Direct *CP* violation in Λ_b decays

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We study direct CP-violating asymmetries (CPAs) in the two-body Λ_b decays of $\Lambda_b \to pM(V)$ with $M(V) = K^-(K^{*-})$ and $\pi^-(\rho^-)$ based on the generalized factorization method in the standard model (SM). After simultaneously explaining the observed decay branching ratios of $\Lambda_b \to (pK^-, p\pi^-)$ with $\mathcal{R}_{\pi K} \equiv \mathcal{B}(\Lambda_b \to p\pi^-)/\mathcal{B}(\Lambda_b \to pK^-)$ being 0.84 ± 0.09, we find that their corresponding direct CPAs are $(5.8 \pm 0.2, -3.9 \pm 0.2)\%$, in comparison with $(-5^{+26}_{-5}, -31^{+43}_{-1})\%$ based on the perturbative QCD calculation and $(-10 \pm 8 \pm 4, 6 \pm 7 \pm 3)\%$ from the CDF experiment, respectively. For $\Lambda_b \to (pK^{*-}, p\rho^-)$, the decay branching ratios and CPVs are predicted to be $(2.5 \pm 0.5, 11.4 \pm 2.1) \times 10^{-6}$ with $\mathcal{R}_{\rho K^*} = 4.6 \pm 0.5$ and $(19.6 \pm 1.6, -3.7 \pm 0.3)\%$, respectively. The uncertainties for the CPAs in these decay modes as well as $\mathcal{R}_{\pi K, \rho K^{*-}}$ mainly arise from the quark mixing elements and nonfactorizable effects, whereas those from the hadronic matrix elements are either totally eliminated or small. We point out that the large CPA for $\Lambda_b \to pK^{*-}$ is promising to be measured by the CDF and LHCb experiments, which is a clean test of the SM.

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I. INTRODUCTION

It is known that one of the main goals in the B meson factories is to confirm the weak CP phase in the standard model (SM) through CP violating effects. However, the direct CP violating asymmetries (CPAs), $\mathcal{A}_{CP}$, in B decays have not been clearly understood yet. In particular, the naive result of $\mathcal{A}_{CP}(\bar{B}^0 \to K^-\pi^+) \simeq \mathcal{A}_{CP}(B^- \to K^-\pi^0)$ in the SM cannot be approved by the experiments [1]. It is known that it is inadequate to calculate the direct CPAs in the two-body mesonic B decays due to the limited knowledge on strong phases [2]. Clearly, one should look for CPV effects in other processes, in which the hadronic effects are well understood.

Unlike the two-body B meson decays, due to the flavor conservation, there is neither color-suppressed nor annihilation contribution in the two-body baryonic modes of $\Lambda_b \to pK^-$ and $\Lambda_b \to p\pi^-$, providing the controllable nonfactorizable effects and traceable strong phases for the CPAs. In fact, their decay branching ratios have been recently observed, given by [3]

$$\mathcal{B}(\Lambda_b \to pK^-) = (4.9 \pm 0.9) \times 10^{-6},$$
$$\mathcal{B}(\Lambda_b \to p\pi^-) = (4.1 \pm 0.8) \times 10^{-6}. \quad (1)$$

Although the two decays have been extensively discussed in the literature [4–6], the measured values in Eq. (1) cannot be simultaneously explained in the studies.

In this paper, we will first examine these two-body baryonic decays based on the configuration of the $\Lambda_b \to p$ transition with a recoiled K or π, and then calculate $\mathcal{A}_{CP}(\Lambda_b \to pK^-, p\pi^-)$, which have been measured by the CDF collaboration [7]. We will also extend our study to the corresponding vector modes of $\Lambda_b \to pV$ with $V = K^{*-}(\rho^-)$ as well as other two-body beauty baryons ($\mathcal{B}_b$) decays, such as Ξ_b.

