New-physics effects on triple-product correlations in Λ_b decays

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We adopt an effective-Lagrangian approach to compute the new-physics contributions to *T*-violating triple-product correlations in charmless Λ_b decays. We use factorization and work to leading order in the heavy-quark expansion. We find that the standard-model (SM) predictions for such correlations can be significantly modified. For example, triple products which are expected to vanish in the SM can be enormous ($\sim 50\%$) in the presence of new physics. By measuring triple products in a variety of Λ_b decays, one can diagnose which new-physics operators are or are not present. Our general results can be applied to any specific model of new physics by simply calculating which operators appear in that model.

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I. INTRODUCTION

The origin of CP violation remains one of the important open questions in particle physics. Within the standard model (SM), CP violation is due to the presence of phases in the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix. The *B* factories BaBar and Belle have been built to test this: if the SM explanation is correct, we expect to observe large CP-violating rate asymmetries in *B* decays [1]. To date, one of the CP phases of the unitarity triangle has been measured: $\sin 2\beta = 0.78 \pm 0.08$ [2], which is consistent with the SM.

Although the main focus has been on rate asymmetries, there is another type of CP-violating signal which could potentially reveal the presence of physics beyond the SM. Triple-product correlations of the form $\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3)$, where each v_i is a spin or momentum, are odd under time reversal (*T*). Therefore, by the CPT theorem, these are also signals of CP violation. A nonzero triple-product correlation is signalled by a nonzero value of the asymmetry

$$A_T \equiv \frac{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) - \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)}{\Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) > 0) + \Gamma(\vec{v}_1 \cdot (\vec{v}_2 \times \vec{v}_3) < 0)}, \quad (1)$$

where Γ is the decay rate for the process in question. However, there is a well-known caveat: strong phases can produce a nonzero value of A_T, even if the weak phases are zero (i.e. CP violation is not really present). Thus, to be sure that one is truly probing *T* and CP violation, one must compare the value of A_T with that of $\bar{A}_T$, which is the *T*-odd asymmetry measured in the CP-conjugate decay process.

Triple-product correlations can be measured in $B \to V_1 V_2$ decays, where V_1 and V_2 are vector mesons [3]. In the rest frame of the *B*, the triple product takes the form $\vec{p} \cdot (\varepsilon_1 \times \varepsilon_2)$, where $\vec{p}$ is the momentum of one of the final-state particles, and ε_i is the polarization of the V_i. One can also consider triple-product correlations in Λ_b decays. Since many such triple products involve the spin of the Λ_b, this means that, in contrast to *B* decays, one is sensitive to the spin of the *b*-quark [4], as it is expected to provide the dominant contribution to the spin of the Λ_b.

In a recent paper [5], we used factorization to study the SM predictions for triple products in charmless two-body Λ_b decays. We considered decays which are generated by the quark-level transitions $b \to s\bar{q}q$ or $b \to d\bar{q}q$. These decays take the form $\Lambda_b \to F_1 F_2$, where F_1 is a light spin-$\frac{1}{2}$ baryon, such as *p*, Λ, etc., and F_2 is a pseudoscalar (*P*) or vector (*V*) meson. There was only one decay in which there was a large effect: the triple-product asymmetry for $\Lambda_b \to pK^-$ was found to be 18%. For all other decays, the asymmetries are found to be at most at the percent level.

The fact that all these triple-product asymmetries are expected to be small in the SM suggests that this is a good area to look for physics beyond the SM. In this paper, we examine the effect of new physics on triple products in charmless Λ_b decays. In order to study this, we adopt an effective-Lagrangian approach: we write down all possible dimension-6 new-physics four-Fermi operators at the quark level. Then, using factorization, we compute their contributions to the various triple-product correlations in Λ_b decays.

There are several advantages to this approach. First, we are able to establish which triple products can be significantly affected by the presence of new physics. Second, we can also determine specifically which new-physics operators contribute to these triple products. Thus, by measuring a number of different triple-product correlations, we may be able to diagnose which operators are or are not present. Finally, these operators include all possible models of new physics. Therefore one can apply our results to a specific model by simply calculating which new-physics operators appear in that model. We will give examples of this procedure.

The paper is organized as follows. In Sec. II, we introduce the new-physics operators used in our analysis. We also give two examples of specific models which generate some of these operators: supersymmetry with *R*-parity breaking, and *Z*- and Z'-mediated flavor-changing neutral currents. We compute the contributions of the new-physics operators to triple-product correlations in Λ_b decays in Sec. III. Here we retain only the leading term in the heavy-quark expansion

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since it is very unlikely that new physics contributes to subleading processes without affecting the leading-order processes. In Sec. IV, we estimate the size of the various triple products in the presence of new physics. By comparing triple products in $\Lambda_b \to F_1 P$ and $\Lambda_b \to F_1 V$ decays, we examine the ''diagnostic power'' of this approach, i.e. the extent to which one can determine which new-physics operators are present. We also show how our results can be applied to the specific models of new physics discussed previously. We conclude in Sec. V.

II. NEW PHYSICS

We are interested in charmless Λ_b decays, which are governed by the quark-level processes $b \to s\bar{q}q$ or $b \to d\bar{q}q$. In what follows we will concentrate on the $b \to s$ transitions; it is straightforward to adapt our analysis to the $b \to d$ case.

Taking into account the two different color structures, as well as all possible Lorentz structures, there are a total of 20 dimension-6 new-physics operators which contribute to each of the $b \to s\bar{q}q$ transitions, $q = u,d,s$. These can be written as

$$\begin{aligned}\mathcal{H}_{NP}^q = \sum_{A,B=L,R} \frac{4G_F}{\sqrt{2}} \{ & f_{q,1}^{AB} \bar{s}_\alpha \gamma_A b_\beta \bar{q}_\beta \gamma_B q_\alpha \\ & + f_{q,2}^{AB} \bar{s} \gamma_A b \bar{q} \gamma_B q + g_{q,1}^{AB} \bar{s}_\alpha \gamma^\mu \gamma_A b_\beta \bar{q}_\beta \gamma_\mu \gamma_B q_\alpha \\ & + g_{q,2}^{AB} \bar{s} \gamma^\mu \gamma_A b \bar{q} \gamma_\mu \gamma_B q \\ & + h_{q,1}^{AB} \bar{s}_\alpha \sigma^{\mu\nu} \gamma_A b_\beta \bar{q}_\beta \sigma_{\mu\nu} \gamma_B q_\alpha \\ & + h_{q,2}^{AB} \bar{s} \sigma^{\mu\nu} \gamma_A b \bar{q} \sigma_{\mu\nu} \gamma_B q \}, \end{aligned} \tag{2}$$

where we have defined $\gamma_{R(L)} = \frac{1}{2}(1 \pm \gamma_5)$. Note: although we have written the tensor operators in the same compact form as the other operators, it should be noted that those with $\gamma_A \neq \gamma_B$ are identically zero. Thus, one can effectively set $h_{q,i}^{LR} = h_{q,i}^{RL} = 0$.

All models of new physics which contribute to $b \to s\bar{q}q$ will generate operators found in the above effective Hamiltonian. These can arise at tree level (e.g. supersymmetry with R-parity breaking, Z- and Z'-mediated flavor-changing neutral currents, models with flavor-changing neutral scalars, etc.) or at loop level (e.g. minimal supersymmetry, left-right symmetric models, four generations, etc.) [6]. In some cases one will obtain operators of the form $\bar{q}\mathcal{O}b\bar{s}\mathcal{O}'q$, but one can perform a Fierz transformation to put them into the form of Eq. (2). Note that, in general, models of new physics do not lead directly to tensor operators ($h_{q,i}^{AB}$), since typically only vector or scalar particles are involved. However, such tensor operators can arise when other operators are Fierz-transformed into the above form, so they must be included in our analysis (the scalar operator $\bar{q}b\bar{s}q$ is such an example).

