

## Lifetime of $B_c^-$ Mesons Constrains Explanations for Anomalies in $B \rightarrow D^{(*)}\tau\nu$

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We investigate a new constraint on new-physics interpretations of the anomalies observed in  $B \rightarrow D^{(*)}\tau\nu$  decays making use of the lifetime of the  $B_c^-$  meson. A constraint is obtained by demanding that the rate for  $B_c^- \rightarrow \tau^- \bar{\nu}$  does not exceed the fraction of the total width that is allowed by the calculation of the lifetime in the standard model. This leads to a very strong bound on new-physics scenarios involving scalar operators since they lift the slight, but not negligible, chiral suppression of the  $B_c^- \rightarrow \tau^- \bar{\nu}$  amplitude in the standard model. The new constraint renders a scalar interpretation of the enhancement measured in  $R_{D^*}$  implausible, including explanations implementing extra Higgs doublets or certain classes of leptoquarks. We also discuss the complementarity of  $R_{D^{(*)}}$  and a measurement of the longitudinal polarization of the  $\tau$  in the  $B \rightarrow D^*\tau\nu$  decay in light of our findings.

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*Introduction.*—The last decade has witnessed a very rapid development of the experimental measurement and theoretical understanding of many  $B$ -meson transitions. Although no observations of new particles have been reported at the LHC, some discrepancies with the predictions of the standard model (SM) have started to appear in semileptonic  $B$  decays [1–8]. For instance, measurements of the charged-current  $b \rightarrow c\tau\nu$  transitions turn out to be enhanced with respect to the SM. These have been measured through the lepton-universality ratios  $R_{D^{(*)}} = \Gamma(B \rightarrow D^{(*)}\tau\nu)/\Gamma(B \rightarrow D^{(*)}\ell\nu)$  (henceforth  $\ell$  stands for the muon or the electron) by two different experiments for the  $B \rightarrow D\tau\nu$  channel, BABAR [4,5] and Belle [6–8], and also by LHCb [9] for the  $B \rightarrow D^*\tau\nu$  one. On the other hand, the theoretical predictions are very accurate as they only rely on parametrizations of the experimental spectra of the  $B \rightarrow D^{(*)}\ell\nu$  decays with form factors described within the heavy-quark expansion and incorporating constraints from unitarity [10–13]. Moreover, calculations of the relevant form factors beyond the zero-recoil limit in lattice QCD have started to appear [14–16]. The global significance of these anomalies currently stands at  $\sim 4\sigma$  and they have been addressed in many different models of new physics (NP) [17–59].

An efficient strategy to analyze the possible NP scenarios is using a bottom-up approach in effective field theory (EFT), which starts with the most-general effective

Lagrangian valid at the energy scale characteristic of the physical decay process,  $\mu \sim m_B$ ,

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & -\frac{4G_F V_{cb}}{\sqrt{2}} [(1 + \epsilon_L)\bar{\tau}\gamma_\mu P_L \nu_\tau \bar{c}\gamma^\mu P_L b \\ & + \epsilon_R \bar{\tau}\gamma_\mu P_L \nu_\tau \bar{c}\gamma^\mu P_R b + \epsilon_T \bar{\tau}\sigma_{\mu\nu} P_L \nu_\tau \bar{c}\sigma^{\mu\nu} P_L b \\ & + \epsilon_{S_L} \bar{\tau} P_L \nu_\tau \bar{c} P_L b + \epsilon_{S_R} \bar{\tau} P_L \nu_\tau \bar{c} P_R b] + \text{H.c.} \quad (1) \end{aligned}$$

In the construction of this low-energy EFT (LEEFT) we have used SM particles only. Furthermore, if there is a mass gap between the NP scale,  $\Lambda$ , and the electroweak-symmetry breaking (EWSB) scale,  $v$ , as the current empirical evidence seems to suggest, then one can express the low-energy operators in terms of  $SU(2)_L \times U(1)_Y$ -invariant ones [60–62]. In the context of this EFT of the SM (SMEFT), it follows that, at leading order in the decoupling limit ( $\sim v^2/\Lambda^2$ ),  $\epsilon_R$  is independent of the charged-lepton flavor [63] and it cannot explain lepton-universality violation in  $R_{D^{(*)}}$  [42] (see Ref. [64] for a discussion of the consequences in the EFT arising with a nonlinear realization of EWSB).