II. FORMALISM

According to the decaying processes depicted in Fig. 1, in the generalized factorization approach [8] the amplitudes of $\Lambda_b \to pM(V)$ with $M(V) = K^-(K^{*-})$ and $\pi^-(\rho^-)$ can be derived as

$$\mathcal{A}(\Lambda_b \to pM) = i\frac{G_F}{\sqrt{2}} m_b f_M [\alpha_M \langle p|\bar{u}b|\Lambda_b\rangle + \beta_M \langle p|\bar{u}\gamma_5 b|\Lambda_b\rangle],$$
$$\mathcal{A}(\Lambda_b \to pV) = \frac{G_F}{\sqrt{2}} m_V f_V \varepsilon^{\mu*} \alpha_V \langle p|\bar{u}\gamma_\mu(1-\gamma_5)b|\Lambda_b\rangle, \quad (2)$$

where G_F is the Fermi constant and the meson decay constants $f_{M(V)}$ are defined by $\langle M|\bar{q}_1\gamma_\mu\gamma_5 q_2|0\rangle = -if_M q_\mu$ and $\langle V|\bar{q}_1\gamma_\mu q_2|0\rangle = m_V f_V \varepsilon^*_\mu$ with the four-momentum q_μ and polarization ε^*_μ, respectively. The constants α_M (β_M) and α_V in Eq. (2) are related to the (pseudo)scalar and vector or axial-vector quark currents, given by

$$\alpha_M(\beta_M) = V_{ub}V^*_{uq}a_1 - V_{tb}V^*_{tq}(a_4 \pm r_M a_6),$$
$$\alpha_V = V_{ub}V^*_{uq}a_1 - V_{tb}V^*_{tq}a_4, \quad (3)$$

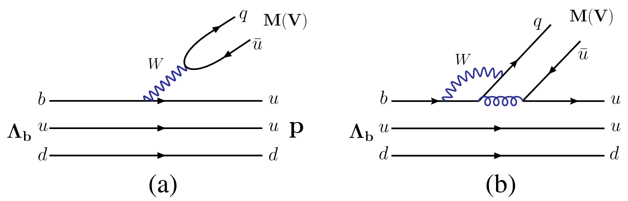

FIG. 1 (color online). Contributions to $\Lambda_b \to pM(V)$ from (a) color-allowed tree-level and (b) penguin diagrams.

where $r_M \equiv 2m_M^2/[m_b(m_q + m_u)]$, V_{ij} are the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, $q = s$ or d, and $a_i \equiv c_i^{\rm eff} + c_{i\pm1}^{\rm eff}/N_c^{\rm (eff)}$ for $i =$ odd (even) are composed of the effective Wilson coefficients $c_i^{\rm eff}$ defined in Ref. [8]. We note that, as seen from Fig. 1, there is no annihilation diagram at the penguin level for $\Lambda_b \to pM(V)$, unlike the cases in the two-body mesonic B decays. In addition, without the color-suppressed tree-level diagram, the nonfactorizable effects in these baryonic decays can be modest. In order to take account of the nonfactorizable effects, we use the generalized factorization method by setting the color number as $N_c^{\rm eff}$, which floats from 2 to ∞. The matrix elements of the $\mathcal{B}_b \to \mathcal{B}$ baryon transition in Eq. (2) have the general forms:

$$\begin{aligned}
\langle\mathcal{B}|\bar{q}\gamma_\mu b|\mathcal{B}_b\rangle &= \bar{u}_\mathcal{B}\left[f_1\gamma_\mu + \frac{f_2}{m_{\mathcal{B}_b}}i\sigma_{\mu\nu}q^\nu + \frac{f_3}{m_{\mathcal{B}_b}}q_\mu\right]u_{\mathcal{B}_b},\\
\langle\mathcal{B}|\bar{q}\gamma_\mu\gamma_5 b|\mathcal{B}_b\rangle &= \bar{u}_\mathcal{B}\left[g_1\gamma_\mu + \frac{g_2}{m_{\mathcal{B}_b}}i\sigma_{\mu\nu}q^\nu + \frac{g_3}{m_{\mathcal{B}_b}}q_\mu\right]\gamma_5 u_{\mathcal{B}_b},\\
\langle\mathcal{B}|\bar{q}b|\mathcal{B}_b\rangle &= f_S\bar{u}_\mathcal{B}u_{\mathcal{B}_b}, \langle\mathcal{B}|\bar{q}\gamma_5 b|\mathcal{B}_b\rangle = f_P\bar{u}_\mathcal{B}\gamma_5 u_{\mathcal{B}_b}, \quad (4)
\end{aligned}$$

