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Possible large *CP* violation in three-body decays of heavy baryon

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A B S T R A C T

We propose a new mechanism which can introduce large *CP* asymmetries in the phase spaces of three-body decays of heavy baryons. In this mechanism, a large *CP* asymmetry is induced by the interference of two intermediate resonances, which subsequently decay into two different combinations of final particles. We apply this mechanism to the decay channel $\Lambda_b^0 \to p\pi^0\pi^-$, and find that the differential *CP* asymmetry can reach as large as 50%, while the regional *CP* asymmetry can reach as large as 16% in the interference region of the phase space.

1. Introduction

CP violation is an important phenomenon in particle physics. Although it has been discovered in the mixing and decay processes of K and B meson systems, including the first discovery of *CP* violation in K system [1], no *CP* violation was established in the baryon sector, except an evidence in the decay channel $\Lambda_b^0 \to pK^-$ [2]. Within the Standard Model, *CP* violation has originated from the weak phase in the Cabibbo–Kobayashi–Maskawa (CKM) matrix [3], along with a strong phase which usually arises from strong interactions. One reason for the smallness of *CP* violation is that the strong phases are usually small, especially when the strong phases come from a scale that is much larger than the QCD scale. However, non-perturbative effects of the strong interaction at low scales provide possibilities for large strong phases, and hence, large *CP* violation.

Three-body decays of heavy hadrons can be dominated by intermediate resonances in certain regions of the phase space. When two resonances decay into two different combinations of final particles, it is possible for them to dominate in the same region of the phase space. As a result, the interference effect together with a possible large strong phase can generate a large *CP* asymmetry.

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2. Differential *CP* asymmetry

It gets more interesting when one applies the aforementioned interference effect to the decay process of heavy baryons. For the decay process $\Lambda_b^0 \to p\pi^0\pi^-$, there is an overlap region in the phase space for resonances $\rho^-(770)$ and $N^+(1440)$, which lies right in the corner of the phase space. The decay amplitude for $\Lambda_b^0 \to p\pi^0\pi^-$ can be expressed as

$$\mathscr{M} = \frac{\langle p\pi^0|\hat{\mathscr{H}}_1|N^+\rangle\langle\pi^- N^+|\hat{\mathscr{H}}_{\text{eff}}|\Lambda_b^0\rangle}{s_0 - m_N^2 + im_N\Gamma_N} + \frac{\langle\pi^0\pi^-|\hat{\mathscr{H}}_2|\rho^-\rangle\langle p\rho^-|\hat{\mathscr{H}}_{\text{eff}}|\Lambda_b^0\rangle}{s - m_\rho^2 + im_\rho\Gamma_\rho}, \quad (1)$$

in the overlap region of the phase space, where $\hat{\mathscr{H}}_{\text{eff}}$ is the effective Hamiltonian for the weak decays, $\hat{\mathscr{H}}_1$ and $\hat{\mathscr{H}}_2$ are the formal Hamiltonian for the strong decays in which the magnitudes of the coupling constants can be determined from experiments, s and s_0 are the invariant mass squares of the systems $\pi^0\pi^-$ and $p\pi^0$, respectively, m_ρ, m_N, Γ_ρ, and Γ_N are the masses and decay widths of $\rho^0(770)$ and $N^+(1440)$, respectively, and the summation over the polarizations of the intermediate particles is understood. The effective Hamiltonian $\hat{\mathscr{H}}_{\text{eff}}$ takes the form [4]

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$$\hat{\mathscr{H}}_{\rm eff}=\frac{G_F}{\sqrt{2}}\Bigg[V_{ub}V_{ud}^*(c_1O_1^u+c_2O_2^u)+V_{cb}V_{cd}^*(c_1O_1^c+c_2O_2^c)-V_{tb}V_{td}^*\sum_{i=3}^{10}c_iO_i\Bigg]+h.c., \tag{2}$$

where G_F is the Fermi constant, V_{ud}, V_{ub}, V_{cd}, V_{cb}, V_{td}, and V_{tb} are the CKM matrix elements, c_i ($c_i=1,\cdots,10$) is the Wilson constant, and O_i is the four-Fermion operator, which takes the form