Because the new-physics operators are of dimension 6, by dimensional analysis we expect them to be suppressed by a factor Λ^2, where Λ is the scale of new physics. However, with the normalization in Eq. (2), the suppression factor is only M_W^2. We therefore expect the size of the coefficients $f_{q,i}^{AB}$, $g_{q,i}^{AB}$ and $h_{q,i}^{AB}$ to be naturally of $O(M_W^2/\Lambda^2) \sim 10^{-2}$ for a new-physics scale of about 1 TeV.

Even so, these new-physics effects can be quite significant. In the SM, one finds only operators of the form $\bar{s}\gamma^\mu \gamma_L b \bar{q} \gamma_\mu \gamma_{L,R} q$, with both color assignments. These operators are typically multiplied by one of two factors: either (i) the CKM matrix elements $V_{tb}V_{ts}^*$ times a Wilson coefficient of $O(10^{-2})$, or (ii) $V_{ub}V_{us}^*$ times a Wilson coefficient of $O(1)$. In either case, new-physics operators with coefficients of $O(10^{-2})$ would actually *dominate* over the SM contributions. (This is, in part, what allows us to put constraints on specific models of new physics.) The bottom line is that the new operators of Eq. (2) can contribute substantially to charmless Λ_b decays.

As noted above, by construction the effective Hamiltonian of Eq. (2) includes all possible models of new physics. Of course, in a particular new-physics model, only a subset of the new operators will appear. Our general analysis can then be applied to that specific model by retaining only the coefficients of the nonzero operators. In order to show explicitly how this works, below we give two examples of such specific models.

A. Supersymmetry with *R*-parity breaking

In supersymmetric models, the R-parity of a field with spin S, baryon number B and lepton number L is defined to be

$$R = (-1)^{2S+3B+L}. \tag{3}$$

R is $+1$ for all the SM particles and -1 for all the supersymmetric particles. R-parity invariance is often imposed on the Lagrangian in order to maintain the separate conservation of baryon number and lepton number. Imposition of R-parity conservation has some important consequences: superparticles must be produced in pairs in collider experiments and the lightest superparticle (LSP) must be absolutely stable. The LSP therefore provides a good candidate for cold dark matter.

Despite the above-mentioned attractive features of R-parity conservation, this conservation is not dictated by any fundamental principle such as gauge invariance, so that there is no compelling theoretical motivation for it. The most general superpotential of the MSSM, consistent with $SU(3) \times SU(2) \times U(1)$ gauge symmetry and supersymmetry, can be written as

$$\mathcal{W} = \mathcal{W}_R + \mathcal{W}_{\not{R}}, \tag{4}$$

where $\mathcal{W}_R$ is the R-parity conserving piece, and $\mathcal{W}_{\not{R}}$ breaks R parity. They are given by

$$\mathcal{W}_R = h_{ij} L_i H_2 E_j^c + h'_{ij} Q_i H_2 D_j^c + h''_{ij} Q_i H_1 U_j^c, \tag{5}$$

$$\mathcal{W}_{\not R}=\frac{1}{2}\lambda_{[ij]k}L_iL_jE_k^c+\lambda'_{ijk}L_iQ_jD_k^c + \frac{1}{2}\lambda''_{i[jk]}U_i^cD_j^cD_k^c+\mu_iL_iH_2. \quad (6)$$

Here $L_i(Q_i)$ and $E_i(U_i,D_i)$ are the left-handed lepton (quark) doublet and lepton (quark) singlet chiral superfields, where i,j,k are generation indices and c denotes a charge conjugate field. $H_{1,2}$ are the chiral superfields representing the two Higgs doublets.

In the R-parity-violating superpotential [Eq. (5)], the λ and λ' couplings violate lepton number conservation, while the λ'' couplings violate baryon number conservation. $\lambda_{[ij]k}$ is antisymmetric in the first two indices and $\lambda''_{i[jk]}$ is antisymmetric in the last two indices. There are therefore 27 λ'-type couplings and 9 each of the λ and λ'' couplings. While it is theoretically possible to have both baryon-number and lepton-number violating terms in the Lagrangian, the nonobservation of proton decay imposes very stringent conditions on their simultaneous presence [7]. One therefore assumes the existence of either L-violating couplings or B-violating couplings, but not both. The terms proportional to λ are not relevant to our present discussion and will not be considered further.

We begin with the B-violating couplings. The transition $b\to s\bar{u}u$ can be generated at tree level through the t-channel exchange of the d-squark, $\widetilde{d_R}$, with strength proportional to $|\lambda''_{112}\lambda^{*''}_{113}|$. However, this product of couplings is already constrained to be $\sim 10^{-8}$ from $n-\bar{n}$ oscillations and double nucleon decay [8]. There are therefore no significant contributions to the new-physics operators of Eq. (2) corresponding to $q=u$.

Similarly, the antisymmetry of the B-violating couplings, $\lambda''_{i[jk]}$ in the last two indices implies that there are no operators that can generate the $b\to s\bar{s}s$ transition, so that all the operators in Eq. (2) vanish for $q=s$.

Finally, the operators that generate the $b\to s\bar{d}d$ transition are given by [9]

$$L_{eff}=\frac{\lambda''_{i12}\lambda''^{*}_{i13}}{4m^2_{\tilde{u}_i}}(\bar{d}_\alpha\gamma_\mu\gamma_R d_\alpha\bar{s}_\beta\gamma_\mu\gamma_R b_\beta - \bar{d}_\alpha\gamma_\mu\gamma_R d_\beta\bar{s}_\beta\gamma_\mu\gamma_R b_\alpha). \quad (7)$$

Hence the only nonvanishing operators in Eq. (2) are

$$g^{RR}_{d,1}=-g^{RR}_{d,2}=-\frac{\sqrt{2}}{G_F}\frac{\lambda''_{i12}\lambda''^{*}_{i13}}{16m^2_{\tilde{u}_i}}. \quad (8)$$

As mentioned above, the constraint on $|\lambda''_{112}\lambda''^{*}_{113}|$ is at the 10^{-8} level. However, the constraint on $|\lambda''_{i12}\lambda''^{*}_{i13}|$, $i\neq 1$, comes only from the nonleptonic decay $B^-\to\bar{K}^0\pi^-$ [9], and is much weaker:

$$|\lambda''_{i12}\lambda''^{*}_{i13}|(i\neq 1)\leq 1.03\times 10^{-2}, \quad (9)$$

where a squark mass $m_{\tilde{f}}=100$ GeV has been assumed. We therefore find

$$|g^{RR}_{d,1}|=|g^{RR}_{d,2}|\leq 7.6\times 10^{-3}. \quad (10)$$

We now turn to the L-violating couplings. In terms of four-component Dirac spinors, these are given by [10]

$$\mathcal{L}_{\lambda'}=-\lambda'_{ijk}[\tilde{\nu}^i_L\bar{d}^k_R d^j_L+\tilde{d}^j_L\bar{d}^k_R\nu^i_L+(\tilde{d}^k_R)^*(\overline{\nu^i_L})^c d^j_L-\tilde{e}^i_L\bar{d}^k_R u^j_L - \tilde{u}^j_L\bar{d}^k_R e^i_L-(\tilde{d}^k_R)^*(\overline{e^i_L})^c u^j_L]+\text{H.c.} \quad (11)$$