where f_j (g_j) ($j = 1, 2, 3, S$ and P) are the form factors. For the $\Lambda_b \to p$ transition, f_j and g_j from different currents can be related by the SU(3) flavor and SU(2) spin symmetries [9,10], giving rise to $f_1 = g_1$ and $f_{2,3} = g_{2,3} = 0$. These relations are also in accordance with the derivations from the heavy-quark and large-energy symmetries in Ref. [11]. Note that the helicity-flip terms of $f_{2,3}$ and $g_{2,3}$ vanish due to the symmetries. Moreover, as shown in Refs. [6,11,12], $f_{2,3}$ ($g_{2,3}$) can only be contributed from the loops, resulting in that they are smaller than $f_1(g_1)$ by one order of magnitude, and can be safely ignored. By the equation of motion, we get $f_S = [(m_{\mathcal{B}_b} - m_\mathcal{B})/(m_b - m_q)]f_1$ and $f_P = [(m_{\mathcal{B}_b} + m_\mathcal{B})/(m_b + m_q)]g_1$. In the double-pole momentum dependences, f_1 and g_1 are in the forms of

$$f_1(q^2) = \frac{C_F}{(1 - q^2/m_{\mathcal{B}_b}^2)^2}, \quad g_1(q^2) = \frac{C_F}{(1 - q^2/m_{\mathcal{B}_b}^2)^2}, \quad (5)$$

with $C_F \equiv f_1(0) = g_1(0)$. To calculate the branching ratio of $\Lambda_b \to pM$ or pV, we take the averaged decay width $\Gamma \equiv (\Gamma_{M(V)} + \Gamma_{\bar{M}(\bar{V})})/2$ with $\Gamma_{M(V)}$ ($\Gamma_{\bar{M}(\bar{V})}$) for $\Lambda_b \to pM(V)$ ($\bar{\Lambda}_b \to \bar{p}\bar{M}(\bar{V})$). From Eq. (2) and Eq. (3), we can derive the ratios

$$\begin{aligned}
\mathcal{R}_{\pi K} &\equiv \frac{\mathcal{B}(\Lambda_b \to p\pi^-)}{\mathcal{B}(\Lambda_b \to pK^-)} = \frac{f_\pi^2}{f_K^2}\frac{|\alpha_\pi|^2 + |\alpha_{\bar{\pi}}|^2}{|\alpha_K|^2 + |\alpha_{\bar{K}}|^2}\frac{1 + \xi_\pi^+}{1 + \xi_K^+},\\
\mathcal{R}_{\rho K^*} &\equiv \frac{\mathcal{B}(\Lambda_b \to p\rho^-)}{\mathcal{B}(\Lambda_b \to pK^{*-})} = \frac{f_\rho^2}{f_{K^*}^2}\frac{|\alpha_\rho|^2 + |\alpha_{\bar{\rho}}|^2}{|\alpha_{K^*}|^2 + |\alpha_{\bar{K}^*}|^2}, \quad (6)
\end{aligned}$$

where ξ_M^+ ($M = \pi, K$) are defined by

$$\xi_M^\pm \equiv \left(\frac{|\beta_M|^2 \pm |\beta_{\bar{M}}|^2}{|\alpha_M|^2 + |\alpha_{\bar{M}}|^2}\right)R_{\Lambda_b\to p}, \quad (7)$$

with $R_{\Lambda_b\to p} = |\langle p|\bar{u}\gamma_5 b|\Lambda_b\rangle|^2/|\langle p|\bar{u}b|\Lambda_b\rangle|^2$, representing the uncertainty from the hadronization. The direct CP asymmetry is defined by

$$\mathcal{A}_{CP} = \frac{\Gamma_{M(V)} - \Gamma_{\bar{M}(\bar{V})}}{\Gamma_{M(V)} + \Gamma_{\bar{M}(\bar{V})}}. \quad (8)$$