Nonetheless, a model that has been very popular in the interpretation of these anomalies is the two-Higgs-doublet model (2HDM) [65], which adds an extra Higgs doublet to the SM and induces effective scalar couplings  $\epsilon_{S_{L,R}}$  mediated by a charged Higgs,  $H^\pm$ . Focusing on implementations with discrete symmetries to avoid flavor-changing neutral currents, the resulting LEEFT of Eq. (1) satisfies the minimal-flavor-violation and minimal-lepton-flavor-violation hypothesis, and as a result  $\epsilon_{L,R}$  are lepton-flavor invariant up to order  $(m_\tau/v)^2$ , while  $\epsilon_{S_{L,R}}^1 \propto m_l/v$  break lepton flavor at leading order in the lepton mass. This naturally explains the characteristic flavor structure required from the interaction, coupling selectively to the

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heavy leptons to explain  $R_{D^{(*)}}$  and secluded from the very stringent tests of lepton universality involving light quarks done with pion, kaon, and nuclear (neutron)  $\beta$  decays [66]. In fact, collider and low-energy flavor data still allow for all the Higgs bosons emerging in the 2HDM to have masses at the EWSB scale [67]. If one only adds this light-Higgs doublet to the particle content of the SM, and still assumes there is a mass gap with respect to any other NP scale, then the resulting 2HDM EFT [68] still cannot produce a nonlepton-universal coupling  $\epsilon_R$  at leading order in the decoupling limit.

Another NP solution to the anomalies that has been discussed extensively in the literature invokes the presence of leptoquarks. As colored particles, the leptoquarks should be produced copiously in  $pp$  collisions at the LHC, which leads to stringent lower bounds on their mass currently reaching the TeV scale. This justifies studying their effects at lower energies by matching first the corresponding model to the SMEFT [42].

Thus, within these assumptions in the EFT framework (encompassing both the SMEFT and the 2HDM EFT) one has a model-independent parametrization of any possible NP contribution in the amplitudes and observables in terms of the four Wilson coefficients  $\epsilon_L$ ,  $\epsilon_{S_L}$ ,  $\epsilon_{S_R}$ , and  $\epsilon_T$ . In order to narrow down the shape of the putative NP effect and discriminate among different models, new observables are necessary that have different sensitivities to the various  $\epsilon_i$ . Examples that have been proposed in the literature include the polarization observables related to the  $\tau$  in the decay [8,29,69,70], the kinematic distributions [5,43], in particular, in terms of the final observable decay products [71–76] or all the very same observables appearing in the equivalent decay modes of other bottom hadrons [77], such as  $B_s$  [78],  $B_c$  [79], and  $\Lambda_b$  [80].

In this work, we investigate a powerful constraint on NP that can be derived from the lifetime of the  $B_c$  meson [53]. This is obtained by demanding that the total rate of the  $B_c^- \rightarrow \tau^- \bar{\nu}_\tau$  decay does not exceed the fraction of the total width that is allowed by the calculation of the lifetime in the SM. Since the chirally flipped operators lift the chiral suppression that the amplitude receives in the SM, one obtains a very strong bound on the scalar operators. We discuss how this new observation puts considerable stress on any “scalar interpretation” of  $R_{D^*}$ , such as those produced by the 2HDM and certain leptoquark models. Finally, we discuss the interplay also between  $R_{D^{(*)}}$  and the longitudinal polarization of the  $\tau$  in the  $B \rightarrow D^* \tau \nu$  decay in light of our findings and the recent measurement of Belle [8]. [For all the calculations related to the  $B \rightarrow D^{(*)} \tau \nu$  decay modes we follow [75].]