There are a variety of sources which bound the above couplings [8,9]. For the sake of brevity we will only quote the bounds and not their sources. Assuming a commom sfermion mass of 100 GeV we find the most stringent bounds are

$$\begin{aligned}
&|\lambda'_{i12}\lambda'^{*}_{i13}|(i\neq 1)\leq 1.7\times 10^{-3},\\
&|\lambda'_{112}\lambda'^{*}_{113}|\leq 4.4\times 10^{-4},\\
&|\lambda'_{111}\lambda'^{*}_{132}|\leq 1.4\times 10^{-4},\\
&|\lambda'^{*}_{211}\lambda'_{232}|\leq 4.7\times 10^{-4},\\
&|\lambda'_{311}\lambda'^{*}_{332}|\leq 4.7\times 10^{-4},\\
&|\lambda'_{111}\lambda'^{*}_{123}|\leq 2.2\times 10^{-5},\\
&|\lambda'_{211}\lambda'^{*}_{223}|\leq 2.2\times 10^{-3},\\
&|\lambda'_{311}\lambda'^{*}_{323}|\leq 2.2\times 10^{-3},\\
&|\lambda'_{131}\lambda'^{*}_{121}|\leq 8.2\times 10^{-4},\\
&|\lambda'_{231}\lambda'^{*}_{221}|\leq 1.3\times 10^{-3},\\
&|\lambda'_{331}\lambda'^{*}_{321}|\leq 1.3\times 10^{-3},\\
&|\lambda'_{132}\lambda'^{*}_{122}|\leq 1.2\times 10^{-2},\\
&|\lambda'_{232}\lambda'^{*}_{222}|\leq 1.2\times 10^{-1},\\
&|\lambda'_{332}\lambda'^{*}_{322}|\leq 2.3\times 10^{-1},\\
&|\lambda'_{122}\lambda'^{*}_{123}|\leq 1.8\times 10^{-3},\\
&|\lambda'_{223}\lambda'^{*}_{222}|\leq 4.4\times 10^{-2},\\
&|\lambda'_{322}\lambda'^{*}_{323}|\leq 2.7\times 10^{-1}.
\end{aligned} \quad (12)$$

There is a single contribution to the $b\to s\bar{u}u$ transition:

$$L_{eff}=-\frac{\lambda'_{i12}\lambda'^{*}_{i13}}{2m^2_{\tilde{e}_i}}\bar{u}_\alpha\gamma_\mu\gamma_L u_\beta\bar{s}_\beta\gamma_\mu\gamma_R b_\alpha. \quad (13)$$

Hence the only nonvanishing operator for $q=u$ in Eq. (2) is

$$g_{u,1}^{RL} = -\frac{\sqrt{2}}{G_F}\frac{\lambda'_{i12}\lambda'^{*}_{i13}}{8m^2_{\tilde{e}_i}}. \quad (14)$$

Using the bounds of Eq. (12), we find

$$|g_{u,1}^{RL}| \leqslant 2.6\times 10^{-3}. \quad (15)$$

Turning now to the $b\to s\bar{d}d$ transition the relevant Lagrangian is

$$L_{eff} = \frac{\lambda'_{i11}\lambda'^{*}_{i23}}{m^2_{\tilde{\nu}_i}}\bar{d}\gamma_L d_\beta \bar{s}\gamma_R b + \frac{\lambda'_{i32}\lambda'^{*}_{i11}}{m^2_{\tilde{\nu}_i}}\bar{d}\gamma_R d\bar{s}\gamma_L b - \frac{\lambda'_{i12}\lambda'^{*}_{i13}}{2m^2_{\tilde{\nu}_i}}\bar{d}_\alpha\gamma_\mu\gamma_L d_\beta \bar{s}_\beta\gamma_\mu\gamma_R b_\alpha - \frac{\lambda'_{i31}\lambda'^{*}_{i21}}{2m^2_{\tilde{\nu}_i}}\bar{d}_\alpha\gamma_\mu\gamma_R d_\beta \bar{s}_\beta\gamma_\mu\gamma_L b_\alpha . \quad (16)$$

The nonvanishing operators in Eq. (2) are then

$$f_{d,2}^{LR} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i32}\lambda'^{*}_{i11}}{4m^2_{\tilde{\nu}_i}}, \quad f_{d,2}^{RL} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i11}\lambda'^{*}_{i23}}{4m^2_{\tilde{\nu}_i}},$$

$$g_{d,1}^{RL} = g_{u,1}^{RL*}, \quad g_{d,1}^{LR} = -\frac{\sqrt{2}}{G_F}\frac{\lambda'_{i31}\lambda'^{*}_{i21}}{8m^2_{\tilde{\nu}_i}}, \quad (17)$$

with

$$|f_{d,2}^{LR}| \leqslant 1.4\times 10^{-3}, \quad |f_{d,2}^{RL}| \leqslant 6.6\times 10^{-3},$$

$$|g_{d,1}^{RL}| \leqslant 2.6\times 10^{-3}, \quad |g_{d,1}^{LR}| \leqslant 2.0\times 10^{-3}. \quad (18)$$

Finally, turning to the $b\to s\bar{s}s$ transition, the relevant Lagrangian is

$$L_{eff} = \frac{\lambda'_{i32}\lambda'^{*}_{i22}}{m^2_{\tilde{\nu}_i}}\bar{s}\gamma_R s\bar{s}\gamma_L b + \frac{\lambda'_{i22}\lambda'^{*}_{i23}}{m^2_{\tilde{\nu}_i}}\bar{s}\gamma_L s\bar{s}\gamma_R b, \quad (19)$$

allowing the identification

$$f_{s,2}^{LR} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i32}\lambda'^{*}_{i22}}{4m^2_{\tilde{\nu}_i}}, \quad f_{s,2}^{RL} = \frac{\sqrt{2}}{G_F}\frac{\lambda'_{i22}\lambda'^{*}_{i23}}{4m^2_{\tilde{\nu}_i}}, \quad (20)$$

with

$$|f_{s,2}^{LR}| \leqslant 0.7, \quad |f_{s,2}^{RL}| \leqslant 0.8. \quad (21)$$

B. Z- and Z'-mediated FCNC's

In these models, one introduces an additional vector-singlet charge $-1/3$ quark h, as is found in E_6 grand unified theories, and allows it to mix with the ordinary down-type quarks d, s, and b. Since the weak isospin of the exotic quark is different from that of the ordinary quarks, flavor-changing neutral currents (FCNC's) involving the Z are induced [11]. The $Zb\bar{s}$ FCNC coupling, which is of interest to us here, is parametrized by the independent parameter U_{sb}^Z:

$$\mathcal{L}_{FCNC}^Z = -\frac{g}{2\cos\theta_W}U_{sb}^Z \bar{s}_L\gamma^\mu b_L Z_\mu . \quad (22)$$

Note that it is only the mixing between the left-handed components of the ordinary and exotic quarks which is responsible for the FCNC: since s_R, b_R and h_R all have the same $SU(2)_L\times U(1)_Y$ quantum numbers, their mixing cannot generate flavor-changing couplings of the Z. Models with Z-mediated FCNC's will therefore generate the $g_{q,2}^{LL}$ and $g_{q,2}^{LR}$ new-physics operators Eq. (2). These are the same operators that appear in the SM, so that this model does not generate new operators. (That is, these are effectively new contributions to the electroweak penguin operators of the SM.)

The strongest constraint on U_{sb}^Z comes from the measurement of $B(B\to\ell^+\ell^- X)$. The most recent result from BELLE gives [12]

$$B(B\to X_s e^+e^-) \leqslant 1.01\times 10^{-5}, \quad (23)$$

leading to the constraint

$$|U_{sb}^Z| \leqslant 7.6\times 10^{-4}. \quad (24)$$

With this constraint, it is straightforward to compute the maximal size of the couplings $g_{q,2}^{LL}$ and $g_{q,2}^{LR}$. We find

$$|g_{u,2}^{LL}| \leqslant 2.7\times 10^{-4}, \quad |g_{u,2}^{LR}| \leqslant 1.1\times 10^{-4},$$

$$|g_{d,2}^{LL}| = |g_{s,2}^{LL}| \leqslant 3.2\times 10^{-4},$$

$$|g_{d,2}^{LR}| = |g_{s,2}^{LR}| \leqslant 6.1\times 10^{-5}. \quad (25)$$

These couplings are therefore comparable in size to those of the SM.

