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Physics Letters B

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Pentaquarks from intrinsic charms in Λ_b decays

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ARTICLE INFO

Article history:

Received 8 September 2015

Received in revised form 19 October 2015

Accepted 7 November 2015

Available online 11 November 2015

Editor: J. Hisano

ABSTRACT

We study the three-body Λ_b decays of $\Lambda_b \rightarrow J/\psi pM$ with $M = K^-$ and π^- . The two new states $\mathcal{P}_{c1} \equiv \mathcal{P}_c(4380)^+$ and $\mathcal{P}_{c2} \equiv \mathcal{P}_c(4450)^+$ observed recently as the resonances in the $J/\psi p$ invariant mass spectrum of $\Lambda_b \rightarrow J/\psi pK^-$ can be identified to consist of five quarks, $uudc\bar{c}$, being consistent with the existence of the pentaquark states. We argue that, in the doubly charmful Λ_b decays of $\Lambda_b \rightarrow J/\psi pK^-$ through $b \rightarrow c\bar{c}s$, apart from those through the non-resonant $\Lambda_b \rightarrow pK^-$ and resonant $\Lambda_b \rightarrow \Lambda^* \rightarrow pK^-$ transitions, the third contribution with the non-factorizable effects is not the dominant part for the resonant $\Lambda_b \rightarrow K^-\mathcal{P}_{c1,c2}, \mathcal{P}_{c1,c2} \rightarrow J/\psi p$ processes, such that we propose that the $\mathcal{P}_{c1,c2}$ productions are mainly from the charmless Λ_b decays through $b \rightarrow \bar{u}s$, in which the $c\bar{c}$ content in $\mathcal{P}_{c1,c2}$ arises from the intrinsic charms within the Λ_b baryon. We hence predict the observables related to the branching ratios and the direct CP violating asymmetries to be $\mathcal{B}(\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p)/\mathcal{B}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p) = 0.58 \pm 0.05$, $\mathcal{A}_{CP}(\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p) = (-7.4 \pm 0.9)\%$, and $\mathcal{A}_{CP}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p) = (+6.3 \pm 0.2)\%$, which can alleviate the inconsistency between the theoretical expectations from the three contributions in the doubly charmful Λ_b decays and the observed data.

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1. Introduction

According to the recent observations of the three-body b -baryon decays of $\Lambda_b \rightarrow J/\psi pM$ with $M = K^-$ and π^- [1–3], apart from the non-resonant $\Lambda_b \rightarrow J/\psi pM$ and resonant $\Lambda_b \rightarrow J/\psi B^*, B^* \rightarrow pM$ ($B^* = \Lambda^* (N^*)$ for $M = K^- (\pi^-)$) contributions, depicted in Figs. 1a and 1b, respectively, there can be another resonant process in $\Lambda_b \rightarrow J/\psi pM$ as shown in Fig. 1c. The LHCb Collaboration has presented the compelling evidence for the new resonant states, being consistent with the existence of the pentaquark states as the five-quark bound states, while the $\mathcal{P}_c(4380)^+$ and $\mathcal{P}_c(4450)^+$ states are observed as the two resonances in the $J/\psi p$ invariant mass spectrum of $\Lambda_b \rightarrow J/\psi K^- p$, with the significance for each state to be more than 9 standard deviations, which can be regarded to be composed of $uudc\bar{c}$. We note that, in the same principle, the two new states should also exist in $\Lambda_b \rightarrow J/\psi p\pi^-$.