Explicitly, from Eqs. (2), (3), and (8), we obtain

$$\begin{aligned}
\mathcal{A}_{CP}(\Lambda_b \to pM) &= \left(\frac{|\alpha_M|^2 - |\alpha_{\bar{M}}|^2}{|\alpha_M|^2 + |\alpha_{\bar{M}}|^2} + \xi_M^-\right)\frac{1}{1 + \xi_M^+},\\
\mathcal{A}_{CP}(\Lambda_b \to pV) &= \frac{|\alpha_V|^2 - |\alpha_{\bar{V}}|^2}{|\alpha_V|^2 + |\alpha_{\bar{V}}|^2}. \quad (9)
\end{aligned}$$

It is interesting to point out that for $\mathcal{R}_{\rho K^*}$ in Eq. (6), there is no uncertainty from the $\Lambda_b \to p$ transition, while both mesonic and baryonic uncertainties are totally eliminated for $\mathcal{A}_{CP}(\Lambda_b \to pV)$ in Eq. (9). Even for $\mathcal{R}_{\pi K}$ and $\mathcal{A}_{CP}(\Lambda_b \to pM)$, we will demonstrate later that the hadron uncertainties can be limited.

III. NUMERICAL RESULTS AND DISCUSSIONS

For the numerical analysis, the theoretical inputs of the meson decay constants and the Wolfenstein parameters for the CKM matrix are taken as [3]

$$\begin{aligned}
(f_\pi, f_K, f_\rho, f_{K^*}) &= (130.4 \pm 0.2, 156.2 \pm 0.7,\\
&\quad 210.6 \pm 0.4, 204.7 \pm 6.1)\ \text{MeV},\\
(\lambda, A, \rho, \eta) &= (0.225, 0.814, 0.120 \pm 0.022,\\
&\quad 0.362 \pm 0.013). \quad (10)
\end{aligned}$$

We note that f_{ρ,K^*} are extracted from the τ decays of $\tau^- \to (\rho^-, K^{*-})\nu_\tau$, and $V_{ub} = A\lambda^3(\rho - i\eta)$ and $V_{td} = A\lambda^3(1 - i\eta - \rho)$ are used to provide the weak phase for CP violation, while the strong phases are coming from the effective Wilson coefficients $c_i^{\rm eff}$ $(i = 1, 2, 3, \ldots, 6)$. Explicitly, at the m_b scale, one has that [8]

$$
\begin{aligned}
c_1^{\rm eff} &= 1.168, \quad c_2^{\rm eff} = -0.365,\\
10^4\epsilon_1 c_3^{\rm eff} &= 64.7 + 182.3\epsilon_1 \mp 20.2\eta - 92.6\rho + 27.9\epsilon_2\\
&\quad + i(44.2 - 16.2\epsilon_1 \mp 36.8\eta - 108.6\rho + 64.4\epsilon_2),\\
10^4\epsilon_1 c_4^{\rm eff} &= -194.1 - 329.8\epsilon_1 \pm 60.7\eta + 277.8\rho - 83.7\epsilon_2\\
&\quad + i(-132.6 + 48.5\epsilon_1 \pm 110.4\eta\\
&\quad + 325.9\rho - 193.3\epsilon_2),\\
10^4\epsilon_1 c_5^{\rm eff} &= 64.7 + 89.8\epsilon_1 \mp 20.2\eta - 92.6\rho + 27.9\epsilon_2\\
&\quad + i(44.2 - 16.2\epsilon_1 \mp 36.8\eta - 108.6\rho + 64.4\epsilon_2),\\
10^4\epsilon_1 c_6^{\rm eff} &= -194.1 - 466.7\epsilon_1 \pm 60.7\eta + 277.8\rho - 83.7\epsilon_2\\
&\quad + i(-132.6 + 48.5\epsilon_1 \pm 110.4\eta\\
&\quad + 325.9\rho - 193.3\epsilon_2), \qquad (11)
\end{aligned}
$$