Of course, since no new operators are generated in this scenario, and since the new-physics effects are about the same size as in the SM, one does not expect large deviations from the SM predictions due to Z-mediated FCNC's. However, models of new physics which contain exotic fermions also predict, in general, the existence of additional neutral Z' gauge bosons. If the s-, b- and h-quarks have different quantum numbers under the new $U(1)$ symmetry, their mixing will induce FCNC's due to Z' exchange [13].

In general, as was the case for Z-mediated FCNC's, such flavor-changing couplings will be constrained by the measurement of $B(B\to\ell^+\ell^- X)$. However, if the Z' is leptophobic, i.e. it does not couple to charged leptons, one can evade the constraints due to Eq. (23). Such models were considered in Ref. [14]. In this case, it is the mixing of the right-handed components of the ordinary and exotic quarks which is most important, and we parametrize the flavor-changing $Z'b\bar{s}$ coupling as

$$\mathcal{L}^{Z'}_{FCNC} = -\frac{g}{2\cos\theta_W} U^{Z'}_{sb} \bar{s}_R \gamma^\mu b_R Z'_\mu. \tag{26}$$

Thus, these models will generate new operators. In particular, the coefficients $g^{RL}_{q,2}$ and $g^{RR}_{q,2}$ will be nonzero.

Even though the Z' is leptophobic, there are constraints on $U^{Z'}_{sb}$ coming from the ALEPH limit $B(b\to s\nu\bar{\nu}) \le 6.4\times 10^{-4}$ [15]. In addition, in realistic models, leptophobia is realized only approximately—there will always be threshold effects which produce a small coupling of the Z' to charged leptons, in which case there are constraints from Eq. (23). The constraints from both of these sources turn out to be similar in size, and lead to [14]

$$|U^{Z'}_{sb}| \frac{M_Z^2}{M_{Z'}^2} \lesssim 6\times 10^{-3}. \tag{27}$$

With this constraint, one can estimate how large the new-physics coefficients can be. One finds

$$|g^{AB}_{q,2}| \le (1-2)\times 10^{-3},$$
$$AB = RR, RL, q = u,d,s, \tag{28}$$

which is about an order of magnitude larger than the coefficients of Eq. (25). Thus, Z'-mediated FCNC's can contribute significantly to charmless hadronic Λ_b decays, and can lead to significant deviations from the SM predictions for triple products in such processes.

III. TRIPLE PRODUCTS

In this section, we compute the contributions from the new-physics operators to triple-product correlations in Λ_b decays. In all cases, we retain only the leading term in the heavy-quark expansion, and neglect terms of order m/m_{Λ_b}, where m is the mass of the light meson. The main reason is that it is very unlikely that new physics will contribute at subleading order, but not at leading order. Indeed, as we will see, this situation can arise only in fine-tuned scenarios. A secondary reason is that the subleading terms are quite a bit smaller, e.g. $m_{K*}/m_{\Lambda_b} \sim 15\%$.

We begin with a review of the results of the SM.

A. SM results

In this subsection, we summarize the predictions of the SM for triple products in Λ_b decays. The discussion is somewhat cursory, and we refer to the reader to Ref. [5] for more details.

We first consider the decay $\Lambda_b \to F_1 P$, whose amplitude can be written generally as

$$\mathcal{M}_P = A(\Lambda_b \to F_1 P) = i\bar{u}_{F_1}(a + b\gamma_5) u_{\Lambda_b}. \tag{29}$$

The calculation of $|\mathcal{M}_P|^2$ yields a single triple-product term:

$$\mathrm{Im}(ab^*)\epsilon_{\mu\nu\rho\sigma} p^\mu_{F_1} s^\nu_{F_1} p^\rho_{\Lambda_b} s^\sigma_{\Lambda_b}, \tag{30}$$

where p_i^μ and s_i^μ are the 4-momentum and polarization of particle i. In the rest frame of the Λ_b, this takes the form $\vec{p}_{F_1}\cdot(\vec{s}_{F_1}\times\vec{s}_{\Lambda_b})$.

Within factorization, one can write

$$\begin{aligned} A(\Lambda_b\to F_1 P) &= \sum_{O,O'} \langle P|O|0\rangle\langle F_1|O'|\Lambda_b\rangle \\ &= i f_P q^\mu \langle F_1|\bar{q}_1\gamma_\mu(1-\gamma_5)b|\Lambda_b\rangle X_P \\ &\quad + i f_P q^\mu \langle F_1|\bar{q}_1\gamma_\mu(1+\gamma_5)b|\Lambda_b\rangle Y_P, \end{aligned} \tag{31}$$

where we have defined the pseudoscalar decay constant f_P as

$$i f_P q^\mu = \langle P|\bar{q}_2\gamma^\mu(1-\gamma_5)q_3|0\rangle, \tag{32}$$

where $q^\mu \equiv p^\mu_{\Lambda_b} - p^\mu_{F_1}$ is the four-momentum transfer. The key point here is that, in order to obtain a nonzero triple-product correlation, one must have two interfering amplitudes, i.e. X_P and Y_P must both be nonzero, and must have a relative weak phase. Furthermore, the triple product will be large only if X_P and Y_P are of similar size.

One can also show that the parameters a and b of Eq. (29) can be written as

$$a = f_P(X_P+Y_P)(m_{\Lambda_b}-m_{F_1}) f_1,$$
$$b = f_P(X_P-Y_P)(m_{\Lambda_b}+m_{F_1}) g_1, \tag{33}$$

where we have dropped terms of $O(m_P/m_{\Lambda_b})$, and f_1 and g_1 are Lorentz-invariant form factors:

$$\langle F_1|\bar{q}_1\gamma^\mu b|\Lambda_b\rangle = \bar{u}_{F_1}\left[f_1\gamma^\mu + i\frac{f_2}{m_{\Lambda_b}}\sigma^{\mu\nu}q_\nu + \frac{f_3}{m_{\Lambda_b}}q^\mu\right]u_{\Lambda_b},$$
$$\langle F_1|\bar{q}_1\gamma^\mu\gamma_5 b|\Lambda_b\rangle = \bar{u}_{F_1}\left[g_1\gamma^\mu + i\frac{g_2}{m_{\Lambda_b}}\sigma^{\mu\nu}q_\nu + \frac{g_3}{m_{\Lambda_b}}q^\mu\right]\gamma_5 u_{\Lambda_b}. \tag{34}$$

From Eqs. (30) and (33), we therefore see explicitly that the triple product in $\Lambda_b\to F_1 P$ is proportional to Im(ab^*) $\sim$Im$(X_P Y_P^*)$.

In the SM, there is only one class of decays which is expected to show a significant effect [5]: the triple-product correlation for $\Lambda_b\to pK^-$ is found to be $\sim 18\%$. For decays such as $\Lambda_b\to\Lambda\eta$, $\Lambda_b\to\Lambda\eta'$ and $\Lambda_b\to n\bar{K}^0$, the triple product is less than 1%. The fundamental reason for this is that $\Lambda_b\to pK^-$ is governed by the quark-level transition $b\to s\bar{u}u$, which has both a tree and a penguin contribution, whereas the other decays are dominated by the $b\to s$ penguin amplitude. Thus, for $\Lambda_b\to\Lambda\eta, \Lambda\eta', n\bar{K}^0$, there is essentially only a single decay amplitude, which precludes any CP- and T-violating effects.

In the above discussion, the size of the triple products has been estimated within factorization. However, it is well-known that nonfactorizable effects can be important in Λ_c decays. For example, the decay $\Lambda_c \to \Sigma^+ \phi$ has been observed [16], and this can only proceed via a (nonfactorizable) *W*-exchange diagram. This then begs the question of whether nonfactorizable effects might be important in Λ_b decays. In fact, the answer is that Λ_b decays are *not* expected to be significantly affected by such effects. In Ref. [17], it was found that the *W*-exchange contributions to inclusive Λ_b decays are suppressed relative to those in Λ_c decays by $O(m_c/m_b)^3$. This implies that even for exclusive decays such nonfactorizable *W*-exchange terms are expected to be small. This was confirmed in Ref. [5]: the *W*-exchange contributions to $\Lambda_b \to pK^-$ were estimated using a pole model, and the ratio of nonfactorizable to factorizable contributions was found to be tiny. For these reasons, here and below we ignore all nonfactorizable effects in Λ_b decays.

