$\bar{\sigma} = \frac{6\pi^2}{m_{B_s^*}^2}\frac{\Gamma_{\ell\ell}}{\Delta E} = \frac{\pi}{4}\frac{\Gamma}{\Delta E}\sigma_0, \tag{12}$$

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].

Another interesting possibility is considering the orbital excitations of the B_s, in particular the lighter mesons in which the s-quark is in the P-wave orbital. This corresponds to two almost degenerate heavy-quark doublets which are predicted to have the quantum numbers $J^P = (0^+,\, 1^+)$ and $(1^+,\, 2^+)$, masses of the order of ~ 5.8 GeV and narrow widths, $\Gamma \sim 0.01 - 1$ MeV [1–4, 99–105]. The last pair of axial-vector and tensor states have been identified with the two observed $B_{s1}(5830)$ and $B_{s2}^*(5840)$ resonances [106–108], where the width of B_{s2}^* has also been measured, $\Gamma(B_{s2}^*) = 1.56(13)(47)$ MeV [108].

These resonances could be produced in resonant $\ell^+\ell^-$ scattering. Their widths are closer to the energy spreads achievable in current and projected accelerators and the scattering would enjoy more luminosity over the resonance region, albeit at the cost of a reduction of the resonant enhancement of the cross section. If the leptonic weak rates for these states were of the same order of magnitude as $\Gamma_{\ell\ell}$, one can see from eq. (9) and (10), that the cross-section for the production would scale as Γ/Γ^*, where Γ^* is the corresponding width. Besides this, studying the leptonic rates and amplitudes for the orbital excitations is interesting because their quantum numbers lead to different independent sensitivities to the short-distance structure of the $b \to s\ell\ell$ weak transition. They are theoretically clean processes provided the relevant decay constants can be calculated accurately and their widths determined. The fact that these states are quite heavier than the $B_s^{(*)}$ could also allow for studying the validity of quark-hadron duality in more detail.

III. THE $B^{*-} \to \ell^- \bar{\nu}$ DECAYS

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 $i = u, c$. Similarly to the B_s^*, the vector nature of these resonances partly compensates for the shorter life-time as compared to the same decays of their pseudoscalar partners. Nevertheless, $B_i \to \ell\nu$ is difficult to observe not only because of the chiral suppression of the rates in the SM but also because the neutrino in the final state. Only the $B_u \to \tau\nu$ has been detected [109, 110] while limits at the level of $\mathcal{B} < 10^{-6}\,(95\%$ C.L.$)$ have been placed on the electronic and muonic modes [111]. The decay channels of the B_c remain unmeasured to a large extent but important progress is expected at the LHC [112, 113].

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,

$$\mathcal{L}_{\text{c.c.}} = -\frac{4G_F}{\sqrt{2}} V_{ib} \left((1 + \epsilon_L^{i\ell})(\bar{u}_i \gamma^\mu P_L b)\,(\bar{\ell}\gamma_\mu P_L \nu) + \epsilon_R^{i\ell}(\bar{u}_i \gamma^\mu P_R b)\,(\bar{\ell}\gamma_\mu P_L \nu) \right), \tag{13}$$

where $\epsilon_{L,R}^{\ell i}$ are Wilson coefficients encoding NP left-handed or right-handed currents that could be lepton dependent. 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:

$$\Gamma_{\nu\ell} = \frac{G_F^2}{8\pi}|V_{ib}|^2(1 + \epsilon_L^{i\ell} - \epsilon_R^{i\ell})^2 m_{B_i}\, f_{B_i}^2 m_\ell^2, \quad \Gamma_{\nu\ell}^* = \frac{G_F^2}{12\pi}|V_{ib}|^2(1 + \epsilon_L^{i\ell} + \epsilon_R^{i\ell})^2 m_{B_i^*}^3\, f_{B_i^*}^2, \tag{14}$$

where we neglect subleading $\mathcal{O}(m_\ell^2/m_{B_i^*}^2)$ corrections. The different quantum numbers of the B_i and B_i^* mesons make their amplitudes sensitive to different and orthogonal combinations of the coefficients $\epsilon_{L,R}^{i\ell}$. This can be exploited better by looking at the ratio of branching fractions:

$$R_{i\ell}^* = \frac{\mathcal{B}(B_i^* \to \ell\nu)}{\mathcal{B}(B_i \to \ell\nu)} = \frac{2}{3}\frac{m_{B_i^*}}{m_{B_i}} \left(\frac{f_{B_i^*}}{f_{B_i}}\right)^2 \frac{\tau_{B_i^*}}{\tau_{B_i}} \left(\frac{m_{B_i^*}}{m_\ell}\right)^2 (1 + 4\epsilon_R^{i\ell}) + \mathcal{O}(\epsilon_{L,R}^{i\ell})^2, \tag{15}$$

which are clean observables sensitive to right-handed currents.