for the $b \to d$ ($\bar{b} \to \bar{d}$) transition, and

$$
\begin{aligned}
c_1^{\rm eff} &= 1.168, \qquad c_2^{\rm eff} = -0.365,\\
10^4 c_3^{\rm eff} &= 241.9 \pm 3.2\eta + 1.4\rho + i(31.3 \mp 1.4\eta + 3.2\rho),\\
10^4 c_4^{\rm eff} &= -508.7 \mp 9.6\eta - 4.2\rho + i(-93.9 \pm 4.2\eta - 9.6\rho),\\
10^4 c_5^{\rm eff} &= 149.4 \pm 3.2\eta + 1.4\rho + i(31.3 \mp 1.4\eta + 3.2\rho),\\
10^4 c_6^{\rm eff} &= -645.5 \mp 9.6\eta - 4.2\rho\\
&\quad + i(-93.9 \pm 4.2\eta - 9.6\rho), \qquad (12)
\end{aligned}
$$

for the $b \to s$ ($\bar{b} \to \bar{s}$) transition, where $\epsilon_1 = (1-\rho)^2 + \eta^2$ and $\epsilon_2 = \rho^2 + \eta^2$. By adopting $C_F = 0.14 \pm 0.03$ from the light-cone sum rules in Ref. [11], with the central value in agreement with those in Refs. [6,12], we find that $\mathcal{B}(\Lambda_b \to pK^-) = (5.1^{+2.4}_{-2.0}) \times 10^{-6}$ and $\mathcal{B}(\Lambda_b \to p\pi^-) = (4.4^{+2.1}_{-1.7}) \times 10^{-6}$, which are consistent with the data in Eq. (1). This is regarded to have a modest nonfactorizable effect, as investigated by the study of $\Lambda_b \to p\pi^-$ in Ref. [11]. Nonetheless, since the uncertainties from the predictions exceed those of the data, we fit C_F with the data in Eq. (1), and obtain

$$C_F = 0.136 \pm 0.009, \qquad (13)$$

which is able to reconcile the theoretical studies of C_F to the data, and to be used in our study. Theoretical inputs in the SM for $R_{\Lambda_b \to p}$ and $\xi_M^\pm$ in Eq. (7) can be evaluated, given by

$$
\begin{aligned}
R_{\Lambda_b \to p} &= 1.008,\\
(\xi_\pi^+, \xi_K^+) &= (1.03 \pm 0.04 \pm 0.00, 0.11 \pm 0.01 \pm 0.02),\\
(\xi_\pi^-, \xi_K^-) &= (-4.0 \pm 0.3 \pm 0.0, -4.0 \pm 0.2 \pm 0.3) \times 10^{-3}, \qquad (14)
\end{aligned}
$$

where the errors for $\xi_M^\pm$ come from the CKM matrix elements and the floating $N_c^{\rm eff}$, respectively.

TABLE I. Ratios of $\mathcal{R}_{\pi K}$ and $\mathcal{R}_{\rho K^*}$ from our calculations, the pQCD and experiments, where the errors of our results are from the CKM matrix elements and nonfactorizable effects, respectively.

	$\mathcal{R}_{\pi K}$	$\mathcal{R}_{\rho K^*}$
Our result	$0.84 \pm 0.09 \pm 0.00$	$4.6 \pm 0.5 \pm 0.1$
pQCD [4]	$2.6^{+2.0}_{-0.5}$	...
CDF [14]	$0.66 \pm 0.14 \pm 0.08$	...
LHCb [15]	$0.86 \pm 0.08 \pm 0.05$	...