The processes of $\mathcal{P}_c \rightarrow J/\psi p$ in Fig. 1c are theoretically known to be dominated by the non-factorizable effects, calculated non-perturbatively with the scatterings of the soft hadrons, such that

the pentaquarks are considered as the molecular states [4–6]. In spite of the non-factorizable diagrams shown in Fig. 1c, which may get enhanced when the strong FSI interactions occur near the threshold to explain the pentaquark productions, we propose another possibility based on the factorizable effects. We note that this type of the processes in Fig. 1c has not been observed in the previous searches of the lighter pentaquarks than \mathcal{P}_c . For example, the resonant $B^+ \rightarrow \bar{p}\Theta(1710)^{++}, \Theta(1710)^{++} \rightarrow pK^+$ decay is measured to be $\mathcal{B}(B^+ \rightarrow \bar{p}\Theta(1710)^{++}, \Theta(1710)^{++} \rightarrow pK^+) < 9.1 \times 10^{-8}$ with the upper bound about 60–70 times smaller than the observed branching ratio of $B^+ \rightarrow p\bar{p}K^+$ [7,8], while $\mathcal{B}(B^0 \rightarrow \bar{p}\Theta(1540)^+, \Theta(1540)^+ \rightarrow pK_s^0) < 5 \times 10^{-8}$ is measured with the upper bound around 50 times smaller than $\mathcal{B}(B^0 \rightarrow p\bar{p}K_s^0)$ [8,9]. Similar to the charmless $B \rightarrow p\bar{p}K$ decays, the upper bound on $\mathcal{B}(B^0 \rightarrow \Theta_c \bar{p}\pi^+, \Theta_c \rightarrow D^{(*)-}p)$ is expected to be about 30–40 times smaller than $\mathcal{B}(B^0 \rightarrow p\bar{p}\pi^+ D^{(*)-})$ [10]. In contrast, since the resonant $\Lambda_b \rightarrow K^-\mathcal{P}_c, \mathcal{P}_c \rightarrow J/\psi p$ decays can contribute to the branching ratio as much as 10%, this leads to the question that if there can be other processes, which are responsible for the resonant $\mathcal{P}_c \rightarrow J/\psi p$ decays, other than the ones in Fig. 1c.

In Ref. [1], the LHCb Collaboration has given

$$\mathcal{R}_{\pi K} \equiv \frac{\mathcal{B}(\Lambda_b \rightarrow J/\psi p\pi^-)}{\mathcal{B}(\Lambda_b \rightarrow J/\psi pK^-)} = 0.0824 \pm 0.0025 \pm 0.0042,$$

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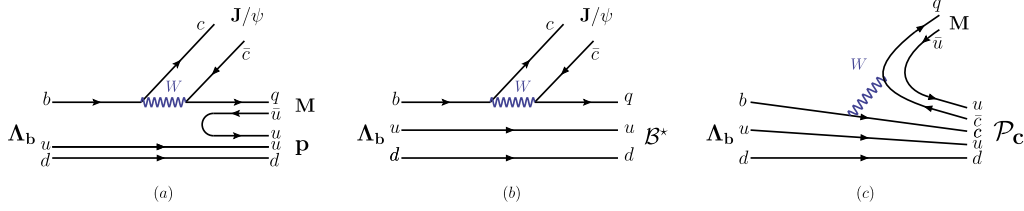


Fig. 1. Doubly charmful Λ_b decays of $\Lambda_b \rightarrow J/\psi pM$ from (a) non-resonant $\Lambda_b \rightarrow J/\psi pM$, (b) resonant $\Lambda_b \rightarrow J/\psi (B^* \rightarrow) pM$ with $B^* \equiv \Lambda^* \rightarrow pK^-$ ($N^* \rightarrow p\pi^-$) for $q = s(d)$, and (c) resonant $\Lambda_b \rightarrow M(\mathcal{P}_c \rightarrow) J/\psi p$ contributions, respectively.