*Interplay between the  $B_c$  lifetime and  $R_{D^*}$ .*—Parity conservation of the strong interactions implies that only the hadronic form factors for the pseudoscalar (scalar) combination  $\epsilon_P = \epsilon_{S_R} - \epsilon_{S_L}$  ( $\epsilon_S = \epsilon_{S_R} + \epsilon_{S_L}$ ) can contribute to the  $R_{D^*}$  ( $R_D$ ) mode (see for instance [81]). This

immediately links the  $R_{D^*}$  anomaly to the tauonic decay of the  $B_c^-$ ,  $B_c^- \rightarrow \tau^- \bar{\nu}_\tau$ , for the latter also receives a contribution from  $\epsilon_P$  with a slight (but non-negligible) enhancement with respect to the SM produced by the chirally flipped nature of the scalar operators. Namely, from the effective Lagrangian in Eq. (1) one obtains [66]

$$\text{Br}(B_c^- \rightarrow \tau \bar{\nu}_\tau) = \tau_{B_c} \frac{m_{B_c} m_\tau^2 f_{B_c}^2 G_F^2 |V_{cb}|^2}{8\pi} \left(1 - \frac{m_\tau^2}{m_{B_c}^2}\right)^2 \times \left|1 + \epsilon_L + \frac{m_{B_c}^2}{m_\tau(m_b + m_c)} \epsilon_P\right|^2, \quad (2)$$

where  $f_{B_c} = 434(15)$  MeV [82] is the  $B_c$  decay constant and  $m_b$  and  $m_c$  are the quark masses in the  $\overline{\text{MS}}$  evaluated at  $\mu = 2$  GeV (and so is  $\epsilon_P$ ). Note that  $\epsilon_T$  does not enter the decay as opposed to  $\epsilon_R$ , which is not included since we study lepton-universality violation. Therefore, any experimental measurement or bound on  $\text{Br}(B_c^- \rightarrow \tau \bar{\nu}_\tau)$  can help constrain  $\epsilon_P$ , and to a lesser extent,  $\epsilon_L$ .

The  $B_c$  enjoys high production rates at the LHC [83] and precise measurements of the lifetime [84,85] or the branching ratios of many of its nonleptonic exclusive decay modes are now available [86,87]. However, a measurement of the tauonic decay branching fraction, giving direct access to the NP operators in Eq. (2), seems to be out of reach [83]. An indirect constraint on NP can be extracted, though, from the experimentally measured  $B_c$  lifetime itself [53], whose PDG value is [85]

$$\tau_{B_c} = 0.507(8) \text{ ps}, \quad (3)$$

and the fact that it should be mainly accounted for by  $b$  and  $c$  decays in the  $B_c$  meson [88,89].

Indeed, the lifetime of the  $B_c$  can be calculated in the SM computing the inclusive  $\sum b \rightarrow c$ ,  $\sum \bar{c} \rightarrow \bar{s}$  and annihilation decay rates using an operator product expansion (OPE) matched to nonrelativistic QCD [89–91], which leads to results consistent with the experimental measurement [90],

$$\tau_{B_c}^{\text{OPE}} = 0.52_{-0.12}^{+0.18} \text{ ps}. \quad (4)$$

For the central value, the width of the  $B_c$ ,  $\Gamma_{B_c}^{\text{OPE}} = 1/\tau_{B_c}^{\text{OPE}}$  is distributed among modes induced by the partonic transitions  $\bar{c} \rightarrow \bar{s} \bar{u} d$  (47%),  $\bar{c} \rightarrow \bar{s} \bar{\ell} \bar{\nu}$  (17%),  $b \rightarrow c \bar{u} d$  (16%),  $b \rightarrow c \bar{\nu} \ell$  (8%), and  $b \rightarrow c \bar{c} s$  (7%). Hence, according to this calculation, only  $\lesssim 5\%$  of the measured experimental width,  $\Gamma_{B_c} = 1/\tau_{B_c}$ , can be explained by (semi)tauonic modes, including those NP effects that would explain the  $R_{D^{(*)}}$  anomalies. None of the NP operators appear enhanced in the semitauonic  $b \rightarrow c \tau \bar{\nu}$  rates, which add up to about 1% of the rate in the SM [90]. (Hence, inclusion of these modes in the analysis improves the constraints derived from the  $B_c$  lifetime only marginally.) Similar conclusions are derived from the other OPE calculations [89,91] as well as from