Turning to $\Lambda_b \to F_1 V$, the general decay amplitude can be written as [18]

$$\mathcal{M}_V = Amp(\Lambda_{F_1} \to BV) = \overline{u}_{F_1} \varepsilon_\mu^* [(p^\mu_{\Lambda_b} + p^\mu_{F_1})(a + b\gamma_5) + \gamma^\mu (x + y\gamma_5)] u_{\Lambda_b}, \tag{35}$$

where ε_μ^* is the polarization of the vector meson. In calculating $|\mathcal{M}_V|^2$, one finds many triple-product terms:

$$\begin{aligned} |\mathcal{M}_V|^2_{t.p.} = {} & 2\,\mathrm{Im}(ab^*) |\varepsilon_V \cdot (p_{\Lambda_b} + p_{F_1})|^2 \epsilon_{\mu\nu\rho\sigma} p^\mu_{F_1} s^\nu_{F_1} p^\rho_{\Lambda_b} s^\sigma_{\Lambda_b} + 2\,\mathrm{Im}(xy^*) \epsilon_{\alpha\beta\mu\nu} [\varepsilon_V \cdot s_{F_1} p^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V - \varepsilon_V \cdot p_{F_1} s^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V \\ & + \varepsilon_V \cdot s_{\Lambda_b} p^\alpha_{F_1} s^\beta_{F_1} \varepsilon^\mu_V p^\nu_{\Lambda_b} - \varepsilon_V \cdot p_{\Lambda_b} p^\alpha_{F_1} s^\beta_{F_1} \varepsilon^\mu_V s^\nu_{\Lambda_b}] + 2\varepsilon_V \cdot (p_{\Lambda_b} + p_{F_1}) \epsilon_{\alpha\beta\mu\nu} [\mathrm{Im}(ax^* + by^*) p^\alpha_{F_1} s^\beta_{F_1} p^\mu_{\Lambda_b} \varepsilon^\nu_V \\ & + m_{\Lambda_b} \mathrm{Im}(bx^* + ay^*) p^\alpha_{F_1} s^\beta_{F_1} s^\mu_{\Lambda_b} \varepsilon^\nu_V - \mathrm{Im}(ax^* - by^*) p^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V - m_{F_1} \mathrm{Im}(ay^* - bx^*) s^\alpha_{F_1} p^\beta_{\Lambda_b} s^\mu_{\Lambda_b} \varepsilon^\nu_V]. \end{aligned} \tag{36}$$

Similar to $\Lambda_b \to F_1 P$ decays, using factorization, one can write

$$\begin{aligned} A(\Lambda_b \to F_1 V) = m_V g_V \Bigg\{ & \varepsilon^*_\mu \langle F_1 | \overline{q}_1 \gamma^\mu (1-\gamma_5) b | \Lambda_b \rangle X_V + \varepsilon^*_\mu \langle F_1 | \overline{q}_1 \gamma^\mu (1+\gamma_5) b | \Lambda_b \rangle Y_V \\ & + \varepsilon \cdot (p_{\Lambda_b} + p_{F_1}) q_\mu \langle F_1 | \overline{q}_1 \gamma^\mu (1-\gamma_5) b | \Lambda_b \rangle \frac{A_V}{m^2_{\Lambda_b}} + \varepsilon \cdot (p_{\Lambda_b} + p_{F_1}) q_\mu \langle F_1 | \overline{q}_1 \gamma^\mu (1+\gamma_5) b | \Lambda_b \rangle \frac{B_V}{m^2_{\Lambda_b}} \Bigg\}, \end{aligned} \tag{37}$$

where the decay constant g_V has been defined as

$$m_V g_V \varepsilon^*_\mu = \langle V | \overline{q}_2 \gamma_\mu q_3 | 0 \rangle. \tag{38}$$

(Above we have included explicit factors of $m^2_{\Lambda_b}$ so that A_V and B_V have the same dimensions as X_V and Y_V. This differs from Ref. [5]. Also, note that, since the magnitudes of p_{Λ_b}, p_{F_1} and q are all of order m_{Λ_b}, the A_V and B_V operators are not *a priori* smaller than the X_V and Y_V operators.) Hence, using factorization, the quantities *a*, *b*, *x* and *y* of Eq. (36) can be expressed as

$$a^\lambda_V = m_V g_V \left[\frac{f_2}{m_{\Lambda_b}} (X^\lambda_V + Y^\lambda_V) + f_1 \frac{m_{\Lambda_b} - m_{F_1}}{m^2_{\Lambda_b}} (A^\lambda_V + B^\lambda_V) \right],$$

$$b^\lambda_V = m_V g_V \left[-\frac{g_2}{m_{\Lambda_b}} (X^\lambda_V - Y^\lambda_V) + g_1 \frac{m_{\Lambda_b} + m_{F_1}}{m^2_{\Lambda_b}} (A^\Lambda_V - B^\lambda_V) \right],$$

$$x^\lambda_V = m_V g_V \left[f_1 - \frac{m_{\Lambda_b} + m_{F_1}}{m_{\Lambda_b}} f_2 \right] [X^\lambda_V + Y^\lambda_V],$$

$$y^\lambda_V = -m_V g_V \left[g_1 + \frac{m_{\Lambda_b} - m_{F_1}}{m_{\Lambda_b}} g_2 \right] [X^\lambda_V - Y^\lambda_V], \tag{39}$$

where λ denotes the polarization of the final-state *V*, and we have again dropped subleading terms of $O(m_V/m_{\Lambda_b})$. If any two of the four terms in Eq. (37) above have a relative weak phase, their interference can lead to triple products.

In the SM, one finds that $A_V \simeq B_V \approx 0$, and that $Y_V \approx 0$ for a longitudinally polarized *V*. For a transversely polarized *V*, Y_V can be nonzero, but is still quite small. Thus, $\Lambda_b \to F_1 V$ decays are dominated by a single amplitude [the X_V term in Eq. (37) above], so that triple products in such decays are expected to be tiny. Specifically, one finds [5] that the triple-product asymmetry in $\Lambda_b \to pK^{*-}$ is $O(1\%)$ for a transversely polarized K^{*-}, while for a longitudinally polarized K^{*-} the asymmetry is $\ll 1\%$. For $\Lambda_b \to \Lambda \phi$ [19] and $\Lambda_b \to n\overline{K}^{*0}$, the asymmetries essentially vanish since these decays are dominated by a single weak decay amplitude (the $b \to s$ penguin).