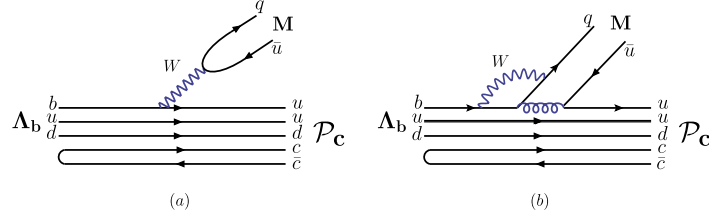


Fig. 2. The new contributions to the resonant $\Lambda_b \rightarrow M(\mathcal{P}_c \rightarrow) J/\psi p$ from the factorizable charmless Λ_b decays, where (a) and (b) are known as the tree and penguin contributions, respectively, while the $c\bar{c}$ contents are coming from the intrinsic charm within the Λ_b baryon.

$$\begin{aligned} \Delta\mathcal{A}_{CP} &\equiv \mathcal{A}_{CP}(\Lambda_b \rightarrow J/\psi p\pi^-) - \mathcal{A}_{CP}(\Lambda_b \rightarrow J/\psi pK^-) \\ &= (5.7 \pm 2.4 \pm 1.2)\%, \end{aligned} \quad (1)$$

where the first and second errors are from the statistical and systematic uncertainties, respectively. The data in Eq. (1) indicate some new $\Lambda_b \rightarrow M\mathcal{P}_c, \mathcal{P}_c \rightarrow J/\psi p$ processes beyond the non-factorizable ones in Fig. 1c with reasons as follows.

First, we note that $\mathcal{B}(\Lambda_b \rightarrow p\pi^-)/\mathcal{B}(\Lambda_b \rightarrow pK^-) = 0.84 \pm 0.09$ has been theoretically reproduced in Ref. [11], both $\mathcal{B}(\Lambda_b \rightarrow D^0 pK^-)/\mathcal{B}(\Lambda_b \rightarrow D^0 p\pi^-) = 0.073 \pm 0.008^{+0.005}_{-0.006}$ and $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c^+ K^-)/\mathcal{B}(\Lambda_b \rightarrow \Lambda_c^+ \pi^-) = 0.0731 \pm 0.0016 \pm 0.0016$ can be understood by the relation of $(V_{us}/V_{ud})^2 (f_K/f_\pi)^2 \simeq 0.075$ [12], and $\mathcal{B}(\Lambda_b \rightarrow \Lambda_c^+ D^-)/\mathcal{B}(\Lambda_b \rightarrow \Lambda_c^+ D_s^-) = 0.042 \pm 0.003 \pm 0.003$ is not far from the relation of $(V_{cd}/V_{cs})^2 (f_D/f_{D_s})^2 \simeq 0.035$ [13]. However, $\mathcal{R}_{\pi K} \simeq 0.08$ in Eq. (1) apparently deviates from $(V_{cd}/V_{cs})^2 \simeq 0.05$ given by the doubly charmful Λ_b decays in Fig. 1. To explain this difference, some new thinking is needed.

Second, $\Delta\mathcal{A}_{CP} \sim 5.7\%$ in Eq. (1) with the significance of 2.2σ suggests that a new contribution must proceed with V_{ub} to provide the weak CP phase, otherwise $\Delta\mathcal{A}_{CP} = 0$ as the case in the doubly charmful Λ_b decays in Fig. 1, in which such a phase is vanishingly small.

We hence propose that the resonant $\Lambda_b \rightarrow M\mathcal{P}_c, \mathcal{P}_c \rightarrow J/\psi p$ processes can be the new contributions to the charmless Λ_b decays as depicted in Fig. 2, where the $c\bar{c}$ content comes from the intrinsic charm (IC) in the Λ_b baryon. In the followings, we will assume that these new processes in Fig. 2 are the dominant ones for $\Lambda_b \rightarrow M\mathcal{P}_c, \mathcal{P}_c \rightarrow J/\psi p$.

