

Polarization fractions in $B \rightarrow V_1 V_2$: U -spin constraints and new physics signatures

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We investigate the decays of B mesons, i.e., B_d , B_s , B^+ , and their antiparticles, to two light vector mesons ($B \rightarrow V_1 V_2$). We use the $SU(2)$ U -spin symmetry, which relates $\Delta S = 0$ and $\Delta S = 1$ decay amplitudes through the interchange $d \leftrightarrow s$ and is an approximate symmetry of the Standard Model (SM), to relate the helicity amplitudes of these decays. Treating all the helicity amplitudes for these decays, and, hence, the reduced matrix elements, as free parameters, we find an acceptable solution within the SM, although this is driven by the fact that the number of observables is smaller than what is needed for a meaningful fit. To reduce the number of free parameters, we then use some apparently reasonable and theoretically motivated approximations, like the dominance of factorizable contributions over the nonfactorizable ones and, hence, a distinct hierarchy between the helicity amplitudes. We find that once the assumption of hierarchy is imposed, there is no solution, both in the exact U -spin limit as well as when substantial U -spin breaking is allowed for. The tension is primarily driven by the longitudinal polarization fractions in almost all $\Delta S = 1$ decays, which are significantly smaller than the corresponding theoretical predictions based on the SM and U -spin symmetry. This is particularly true for $B_s \rightarrow K^{*0} \overline{K}^{*0}$, for which the individual disagreement with U -spin-based expectation is more than 7σ . Within the SM framework, the only effective resolution would be to entirely disregard the hierarchy between the longitudinal and transverse helicity amplitudes in the heavy-quark limit, as dictated by naive factorization, implying that there must be large nonfactorizable contributions to all these decay amplitudes. We also explore whether some new physics in the $b \rightarrow s$ sector that does not respect the hierarchy among the helicity amplitudes can reduce the tension for all the $\Delta S = 1$ modes. While the answer is partially in the affirmative, we find that, for simplistic new physics scenarios, the tension still exists and the fit remains poor enough, if the hierarchy exists among the SM amplitudes. Some possible scenarios for a complete solution of the puzzle are also suggested.

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I. INTRODUCTION

Over the past couple of decades, several cases of apparent tension between the Standard Model (SM) predictions and the experimental data on hadronic B -meson decays have surfaced, and some of them are yet to be resolved. One of the long-standing examples is the so-called $B \rightarrow \pi K$ puzzle (for details, see, e.g., Refs. [1,2], and references therein). More recently, tensions have also been reported in B decays to two

pseudoscalar mesons [3–10], which are related by the $SU(3)_F$ flavor symmetry. A common conclusion of these studies is that the measured branching ratios (BRs) and CP asymmetries (A_{CP}) show significant inconsistencies with their SM predictions. These discrepancies could be indicative of either large (and unexpected) $SU(3)_F$ -breaking effects or possible short-distance new physics (NP) contributions. Similar tensions have also been observed in the two-body hadronic decays of the D mesons [11–14].

In this paper, we will focus on a tension that has surfaced rather recently in $B \rightarrow V_1 V_2$, i.e., B decays to two vector mesons [15], particularly in the measurement of the longitudinal polarization fractions (LPFs) in $B_{s,d} \rightarrow K^{*0} \overline{K}^{*0}$ decays.¹ The recent results from the LHCb Collaboration

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¹Some authors, like Ref. [16], argue that nonfactorizable effects like large final-state rescattering are non-negligible, but their contribution is still way too small to explain the unexpectedly low LPF in $B_s \rightarrow K^{*0} \overline{K}^{*0}$.

[17], which we will talk about a lot later in this paper, have only helped in strengthening the tension to an almost unacceptable value. These two decays, with identical final states, are related by the SU(2) U -spin symmetry, characterized by the $d \leftrightarrow s$ interchange, and, hence, relate the $\Delta S = 0$ ($b \rightarrow d$) and $\Delta S = 1$ ($b \rightarrow s$) amplitudes (see Refs. [18,19]). Further detailed analysis of these two particular decay modes [20] as well as possible NP explanations [21,22] are also available in the literature. Here, though, we adopt a broader outlook and try to encompass a multitude of final states (with two vector mesons) that, pairwise, are related by U spin. The advantages of using U spin over conventional quark-diagrammatic approaches may be found in detail in Ref. [23]. For us, the rationale for using the U -spin symmetry, of course, is to relate the $\Delta S = 0$ and $\Delta S = 1$ amplitudes, so that they need not be treated as independent variables.²

The vector mesons that we consider in the final state are ρ and/or K^* , both charged and neutral. The SM amplitudes are dominated by either $b \rightarrow d$ or $b \rightarrow s$ penguins; tree amplitudes, wherever they exist, are suppressed. Thus, this is a fertile ground to look for NP effects. Factorization of the hadronic currents leads to the fact that, in the infinite b -quark mass limit, both the final state vector mesons must be produced in the longitudinally polarized state, i.e., with spin projection zero. On the inclusion of finite- m_b effects, the final state mesons can also be transversely polarized, with amplitudes relatively suppressed by factors of the order of Λ_{QCD}/m_b , with the QCD scale being of the order of a few hundred MeV. This hierarchy between the longitudinal and the transverse polarization amplitudes holds within the context of the SM and may change if one considers a different Lorentz structure for the relevant operators [24]. Nonfactorizable effects too may affect the hierarchy, but there is no *a priori* way to reliably ascertain their relative importance. However, that they are suppressed compared to the factorizable part is more or less an accepted assumption.

The data, as shown in Table I, show some interesting features. The LPFs, denoted by f_L , are consistently smaller than the corresponding theoretical expectations for all the $\Delta S = 1$ modes, although the discrepancy is most severe for $B_s \rightarrow K^{*0} \bar{K}^{*0}$. The tension exists even for $B_s \rightarrow \phi\phi$, a mode without a U -spin partner channel, for which $f_L = 0.379 \pm 0.008$ [25,26].³ One may try to argue that these decays are dominated by penguins in the SM and so may not be as clean as the tree-dominated decays; in particular, that the nonfactorizable contributions can be large. The counterargument would be the numbers for the $\Delta S = 0$

TABLE I. Experimental values of $B \rightarrow V_1 V_2$ observables such as CP -averaged branching ratios ($\mathcal{B}_{CP}$), direct (A_{CP}) and indirect (S_{CP}) CP asymmetries, and longitudinal (f_L) and transverse ($f_\perp$) polarization fractions. There are a total 18 measurements. These values are taken from Ref. [28], unless indicated otherwise. The entries marked with ♠ for the yet-to-be-observed modes are our predictions, made from the best-fit values of the RMEs (see the text for details).

$\Delta S = 0$	$\Delta S = 1$
$B^+ \rightarrow \bar{K}^{*0} K^{*+}$	$B^+ \rightarrow K^{*0} \rho^+$
$\mathcal{B}_{CP} = (0.91 \pm 0.29) \times 10^{-6}$	$\mathcal{B}_{CP} = (9.2 \pm 1.5) \times 10^{-6}$
$f_L = 0.82^{+0.15}_{-0.21}$	$A_{CP} = -0.01 \pm 0.16$
	$f_L = 0.48 \pm 0.08$
$B_d \rightarrow K^{*0} \bar{K}^{*0}$	$B_s \rightarrow K^{*0} \bar{K}^{*0}$
$\mathcal{B}_{CP} = (0.83 \pm 0.24) \times 10^{-6}$	$\mathcal{B}_{CP} = (11.1 \pm 2.7) \times 10^{-6}$
$f_L = 0.600 \pm 0.028$ [17]	$f_L = 0.159 \pm 0.012$ [17]
$f_\perp = 0.24 \pm 0.03$ [17]	$f_\perp = 0.500 \pm 0.014$ [17]
$B_d \rightarrow \rho^+ \rho^-$	$B_s \rightarrow K^{*-} K^{*+}$ (not yet observed)
$\mathcal{B}_{CP} = (27.7 \pm 1.9) \times 10^{-6}$	♠ $\mathcal{B}_{CP} = (1.27 \pm 0.28) \times 10^{-5}$
$A_{CP} = 0.00 \pm 0.09$	
$S_{CP} = -0.14 \pm 0.13$	
$f_L = 0.990^{+0.021}_{-0.019}$	♠ $f_L = 0.926 \pm 0.063$
$B_s \rightarrow K^{*-} \rho^+$ (not yet observed)	$B_d \rightarrow K^{*+} \rho^-$
♠ $\mathcal{B}_{CP} = (2.96 \pm 0.21) \times 10^{-5}$	$\mathcal{B}_{CP} = (10.3 \pm 2.6) \times 10^{-6}$
	$A_{CP} = 0.21 \pm 0.15$
♠ $f_L = 0.972 \pm 0.024$	$f_L = 0.38 \pm 0.13$

($b \rightarrow d$) decays, which are also penguin dominated but for which the LPFs are closer to the SM predictions, except for the $B_d \rightarrow K^{*0} \bar{K}^{*0}$ channel.⁴

Our goal in this paper will be to identify, in a quantitative way, how serious the tension is between the LPF measurements for the $\Delta S = 1$ channels, in particular, for $B_{s,d} \rightarrow K^{*0} \bar{K}^{*0}$, and the corresponding SM expectations. As the number of observables is limited, one needs to make certain reasonable assumptions to check the robustness of the SM dynamics. We will use two such assumptions, namely,

- (1) U spin is an approximate symmetry of the SM. The amount of U -spin breaking should be of the order of the ratio of the meson decay constants, f_K/f_π , which is approximately 1.19. Thus, we expect U spin to be broken by about 20% or, to be even more conservative, 25%–30% at the most.