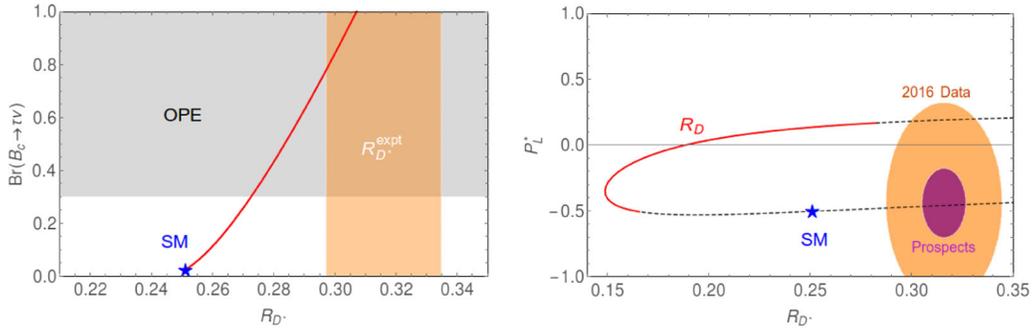


FIG. 1. Left panel: Parametric dependence of  $R_{D^*}$  and  $\text{Br}(B_c \rightarrow \tau\nu)$  for a pseudoscalar NP interaction (red line). The shaded areas are the  $1\sigma$  band corresponding to the measurement of  $R_{D^*}$  (vertical orange) and to the bound on the NP contribution to the lifetime of the  $B_c$  assuming that the SM accounts for 70% of it (gray horizontal). Right panel: The (black) dashed line represents the parametric  $\epsilon_T$ -dependence of  $R_{D^*}$  and  $P_L^*$ . The overlaid (red) solid line corresponds to values of  $\epsilon_T$  for which  $R_D$  is consistent with the experimental measurement at  $1\sigma$ . The shaded areas are the  $1\sigma$  bands corresponding to the current data set (orange) and naive experimental prospects discussed in the main text (purple).

those obtained using potential models [88] or QCD sum rules [92] (see Ref. [83] for a comparison among the different predictions). This constraint can be relaxed up to a  $\lesssim 30\%$  of  $\Gamma_{B_c}$  if the longer lifetime given by the upper limit in Eq. (4) is taken as an input for the SM calculation in the OPE.

The discussion above is immaterial for a NP scenario involving the left-handed contribution  $\epsilon_L$ , which explains  $R_{D^{(*)}}$  by an overall  $\sim 15\%$  enhancement of the SM at the amplitude level (see below) and modifies the corresponding branching fraction for  $B_c^- \rightarrow \tau^- \bar{\nu}$  ( $\sim 2\%$ ) well below these limits. On the other hand, the  $B_c$  lifetime becomes a very severe constraint for scenarios aiming to explain  $R_{D^*}$  with pseudoscalar contributions. Indeed, fitting the average of the experimental determinations [93],

$$R_{D^*}^{\text{expt}} = 0.316(16)(10), \quad (5)$$

to  $\epsilon_P$  we find the following  $1\sigma$  range,

$$\epsilon_P = 1.48(34), \quad (6)$$

which is consistent with the fit in the very same scenario reported in [27] and where we have neglected another solution that interferes destructively with the SM,  $\epsilon_P \simeq -4.3$ . The value in Eq. (6), together with the enhancement factor  $m_{B_c}^2/[m_\tau(m_b + m_c)] \simeq 4$ , leads to a boost of the  $B_c^- \rightarrow \tau^- \bar{\nu}$  branching fraction by a factor of no less than  $\sim 25$ , which not only exceeds the limits discussed above but can also be larger than the experimental width of the  $B_c$ .