TABLE II. Results for the branching fractions of the different charged-current leptonic $B_i^{(*)}$ decays considered in this work ($i = u,\, c$ and $\ell = e,\, \mu$). The uncertainties in the B_i^* decays are dominated by the error of their widths whereas those of the B_i are of a few percent relative to the 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}

In order to know the practical interest of these modes we need to know the width (or lifetime) of the B_i^* mesons, induced by their electromagnetic decay, which can be again estimated as explained in the Appendix, giving $\Gamma_u = 0.50(25)$ KeV and $\Gamma_c = 0.030(7)$ KeV. In Table II we show the subsequent predictions for the branching fractions of the $B_i^* \to \ell\nu$ decays as compared to the electronic and muonic modes of the decays of the B_i mesons. [3] We observe that for the B_u^* state the branching fraction is very small although it is larger than the one of the $B_u \to e\nu$ mode. On the other hand, the branching fraction of the B_c^* state is not unreasonably small. It is only an order of magnitude smaller than the $B_c \to \mu\nu$ mode and still much larger than $B_c \to e\nu$.

IV. 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 could provide a clean window to a class of semileptonic $b \to s\ell\ell$ operators, in particular $\mathcal{O}_9$, that could contain information of new physics at the TeV scales.

The advantage of $B_s^* \to \ell\ell$ over other decays (e.g. semileptonic rare decays) is its theoretical cleanness since *(i)* the amplitude only depends on decay constants which are determined accurately

[3] For the masses of the B_i and B_i^* we take the PDG averages [71] and we use lattice calculations for the rest of the inputs, $f_B = 190.5(4.2)$ MeV [70], $f_B^*/f_B = 0.941(26)$ [78], $m_{B_c^*} = 6315(8)$ MeV, $f_{B_c} = 489$ MeV [114] and assume $f^*_{B_c}/f_{B_c} = 1$. For the CKM matrix elements we use the inclusive determinations $|V_{ub}| = 4.13(49) \times 10^{-3}$ [71] and $|V_{cb}| = 0.0424(9)$ [115].

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. In fact, the cross-section at the mass of the B_s^* is of the same order of magnitude as, for example, the one for neutrino-nucleon scattering. Taking into account the energy spread of the beams reduces the effective cross-section and we estimated that this 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.

The same type of analysis can be extended to the leptonic charged-current decays of the $B_{u,c}^{*\pm}$ mesons studying their complementarity with those of the $B_{u,c}^{\pm}$ mesons. For instance, the sensitivities induced by their different quantum numbers could be used to test for the left-handedness of the charged-current transitions. For the B_u^* state the branching fraction results to be $\sim 10^{-9}$ whereas for the B_c^* state is of the order of 10^{-5} and only an order of magnitude smaller than the $B_c \to \mu\nu$ mode.

Notwithstanding the entertainment value (at least to the present authors) of this investigation, we maintain that the rates of the proposed experiments are not ridiculously small; the creativity and prowess of the experimenter should not be discounted. [4]

[4] During a seminar right after publication of [62] one of the authors was ridiculed for the preposterous notion that $B \to X_s\ell\ell$ would ever be measured.

V. ACKNOWLEDGMENTS

We thank U. Egede for encouragement and discussions. This work was supported in part by DOE grant 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.

Appendix A: $B_s^* \to B_s\gamma$ and the B_s^* width

The $B_s^* \to B_s\gamma$ decay rate (and B_s^* width) can be estimated in a model-independent way using heavy-quark and chiral effective theories [85–87]. The amplitude of this transition is:

$$\mathcal{M}_\gamma = \langle B_s(q-k)|j^\mu_{\text{e.m.}}|B_s^*(q,\,\varepsilon)\rangle\eta^*_\mu = e\,\mu_{bs}\,\epsilon^{\mu\nu\rho\sigma}\eta^*_\mu q_\nu k_\rho \varepsilon_\sigma, \tag{A1}$$

where e is the electric charge, η (k) and ε (q) are the polarization vectors (four-momenta) of the photon and the B_s^* respectively, and with μ_{bs} a nonperturbative magnetic moment. The electromagnetic decay rate then follows as:

$$\Gamma_\gamma = \frac{\alpha_{\text{em}}}{3}\mu_{bs}^2\,|\vec{k}|^3. \tag{A2}$$