²One could have used the method of topological amplitudes. However, there would have been more independent variables than what we use in our approach.

³The only exception is $B^+ \rightarrow \rho^0 K^{*+}$, for which $f_L = 0.720 \pm 0.029$ [27], but this mode is also without a U -spin partner.

⁴It is hard to envisage a scenario where the final state interaction is significant only for this final state, so we do not discuss such a possibility.

- (2) The factorizable parts of the amplitudes dominate over the nonfactorizable parts. U -spin amplitudes include both of them, but if factorization dominates, one expects a hierarchy among the magnitudes of the helicity amplitudes, with the longitudinal helicity amplitude dominating over the transverse ones. If nonfactorizable contributions are small, one may reasonably expect this hierarchy to be reflected in the U -spin amplitudes, too, expressed in the helicity basis. We will quantify later what we mean by this statement.

In the first part of this paper, we will try to see how good these assumptions are, or whether the tension can be explained within the framework of the SM, by relaxing some of these underlying assumptions. Such a relaxation, although not favored in models like QCD factorization, may not be too drastic, as nonperturbative QCD dynamics can still hold some surprises. While including all the relevant modes (such as $B \rightarrow V_1 V_2$, with $V_{1,2} = \{\rho, K^*\}$) in our analysis, we focus mainly on those decays that are related by the U -spin symmetry, with a moderate breaking. Thus, our analysis includes a total of eight decay modes, namely, four with $\Delta S = 0$ and four with $\Delta S = 1$. Among these, two decays, viz., $B_s \rightarrow K^{*-} \rho^+$ and $B_s \rightarrow K^{*-} K^{*+}$, have not yet been observed experimentally, but we will predict the expected LPFs for them. The decay amplitudes, in terms of the U -spin reduced matrix elements (RMEs), are treated as free parameters and can be extracted from the available data. Note that in any other approach, like the one using topological diagrams, the number of free parameters is much larger, and the fit does not make much sense.

The observables that we use are, wherever available, CP -averaged BRs, CP asymmetries (A_{CP}), and polarization fractions, mostly longitudinal, but a couple of transverse fractions, too. We find that, in the exact U -spin symmetric limit, the fit is poor, with a very high χ^2 , indicating a large tension with the SM. The main source of the tension is the LPF of $B_s \rightarrow K^{*0} \bar{K}^{*0}$, which is only about one-fourth of that of its U -spin partner channel, $B_d \rightarrow K^{*0} \bar{K}^{*0}$. Consequently, whereas U -spin symmetry is expected to be broken in the SM at most at the level of 25%–30%, even a much larger breaking is insufficient to significantly alleviate the tension. We find that the root cause of the tension is the assumed hierarchy among the helicity amplitudes, if only the SM dynamics is at play.

In the second part, we investigate whether this discrepancy could be a signal of NP beyond the SM. The NP, of course, need not respect U spin. In fact, it can be very much flavor-specific. As an instructive exercise, we perform a model-independent analysis, introducing some four-fermion effective operators with various possible Lorentz structures. As the tension is mostly confined to the $\Delta S = 1$ sector, the NP operators that we introduce are those that mediate $b \rightarrow s \bar{q} q$ decays, $q \in \{u, d\}$, leaving $\Delta S = 0$ channels essentially uncontaminated by NP and, hence,

explicitly breaking U spin for NP.⁵ What underlying theory leads to such NP structure is, however, a moot question; the only motivation for such a choice is the data—this is entirely a data-driven analysis. One may note that, only in SM and theories with minimal flavor violation, the $b \rightarrow d$ transition strength is suppressed compared to the $b \rightarrow s$ one; this need not be true for NP models, in general. For the SM part of the amplitudes, we stick to the two assumptions enlisted before. While the domination of the LPF is expected for the SM part of the $\Delta S = 1$ amplitudes, the NP operators, particularly scalar and tensor types, lead to a significant transverse polarization fraction. We show that this leads to a softening of the tension but not to a complete explanation. We also discuss qualitatively what happens when one introduces NP in $b \rightarrow d$ transitions, too.

The paper is organized as follows. In Sec. II, we briefly discuss how U spin relates various $B \rightarrow V_1 V_2$ amplitudes in the SM. This is followed by an analysis within the framework of SM, where we show quantitatively how large the tension is and pinpoint its source. In Sec. III, we introduce various NP operators and again do the analysis, taking the NP only in the $\Delta S = 1$ sector. However, both $\Delta S = 0$ and $\Delta S = 1$ amplitudes related by U spin are taken into account for the SM part of the decay channels. We find that several possible NP operators can help in partially resolving the tension. We summarize and conclude in Sec. IV and relegate a few necessary details to the Appendix.

II. U -SPIN ANALYSIS

The U -spin group $SU(2)_U$ is a subgroup of the full $SU(3)_F$ symmetry, with (d, s) and $(\bar{s}, -\bar{d})$ in the fundamental representation and all others quark fields (u, c, b, t) as singlets. Of the spin-1 $SU(3)_F$ octet given by

$$\begin{aligned} K^{*+} &= u\bar{s}, & K^{*-} &= \bar{u}s, & \rho^+ &= -u\bar{d}, & \rho^- &= \bar{u}d, \\ K^{*0} &= d\bar{s}, & \bar{K}^{*0} &= -s\bar{d}, & U^0 &= \frac{1}{\sqrt{2}}(s\bar{s} - d\bar{d}), \\ U_8^* &= \frac{1}{\sqrt{6}}(2u\bar{u} - d\bar{d} - s\bar{s}), \end{aligned} \quad (1)$$

under $SU(2)_U$, (K^{*+}, ρ^+) and (ρ^-, K^{*-}) are doublets, and $(K^{*0}, U^0, \bar{K}^{*0})$ form a triplet, while U_8^* is a singlet. The ninth vector meson $U_1^* = \frac{1}{\sqrt{6}}(u\bar{u} + d\bar{d} + s\bar{s})$ is, of course, a singlet under the entire $SU(3)_F$. The states U^{*0} , U_8^* , and U_1^* are linear combinations of the vector mesons ρ^0 , ϕ , and ω . Similarly, B^+ is a U -spin singlet, while (B_d, B_s) form a U -spin doublet.

⁵Note that while small Wilson coefficients (WCs) can be accommodated for $\Delta S = 0$ NP operators, these neither add anything to the global fit nor are necessary. For such flavor-specific NP, see, e.g., Ref. [29].

The decays $B \rightarrow V_1 V_2$, where $B \in \{B^+, B_d, B_s\}$ and $V_1, V_2 \in \{\rho, K^*\}$, including all charged and neutral states, can be classified into two broad classes, namely, strangeness-conserving $\Delta S = 0$ ($\bar{b} \rightarrow \bar{d} q \bar{q}$) and strangeness-violating $\Delta S = 1$ ($\bar{b} \rightarrow \bar{s} q \bar{q}$), with $q \in \{u, d, s\}$. The effective weak Hamiltonian for these decays can be written as a sum of local operators $\mathcal{O}_i$ multiplied by short-distance WCs C_i and products of relevant Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, $\lambda_p^D = V_{pb}^* V_{pD}$, where $p \in \{u, c, t\}$ and $D \in \{d, s\}$, depending on the decay mode under consideration. C_i and the expectation values of $\mathcal{O}_i$ are functions of the renormalization scale μ , for which we adopt the natural choice⁶ $\mu = m_b$. Exploiting the unitarity of the CKM matrix, we can write the Hamiltonian as⁷

$$\mathcal{H}_{\text{eff}}^{\bar{b} \rightarrow \bar{D}} = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^D \left(C_1 \mathcal{O}_{1D}^p + C_2 \mathcal{O}_{2D}^p + \sum_{i=3}^{10} C_i \mathcal{O}_{iD} \right) + \text{H.c.} \quad (2)$$

Here, $\mathcal{O}_{(1,2)D}^p$ are the left-handed current-current operators (generated by tree-level W boson exchange) that depend on the identity of the up-type quark p , whereas $\mathcal{O}_{(3-6)D}$ and $\mathcal{O}_{(7-10)D}$ are the QCD and electroweak penguin operators respectively, as defined in Ref. [30].⁸ From now on, we will absorb the SM WCs into the definition of the operators, i.e., define $C_i \mathcal{O}_i$ as just $\mathcal{O}_i$, so that the WCs will not appear explicitly any further.