This is illustrated in Fig. 1 where we plot the parametric dependence of  $R_{D^*}$  and  $\text{Br}(B_c \rightarrow \tau\nu)$  with respect to the pseudoscalar operator. The maximum value of  $R_{D^*}$  that can be achieved without violating  $\text{Br}(B_c \rightarrow \tau\nu) \lesssim 30\%$  is  $R_{D^*} \simeq 0.27$ , corresponding to the value  $\epsilon_P^{\text{max}} = 0.61$ , and describing the experimental value for  $R_{D^*}$  in Eq. (5) would

saturate (or even exceed) the experimental width of the  $B_c$  with the rate of the tauonic decay. The conclusion is that the  $B_c$  lifetime makes highly implausible any explanation of  $R_{D^*}$  induced by  $\epsilon_P$ .

In fact, this type of scenario encompasses models of NP that have been popular in the interpretation of the  $R_{D^{(*)}}$ , such as the 2HDM [28–38]. In light of the constraint from the  $B_c$  lifetime, any of the 2HDM interpretations of  $R_{D^*}$  is ruled out, in particular, those realizations of the model based on the so-called type III, with general Yukawa couplings to the different fermions [65], which can explain simultaneously  $R_D$  and  $R_{D^*}$  [32,34–38] unlike the 2HDM tested by BABAR [4,5] and Belle [6]. Model-building possibilities invoking leptoquarks are also subject to the bound from the lifetime of the  $B_c$  meson if they generate sizable scalar operators (see for instance Ref. [53] for a first discussion of this).

*Interplay between  $P_L^*$  and  $R_{D^{(*)}}$ .*—The analysis of the  $B_c$  lifetime leaves only  $\epsilon_L$  and  $\epsilon_T$  as possible mechanisms to explain the  $R_{D^*}$  anomaly. As discussed succinctly above, the left-handed scenario leads to a universal enhancement with respect to the SM of all the  $b \rightarrow c\tau\nu$  decay rates by the very same  $\sim 30\%$  that is required to simultaneously explain  $R_D$  and  $R_{D^*}$ . Moreover, observables of a given semitaonic decay, normalized by the corresponding total rate of the same decay, must be insensitive to  $\epsilon_L$  because the enhancements cancel in the ratio. The tensor (and scalar) contribution, on the other hand, has a different structure that is manifest not only in the total rate but also in the kinematic distributions of the decay [17,22,23,66,75]. Therefore, normalized observables together with  $R_{D^{(*)}}$  can discriminate between the two NP scenarios.

To illustrate our argument let us start by fitting the tensor contribution to  $R_{D^*}^{\text{expt}}$  together with the experimental average of  $R_D$  [93],

$$R_D^{\text{expt}} = 0.397(40)(28). \quad (7)$$

Adding the statistical and systematic error of the experimental results in quadratures and taking into account their relative negative correlation,  $\rho = -0.21$  [93], we find the solution

$$\epsilon_T = 0.377(12), \quad (8)$$

with a  $\chi^2_{\text{d.o.f.}} = 1.49$ . In contrast, in the left-handed scenario we find  $\epsilon_L = 0.13(3)$  at the minimum  $\chi^2_{\text{d.o.f.}} = 0.013$ .

One example of the normalized observables defined above is the longitudinal polarization of the  $\tau$  in the decay, which is defined as [29]

$$dP_L^{(*)} = \frac{d\Gamma_+ - d\Gamma_-}{d\Gamma_+ + d\Gamma_-}, \quad (9)$$

where  $d\Gamma_{\pm}$  denotes generically a differential decay rate (e.g., in  $q^2$ ) of  $B \rightarrow D^{(*)}\tau\nu$  and with the helicity of the  $\tau$  being  $\lambda_{\tau} = \pm$ . The longitudinal polarization can be measured using the kinematic distributions of the detected decay products of the  $\tau$  as polarimeters [69,70] and a very first measurement of the integrated longitudinal polarization for the  $BD^*$  channel,  $P_L^*$ , has been reported using the  $\tau \rightarrow \pi\nu$  and  $\tau \rightarrow \rho\nu$  modes in the full Belle data set [8],