The magnetic moment can be separated into two components, $\mu_{bs} = \mu_b + \mu_s$. The first one, μ_b, is obtained in heavy-quark effective theory simply as the magnetic moment of the b-quark appearing in the effective Lagrangian at $\mathcal{O}(\Lambda_{\text{QCD}}/m_b)$ [85]:

$$\mu_b = -\frac{1}{3\,m_b}. \tag{A3}$$

The light component, μ_s, involves the long-distance, heavy-quark spin conserving contributions of the light-quarks which are described by pion and kaon fluctuations coupled to the heavy hadron within the framework of chiral perturbation theory [86]. At leading order in the chiral expansion, μ_s, μ_u and μ_d are related by $SU(3)$ flavor symmetry to a single nonperturbative parameter, $\mu_l = Q_l\beta$ where Q_l is the electric charge of the light quark l. Leading $SU(3)$-breaking corrections are given at the next-to-leading order by pion and kaon loops [86]:

$$\delta\mu_u = -\frac{g_1^2\,m_\pi}{4\pi f_\pi^2} - \frac{g_1^2\,m_K}{4\pi f_\pi^2},\quad \delta\mu_d = \frac{g_1^2\,m_\pi}{4\pi f_\pi^2},\quad \delta\mu_s = \frac{g_1^2\,m_K}{4\pi f_\pi^2}, \tag{A4}$$

where $f_\pi \simeq 131$ MeV is the pion semileptonic decay constant, and $m_\pi \simeq 139$ MeV and $m_K \simeq 495$ MeV are the meson masses. The g_1 is the effective coupling of the pseudoscalar and heavy mesons and which has been obtained from lattice calculations, $g_1 \simeq 0.50$ [116–118].

In the heavy quark-limit β relates the magnetic transitions of the B^* mesons to those in the charm sector, where experimental information is available. In particular:

$$\Gamma(D^{*\pm} \to D^\pm\gamma) = \Gamma(D^{*\pm}) \times \mathcal{B}(D^{*\pm} \to D^\pm\gamma) = 1.33(33)\ \text{ KeV}, \tag{A5}$$

where we have used the results obtained by BaBar [119] on the $D^{*\pm}$ width and the CLEO results for the branching fraction [120]. Equating eq. (A2) to this experimental result one obtains that $\mu_{cd} = -0.46(5)$ GeV^{-1}, and then, $\beta = 3.41(16)$ GeV^{-1}. This value for μ_{cd} together with the one readily obtained as a prediction for $\mu_{cu} = 2.1(1)$ GeV^{-1}, compare well with the results of the recent lattice calculation of ref. [91], $\mu_{cd} = -0.2(3)$ GeV^{-1} and $\mu_{cu} = 2.0(6)$ GeV^{-1}. With this, we are now ready to predict $\mu_{bs} = -0.64(5)$ GeV^{-1}, which translates into $\Gamma_\gamma \simeq 0.11(1)$ KeV.

This value is consistent with the predictions obtained using a similar formalism and older data [85–87] and with those in various quark models [88–90]. Nevertheless, and beyond the experimental uncertainties, we expect this result to receive sizable corrections from the chiral expansion or breaking the heavy-quark symmetry (e.g. the heavy-quark dependence of the constant β or through recoil corrections of the heavy mesons). These scale like $\mathcal{O}(m_K^2/\Lambda_{\chi\text{SB}}^2)$ and $\mathcal{O}(\Lambda_{\text{QCD}}/m_c)$ respectively, each of which could be as large as a 25% correction. A manifestation of this problem is the large size of the kaon loops (about 1/2 of the total contribution) which makes our results very sensitivity to the exact value of g_1 [116–118], or to whether one implements higher $SU(3)$-breaking corrections phenomenologically by using $f_K = 1.22 f_\pi$ in the kaon loops [86] or not. With all this in mind, we will use,

$$\Gamma_\gamma = 0.10(5)\ \text{ KeV}, \tag{A6}$$

for the phenomenological discussion of this paper.

This analysis can be extended to calculate the electromagnetic decay rates of the $B_{u,c}^{*\pm}$ mesons. In case of the B_u^* one simply implements the contribution of the light quark $\mu_u = 2/3\,\beta + \delta\mu_u$, giving $\Gamma_\gamma = 0.50(25)$ KeV. For the case of the B_c^*, the contributions from the two heavy quarks are of the type induced by eq. (A3), leading to $\Gamma_\gamma = 0.030(7)$ KeV, where the error has been estimated as of $\mathcal{O}(\Lambda_{\rm QCD}/m_c)$.

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