In the $|UU_3\rangle$ basis, the $\bar{b} \rightarrow \bar{d}$ Hamiltonian transforms like $\bar{d} \equiv -|\frac{1}{2} - \frac{1}{2}\rangle$, whereas the $\bar{b} \rightarrow \bar{s}$ Hamiltonian transforms like $\bar{s} \equiv |\frac{1}{2} \frac{1}{2}\rangle$. Using the Wigner-Eckart theorem, it is straightforward to express the amplitudes for the $B \rightarrow V_1 V_2$ decays in terms of the U -spin RMEs [23]. From now on, we will write these amplitudes, derived using the Hamiltonian in Eq. (2), in terms of the final state vector meson helicity $h \in \{0, -, +\}$.

In terms of the U -spin RMEs defined as

$$\begin{aligned} \mathcal{A}_{\frac{1}{2}D}^{p,h} &= \left\langle \frac{1}{2}; h \left| [\mathcal{O}_D^p]^{\frac{1}{2}} \right| 0 \right\rangle, & \mathcal{A}_{0D}^{p,h} &= \left\langle 0; h \left| [\mathcal{O}_D^p]^{\frac{1}{2}} \right| \frac{1}{2} \right\rangle, \\ \mathcal{A}_{1q}^{p,h} &= \left\langle 1; h \left| [\mathcal{O}_D^p]^{\frac{1}{2}} \right| \frac{1}{2} \right\rangle, \end{aligned} \quad (3)$$

the $\Delta S = 0$ decay amplitudes, in units of $G_F/\sqrt{2}$, are

$$\begin{aligned} A_h(B^+ \rightarrow K^{*+} \bar{K}^{*0}) &= -\sqrt{\frac{2}{3}} \left[\lambda_{bd}^u \mathcal{A}_{\frac{1}{2}d}^{u,h} + \lambda_{bd}^c \mathcal{A}_{\frac{1}{2}d}^{c,h} \right], \\ A_h(B_d \rightarrow K^{*0} \bar{K}^{*0}) &= -\frac{1}{\sqrt{6}} \left[\lambda_{bd}^u \mathcal{A}_{0d}^{u,h} + \lambda_{bd}^c \mathcal{A}_{0d}^{c,h} \right] \\ &\quad - \frac{1}{2} \left[\lambda_{bd}^u \mathcal{A}_{1d}^{u,h} + \lambda_{bd}^c \mathcal{A}_{1d}^{c,h} \right], \\ A_h(B_d \rightarrow \rho^+ \rho^-) &= \frac{1}{2} \left[\lambda_{bd}^u \mathcal{A}_{0d}^{u,h} + \lambda_{bd}^c \mathcal{A}_{0d}^{c,h} \right] \\ &\quad - \frac{1}{2} \left[\lambda_{bd}^u \mathcal{A}_{1d}^{u,h} + \lambda_{bd}^c \mathcal{A}_{1d}^{c,h} \right], \\ A_h(B_s \rightarrow \rho^+ K^{*-}) &= \lambda_{bd}^u \mathcal{A}_{1d}^{u,h} + \lambda_{bd}^c \mathcal{A}_{1d}^{c,h}, \end{aligned} \quad (4)$$

with⁹ $p \in \{u, c\}$ and $D \in \{d, s\}$. The labels 0 and $\frac{1}{2}$ in the initial and final states denote their total U -spin quantum numbers. With the initial state being spin 0, the helicities of the daughters are equal and are denoted by a single label h .

Similarly, the $\Delta S = 1$ amplitudes (again, in units of $G_F/\sqrt{2}$) are

$$\begin{aligned} A_h(B^+ \rightarrow K^{*0} \rho^+) &= -\sqrt{\frac{2}{3}} \left[\lambda_{bs}^u \mathcal{A}_{\frac{1}{2}s}^{u,h} + \lambda_{bs}^c \mathcal{A}_{\frac{1}{2}s}^{c,h} \right], \\ A_h(B_s \rightarrow K^{*0} \bar{K}^{*0}) &= -\frac{1}{\sqrt{6}} \left[\lambda_{bs}^u \mathcal{A}_{0s}^{u,h} + \lambda_{bs}^c \mathcal{A}_{0s}^{c,h} \right] \\ &\quad - \frac{1}{2} \left[\lambda_{bs}^u \mathcal{A}_{1s}^{u,h} + \lambda_{bs}^c \mathcal{A}_{1s}^{c,h} \right], \\ A_h(B_s \rightarrow K^{*-} K^{*+}) &= \frac{1}{2} \left[\lambda_{bs}^u \mathcal{A}_{0s}^{u,h} + \lambda_{bs}^c \mathcal{A}_{0s}^{c,h} \right] \\ &\quad - \frac{1}{2} \left[\lambda_{bs}^u \mathcal{A}_{1s}^{u,h} + \lambda_{bs}^c \mathcal{A}_{1s}^{c,h} \right], \\ A_h(B_d \rightarrow K^{*+} \rho^-) &= \lambda_{bs}^u \mathcal{A}_{1s}^{u,h} + \lambda_{bs}^c \mathcal{A}_{1s}^{c,h}. \end{aligned} \quad (5)$$

However, one may note that not all the RMEs are independent. In the limit of exact U spin [23], the RMEs $\mathcal{A}_{nD}^{p,h}$, with $n \in \{0, \frac{1}{2}, 1\}$, are related:

$$\mathcal{A}_{\frac{1}{2}d}^{p,h} = \mathcal{A}_{\frac{1}{2}s}^{p,h}, \quad \mathcal{A}_{0d}^{p,h} = \mathcal{A}_{0s}^{p,h}, \quad \mathcal{A}_{1d}^{p,h} = \mathcal{A}_{1s}^{p,h}, \quad \forall p, h, \quad (6)$$

rendering the D label irrelevant.

The heavy-quark symmetry and the absence of large nonfactorizable contributions imply a certain hierarchy among the moduli of the helicity amplitudes, which is, in turn, reflected in the hierarchy among the RMEs [24]:

$$A_0 : A_- : A_+ \sim 1 : \frac{\Lambda_{\text{QCD}}}{m_b} : \left(\frac{\Lambda_{\text{QCD}}}{m_b} \right)^2. \quad (7)$$

It should be noted that Eq. (7) is a joint consequence of the $(V-A)$ structure of the weak interaction as well as the

⁶Varying μ in the window $(\frac{1}{2}m_b, 2m_b)$ brings in only relatively minor quantitative changes in our conclusions.

⁷Note that the CKM factors, $\lambda_u^D = V_{ub}^* V_{uD}$ and $\lambda_c^D = V_{cb}^* V_{cD}$, involve only the u and c quarks, with the top-quark contribution having been appropriately subsumed using the unitarity condition $\lambda_u^D + \lambda_c^D + \lambda_t^D = 0$.

⁸We do not consider here the electromagnetic and chromomagnetic penguin operators, commonly denoted by $\mathcal{O}_{7\gamma}$ and $\mathcal{O}_{8g}$, respectively, as they are not relevant for our discussion.

⁹The top contribution has already been absorbed in the definition of the operators.

assumption of naive factorization being a good approximation. As stressed before, we assume nonfactorizable contributions to the amplitudes to be subdominant. Under this assumption, the matrix elements of the four-fermion operators in $\mathcal{H}_{\text{eff}}$ are sought to be approximated by products of two current matrix elements. In other words, the helicity amplitudes $A_h^{V_1 V_2}$ for the decay $B \rightarrow V_1 V_2$ are proportional to

$$A_h^{V_1 V_2} \propto \langle V_2^h | (\bar{D}b)_{V-A} | \bar{B} \rangle \langle V_1^h | (\bar{q}q')_V | 0 \rangle. \quad (8)$$

Using the form factors defined in the Appendix, we obtain

$$\begin{aligned} A_0^{V_1 V_2} &\propto i m_B^2 f_{V_1} \mathfrak{A}_0^{B \rightarrow V_2}(0), \\ A_{\pm}^{V_1 V_2} &\propto i m_B m_{V_1} f_{V_1} \mathfrak{F}_{\pm}^{B \rightarrow V_2}(0), \end{aligned} \quad (9)$$

where

$$\begin{aligned} \mathfrak{F}_{\pm}^{B \rightarrow V_2}(q^2) &= \left(1 + \frac{m_{V_2}}{m_B}\right) \mathfrak{A}_1^{B \rightarrow V_2}(q^2) \\ &\mp \left(1 - \frac{m_{V_2}}{m_B}\right) \mathfrak{B}^{B \rightarrow V_2}(q^2). \end{aligned} \quad (10)$$

The transverse amplitudes $A_{\pm}^{V_1 V_2}$ are suppressed by a factor of m_{V_2}/m_B relative to the longitudinal amplitude $A_0^{V_1 V_2}$. Moreover, in the heavy-quark limit, the vector and axial-vector contributions to $\mathfrak{F}_{\pm}^{B \rightarrow V_2}(0)$ cancel due to an exact form-factor relation [24]. This suppression pattern leads to Eq. (7).