B. $\Lambda_b \to F_1 P$: New-physics

We begin by considering the new-physics contributions to triple-product correlations in $\Lambda_b \to pK^-$ decays. Although this process is governed by the quark transition $b\to s\bar{u}u$, one still has to perform Fierz transformations on the operators in Eq. (2) to put them in a form appropriate for this decay. Using the relations

$$
\begin{aligned}
if_K q^\mu &= \langle K|\bar{s}\gamma^\mu(1-\gamma_5)u|0\rangle,\\
\langle K|\bar{s}(1\pm\gamma_5)u|0\rangle &= \mp\frac{if_K m_K^2}{m_s+m_u},\\
\langle p|\bar{u}(1\pm\gamma_5)b|\Lambda_b\rangle &= \frac{q^\mu}{m_b}\langle p|\bar{u}\gamma_\mu(1\mp\gamma_5)b|\Lambda_b\rangle,
\end{aligned}
\tag{40}
$$

we find that the new-physics contributions to X_K and Y_K [Eq. (31)] are

$$
\begin{aligned}
X_K^{NP} &= \frac{G_F}{\sqrt{2}}\left[\frac{1}{4}a_{u,1}^{RR}\chi_K + \frac{1}{2}a_{u,1}^{LR} + b_{u,1}^{LL} - b_{u,1}^{RL}\chi_K + 3c_{u,1}^{RR}\chi_K\right],\\
Y_K^{NP} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{4}a_{u,1}^{LL}\chi_K - \frac{1}{2}a_{u,1}^{RL} - b_{u,1}^{RR} + b_{u,1}^{LR}\chi_K - 3c_{u,1}^{LL}\chi_K\right],
\end{aligned}
\tag{41}
$$

where

$$
\chi_K \equiv \frac{2m_K^2}{(m_s+m_u)m_b}, \tag{42}
$$

and we have defined

$$
\begin{aligned}
a_{q,1}^{AB} &\equiv f_{q,1}^{AB} + \frac{1}{N_c}f_{q,2}^{AB},\\
b_{q,1}^{AB} &\equiv g_{q,1}^{AB} + \frac{1}{N_c}g_{q,2}^{AB},\\
c_{q,1}^{AB} &\equiv h_{q,1}^{AB} + \frac{1}{N_c}h_{q,2}^{AB}.
\end{aligned}
\tag{43}
$$

Note that we can obtain Y_K^{NP} from X_K^{NP}, up to an overall minus sign, simply by changing the chiralities $L\leftrightarrow R$.

As discussed in the previous subsection, within the SM the triple-product correlation for $\Lambda_b\to pK^-$ is expected to be large, $\sim 18\%$. However, since the new-physics operators of Eq. (2) can contribute to both X_K and Y_K, this prediction can easily be modified.

We now turn to the decay $\Lambda_b\to\Lambda\eta(\eta')$, which receives contributions from all three quark-level processes $b\to s\bar{q}q$, $q=u,d,s$. The calculation is similar to that above. For $\Lambda_b\to\Lambda\eta$ we find

$$
\begin{aligned}
X_\eta^{NP} &= \frac{G_F}{\sqrt{2}}[x_u + x_d + x_s + 3c_{s,1}^{RR}\chi_{\eta_s}],\\
x_{u(d)} &= r_{u(d)}\left[\frac{1}{2}(a_{u(d),2}^{RL} - a_{u(d),2}^{RR})\chi_{\eta_{u(d)}} + (b_{u(d),2}^{LL} - b_{u(d),2}^{LR})\right],\\
x_s &= r_s\left[\frac{1}{2}\left(a_{s,2}^{RL} - a_{s,2}^{RR} + \frac{1}{2}a_{s,1}^{RR} - 2b_{s,1}^{RL}\right)\chi_{\eta_s}\right.\\
&\qquad\left. + \left(b_{s,1}^{LL} + b_{s,2}^{LL} - b_{s,2}^{LR} + \frac{1}{2}a_{s,1}^{LR}\right)\right],\\
Y_\eta^{NP} &= \frac{G_F}{\sqrt{2}}[y_u + y_d + y_s - 3c_{s,1}^{LL}\chi_{\eta_s}],\\
y_{u(d)} &= r_{u(d)}\left[\frac{1}{2}(-a_{u(d),2}^{LR} + a_{u(d),2}^{LL})\chi_{\eta_{u(d)}}\right.\\
&\qquad\left. + (-b_{u(d),2}^{RR} + b_{u(d),2}^{RL})\right],\\
y_s &= r_s\left[\frac{1}{2}\left(-a_{s,2}^{LR} + a_{s,2}^{LL} - \frac{1}{2}a_{s,1}^{LL} + 2b_{s,1}^{LR}\right)\chi_{\eta_s}\right.\\
&\qquad\left. + \left(-b_{s,1}^{RR} - b_{s,2}^{RR} + b_{s,2}^{RL} - \frac{1}{2}a_{s,1}^{RL}\right)\right],
\end{aligned}
\tag{44}
$$

with

$$
\chi_{\eta_{u,d,s}} = \frac{m_\eta^2}{m_{u,d,s}m_b}, \tag{45}
$$

and

$$
a_{q,2}^{AB} \equiv f_{q,2}^{AB} + \frac{1}{N_c}f_{q,1}^{AB}, \quad b_{q,2}^{AB} \equiv g_{q,2}^{AB} + \frac{1}{N_c}g_{q,1}^{AB}. \tag{46}
$$

In the above, we have defined $r_{u,d,s}\equiv f_\eta^{u,d,s}/f_\pi$, where

$$
\begin{aligned}
if_\eta^u p_\eta^\mu &= \langle\eta|\bar{u}\gamma^\mu(1-\gamma_5)u|0\rangle\\
&= \langle\eta|\bar{d}\gamma^\mu(1-\gamma_5)d|0\rangle,\\
if_\eta^s p_\eta^\mu &= \langle\eta|\bar{s}\gamma^\mu(1-\gamma_5)s|0\rangle.
\end{aligned}
\tag{47}
$$

The amplitude for $\Lambda_b\to\Lambda\eta'$ has the same form as Eq. (44) with the replacement $\eta\to\eta'$.

Finally, we consider the decay $\Lambda_b\to n\bar{K}^0$, which is related by isospin to $\Lambda_b\to pK^-$. This is a pure penguin decay, with $b\to s\bar{d}d$. This decay will be very difficult to detect experimentally. Nevertheless, we include it here for completeness. We find

$$
\begin{aligned}
X_{\bar{K}}^{NP} &= \frac{G_F}{\sqrt{2}}\left[\frac{1}{4}a_{d,1}^{RR}\chi_{\bar{K}} + \frac{1}{2}a_{d,1}^{LR} + b_{d,1}^{LL} - b_{d,1}^{RL}\chi_{\bar{K}} + 3c_{d,1}^{RR}\chi_{\bar{K}}\right],\\
Y_{\bar{K}}^{NP} &= \frac{G_F}{\sqrt{2}}\left[-\frac{1}{4}a_{d,1}^{LL}\chi_{\bar{K}} - \frac{1}{2}a_{d,1}^{RL} - b_{d,1}^{RR} + b_{d,1}^{LR}\chi_{\bar{K}} - 3c_{d,1}^{LL}\chi_{\bar{K}}\right],
\end{aligned}
\tag{48}
$$

where

$$\chi_{\bar{K}} \equiv \frac{2m_K^2}{(m_s+m_d)m_b}. \tag{49}$$

For each of the decays $\Lambda_b \to \Lambda\eta$, $\Lambda_b \to \Lambda\eta'$ and $\Lambda_b \to n\bar{K}^0$, the triple product is tiny in the SM. This is due essentially to the fact that these decays are dominated by a single weak decay amplitude (the $b\to s$ penguin). However, this is no longer true in the presence of new physics; on the contrary, there may be several decay amplitudes. The new-physics operators may therefore lead to sizeable triple products in these decays.

C. $\Lambda_b \to F_1 V$: New-physics

We now examine the new-physics contributions to triple products in $\Lambda_b \to F_1 V$ decays. Before turning to specific decays, one can make some very general observations.

First, the amplitude for the production of a transversely polarized vector boson V is suppressed relative to that for a longitudinally polarized V by a factor m_V/E_V. Since $E_V \sim m_{\Lambda_b}/2$, this means that this production amplitude is subleading in the heavy-quark expansion, and can be neglected. In other words, in our analysis, we will assume the vector meson in the decay $\Lambda_b \to F_1 V$ to be essentially longitudinally polarized. As explained earlier, this is justified by the fact that it is very unlikely that the new physics will affect the production of a transversely polarized V without also affecting that of a longitudinally polarized V.