$$P_L^{*\text{expt}} = -0.44(47)_{-0.17}^{+0.20}, \quad (10)$$

which is consistent with the SM prediction,  $P_L^{*\text{SM}} = -0.504(24)$ . Using the result of the fit obtained in Eq. (8), we obtain

$$P_L^* = 0.190(10). \quad (11)$$

This prediction of the tensor scenario leads to a  $P_L^*$  with the opposite sign compared to the SM and stands at about  $1\sigma$  from the experimental measurement in Eq. (10). To be more quantitative, we can add  $P_L^*$  to the fit of  $\epsilon_T$  to the experimental averages of  $R_{D^{(*)}}$  performed above. Assuming no further correlation and adding the errors in Eq. (10) in quadratures, neither the position of the minima in Eq. (8) nor the  $\chi^2_{\text{d.o.f.}}$  quoted below for the tensor case change. This suggests that this scenario is disfavored with respect to the left-handed one, although more precise measurements of  $P_L^*$  (or other normalized observables) would be essential to confirm or refute this conclusion with a higher statistic significance. This is illustrated on the right-hand panel of Fig. 1, where we display the parametric dependence of  $R_{D^*}$  and  $P_L^*$  as functions of  $\epsilon_T$  compared to the current experimental measurements. We also show prospects estimated naively assuming an order-of-magnitude increase in statistics and improving the systematics by a factor of 2.

We end the discussion by pointing out a caveat in our latter conclusions. The tension between the ‘‘tensor interpretation’’ of the anomalies and data can be relaxed if one adds the coefficient  $\epsilon_S$ , which exclusively contributes to the

$B \rightarrow D\tau\nu$  channel, to the fit. One can, indeed, see that in this case  $R_{D^*}$  can be explained by a small negative value of  $\epsilon_T$  that interferes constructively with the SM in the  $B \rightarrow D^*\tau\nu$  channel and does not contribute to  $P_L^*$  significantly. This solution interferes destructively in the  $B \rightarrow D\tau\nu$  channel, producing a deficit in  $R_D$  that needs a moderate value of  $\epsilon_S$  to make it consistent with the experimental value. The effect of this combined  $\epsilon_T - \epsilon_S$  scenario is represented by the dashed-line that develops to the right of the SM in the right panel of Fig. 1. It can be distinguished from the other scenarios above efficiently only by looking for the effects of  $\epsilon_S$  in precise measurements of normalized observables of the  $B \rightarrow D\tau\nu$  channel such as  $P_L$ , or the kinematic distributions like the  $q^2$  spectrum.

*Conclusions.*—We have discussed a powerful constraint on possible new-physics interpretations of the  $R_{D^{(*)}}$  anomalies that is derived from the lifetime of the  $B_c$  meson. This is obtained by demanding that the contribution to the rate of the  $B_c \rightarrow \tau\nu$  decay channel does not exceed the fraction of the total width that is allowed by the calculation of the lifetime in the standard model. This leads to a very strong bound on scenarios manifest at low energies as four-fermion scalar contact operators, since they lift the slight (but not negligible) chiral suppression of the  $B_c^- \rightarrow \tau^-\bar{\nu}$  amplitude in the standard model. The new constraint is sufficient to rule out, model independently, explanations of the enhancement in  $R_{D^*}$  based on scalar operators because the  $B \rightarrow D^*\tau\nu$  decay amplitude depends on exactly the same pseudoscalar combination of the corresponding Wilson coefficients as the one contributing to the pure tauonic  $B_c$  decay. Popular new-physics models implementing extra Higgs doublets or certain classes of leptoquarks enter into this category. Updated calculations of the  $B_c$  lifetime in QCD and measurements of the branching fractions of its decay channels at the LHC should refine the bound in the future.

Taking into account the constraint of the lifetime of the  $B_c$  meson simplifies the new-physics analyses of the  $R_{D^{(*)}}$  anomaly. For instance, only two effective operators, left handed or tensorial, remain at low energies to account for the enhancement in  $R_{D^*}$ . We argued that the precise measurement of normalized observables, such as the longitudinal polarizations, could become instrumental to distinguish between the different scenarios.

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