These helicity amplitudes can also be expressed in terms of transversity amplitudes such as

$$A_L = A_0, \quad A_{\parallel} = \frac{A_+ + A_-}{\sqrt{2}}, \quad A_{\perp} = \frac{A_+ - A_-}{\sqrt{2}}. \quad (11)$$

In general, the amplitude for $B(p) \rightarrow V_1(q, \epsilon_1) V_2(k, \epsilon_2)$ decay can be written in the form [31,32]

$$\mathcal{M} = a \epsilon_1^* \cdot \epsilon_2^* + \frac{b}{m_B^2} (p \cdot \epsilon_1^*) (p \cdot \epsilon_2^*) + i \frac{c}{m_B^2} \epsilon_{\mu\nu\alpha\beta} p^\mu q^\nu \epsilon_1^\alpha \epsilon_2^\beta, \quad (12)$$

where the form factors appear in the constants a , b , and c . The transversity amplitudes are related to the coefficients a , b , and c by

$$\begin{aligned} A_{\parallel} &= \sqrt{2}a, \quad A_L = -ax - \frac{m_1 m_2}{m_B^2} b(x^2 - 1), \\ A_{\perp} &= 2\sqrt{2} \frac{m_1 m_2}{m_B^2} c \sqrt{x^2 - 1}, \end{aligned} \quad (13)$$

where

$$x = \frac{q \cdot k}{m_1 m_2} = \frac{m_B^2 - m_1^2 - m_2^2}{2m_1 m_2}. \quad (14)$$

Equation (13) shows why, in the limit $m_B \gg m_1, m_2$, A_L dominates over $A_{\perp}$ or $A_{\parallel}$. The expressions are true if factorization of the current-current operators holds, although, even in the presence of nonfactorizable contributions, A_L still dominates.

In the naive factorization limit, and using the values for the form factors and masses, we find that for the modes under discussion, $A_0 : A_- : A_+ \simeq 1 : 0.27 : 0.01$ to a good approximation,¹⁰ subject to the caveat that this need not be valid [24] in the presence of large nonfactorizable corrections. Therefore, one should be able to safely ignore A_+ for the SM analysis, but, at the same time, it is worthwhile to test this hierarchy experimentally. Note that, for operators with other Lorentz structures, this hierarchy need not hold (even A_+ may not be neglected), and this fact will be important in our discussion of possible NP scenarios to explain the tension.

The number of independent RMEs, of the form $\mathcal{A}_{nD}^{p,h}$, with $p \in \{u, c\}$, $h \in \{0, -\}$, and $n \in \{0, \frac{1}{2}, 1\}$ are as follows:

- (i) If all the helicity amplitudes are taken to be independent and U -spin symmetry is not assumed, there are a total of 36 RMEs, 18 each for $\Delta S = 1$ and $\Delta S = 0$ decays.
- (ii) If even an approximate U -spin symmetry is assumed, the $\Delta S = 0$ and $\Delta S = 1$ amplitudes get related, and the number of independent amplitudes drops to 18.
- (iii) In the limit $A_+ = 0$, one has 24 independent RMEs without U spin and 12 in the limit of exact $SU(2)_U$.

At this point, we would like to stress again that U -spin amplitudes include both factorizable and nonfactorizable contributions and are not decomposed explicitly in terms of the form factors; they depend only on the initial and final U -spin quantum numbers and not the diagram topologies. The hierarchy among the amplitudes is not inevitable in the U -spin approach. In the heavy-quark limit, one has $\mathfrak{F}_-/\mathfrak{A}_0 \approx 1$ and $\mathfrak{F}_+/\mathfrak{A}_0 \approx \Lambda_{\text{QCD}}/m_B$. These relations indicate that the hierarchical structure of the amplitudes originates from the relative sizes of the form factors contributing to them. Thus, any expectation of hierarchy among the U -spin RMEs essentially means that the nonfactorizable contributions are assumed to be much smaller than the corresponding factorizable ones.

At the same time, U spin is not an exact symmetry of nature, so one should also check whether the polarization puzzle is due to an imperfect $SU(2)_U$ symmetry. One way to incorporate U -spin breaking is to add an explicit breaking term in the $\Delta S = 1$ Hamiltonian. For the SM-based analysis that follows, we take a different but equivalent route. We propose a possible mismatch of at most 30%—as the breaking is expected to be moderate—between any $\Delta S = 0$

¹⁰One may note that the hierarchy given in Eq. (7) is much stronger, even if we take $\Lambda_{\text{QCD}} \sim 300$ MeV. However, that is only an order-of-magnitude estimate.

RME and its $\Delta S = 1$ counterpart, as shown later in Eq. (21). The mismatch need not be the same for all U -spin-related pairs. While we will show explicitly that any such possible mismatch is not enough to address the polarization puzzle, the full analysis is detailed in Sec. II B.

Before moving further, let us note that one could have relaxed the assumption of $A_+ = 0$. In that case, there would have been 18 variables in the unbroken U -spin limit. With 18 observables, one could try for an exact solution. We found that such an exact solution indeed exists—to be precise, very close to an exact solution, because all the RMEs are taken to be real, and $A_{CP}(B_d \rightarrow K^{*+}\rho^-)$ is not consistent with zero at the 1σ level—but none of the six A_+ 's are negligibly small. In other words, the hierarchy does not exist, and an SM-only explanation desperately needs a large nonfactorizable contribution.

A. Observables

To begin with, we list the observables, pertaining to various $\Delta S = 0$ and $\Delta S = 1$ decays (into a vector boson pair) related by $SU(2)_U$, that we would use in our analysis. These comprise (wherever available) the CP -averaged branching fraction $\mathcal{B}_{CP}$, the direct CP asymmetry A_{CP} , the mixing-induced CP asymmetry S_{CP} and, most importantly, the polarization fractions, f_L (longitudinal), $f_\perp$ (transverse), and $f_\parallel$ (parallel) and are defined as¹¹

$$\begin{aligned} \mathcal{B}_{CP} &= \frac{\sqrt{m_B^2 - (m_{V_1} + m_{V_2})^2} \sqrt{m_B^2 - (m_{V_1} - m_{V_2})^2}}{32\pi m_B^3 \Gamma_B} \\ &\quad \times \sum_{h=0,\pm} (|A_h|^2 + |\bar{A}_h|^2), \\ A_{CP} &= \frac{\sum_{h=0,\pm} (|\bar{A}_h|^2 - |A_h|^2)}{\sum_{h=0,\pm} (|\bar{A}_h|^2 + |A_h|^2)}, \\ S_{CP} &= 2\text{Im} \left[e^{-2i\phi_M} \frac{\sum_{h=0,\pm} \bar{A}_h A_h^*}{\sum_{h=0,\pm} (|\bar{A}_h|^2 + |A_h|^2)} \right], \\ f_{L,\perp,\parallel} &= \frac{1}{2} \left(\frac{|A_{L,\perp,\parallel}|^2}{|A_L|^2 + |A_\perp|^2 + |A_\parallel|^2} + \frac{|\bar{A}_{L,\perp,\parallel}|^2}{|\bar{A}_L|^2 + |\bar{A}_\perp|^2 + |\bar{A}_\parallel|^2} \right). \end{aligned} \quad (15)$$

Here, A_h and $\bar{A}_h$ are the helicity amplitudes for $B \rightarrow V_1 V_2$ and its CP conjugate decays, and ϕ_M is the weak phase of $B_d - \bar{B}_d$ mixing. The experimental data are summarized in Table I. Among the 18 measurements, only one shows a striking tension with the SM, namely, $f_L(B_s \rightarrow K^{*0}\bar{K}^{*0})$, but f_L for all $\Delta S = 1$ modes are smaller than the corresponding theoretical expectations. While Eq. (7) predicts $f_L \gg f_\perp$, one finds, instead, $f_L = 0.159 \pm 0.012$ and

¹¹Only two of the polarization fractions are independent, as their sum must be unity. Thus, we take only two of them as independent observables (with correlation), even when data are available on all three.

$f_\perp = 0.500 \pm 0.014$ [17] for $B_s \rightarrow K^{*0}\bar{K}^{*0}$, in contradiction with Eq. (7). This is also in contradiction with U -spin symmetry, as its U -spin partner $B_d \rightarrow K^{*0}\bar{K}^{*0}$ has $f_L = 0.600 \pm 0.028$. In the exact U -spin limit, these two f_L 's should be equal to each other, as predicted by QCD factorization [18]:

$$\left[\frac{f_L(B_d \rightarrow K^{*0}\bar{K}^{*0})}{f_L(B_s \rightarrow K^{*0}\bar{K}^{*0})} \right]_{\text{QCDF}} = 1.09^{+0.19}_{-0.08}, \quad (16)$$

while we find

$$\left[\frac{f_L(B_d \rightarrow K^{*0}\bar{K}^{*0})}{f_L(B_s \rightarrow K^{*0}\bar{K}^{*0})} \right]_{\text{exp}} = 3.77 \pm 0.33, \quad (17)$$

corresponding to a tension of approximately 7.1σ between data and theory. This is the polarization puzzle in the $B_{d,s} \rightarrow K^{*0}\bar{K}^{*0}$ decays.