In Refs. [7,13], it is pointed out that the ratio of $\mathcal{R}_{\pi K}$ observed by CDF [14] or LHCb [15] has not been realized theoretically, as shown in Table I. In particular, we note that $\mathcal{R}_{\pi K} = 2.6^{+2.0}_{-0.5}$ in the pQCD prediction [4] is about 3 times larger than the data, but better than other QCD calculations, such as $\mathcal{R}_{\pi K} = 10.7$ in Ref. [6]. However, in Table I our result of this study shows that $\mathcal{R}_{\pi K} = 0.84 \pm 0.09$, which agrees well with the combined experimental value of 0.84 ± 0.22 by CDF and LHCb. Clearly, our result justifies the theoretical approach based on the factorization in the two-body Λ_b decays. We emphasize that the ratio of $\mathcal{R}_{\rho K^*}$ for the vector meson modes, which is predicted to be around 4.6, is an interesting physical observable as it is free of the hadronic uncertainties from the baryon sectors. A measurement for this ratio will be a firm test of the factorization approach in these baryonic decays.

In Table II, we present the branching ratios and direct CP asymmetries of $\Lambda_b \to pM(V)$ with $M(V) = K^-(K^{*-})$ and $\pi^-(\rho^-)$. For the decays of $\Lambda_b \to (pK^{*-}, p\rho^-)$, the predictions of the branching ratios in Table II are accessible to the experiments by CDF and LHCb. Note that our results of $\mathcal{B}(\Lambda_b \to pK^{*-}, p\rho^-) \simeq (2.5, 11.4) \times 10^{-6}$ in Table II are larger than those of $(0.3, 6.1) \times 10^{-6}$ [6] and $(0.8, 1.9) \times 10^{-6}$ [16] in other theoretical calculations.

As shown in Table II, for the first time, the theoretical values of $\mathcal{B}(\Lambda_b \to pK^-)$ and $\mathcal{B}(\Lambda_b \to p\pi^-)$ are found to be simultaneously in agreement with the data. Moreover, we demonstrate that the uncertainties from the form factors, the CKM matrix elements, and the nonfactorizable effects are small and well controlled.

For CP violation, from Eqs. (9) and (14), one can use the reduced forms of $\mathcal{A}_{CP}(\Lambda_b \to pM) \propto (|\alpha_M|^2 - |\alpha_{\bar{M}}|^2)/(|\alpha_M|^2 + |\alpha_{\bar{M}}|^2)$ similar to $\mathcal{A}_{CP}(\Lambda_b \to pV)$, which indeed present the limited hadron uncertainties, except for the factor of $1/2$ for $\mathcal{A}_{CP}(\Lambda_b \to p\pi^-)$. As shown in Table II, our predictions of $\mathcal{A}_{CP}(\Lambda_b \to p\pi, pK^-)$ are around $(-3.9, 5.8)\%$ with the errors less than 0.2%, while the results from the data [7] as well as the pQCD calculations are given to be consistent with zero.

For the vector modes, as the uncertainties from the hadronizations have been totally eliminated in Eq. (9), we are able to obtain reliable theoretical predictions for $\mathcal{A}_{CP}$, which should be helpful for experimental searches. In particular, it is worthwhile to note that $\mathcal{A}_{CP}(\Lambda_b \to pK^{*-}) = (19.6 \pm 1.6)\%$ is another example of the large and clean CP

TABLE II. Decay branching ratios and direct CP asymmetries of $\Lambda_b \to pM(V)$, where the errors for $\mathcal{B}(\Lambda_b \to pM(V))$ arise from $f_{M(V)}$ and $f_1(g_1)$, the CKM matrix elements and nonfactorizable effects, while those for $\mathcal{A}_{CP}(\Lambda_b \to pM(V))$ are from the CKM matrix elements and nonfactorizable effects, respectively.