It is not surprising that the Λ_b baryon contains the ICs, which are presented in the Fock state decomposition [14,15] as $|\Lambda_b\rangle = \Psi_{bud}|bud\rangle + \Psi_{bud\bar{c}}|bud\bar{c}\rangle + \dots$. In fact, the existence of the IC was first suggested in the proton to explain the large D^+ and Λ_c^+ productions at large energies in the proton–proton scattering [14,15]. In addition, as a possible solution to the so-called ρ - π puzzle [16], the IC in ρ for $J/\psi \rightarrow \rho^+\pi^-$ allows a strong decay not through the J/ψ annihilation suppressed by the Okubo–Zweig–Iizuka (OZI) rule. For a heavier hadron, since the gluon fluctuation, such as $gg \rightarrow c\bar{c}$, can easily occur without costing a large energy [21], it is expected that [22] the IC component in Λ_b ($m_{\Lambda_b} > m_B > m_p$) can be larger than the proton and the B mesons, estimated to be 1% and 4%, respectively. Consequently, in the $\Lambda_b \rightarrow p$ transition, we only consider the ICs in Λ_b since the heavier baryon would contribute a larger $c\bar{c}$ production. Note that, to distinguish the IC

in the proton from that in Λ_b , the J/ψ photoproduction can be useful [17–19], which is in accordance with Ref. [20]. While the study of the ICs in the B decays has been done extensively in the literature [21–24], it is not well examined in Λ_b , which should be a suitable scenario.

In this paper, since we propose that the two new resonant \mathcal{P}_c states, i.e. the pentaquark states, in the $m_{J/\psi p}$ spectrum of the $\Lambda_b \rightarrow J/\psi pK^-$ decay can be traced back to the charmless Λ_b decays from $b \rightarrow u\bar{u}s$, while the $c\bar{c}$ content in J/ψ is from the IC in the Λ_b baryon, we will study the branching ratios and the direct CP violating asymmetries, and check if our results will be able to understand the inconsistency between the theoretical estimations in the doubly charmful Λ_b decays and the observed data in Eq. (1).

2. Formalism

In terms of the effective Hamiltonian for $b \rightarrow c\bar{c}q$ at the quark level, the amplitude of $\Lambda_b \rightarrow J/\psi pM$ from Figs. 1a and 1b is given by

$$\begin{aligned} \mathcal{A}_{c\bar{c}q}(\Lambda_b \rightarrow J/\psi pM) &= \frac{G_F}{\sqrt{2}} V_{cb} V_{cq}^* a_2 \langle J/\psi | \bar{c}\gamma^\mu (1 - \gamma_5) c | 0 \rangle \\ &\quad \times \langle pM | \bar{q}\gamma_\mu (1 - \gamma_5) b | \Lambda_b \rangle, \end{aligned} \quad (2)$$

where G_F is the Fermi constant, V stands for the CKM mixing matrix, $q = s(d)$ corresponds to $M = K^-(\pi^-)$, $\langle pK | \bar{s}\gamma_\mu (1 - \gamma_5) b | \Lambda_b \rangle$ contains the contributions from the non-resonant $\Lambda_b \rightarrow pM$ and resonant $\Lambda_b \rightarrow B^* \rightarrow pM$ transitions, and $\langle J/\psi | \bar{c}\gamma^\mu (1 - \gamma_5) c | 0 \rangle = m_{J/\psi} f_{J/\psi} \varepsilon^{\mu*}$ with $m_{J/\psi}$, $f_{J/\psi}$ and $\varepsilon^{\mu*}$ being the mass, the decay constant and the polarization of J/ψ , respectively. Subsequently, the matrix elements of the combined $\Lambda_b \rightarrow pM$ transition can be parameterized as

$$\langle pM | \bar{q}\gamma_\mu (1 - \gamma_5) b | \Lambda_b \rangle \simeq F_M e^{i\delta_1} \bar{u}_p \gamma_\mu (1 - \gamma_5) u_{\Lambda_b}, \quad (3)$$

where δ_1 is the strong phase from the on-shell resonant $B^* \rightarrow pM$ decay and F_M is the parameter with $F_\pi/F_K \simeq (f_\pi/f_K)$ representing the flavor $SU(3)$ symmetry breaking. As a result, we rewrite the amplitude in Eq. (2) as