B. Analysis within the SM

We have made two assumptions so far: first, that $SU(2)_U$ is exact and, second, that the hierarchy of Eq. (7) is valid. Let us first check how far these assumptions are justified. First, let us assume that the U spin is an exact symmetry but relax the hierarchy of Eq. (7). Thus, there are 18 independent RMEs, and, as can be seen from Table I, there are 18 observables. So, one expects an exact solution—at least almost an exact solution, as we have taken all RMEs to be real, but $A_{CP}(B_d \rightarrow K^{*+}\rho^-)$ is not consistent with zero at the 1σ level. We perform a fit to the data, taking all the error bars at the 1σ level, and that is precisely what we find, even when U spin is taken to be unbroken, underlining that it is indeed a good symmetry of the flavor sector. We also predict the values of several observables, based on the fit, which are shown in Table I.

This, however, does not mean that there is no tension with the SM, because the magnitudes of the helicity amplitudes do not even come close to respecting the hierarchy as shown in Eq. (7). For example,¹²

$$\begin{aligned} B_d \rightarrow K^{*0}\bar{K}^{*0}: \quad & A_0 = (8.03 \pm 1.65) \times 10^{-9}, \\ & A_- = (0.07 \pm 1.80) \times 10^{-8}, \\ & A_+ = (0.67 \pm 1.55) \times 10^{-8}, \\ B_s \rightarrow K^{*0}\bar{K}^{*0}: \quad & A_0 = (1.42 \pm 0.37) \times 10^{-8}, \\ & A_- = (3.38 \pm 7.54) \times 10^{-9}, \\ & A_+ = (3.30 \pm 0.80) \times 10^{-8}. \end{aligned} \quad (18)$$

Next, we try to see how badly the hierarchy approximation is violated. For that, we still work in the exact

¹²The helicity amplitudes also include the CKM elements, and so, even though these two decays are related by U spin, the helicity amplitudes are not the same.

TABLE II. The magnitudes of the CKM matrix elements [28] and the CP -violating phases [35] used in the analysis.

$V_{ud} = 0.97367 \pm 0.00032$	$V_{us} = 0.22431 \pm 0.00085$	$V_{ub} = 0.00382 \pm 0.00020$
$V_{cd} = 0.221 \pm 0.004$	$V_{cs} = 0.975 \pm 0.006$	$V_{cb} = 0.0411 \pm 0.0012$
$\gamma = (66.4^{+2.8}_{-3.0})^\circ$	$\beta = (22.6^{+0.5}_{-0.4})^\circ$	

U -spin limit but take all $A_+ = 0$. This does not appear to be a drastic approximation, at least within the purview of the SM, as A_+ is anyway suppressed by $(\Lambda_{\text{QCD}}/m_B)^2$. This gives us altogether 12 RMEs, following Eq. (6). We assume these RMEs to be real and perform a χ^2 fit to the 18 measurements to extract the best-fit values of the 12 RMEs. For the minimization, we use the MINUIT library [33,34]. The CP -violating phases γ and β are taken as external constraints, with current averages given in Table II along with the CKM matrix elements.

We find a rather poor fit with $\chi^2/\text{d.o.f.} = 37.8/6$ and the p value of 1.23×10^{-6} . The fit being so poor, we do not even display the fit results. However, one may immediately see that even this best fit does not respect the hierarchy between A_0 and A_- as shown in Eq. (7). To demonstrate this, we have computed the ratios of the helicity amplitudes, $|A_-/A_0|$, using the best-fit values. For $\Delta S = 0$ decays, we find

$$\begin{aligned} \left| \frac{A_-}{A_0} \right|_{B^+ \rightarrow K^{*0} K^{*+}} &= 0.52 \pm 0.87, \\ \left| \frac{A_-}{A_0} \right|_{B_d \rightarrow K^{*0} \bar{K}^{*0}} &= 0.81 \pm 0.64, \\ \left| \frac{A_-}{A_0} \right|_{B_d \rightarrow \rho^+ \rho^-} &= 0.10 \pm 0.09. \end{aligned} \quad (19)$$

The hierarchy looks well respected for the last one, and, given the large error margins, the first two numbers are also consistent with the hierarchy. The situation changes when we look at the $\Delta S = 1$ decays:

$$\begin{aligned} \left| \frac{A_-}{A_0} \right|_{B^+ \rightarrow K^{*0} \rho^+} &= 1.04 \pm 0.19, \\ \left| \frac{A_-}{A_0} \right|_{B_s \rightarrow K^{*0} \bar{K}^{*0}} &= 2.39 \pm 0.80, \\ \left| \frac{A_-}{A_0} \right|_{B_d \rightarrow K^{*+} \rho^-} &= 1.30 \pm 0.40. \end{aligned} \quad (20)$$

To ascertain that the poor fit is not due to the imposition of an exact $SU(2)_U$, we insert a reasonable breaking:

$$0.7 \leq \frac{\mathcal{A}_{nd}^{p,h}}{\mathcal{A}_{ns}^{p,h}} \leq 1.3 \quad \forall \quad p, h, n, \quad (21)$$

where we refer to Eqs. (3) and (6) for the notation. This means a breaking of 30% on either side of equality, but the breaking need not be the same for different p , h , and n . This

feature has been incorporated in our analysis by randomizing the breaking within $\pm 30\%$ for different pair of U -spin-related RMEs. Thus, to take extreme values, while the ratio can be as low as 0.7 for a certain pair, it can be as high as 1.3 for another pair. What we find is that the fit result remains almost equally poor; the improvement is so marginal as to be almost imperceptible.

We can even perform an analysis of $\Delta S = 0$ and $\Delta S = 1$ decays separately. This does not mean disregarding U spin, as the RMEs are still related by the symmetry, e.g., in $B_d \rightarrow \rho^+ \rho^-$ and $B_s \rightarrow \rho^+ K^{*-}$. To ascertain whether the hierarchy is broken only for the $\Delta S = 1$ modes, we impose $A_+ = 0$ and $A_- = 0.3A_0$. For $\Delta S = 0$ decays, we have six parameters and nine measurements and get a rather poor fit with $\chi^2/\text{d.o.f.} = 144.1/3$, with $f_L(B_d \rightarrow K^{*0} \bar{K}^{*0})$ being the most responsible one for the quality of the fit, contributing about 125 or 87% to the total χ^2 . There is no minimum for the $\Delta S = 1$ modes within the allowed ranges of the observables, as expected (which, in other words, means a disastrously poor χ^2).

This just reaffirms what we have stated before: The tension is real and rather severe if we stick to the hierarchy among the helicity amplitudes. In other words, if there is no new physics beyond the SM, there must be a large non-factorizable contribution for $B_{s,d} \rightarrow K^{*0} \bar{K}^{*0}$. In the next section, we will see whether NP can help.

III. NEW PHYSICS INTERPRETATION

Let us now entertain the possibility that NP manifests itself in $b \rightarrow s$ decays alone, a choice, while not forced upon us, that is motivated by the fact that the tension is more severe for $\Delta S = 1$ decays. This simplifying assumption of a flavor-specific NP is, of course, not U -spin invariant. However, as we will soon see (and as could be gleaned from the preceding section), this lack of U -spin symmetry is not the key ingredient. Rather, the resolution of the tension would depend on the introduction of specific $\Delta S = 1$ operators that may not respect the hierarchy of Eq. (7).

The aforementioned U -spin violation could be motivated by postulating a heavy mediator X with a tree-level (or effective) $\bar{b}sX$ vertex and also coupling to a light-quark current. It needs to be borne in mind that such a flavor-violating coupling is constrained by $B_s - \bar{B}_s$ mixing. However, rather than committing ourselves to a specific ultraviolet completion, we consider, instead, to a NP effective Hamiltonian at the scale m_b , written as a sum

of possible vector, scalar, and tensor operators:

$$\mathcal{H}_{\text{NP}} = \mathcal{H}_V + \mathcal{H}_S + \mathcal{H}_T + \text{H.c.}, \quad (22)$$

where

$$\begin{aligned} \text{Vector: } \mathcal{H}_V &= h_v e^{i\zeta_v} (\bar{q}_\alpha \gamma_\mu (c_1 + c_2 \gamma_5) q_\alpha) \\ &\quad \times (\bar{s}_\beta \gamma^\mu (c_3 + c_4 \gamma_5) b_\beta), \\ \text{Scalar: } \mathcal{H}_S &= h_s e^{i\zeta_s} (\bar{q}_\alpha (c_1 + c_2 \gamma_5) q_\alpha) (\bar{s}_\beta (c_3 + c_4 \gamma_5) b_\beta), \\ \text{Tensor: } \mathcal{H}_T &= h_t e^{i\zeta_t} (\bar{q}_\alpha \sigma_{\mu\nu} (c_1 + c_2 \gamma_5) q_\alpha) \\ &\quad \times (\bar{s}_\beta \sigma^{\mu\nu} (c_3 + c_4 \gamma_5) b_\beta). \end{aligned} \quad (23)$$