	Our result	pQCD [4]	Data
$10^6\mathcal{B}(\Lambda_b \to pK^-)$	$4.8 \pm 0.7 \pm 0.1 \pm 0.3$	$2.0^{+1.0}_{-1.3}$	4.9 ± 0.9 [3]
$10^6\mathcal{B}(\Lambda_b \to p\pi^-)$	$4.2 \pm 0.6 \pm 0.4 \pm 0.2$	$5.2^{+2.5}_{-1.9}$	4.1 ± 0.8 [3]
$10^6\mathcal{B}(\Lambda_b \to pK^{*-})$	$2.5 \pm 0.3 \pm 0.2 \pm 0.3$	...	...
$10^6\mathcal{B}(\Lambda_b \to p\rho^-)$	$11.4 \pm 1.6 \pm 1.2 \pm 0.6$	...	...
$10^2\mathcal{A}_{CP}(\Lambda_b \to pK^-)$	$5.8 \pm 0.2 \pm 0.1$	-5^{+26}_{-5}	$-10 \pm 8 \pm 4$ [7]
$10^2\mathcal{A}_{CP}(\Lambda_b \to p\pi^-)$	$-3.9 \pm 0.2 \pm 0.0$	-31^{+43}_{-1}	$6 \pm 7 \pm 3$ [7]
$10^2\mathcal{A}_{CP}(\Lambda_b \to pK^{*-})$	$19.6 \pm 1.3 \pm 1.0$	...	...
$10^2\mathcal{A}_{CP}(\Lambda_b \to p\rho^-)$	$-3.7 \pm 0.3 \pm 0.0$	...	...

violating effects without hadronic uncertainties as the process in the baryonic B decays of $B^\pm \to K^{*\pm}\bar{p}p$ [17].

Interestingly, one would ask why the CP symmetry in $\Lambda_b \to pK^{*-}$ is larger than those in the other baryonic decay modes. The reason is that the term related to a_4 from the penguin diagram in Eq. (3) can be the primary contribution to $\Lambda_b \to pK^{*-}$ in Eq. (2), while allowing the certain contribution to the a_1 term from the tree diagram, such that the apparent large interference is able to take place. In contrast, in $\Lambda_b \to p\pi^-(\rho^-)$ and $\Lambda_b \to pK^-$, the a_1 and $(a_4 + r_M a_6)$ terms are dominating the branching ratios, respectively, leaving less room for the interferences. Clearly, $\mathcal{A}_{CP}(\Lambda_b \to pK^{*-})$ as well as the CPAs in other modes should receive more attention, which have also been emphasized in Ref. [18]. Finally, we remark that our approach can be extended to the two-body decay modes of other beauty baryons ($\mathcal{B}_b$), such as Ξ_b. For example, the branching ratio and CPA of $\Xi_b \to \Sigma^+K^{*-}$ can be estimated to be around 2.8×10^{-6} and 20%, respectively.

IV. CONCLUSIONS

Based on the generalized factorization method and $SU(3)$ flavor and $SU(2)$ spin symmetries, we have simultaneously explained the recent observed decay branching ratios in $\Lambda_b \to pK^-$ and $\Lambda_b \to p\pi^-$ and obtained the ratio of $\mathcal{R}_{\pi K}$ being 0.84 ± 0.09, demonstrating a reliable theoretical approach to study the two-body Λ_b decays. We have also predicted that $\mathcal{A}_{CP}(\Lambda_b \to pK^-) = (5.8 \pm 0.2)\%$ and $\mathcal{A}_{CP}(\Lambda_b \to p\pi^-) = (-3.9 \pm 0.2)\%$ with well-controlled uncertainties, whereas the current data for these CPAs are consistent with zero. We have used this approach to study the corresponding vector modes. Explicitly, we have found that $\mathcal{B}(\Lambda_b \to pK^{*-}, p\rho^-) = (2.5 \pm 0.5, 11.4 \pm 2.1) \times 10^{-6}$ with $\mathcal{R}_{\rho K^*} = 4.6 \pm 0.5$ and $\mathcal{A}_{CP}(\Lambda_b \to pK^{*-}, p\rho^-) = (19.6 \pm 1.6, -3.7 \pm 0.3)\%$. Since our prediction for $\mathcal{A}_{CP}(\Lambda_b \to pK^{*-})$ is large and free of both mesonic and baryonic uncertainties from the hadron sector, it would be the most promised direct CP asymmetry to be measured by the experiments at the CDF and LHCb to test the SM.

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