$$\begin{aligned} \mathcal{A}_{c\bar{c}q}(\Lambda_b \rightarrow J/\psi pM) &= \frac{G_F}{\sqrt{2}} V_{cb} V_{cq}^* a_2 m_{J/\psi} f_{J/\psi} F_M e^{i\delta_1} u_p \not{\varepsilon} (1 - \gamma_5) u_{\Lambda_b}, \end{aligned} \quad (4)$$

with $\not{g} = \varepsilon^{\mu*} \cdot \gamma_\mu$. From Fig. 2, which depicts the charmless Λ_b decays of $\Lambda_b \rightarrow \mathcal{P}_c M$, $\mathcal{P}_c \rightarrow J/\psi M$ decays, with $c\bar{c}$ in \mathcal{P}_c coming from the IC in Λ_b , the amplitudes of $\Lambda_b \rightarrow M\mathcal{P}_c$, $\mathcal{P}_c \rightarrow J/\psi p$ can be derived as [25]

$$\mathcal{A}_{\mathcal{P}_c} = i \frac{G_F}{\sqrt{2}} m_b f_M \left[\alpha_M \langle J/\psi p | \bar{u} b | \Lambda_b \rangle + \beta_M \langle J/\psi p | \bar{u} \gamma_5 b | \Lambda_b \rangle \right], \quad (5)$$

where f_M is the meson decay constant, defined by $\langle M | \bar{q}_1 \gamma_\mu \gamma_5 q_2 | 0 \rangle = -i f_M q_\mu$ with the four-momentum q_μ . The constant α_M (β_M) in Eq. (5) is related to the (pseudo)scalar quark current, given by

$$\alpha_M (\beta_M) = V_{ub} V_{uq}^* a_1 - V_{tb} V_{tq}^* (a_4 \pm r_M a_6), \quad (6)$$

where $r_M \equiv 2m_M^2/[m_b(m_q + m_u)]$ and $a_i \equiv c_i^{\text{eff}} + c_{i\pm 1}^{\text{eff}}/N_c^{\text{(eff)}}$ for $i = \text{odd}$ (even) are composed of the effective Wilson coefficients c_i^{eff} defined in Ref. [25]. In Eq. (5), the matrix elements for the resonant $\Lambda_b \rightarrow \mathcal{P}_c$, $\mathcal{P}_c \rightarrow J/\psi p$ transition can be given as

$$\langle J/\psi p | \bar{u} (\gamma_5) b | \Lambda_b \rangle = \langle J/\psi p | \mathcal{P}_c \rangle \mathcal{R}_{\mathcal{P}_c} \langle \mathcal{P}_c | \bar{u} (\gamma_5) b | \Lambda_b \rangle, \quad (7)$$

where the Breit–Wigner factor $\mathcal{R}_{\mathcal{P}_c}$ for \mathcal{P}_c is described as an intermediate state, given by

$$\mathcal{R}_{\mathcal{P}_c} = \frac{i}{(t - m_{\mathcal{P}_c}^2) + i m_{\mathcal{P}_c} \Gamma_{\mathcal{P}_c}}, \quad (8)$$

with $m_{\mathcal{P}_c}$ and $\Gamma_{\mathcal{P}_c}$ the mass and the decay width for the \mathcal{P}_c state, respectively. Despite the fact that there is no sufficient information for the detailed parameterization of $\langle J/\psi p | \mathcal{P}_c \rangle \langle \mathcal{P}_c | \bar{u} (\gamma_5) b | \Lambda_b \rangle$, the matrix elements of $\langle J/\psi p | \bar{u} (\gamma_5) b | \Lambda_b \rangle$ in Eq. (7) can still be reduced as

$$\begin{aligned} \langle J/\psi p | \bar{u} b | \Lambda_b \rangle &= \mathcal{R}_{\mathcal{P}_c} (\varepsilon \cdot q) F_S \bar{u}_p u_{\Lambda_b}, \\ \langle J/\psi p | \bar{u} \gamma_5 b | \Lambda_b \rangle &= \mathcal{R}_{\mathcal{P}_c} (\varepsilon \cdot q) F_P \bar{u}_p \gamma_5 u_{\Lambda_b}. \end{aligned} \quad (9)$$