Here, $\zeta_{v,s,t}$ are the NP weak phases, α and β are color indices, and $q \in \{u, d\}$. The WCs $h_{v,s,t}$ encapsulate the NP coupling constants as well the suppression Λ^{-2} , with Λ being the scale at which the NP fields have been integrated out (and constrained to be larger than the mass of the heaviest such field). Nominally, we consider $\Lambda = 1$ TeV, with the fit values of $h_{v,s,t}$ scaling as $(\Lambda/1 \text{ TeV})^2$. As for the coefficients c_i , they determine only the Lorentz structure of the current-current interactions, and we would consider only four possible combinations, viz. $(c_1, c_2, c_3, c_4) = (1, 1, 1, 1)$, $(1, -1, 1, 1)$, $(1, 1, 1, -1)$, and $(1, -1, 1, -1)$ for vector and scalar NP and $(c_1, c_2, c_3, c_4) = (1, 1, 1, 1)$ and $(1, -1, 1, -1)$ for tensor NP. Such a choice ensures that the relevant fermion fields are chiral. Furthermore, to keep the analysis tractable, we restrict ourselves to only one type of NP operator at a time, viz. only a single set of the c_i 's. Only with more data would the relaxation of these simplifying assumptions still be amenable to a meaningful fit.

This is how we propose to do the analysis.

- (i) As already said, we will consider only one type of NP at a time. Furthermore, we assume the corresponding h_i to be real (i.e., the weak phase $\zeta_i = 0$), a choice motivated by the twin facts of the SM phase (arising almost entirely coming from $V_{tb}^* V_{ts}$) being very small and that there is hardly any tension in A_{CP} for the $b \rightarrow s$ decays. In other words, the present data are completely consistent with the choice $\zeta_i = 0$. This simplifying choice implies that, for a given combination of c_i (a discrete set of choices), the number of free parameters has increased just by one.
- (ii) We will first concentrate on the three $\Delta S = 1$ modes listed in Table I. Since the SM amplitudes are supposed to respect U spin, we would relate the SM contributions to the $\Delta S = 1$ amplitudes through an exact U -spin symmetry to those for the corresponding $\Delta S = 0$ modes which, by definition, do not feel NP.
- (iii) The $\Delta S = 1$ helicity amplitudes may be expressed as sum of the SM part and the NP part. For the SM part, we follow the hierarchy of Eq. (7), assume $A_+ = 0$, and take $|A_-/A_0| = 0.3$. For the NP part,

we assume naive factorization and use Eqs. (12) and (13) to obtain the NP parts of the helicity amplitudes A_0 and $A_\pm$ from the a , b , and c coefficients which are functions of the NP coupling h_i . These are rather simplistic assumptions, but our aim is to check whether the theoretically motivated SM dynamics, in the presence of NP, can soften the tension.

- (iv) As we have taken the NP to be confined to the $b \rightarrow s$ transition only, it is reasonable to redo the analysis with only the $\Delta S = 1$ modes. The U -spin symmetry is no longer relevant, but we keep the hierarchy between A_- and A_0 for the SM part of the amplitudes. Again, we find a very distinct improvement of the fit compared to the SM case.¹³
- (v) We also predict the LPF f_L for the modes $B_s \rightarrow K^{*+} K^{*-}$ and $B_s \rightarrow K^{*-} \rho^+$. For the latter, NP does not play a role, but the best-fit values for the SM amplitudes change in the presence of NP and exact $SU(2)_U$.

Let us first calculate the NP contributions to the following $\Delta S = 1$ decays: $B^+ \rightarrow K^{*0} \rho^+$, $B_s \rightarrow K^{*0} \bar{K}^{*0}$, and $B_d \rightarrow K^{*+} \rho^-$. The quark-level processes are $b \rightarrow s d \bar{d}$ and $b \rightarrow s \bar{u} u$. For the first channel, we get the following expressions with Eqs. (12) and (13) using naive factorization.

(i) Scalar NP:

$$\begin{aligned} a &= \frac{h_s e^{i\zeta_s}}{4N_c} [(c_1 c_4 - c_2 c_3) f_{K^{*0} m_{K^{*0}}} (m_B + m_\rho) \mathfrak{A}_1(q^2) \\ &\quad + c_1 c_4 f_{K^{*0}}^T \{ (m_B^2 - m_\rho^2) g_+(q^2) + q^2 g_-(q^2) \}], \\ \frac{b}{m_B^2} &= -\frac{h_s e^{i\zeta_s}}{4N_c} \left[(c_1 c_4 - c_2 c_3) f_{K^{*0} m_{K^{*0}}} \frac{2\mathfrak{A}_2(q^2)}{(m_B + m_\rho)} \right. \\ &\quad \left. + c_1 c_4 f_{K^{*0}}^T \{ 2g_+(q^2) + q^2 g_0(q^2) \} \right], \\ \frac{c}{m_B^2} &= \frac{h_s e^{i\zeta_s}}{4N_c} \left[(c_1 c_3 - c_2 c_4) f_{K^{*0} m_{K^{*0}}} \frac{2\mathfrak{B}(q^2)}{(m_B + m_\rho)} \right. \\ &\quad \left. - 2(c_2 c_4 + c_1 c_3) f_{K^{*0}}^T g_+(q^2) \right]. \end{aligned} \quad (24)$$

(ii) Vector NP:

$$\begin{aligned} a &= -\frac{h_v e^{i\zeta_v}}{2N_c} (c_1 c_4 + c_2 c_3) f_{K^{*0} m_{K^{*0}}} (m_B + m_\rho) \\ &\quad \times \mathfrak{A}_1(q^2), \\ \frac{b}{m_B^2} &= \frac{h_v e^{i\zeta_v}}{N_c} (c_1 c_4 + c_2 c_3) f_{K^{*0} m_{K^{*0}}} \frac{\mathfrak{A}_2(q^2)}{(m_B + m_\rho)}, \\ \frac{c}{m_B^2} &= \frac{h_v e^{i\zeta_v}}{N_c} (c_2 c_4 - c_1 c_3) f_{K^{*0} m_{K^{*0}}} \frac{\mathfrak{B}(q^2)}{(m_B + m_\rho)}. \end{aligned} \quad (25)$$

¹³One has to compare with the SM fit where the hierarchy between A_- and A_0 has been imposed.

TABLE III. Fit results for different NP Lorentz structures taking NP only in $\Delta S = 1$ decay assuming exact U -spin and $A_-/A_0 = 0.3$. The NP WCs h_i are in TeV^{-2} unit.

NP type	Lorentz structure	Best-fit value (h_i)	$\chi^2/\text{d.o.f.}$
Scalar	$(S - P) \otimes (S - P)$	0.066 ± 0.005	340.0/11
	$(S - P) \otimes (S + P)$	-0.168 ± 0.015	198.4/11
	$(S + P) \otimes (S - P)$	0.132 ± 0.009	484.2/11
	$(S + P) \otimes (S + P)$	0.068 ± 0.005	400.9/11
Vector	$(V - A) \otimes (V - A)$	0.051 ± 0.006	622.6/11
	$(V - A) \otimes (V + A)$	0.107 ± 0.008	565.7/11
	$(V + A) \otimes (V - A)$	-0.107 ± 0.008	565.7/11
	$(V + A) \otimes (V + A)$	-0.051 ± 0.006	622.6/11
Tensor	$(T \pm PT) \otimes (T - PT)$	-0.029 ± 0.002	163.6/11
	$(T \pm PT) \otimes (T + PT)$	0.032 ± 0.002	224.4/11

(iii) Tensor NP:

$$\begin{aligned}
 a &= -\frac{h_t e^{i\zeta_t}}{N_c} c_1 c_4 f_{K^{*0}}^T [(m_B^2 - m_\rho^2) g_+(q^2) + q^2 g_-(q^2)], \\
 \frac{b}{m_B^2} &= \frac{h_t e^{i\zeta_t}}{N_c} c_1 c_4 f_{K^{*0}}^T [2g_+(q^2) + q^2 g_0(q^2)], \\
 \frac{c}{m_B^2} &= \frac{h_t e^{i\zeta_t}}{N_c} 2c_1 c_3 f_{K^{*0}}^T g_+(q^2), \quad (26)
 \end{aligned}$$

where the a , b , and c terms are contributions from NP only. Here, N_c is the number of colors, $f_{K^{*0}}$ and $f_{K^{*0}}^T$ are the decay constants, and $\mathfrak{B}$, $\mathfrak{A}_1$, $\mathfrak{A}_2$, g_+ , g_0 , and g_- are the $B \rightarrow \rho$ form factors, the definitions of which can be found in the Appendix. For our analysis, we take $f_{K^*} = 0.204 \pm 0.007 \text{ GeV}$ and $f_{K^*}^T/f_{K^*} = 0.712 \pm 0.012$ [36]. Similar expressions for other decay channels may be obtained by suitable replacement of the daughter mesons and appropriate changes for the form factors, decay constants, and masses. Our parametrization of the NP operators, consistent with a hypothetical flavor-changing charge-neutral color-singlet mediator that has been integrated out, leads to all the NP amplitudes being color suppressed.¹⁴

When we take both $\Delta S = 0$ and $\Delta S = 1$ modes into account and impose (i) exact $\text{SU}(2)_U$ for the SM part of the amplitudes and (ii) the hierarchy $|A_-/A_0| = 0.3$ for all the SM amplitudes, we have 18 measurements and only seven free parameters: six SM RMEs plus one coming from the NP, and characterized by h_i , with $i \in \{s, v, t\}$. The fit results are shown in Tables III and IV. The Lorentz structure of NP is shown as the product of two currents; the first one is the $\bar{b}s$ current, and the second one is the $\bar{q}q$ current.