Second, in the rest frame of the Λ_b, we can write the 4-momentum of the final state vector meson as $q_\mu = (E_V, 0, 0, |\vec{p}_V|)$, so that the longitudinal polarization vector takes the form $\varepsilon_\mu^{\lambda=0} = (1/m_V)(|\vec{p}_V|, 0, 0, E_V)$. In the heavy-quark limit, $E_V \gg m_V$. Thus, in this limit, the longitudinal polarization vector can be written approximately as

$$\varepsilon_\mu^{\lambda=0} \simeq \frac{1}{m_V}\left(q_\mu + \frac{m_V^2}{2E_V} n_\mu\right), \tag{50}$$

with $n_\mu = (-1,0,0,1)$. In other words, to leading order in the heavy-quark expansion, $\varepsilon_\mu^{\lambda=0}$ is proportional to q_μ. This has two important consequences.

Consider first the A_V amplitude of Eq. (37), which is one of the four amplitudes describing $\Lambda_b \to F_1 V$ decays:

$$m_V g_V \varepsilon\cdot(p_{\Lambda_b}+p_{F_1}) q_\mu \langle F_1|\bar{q}_1\gamma^\mu(1-\gamma_5)b|\Lambda_b\rangle \frac{A_V}{m_{\Lambda_b}^2}. \tag{51}$$

Since $p_{F_1}^\mu = (E_{F_1}, 0, 0, -|\vec{p}|)$, one sees that $\varepsilon_V^* \cdot (p_{\Lambda_b}+p_{F_1})$ will be nonzero only for a longitudinally polarized V. Now, writing the quark content of the V as $\bar{q}_2 q_3$, the operators which correspond to the V take the form $\bar{q}_2(1\pm\gamma_5)q_3$, $\bar{q}_2\gamma^\mu(1\pm\gamma_5)q_3$ or $\bar{q}_2\sigma^{\mu\nu}(1\pm\gamma_5)q_3$. In calculating the V matrix elements, these yield

$$\langle V|\bar{q}_2(1\pm\gamma_5)q_3|0\rangle = 0,$$

$$\langle V|\bar{q}_2\gamma_\mu(1\pm\gamma_5)q_3|0\rangle = m_V g_V \varepsilon_\mu^*,$$

$$\langle V|\bar{q}_2\sigma_{\mu\nu}q_3|0\rangle = -i g_V^T[\varepsilon_\mu^* q_\nu - \varepsilon_\nu^* q_\mu]. \tag{52}$$

Thus, we see that it is only the tensor matrix element which could potentially contribute to A_V. However, to leading order in the heavy-quark expansion, $\varepsilon_\mu^{\lambda=0} \sim q_\mu$, so that the tensor matrix element vanishes. Thus, we have $A_V = O(m_V/m_{\Lambda_b})$, even in the presence of new-physics operators, and we neglect it. This argument applies also to the B_V amplitude of Eq. (37). [By comparison, X_V and Y_V are expected to be $O(1)$, i.e., leading order in the heavy-quark expansion.]

Neglecting the A_V and B_V terms, Eq. (39) reduces to

$$a_V^\lambda = m_V g_V \frac{f_2}{m_{\Lambda_b}}[X_V^\lambda + Y_V^\lambda],$$

$$b_V^\lambda = -m_V g_V \frac{g_2}{m_{\Lambda_b}}[X_V^\lambda - Y_V^\lambda],$$

$$x_V^\lambda = m_V g_V\left[f_1 - \frac{m_{F_1}+m_{\Lambda_b}}{m_{\Lambda_b}} f_2\right][X_V^\lambda + Y_V^\lambda],$$

$$y_V^\lambda = -m_V g_V\left[g_1 + \frac{m_{\Lambda_b}-m_{F_1}}{m_{\Lambda_b}} g_2\right][X_V^\lambda - Y_V^\lambda]. \tag{53}$$

Note that a_V^λ and x_V^λ now have the same weak phase, as do b_V^λ and y_V^λ.

Now consider again the triple-product terms of Eq. (36). As discussed above, to leading order in the heavy-quark expansion, only longitudinally polarized vector mesons need be considered, and $\varepsilon_\mu^{\lambda=0} \sim q_\mu$ in this limit. Thus, we see that triple products of the form $\epsilon_{\alpha\beta\mu\nu} p_{F_1}^\alpha p_{\Lambda_b}^\beta s_{\Lambda_b}^\mu \varepsilon_V^\nu$ are of subleading order, and we neglect them. In fact, to leading order, there is only a single triple product which remains:

$$|\mathcal{M}_V|^2_{t.p.} \simeq \frac{4}{m_V^2}\epsilon_{\alpha\beta\mu\nu} p_{\Lambda_b}^\alpha s_{\Lambda_b}^\beta q^\mu s_{F_1}^\nu \quad \{-2\,\mathrm{Im}(ab^*)|q\cdot p_{\Lambda_b}|^2 + \mathrm{Im}(xy^*) q\cdot p_{\Lambda_b} + q\cdot p_{\Lambda_b}[(m_{\Lambda_b}+m_{F_1})\mathrm{Im}(ay^*) + (m_{\Lambda_b}-m_{F_1})\mathrm{Im}(bx^*)]\}. \tag{54}$$

All other triple products are expected to be smaller, by a factor of order m_V/m_{Λ_b}.

We now turn to specific decays, and start with $\Lambda_b \to pK^{*-}$. First, for the tensor operators, one needs to evaluate matrix elements of the form

$$\langle V|\bar{q}_2\sigma_{\mu\nu}(1\pm\gamma_5)q_3|0\rangle\langle F|\bar{q}_3\sigma_{\mu\nu}(1\pm\gamma_5)b|\Lambda_b\rangle. \tag{55}$$

However, as we have argued above, the tensor matrix element vanishes to leading order in the heavy-quark expansion. Therefore the tensor operators will not contribute to this decay. The same does not hold true for the scalar/pseudoscalar and vector/axial vector new-physics operators, and we find

$$X_{K*}^{NP,\lambda}=\frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{u,1}^{LR}+b_{u,1}^{LL}\right],$$

$$Y_{K*}^{NP,\lambda}=\frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{u,1}^{RL}+b_{u,1}^{RR}\right]. \tag{56}$$

Note that, as expected, the new-physics operators contribute equally to longitudinally and transversely polarized V's. It is therefore reasonable to concentrate on the longitudinal V's, which dominate the decay $\Lambda_b\to pK^{*-}$.

The expressions for the decay $\Lambda_b\to n\overline{K^{0*}}$ can be easily obtained from those above by the replacement $u\to d$.

Finally, for $\Lambda_b\to\Lambda\phi$, we have

$$X_{\phi}^{NP,\lambda}=\frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{s,1}^{LR}+b_{s,1}^{LL}+b_{s,2}^{LL}+b_{s,2}^{LR}\right],$$

$$X_{\phi}^{SM,\lambda}=-\frac{G_F}{\sqrt{2}}V_{tb}V_{ts}^{*}\left[a_3^t+a_4^t+a_5^t-\frac{1}{2}a_7^t-\frac{1}{2}a_9^t-\frac{1}{2}a_{10}^t-a_3^c-a_4^c-a_5^c+\frac{1}{2}a_7^c+\frac{1}{2}a_9^c+\frac{1}{2}a_{10}^c\right],$$

$$Y_{\phi}^{NP,\lambda}=\frac{G_F}{\sqrt{2}}\left[-\frac{1}{2}a_{s,1}^{RL}+b_{s,1}^{RR}+b_{s,2}^{RR}+b_{s,2}^{RL}\right],$$

$$Y_{\phi}^{SM,\lambda}\simeq 0, \tag{57}$$

where we have included the standard model contribution without the tiny dipole contribution. The definitions of the various coefficients a_i^q, as well as their values, can be found in Ref. [5].