This is due to the fact that, after the summations of the intermediate \mathcal{P}_c spins with spin = 3/2 or 5/2, all Lorentz indices are in fact coupled to be a scalar quantity, which can be parameterized as F_S and F_P . In general, $F_{S,P}$ are momentum dependent, but they can be taken as nearly constants around the threshold area of $t \simeq m_{\mathcal{P}_c}^2$, at which the threshold effect dominates the decay branching ratio. Besides, we take $F_S = F_P \equiv F_{\mathcal{P}_c}$ as a consequence of the Λ_b transition [11]. We hence obtain $\mathcal{A}_{\mathcal{P}_c} \simeq i \frac{G_F}{\sqrt{2}} m_b f_M \mathcal{R}_{\mathcal{P}_c} F_{\mathcal{P}_c} \bar{u}_p (\alpha_M + \beta_M \gamma_5) u_{\Lambda_b}$, such that the total amplitude for the two resonant \mathcal{P}_c states is in the form of

$$\begin{aligned} \mathcal{A}(\Lambda_b \rightarrow M(\mathcal{P}_{c1}, \mathcal{P}_{c2} \rightarrow) J/\psi p) \\ = \mathcal{A}_{\mathcal{P}_{c1}} + \mathcal{A}_{\mathcal{P}_{c2}} \simeq i \frac{G_F}{\sqrt{2}} m_b f_M F_2 e^{i\delta_2} \bar{u}_p (\alpha_M + \beta_M \gamma_5) u_{\Lambda_b}, \end{aligned} \quad (10)$$

with $F_2 e^{i\delta_2} = \mathcal{R}_{\mathcal{P}_{c1}} F_{\mathcal{P}_{c1}} + \mathcal{R}_{\mathcal{P}_{c2}} F_{\mathcal{P}_{c2}}$, where δ_2 is the strong phase from the on-shell $\mathcal{P}_c \rightarrow J/\psi p$ decays, and \mathcal{P}_{c1} and \mathcal{P}_{c2} denote $\mathcal{P}_c(4380)^+$ and $\mathcal{P}_c(4380)^+$, respectively. Note that $\mathcal{P}_{c1,c2}$ have been observed to have the masses and the decay widths as $(m, \Gamma) = (4380 \pm 8 \pm 29, 205 \pm 18 \pm 86)$ MeV and $(4449.8 \pm 1.7 \pm 2.5, 39 \pm 5 \pm 19)$ MeV, respectively, while their quantum numbers for J^P can be $(3/2^-, 5/2^+)$ or $(3/2^+, 5/2^-)$. However, the information of $\mathcal{P}_{c1,c2}$ can be cast into the to-be-determined parameters $F_2 e^{i\delta_2}$, without losing generality.

Table 1

The experimental inputs for the fitting, where the numbers are taken from [1,2].

	Data	Fitting results
$\mathcal{R}_{\pi K}$	$(8.24 \pm 0.49)\%$	$(8.38 \pm 0.77)\%$
$\Delta \mathcal{A}_{CP}$	$(5.7 \pm 2.7)\%$	$(2.9 \pm 1.4)\%$
$10^4 \mathcal{B}(\Lambda_b \rightarrow K^- J/\psi p)$	3.04 ± 0.55	3.21 ± 0.44
$10^6 \mathcal{B}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1} \rightarrow) J/\psi p)$	25.6 ± 13.8	10.3 ± 3.9
$10^6 \mathcal{B}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c2} \rightarrow) J/\psi p)$	12.5 ± 4.2	10.9 ± 2.7