A look at the last column in Table III immediately gives the impression that the fit, even with an order-of-magnitude

improvement in χ^2 from the corresponding SM fit just with the introduction of one single parameter, namely, the NP coupling h_i , still remains very poor. While this impression is not wrong, we would like to analyze the individual contributions of different observables to the final χ^2 to pinpoint the source of the tension. We have explicitly checked that a moderate breaking of $\text{SU}(2)_U$, similar to what we have used before, hardly changes either the best-fit values or the χ^2 .

From Table III, the most promising NP combinations are either scalar type: $(S - P) \otimes (S + P)$, with $\{c_1, c_2, c_3, c_4\} = \{1, -1, 1, 1\}$, or tensor type: $(T \pm PT) \otimes (T - PT)$, with $\{c_1, c_2, c_3, c_4\} = \{1, \pm 1, 1, 1\}$. For tensor operators, c_2 does not play role, as can be seen from Eq. (26). Let us concentrate only on these combinations.

It is obvious that the more an operator can raise the transverse polarization fraction, the more successful it will be in bringing down the χ^2 . However, note that the NP operators are only of $\Delta S = 1$ type. There is one observable in the $\Delta S = 0$ channel—the LPF of $B_d \rightarrow K^{*0} \bar{K}^{*0}$ —that contributes ~ 125 to χ^2 . This was so for the SM fit, and this remains true for the SM + NP fit, too.¹⁵ To highlight the improvement of the fit for the $\Delta S = 1$ modes, we perform the same exercise taking only these channels into account. This naturally takes the U spin out of consideration. There are only nine observables and seven free parameters, so we expect the fit to improve significantly. This is indeed the case, and the total χ^2 values for both the cases drop below 20. We have also checked explicitly that introduction of a nonzero NP weak phase ζ_i hardly changes the fit. There is no significant tension with the CP asymmetries for these modes, all of them being consistent with zero. Even if we introduce a NP weak phase ζ_i , the direct CP asymmetries will also depend on the strong phase difference of the relevant amplitudes, which are, *a priori*, unknown, and, therefore, this hardly puts any significant constraint on the NP weak phase ζ_i . The mixing-induced CP asymmetry is expected to be tiny, as the $B_s - \bar{B}_s$ mixing amplitude is almost real within the SM. Thus, we do not lose anything by keeping all $\zeta_i = 0$; it only simplifies the analysis.

Table IV also shows the ratios of the transverse (A_-) and longitudinal (A_0) amplitudes for the NP and the SM. One may note that the ratio is consistently larger than unity for the transverse amplitudes and much smaller than unity for the longitudinal amplitudes; in other words, there is no hierarchy between A_- and A_0 for the NP amplitudes. This leads to smaller f_L values, in agreement with the data. Also note that $|A_0^{\text{NP}}| \ll |A_0^{\text{SM}}|$. With the BRs being dominated by A_0^{SM} , these are not significantly affected by the NP and remain consistent with the experimental numbers.

¹⁴We could have started with a NP Hamiltonian that is a product of two color-octet currents, and then all our amplitudes would have been color allowed.

¹⁵Thus, one may reasonably argue that NP should also occur in $b \rightarrow d$ transitions. However, in a data-driven analysis like this, there is no way one can relate the strengths of $b \rightarrow s$ and $b \rightarrow d$ transitions.

TABLE IV. The best-fit polarization fractions for various $\Delta S = 1$ modes and for some select Lorentz structures of the NP Hamiltonian. The two numbers shown immediately below f_L (and $f_\perp$) are the NP and SM transverse (A_-) and longitudinal (A_0) amplitude ratios. Note that A_+ (SM) has been taken to be zero in the analysis, but A_+ (NP) is nonzero. We also show the prediction for $f_L(B_s \rightarrow K^{*+} K^{*-})$, indicated by ♠.

Mode	$(S - P) \otimes (S + P)$		$(T \pm PT) \otimes (T - PT)$	
	$ A_-^{\text{NP}} / A_-^{\text{SM}} $	$ A_0^{\text{NP}} / A_0^{\text{SM}} $	$ A_-^{\text{NP}} / A_-^{\text{SM}} $	$ A_0^{\text{NP}} / A_0^{\text{SM}} $
$B_s \rightarrow K^{*0} \overline{K}^{*0}$	$f_L = 0.18 \pm 0.01$ $f_\perp = 0.53 \pm 0.01$ 1.08 ± 0.06	0.92 ± 0.05	$f_L = 0.18 \pm 0.01$ $f_\perp = 0.53 \pm 0.01$ 1.94 ± 0.63	0.15 ± 0.05
$B^+ \rightarrow K^{*0} \rho^+$	$f_L = 0.72 \pm 0.04$ 2.16 ± 0.30	1.79 ± 0.25	$f_L = 0.42 \pm 0.07$ 1.14 ± 0.15	0.09 ± 0.01
$B_d \rightarrow K^{*+} \rho^-$	$f_L = 0.78 \pm 0.02$ 2.31 ± 0.26	1.94 ± 0.22	$f_L = 0.39 \pm 0.01$ 1.27 ± 0.33	0.05 ± 0.01
♠ $B_s \rightarrow K^{*+} K^{*-}$	$f_L = 0.90 \pm 0.01$		$f_L = 0.63 \pm 0.05$	

TABLE V. BRs, multiplied by 10^6 , for several modes affected by the NP operators discussed in the text. The numbers for the experimental data are taken from Ref. [37].

Mode	Data ($\times 10^6$)	$(S - P) \otimes (S + P) (\times 10^6)$	$(T \pm PT) \otimes (T - PT) (\times 10^6)$
$B_s \rightarrow K^+ K^-$	26.1 ± 1.6	28.1 ± 3.2	19.2 ± 3.8
$B_s \rightarrow K^0 \overline{K}^0$	17.6 ± 3.1	25.2 ± 4.1	17.0 ± 3.4
$B_s \rightarrow K^{*0} \overline{K}^{*0}$	11.1 ± 2.7	0.93 ± 0.06	5.61 ± 0.93
$B_d \rightarrow K^{*+} \rho^-$	10.3 ± 2.0	9.93 ± 2.00	9.55 ± 1.80
$B_d \rightarrow \pi^- K^+$	20.0 ± 0.4	18.5 ± 1.2	18.8 ± 0.9
$B_d \rightarrow \pi^0 K^0$	10.1 ± 0.4	9.7 ± 0.7	9.1 ± 0.9
$B^+ \rightarrow K^{*0} \rho^+$	9.2 ± 1.5	7.32 ± 1.80	10.34 ± 1.33
$B^+ \rightarrow \pi^+ K^0$	23.9 ± 0.6	20.5 ± 1.8	21.2 ± 0.8
$B^+ \rightarrow \pi^0 K^+$	13.2 ± 0.4	11.2 ± 0.9	12.1 ± 0.5

The BRs for the $B \rightarrow V_1 2V_2$ modes are used as inputs to the fit. With the best-fit values for the NP couplings h_i —for both $(S - P) \otimes (S + P)$ and $(T \pm PT) \otimes (T - PT)$ —we have checked how far the BRs are affected even with final states that include pseudoscalar mesons. Our results are shown in Table V; all the modes, including those that are not shown explicitly, are consistent with the data within 1σ , except $B_s \rightarrow K^{*0} \overline{K}^{*0}$, which shows some tension.

One may improve the consistency marginally with various measurements of A_{CP} if h_i is taken to be complex (which, in turn, reduces χ^2 further). This, however, is not needed for the study of LPFs, and we have refrained from doing so. It has been checked that all other observables, e.g., mass and width differences in the $B_s - \overline{B}_s$ system and $\text{BR}(B \rightarrow X_s \gamma)$, remain consistent with data even after NP effects are included.

Our assumption of the existence of NP only for the $\Delta S = 1$ channels is rather simplistic. There is no symmetry that prevents such NP for the $\Delta S = 0$ channels, i.e., for the $b \rightarrow d$ transitions. Even a basis rotation in the down-quark sector can induce such NP for $b \rightarrow d$. While we do not

perform any detailed analysis of this scenario here, it is immediately obvious that this NP can reduce the tension in $f_L(B_d \rightarrow K^{*0} \overline{K}^{*0})$ and may lead to an acceptable fit.