In all of the above decays, Y_V is expected to be very small in the SM, so that the triple products in $\Lambda_b\to F_1V$ are predicted to be at most $O(1\%)$. However, this can change significantly in the presence of new physics—from the above expressions one sees that the new-physics operators can easily produce a nonzero Y_V. The triple products in $\Lambda_b\to F_1V$ may well be sizeable in the presence of new physics.

IV. DIAGNOSTIC POWER

In the previous section, we saw that the presence of new-physics operators can significantly modify the SM predictions for triple-product correlations in $b\to s$ Λ_b decays. In particular, triple products which were expected to be tiny in the SM may now be sizeable. This is not at all surprising: most of those triple products are vanishingly small because the decays are dominated by a single weak $b\to s$ penguin decay amplitude. However, in the presence of new physics, one can have several decay amplitudes and, consequently, large triple-product asymmetries.

Although this particular result is entirely expected, the previous exercise is still useful for several reasons. First, the pattern of nonzero triple products provides information about the type of new-physics operators which may be present. And second, one can apply the above general analysis to specific models of new physics to obtain model-dependent predictions. These are the issues we explore in this section.

We begin with the model-independent analysis. The first observation is simple: if one sees no new effect in a particular decay, this implies that certain new-physics operators are absent (barring fine-tuned cancellations among these operators). For example, suppose that the triple-product asymmetry in $\Lambda_b\to pK^{*-}$ is found to be tiny, as in the SM. This means that $Y_{K*}^{NP}=0$ [Eq. (56)], so that $a_{u,1}^{RL}=b_{u,1}^{RR}=0$. (Note: since $Y_{K*}^{SM}\simeq 0$, X_{K*}^{NP} could still be nonzero, since the triple product is proportional to the product of these two quantities.) This in turn suggests that each of $f_{u,1}^{RL}$, $f_{u,2}^{RL}$, $g_{u,1}^{RR}$ and $g_{u,2}^{RR}$ vanish, since they make up $a_{u,1}^{RL}$ and $b_{u,1}^{RR}$. Similarly, should no new effects be seen in $\Lambda_b\to\Lambda\eta$, each of the 30 operators in Y_{η}^{NP} [Eq. (44)] must vanish.

Of course, one can obtain more information by combining measurements, since the same operators can contribute to more than one decay. In fact, one can even partially test the assumption that there are no fine-tuned cancellations. For example, suppose that the triple-product asymmetry in $\Lambda_b\to pK^-$ is found to agree with the SM, but that in $\Lambda_b\to pK^{*-}$ does not. The latter result implies that $a_{u,1}^{RL}$ and/or $b_{u,1}^{RR}$ are nonzero. However, these operators also contribute to Y_K^{NP} [Eq. (41)]. Thus, in order to obtain $Y_K^{NP}=0$, there must be cancellations among the various operators. Should such a result be found, it would be necessary to explain these cancellations, either via a symmetry, or by construction within a given model.

We now turn to the model-dependent analysis. The very general results of the previous section can be applied to specific models of new physics. Of course, in a given model, not all the operators of Eq. (2) will appear. In addition, it may be that the coefficients of those operators which do appear are related in some way. As examples of this behavior, we examine those models described in Sec. II, but this analysis can be applied to any models of new physics (e.g., supersymmetry, left-right symmetric models, etc.).

Consider first supersymmetric models with R-parity breaking (Sec. II A). If only B-violating couplings are present, then the only new-physics operators are vector operators contributing to $b\to s\overline{d}d$ [Eq. (8)]. This leads to a clear pattern of predictions: no new-physics effects are expected in the decays of a Λ_b to pK^-, pK^{*-} and $\Lambda\phi$. Indeed, if measurements of these triple-product asymmetries disagree with the SM predictions, this particular model is ruled out.

On the other hand, the decays $\Lambda_b\to\Lambda\eta$, $n\overline{K^0}$ and $n\overline{K}^{*0}$ can be affected in this model. How big can these effects be? In general, they can be enormous. As we have already noted, the new-physics contributions to these rare decays are still allowed by data to be comparable to, or even larger than, the SM contributions. If the two interfering amplitudes are of similar size, the triple-product asymmetry can be as large as $\sim 50\%$ (to be contrasted with the SM prediction of $\simeq 0$).

This also holds for the other models discussed below.

Turning to the *L*-violating couplings, one sees that more operators may be present [Eqs. (14), (17), and (20)]. In this case, all decays may be affected, except Λ_b to pK^-. This is a quite distinctive signature for this model.

Finally, we consider leptophobic Z'-mediated FCNC's (Sec. II B). There are only six nonzero new-physics coefficients, given in Eq. (28), and these all depend on the parameters $|U_{sb}^{Z'}|$ and $M_{Z'}$ [Eq. (27)]. In this case, all Λ_b decays will be affected. However, note that, within this model, there are more observables (6) than there are theoretical parameters (2). This means that if deviations from the SM predictions are measured, we will be able to get a handle on $|U_{sb}^{Z'}|$ and $M_{Z'}$. Conversely, if no new-physics effects are observed, we will be able to place strong constraints on these quantities.

V. CONCLUSIONS

In the standard model (SM), (almost) all *T*-violating triple-product correlations in charmless Λ_b decays are expected to be tiny. (The one exception is the decay $\Lambda_b \to pK^-$, for which the asymmetry is 18%.) This is therefore a good place to look for physics beyond the standard model.

In this paper, using an effective-Lagrangian approach, we have computed the effects of new physics on such triple products. This approach has the advantage of indicating which specific new-physics operators affect each of the Λ_b triple-product correlations. Thus, the measurement of a number of different triple products permits us to determine which new-physics operators are or are not present. Furthermore, the approach is completely general—the effects of any specific model can be obtained by simply calculating which operators appear in that model.

The new-physics effects on triple products are calculated using factorization. In addition, we work only to leading order in the heavy-quark expansion, neglecting terms of order m/m_{Λ_b}, where m is the mass of the light final-state meson. The justification for this is that it is only in fine-tuned scenarios that the new physics contributes at subleading order, but not at leading order. (Also, the subleading terms are quite a bit smaller, e.g., $m_{K*}/m_{\Lambda_b} \sim 15\%$.)

We have found that all Λ_b triple products can be significantly modified by new physics. Of course, this to be expected. Most of the triple products are vanishingly small in the SM because the decays are dominated by a single weak $b \to s$ penguin decay amplitude. Thus, in the presence of new physics, there may be several decay amplitudes which can interfere with the SM amplitude. However, in order to obtain a sizeable asymmetry, the interfering amplitudes must be of similar size. We note that the constraints on the new-physics operators are sufficiently weak that they can be comparable to, or even larger than, the SM contributions. Thus, triple products which vanish in the SM can be as large as $\sim 50\%$ with new physics.

We have demonstrated how the measurement of triple-product asymmetries provides diagnostic information about the new-physics operators present. For example, all operators which affect $\Lambda_b \to pK^{*-}$ also affect $\Lambda_b \to pK^-$, but not vice versa. Thus, if the triple product in $\Lambda_b \to pK^-$ is found to agree with the SM, we would also expect no new effects in $\Lambda_b \to pK^{*-}$. If this were found not to hold, then we would conclude that there must be cancellations among the operators in $\Lambda_b \to pK^-$, and this would have to be explained in some way (e.g., symmetry, specific model, etc.).

Finally, we have also applied this general approach to two specific models: supersymmetry with *R*-parity breaking, and leptophobic Z'-mediated flavor-changing neutral currents. In both cases, we have worked out the new-physics operators which appear in those models, and used the previous formalism to calculate which Λ_b triple products can be affected. For example, in the case of *R*-parity breaking models, there is a clear pattern of effects. One such model predicts significant new effects in the decays $\Lambda_b \to \Lambda \eta$, $n\bar{K}^0$ and $n\bar{K}^{*0}$, but not in Λ_b to pK^-, pK^{*-} and $\Lambda \phi$. Any deviation from this pattern would rule out this model. Other models of new physics can be treated similarly.

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