3. Numerical analysis and discussions

For the numerical analysis, the theoretical inputs of the meson decay constants and Wolfenstein parameters in the CKM matrix are taken as [26,27]

$$\begin{aligned} (f_{J/\psi}, f_\pi, f_K) &= (418 \pm 9, 130.4 \pm 0.2, 156.2 \pm 0.7) \text{ MeV}, \\ (\lambda, A, \rho, \eta) \\ &= (0.225, 0.814, 0.120 \pm 0.022, 0.362 \pm 0.013), \end{aligned} \quad (11)$$

while the parameters $a_{1,4,6}$ can be adopted from Refs. [11,28], along with $a_2 = 0.2$ [29]. The data in the fitting are given in Table 1. As a result, we obtain

$$\begin{aligned} F_K &= 2.8 \pm 0.2, \quad F_{\mathcal{P}_{c1}} = 19.6 \pm 3.1, \\ F_{\mathcal{P}_{c2}} &= 5.5 \pm 1.0, \quad \delta_1 = (54.8 \pm 31.9)^\circ, \end{aligned} \quad (12)$$

with the fitted numbers in column 2 of Table 1 to be consistent with the data. First, for the three-body Λ_b decays only from the resonant $\Lambda_b \rightarrow M\mathcal{P}_c$, $\mathcal{P}_c \rightarrow J/\psi p$ contributions in Fig. 2, we obtain

$$\frac{\mathcal{B}(\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p)}{\mathcal{B}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p)} = 0.58 \pm 0.05, \quad (13)$$

where the parameters $F_2 e^{i\delta_2}$ in Eq. (10) have been canceled by the ratio. In the doubly charmful Λ_b decays, since the three contributions are all through $b \rightarrow c\bar{c}q$ at the quark level (see Fig. 1), the ratio of $\mathcal{R}_{\pi K}$ defined in Eq. (1) should be $(V_{cd}/V_{cs})^2 \simeq 0.05$, which is not approved by the data in Eq. (1). However, by adding the contributions from the charmless decays of $\Lambda_b \rightarrow K^-(\mathcal{P}_c \rightarrow) J/\psi p$, the doubly charmful $\mathcal{B}(\Lambda_b \rightarrow J/\psi p \pi^-)$ is getting close to the charmless $\mathcal{B}(\Lambda_b \rightarrow J/\psi(\mathcal{P}_c \rightarrow) p \pi^-)$, so that the value of $\mathcal{R}_{\pi K}$ is able to increase from 0.05 to a larger one to meet the data in Eq. (1), as the fitted result of $(8.38 \pm 0.77)\%$ in Table 1.

Second, the direct CP violating asymmetries from the resonant $\Lambda_b \rightarrow M\mathcal{P}_c$, $\mathcal{P}_c \rightarrow J/\psi p$ parts are evaluated to be

$$\begin{aligned} \mathcal{A}_{CP}(\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p) &= (-7.4 \pm 0.9)\%, \\ \mathcal{A}_{CP}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p) &= (+6.3 \pm 0.2)\%. \end{aligned} \quad (14)$$

However, since the measurement by the LHCb in Ref. [2] has suggested that the doubly charmful $\Lambda_b \rightarrow J/\psi p K^-$ mode dominates the corresponding decay, it leaves little room for the interference effects with the charmless ones of $\Lambda_b \rightarrow K^-(\mathcal{P}_{c1}, \mathcal{P}_{c2} \rightarrow) J/\psi p$ that provide the weak CP phase, of which $\mathcal{A}_{CP}^{\text{aw}}(\Lambda_b \rightarrow J/\psi p K^-) = (1.1 \pm 0.9)\%$ from the LHCb [1] agrees with the fitted result of $\mathcal{A}_{CP}(\Lambda_b \rightarrow J/\psi p K^-) = (-0.22 \pm 0.16)\%$. Note that $\Delta \mathcal{A}_{CP} = 0$ from $b \rightarrow c\bar{c}q$ to be different from the data of $\Delta \mathcal{A}_{CP} = 5.7\%$ in Eq. (1) requires the interference between the two compatible $\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow) J/\psi p$ and $\Lambda_b \rightarrow J/\psi(N^*(1440), N^*(1520) \rightarrow) p \pi^-$ channels. It is found that the contributions from $b \rightarrow c\bar{c}d$ with the strong phase $\delta_1 = 54.8^\circ$ and the contributions from $b \rightarrow u\bar{u}d$ with the weak phase by V_{ub} gives $\Delta \mathcal{A}_{CP} = (2.9 \pm 1.4)\%$, which is in good agreement with the data. Finally, the branching ratio for $\Lambda_b \rightarrow J/\psi p \pi^-$ is predicted as