IV. CONCLUSIONS

In this work, we perform a U -spin analysis of B meson decays into two light vector mesons $V_1 V_2$, where $V_1, V_2 \in \{\rho, K^*\}$. These decays form a set of eight processes related by the $d \leftrightarrow s$ exchange symmetry of U spin. Out of these, we have data on six modes. The available observables include CP -averaged branching ratios, CP asymmetries, and polarization fractions (mostly longitudinal, but in a few cases transverse, too).

Our motivation has been to probe the systematically low LPFs for $\Delta S = 1$ decays, which is most aggravated for the decay $B_s \rightarrow K^{*0} \overline{K}^{*0}$. On the other hand, LPFs for the $\Delta S = 0$ decays are consistent with theoretical expectation based on the SM dynamics, which predicts LPFs to be close to unity if nonfactorizable effects are neglected.

The rationale for using U spin is as follows. Without U spin, we would have lost the ability to express the amplitudes in a compact and structured form. In such cases, we would need to parametrize the various contributions based on topological classifications or adopt factorization approaches, systematically accounting for nonfactorizable corrections. In principle, nonfactorizable contributions, such as annihilation or other subleading topologies, can also be interpreted in terms of U -spin-symmetric or U -spin-broken RMEs. This interpretation reinforces the central role of U -spin symmetry, even when considering corrections and breaking effects.

To summarize the chain of reasoning, the decomposition of the amplitudes is carried out using U -spin symmetry, which forms the backbone of our analysis. While the U -spin-symmetric approximation alone cannot resolve the polarization puzzle, it provides the essential framework. However, as U spin is not an exact symmetry of nature, we have to incorporate the U -spin-breaking effects. We found that even a sizable breaking (larger than the expected level) is not enough to satisfactorily explain the data. Then we tried some flavor-specific NP to address the puzzle and found a much improved solution.

We take all the six modes, three each for $\Delta S = 0$ and $\Delta S = 1$, and perform an analysis based on the SM only. We find that (i) the quality of the fit does not at all depend on whether the U spin is exact or moderately broken, but (ii) it severely depends on how seriously we take the hierarchy among the polarization fractions. If we consider LPF dominance as suggested by theoretical models, the fit becomes abysmally poor. On the other hand, if we keep polarization fractions as free parameters, the fit becomes excellent, but the hierarchy is no longer respected. The transverse polarization amplitudes become comparable to, or even bigger than, the LPFs, which is expected from the data.

Given this conundrum, we entertain the idea that there might be some new physics that lowers the LPFs for the $\Delta S = 1$ modes. We take a bottom-up approach and use a simplified effective theory where NP appears in $b \rightarrow s$ transitions but not in $b \rightarrow d$ transitions. To make the analysis as simple as possible, we assume exact U -spin symmetry among the SM parts of the amplitudes, as well as the hierarchy $|A_-/A_0| = 0.3$ between the polarization

fractions as dictated by naive factorization within SM. No such hierarchy has been assumed for the NP part, which is natural, as Lorentz structures other than the SM one do not reproduce the hierarchy.

We find a very significant improvement in the fit for almost all types of NP, even without considering a nonzero weak phase in the NP Hamiltonian. The reason is obvious: The new Lorentz structures do not respect the hierarchy and, hence, can lower the LPFs. We have also checked that the fit results to other observables, namely, the CP -averaged branching ratios and CP asymmetries, are consistent with the data. Finally, we make predictions for the LPF of the yet-to-be-observed mode $B_s \rightarrow K^{*+} K^{*-}$, which has the power to differentiate among various NP options, as can be seen in Table IV. We expect our experimental colleagues to observe and measure this mode in the near future.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX: FORM FACTORS

The vector and axial-vector form factors for $B \rightarrow V_2$ are defined as

$$\begin{aligned} \langle V_2(k, \epsilon_2) | \bar{q} \gamma_\mu b | \bar{B}(p) \rangle &= -i \epsilon_{\mu\alpha\beta} \epsilon_2^{*\nu} p^\alpha k^\beta \frac{2\mathfrak{V}(q^2)}{m_B + m_{V_2}}, \\ \langle V_2(k, \epsilon_2) | \bar{q} \gamma_\mu \gamma_5 b | \bar{B}(p) \rangle &= \epsilon_2^*(m_B + m_{V_2}) \mathfrak{A}_1(q^2) - (p + k)_\mu (\epsilon_2^* \cdot q) \frac{\mathfrak{A}_2(q^2)}{m_B + m_{V_2}} \\ &\quad - q_\mu (\epsilon_2^* \cdot q) \frac{2m_{V_2}}{q^2} [\mathfrak{A}_3(q^2) - \mathfrak{A}_0(q^2)], \end{aligned} \quad (\text{A1})$$

with

$$\mathfrak{A}_3(q^2) = \frac{m_B + m_{V_2}}{2m_V} \mathfrak{A}_1(q^2) - \frac{m_B - m_{V_2}}{2m_V} \mathfrak{A}_2(q^2). \quad (\text{A2})$$

The scalar matrix element is zero, whereas the pseudoscalar matrix element is given by

$$\langle V_2(k, \epsilon_2) | \bar{q} \gamma_5 b | \bar{B}(p) \rangle = -(\epsilon_2^* \cdot q) \frac{2m_{V_2}}{m_b + m_q} \mathfrak{A}_0(q^2). \quad (\text{A3})$$

The tensor contributions can be parametrized as

$$\begin{aligned} \langle V_2(k, \epsilon_2) | \bar{q} \sigma_{\mu\nu} b | \bar{B}(p) \rangle &= \epsilon_{\mu\nu\alpha\beta} [\epsilon_2^{*\alpha} (p+k)^\alpha g_+(q^2) + \epsilon_2^{*\alpha} q^\beta g_-(q^2) + (\epsilon_2^* \cdot q) p^\alpha k^\beta g_0(q^2)], \\ \langle V_2(k, \epsilon_2) | \bar{q} \sigma_{\mu\nu} \gamma_5 b | \bar{B}(p) \rangle &= i [\{ \epsilon_{2\mu}^* (p+k)_\nu - (p+k)_\mu \epsilon_{2\nu}^* \} g_+(q^2) + \{ \epsilon_{2\mu}^* q_\nu - q_\mu \epsilon_{2\nu}^* \} g_-(q^2) \\ &\quad + (\epsilon_2^* \cdot q) \{ p_\mu k_\nu - k_\mu p_\nu \} g_0(q^2)], \end{aligned} \quad (\text{A4})$$

where $g_{\pm,0}$ are related to commonly used $T_{1,2,3}$ form factors by following

$$\begin{aligned} g_+(q^2) &= -T_1(q^2), \\ g_-(q^2) &= \frac{m_B^2 - m_{V_2}^2}{q^2} [T_1(q^2) - T_2(q^2)], \\ g_0(q^2) &= \frac{2}{q^2} \left[T_1(q^2) - T_2(q^2) - \frac{q^2}{m_B^2 - m_{V_2}^2} T_3(q^2) \right]. \end{aligned} \quad (\text{A5})$$

We use the light cone sum rule prescription to evaluate the form factors. We use the expansion in the exact same form as described in Ref. [38]:

$$\mathfrak{F}_i(q^2) = \frac{1}{1 - q^2/m_{R,i}^2} \sum_{n=0}^N \alpha_n^i [z(q^2) - z(0)]^n, \quad (\text{A6})$$

where

$$\begin{aligned} z(q^2) &= \frac{\sqrt{t_+ - q^2} - \sqrt{t_+ - t_0}}{\sqrt{t_+ - q^2} + \sqrt{t_+ - t_0}} \quad \text{and} \\ t_0 &= t_+ \left(1 - \sqrt{1 - \frac{t_-}{t_+}} \right). \end{aligned} \quad (\text{A7})$$

Here, $t_\pm = (m_B \pm m_{V_2})^2$; the masses of resonances $m_{R,i}$ and the expansion parameters α_n^i can be found in Ref. [38]. We calculate the $B \rightarrow \rho$ and $B \rightarrow K^{*0}$ form factors at $q^2 = m_{K^0}^2$ and $q^2 = m_{\bar{K}^{*0}}^2$, respectively. The matrix elements from a vacuum to a vector meson are parametrized as follows:

$$\begin{aligned} \langle V_1(q, \epsilon_1) | \bar{q}' \gamma^\mu q | 0 \rangle &= f_{V_1} m_{V_1} \epsilon_1^\mu, \\ \langle V_1(q, \epsilon_1) | \bar{q}' \sigma^{\mu\nu} q | 0 \rangle &= i f_{V_1}^T (q^\mu \epsilon_1^\nu - \epsilon_1^\mu q^\nu), \end{aligned} \quad (\text{A8})$$

where f_{V_1} and $f_{V_1}^T$ are the decay constants of the V_1 meson. The matrix elements for the remaining Lorentz structures are zero.

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