$$\mathcal{B}(\Lambda_b \rightarrow J/\psi p\pi^-) = (2.68 \pm 0.34) \times 10^{-5}, \quad (15)$$

which includes the compatible contribution from $\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow)J/\psi p$ to agree well with $\mathcal{B}(\Lambda_b \rightarrow J/\psi p\pi^-) = (2.51 \pm 0.08 \pm 0.13_{-0.35}^{+0.45}) \times 10^{-5}$ measured by the LHCb [3], whereas the contributions only from the doubly charmful Λ_b decays give $\mathcal{B}(\Lambda_b \rightarrow J/\psi p\pi^-) = (1.68 \pm 0.24) \times 10^{-5}$, which is around $0.05\mathcal{B}(\Lambda_b \rightarrow J/\psi pK^-)$, borne by the relation of $|V_{cd}/V_{cs}|^2$.

In sum, the charmless processes of $\Lambda_b \rightarrow M(\mathcal{P}_{c1}, \mathcal{P}_{c2} \rightarrow)J/\psi$ provide us with a possible way to understand the CP asymmetry in Eq. (1) due to the origin of the weak phase from V_{ub} . Furthermore, to realize the ratio of $\mathcal{R}_{\pi K}$ in Eq. (1), which is unable to be explained from $b \rightarrow c\bar{c}q$, the contributions apart from the non-perturbative processes in Fig. 1c have to be included.

4. Conclusions

Since the non-factorizable effects for the doubly charmful Λ_b decays through $b \rightarrow c\bar{c}s$ may not be suitable to understand the resonant $\Lambda_b \rightarrow K^-\mathcal{P}_c, \mathcal{P}_c \rightarrow J/\psi p$ decays, while the new \mathcal{P}_c states observed in $m_{J/\psi p}$ spectrum can be identified as the pentaquark states with five quarks, $uudc\bar{c}$, we have proposed that these resonant processes could proceed as the charmless Λ_b decays through $b \rightarrow u\bar{u}s$, while the $c\bar{c}$ content in the \mathcal{P}_c states is from the intrinsic charms in the Λ_b baryon. As a result, we predicted that $\mathcal{B}(\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow)J/\psi p)/\mathcal{B}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1,c2} \rightarrow)J/\psi p) = 0.58 \pm 0.04$, $\mathcal{A}_{CP}(\Lambda_b \rightarrow \pi^-(\mathcal{P}_{c1,c2} \rightarrow)J/\psi p) = (-7.4 \pm 0.9)\%$, and $\mathcal{A}_{CP}(\Lambda_b \rightarrow K^-(\mathcal{P}_{c1,c2} \rightarrow)J/\psi p) = (+6.3 \pm 0.2)\%$, which could alleviate the inconsistency between the theoretical expectations of $(\mathcal{R}_{\pi K}, \Delta A_{CP}) = (0.05, 0)$ in the doubly charmful Λ_b decays and the observed data of $(\mathcal{R}_{\pi K}, \Delta A_{CP}) \sim (0.08, 5.7\%)$.

Acknowledgements

The authors would like to thank Yuan-Ning Gao and Qiang Zhao for useful discussions. The work was supported in part

by the National Center for Theoretical Sciences, National Science Council Taiwan (NSC-101-2112-M-007-006-MY3), MoST (MoST-104-2112-M-007-003-MY3) and National Tsing Hua University (104N2724E1).

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