$$\begin{aligned}
O_1^q&=\bar{d}_\alpha\gamma_\mu(1-\gamma_5)q_\beta\bar{q}_\beta\gamma^\mu(1-\gamma_5)b_\alpha,\\
O_2^q&=\bar{d}\gamma_\mu(1-\gamma_5)q\bar{q}\gamma^\mu(1-\gamma_5)b,\\
O_3&=\bar{d}\gamma_\mu(1-\gamma_5)b\sum_{q'}\bar{q}'\gamma^\mu(1-\gamma_5)q',\\
O_4&=\bar{d}_\alpha\gamma_\mu(1-\gamma_5)b_\beta\sum_{q'}\bar{q}'_\beta\gamma^\mu(1-\gamma_5)q'_\alpha,\\
O_5&=\bar{d}\gamma_\mu(1-\gamma_5)b\sum_{q'}\bar{q}'\gamma^\mu(1+\gamma_5)q',\\
O_6&=\bar{d}_\alpha\gamma_\mu(1-\gamma_5)b_\beta\sum_{q'}\bar{q}'_\beta\gamma^\mu(1+\gamma_5)q'_\alpha,\\
O_7&=\frac{3}{2}\bar{d}\gamma_\mu(1-\gamma_5)b\sum_{q'}e_{q'}\bar{q}'\gamma^\mu(1+\gamma_5)q',\\
O_8&=\frac{3}{2}\bar{d}_\alpha\gamma_\mu(1-\gamma_5)b_\beta\sum_{q'}e_{q'}\bar{q}'_\beta\gamma^\mu(1+\gamma_5)q'_\alpha,\\
O_9&=\frac{3}{2}\bar{d}\gamma_\mu(1-\gamma_5)b\sum_{q'}e_{q'}\bar{q}'\gamma^\mu(1-\gamma_5)q',\\
O_{10}&=\frac{3}{2}\bar{d}_\alpha\gamma_\mu(1-\gamma_5)b_\beta\sum_{q'}e_{q'}\bar{q}'_\beta\gamma^\mu(1-\gamma_5)q'_\alpha,
\end{aligned}\tag{3}$$

with d, b, q, and q' being quark fields and α and β being colour indices.

Under the factorization hypothesis, the weak decay amplitudes can be expressed as

$$\langle\pi^-N^+|\mathscr{H}_{\rm eff}|\Lambda_b^0\rangle=i\eta^N\overline{u}_N\not{p}_{\pi^-}(1-\gamma_5)u_{\Lambda_b}, \tag{4}$$

$$\langle\rho^-p|\mathscr{H}_{\rm eff}|\Lambda_b^0\rangle=\eta^p m_\rho\overline{u}_N\not{\epsilon}_{\rho^-}(1-\gamma_5)u_{\Lambda_b}, \tag{5}$$

where ϵ_{ρ^-} is the polarization vector of ρ^-, u_N and u_{Λ_b} are the spinors for $N^+(1440)$ and Λ_b, respectively,

$$\eta^N=\frac{G_F}{\sqrt{2}}f_\pi F^{\Lambda_b\to N^+}\Bigg\{a_2V_{ub}V_{ud}^*-V_{tb}V_{td}^*\Bigg[(a_4+a_{10})-\frac{2m_\pi^2(a_6+a_8)}{(m_u+m_d)m_b}\Bigg]\Bigg\}, \tag{6}$$

$$\eta^p=\frac{G_F}{\sqrt{2}}f_\rho F^{\Lambda_b\to p}\left\{a_2V_{ub}V_{ud}^*-V_{tb}V_{td}^*\left[(a_4+a_{10})\right]\right\}, \tag{7}$$

with f_π being the decay constant of the pion, $F^{\Lambda_b\to N^+}$ and $F^{\Lambda_b\to p}$ being the form factors for the transition $\Lambda_b\to N^+(1440)$ and $\Lambda_b\to p$, respectively, and $a_i=c_i+c_{i-1}/N_c$ for even i.

Because of the non-perturbative effects of strong interactions, there can be a relative strong phase between the coupling constants of $\hat{\mathscr{H}}_1$ and $\hat{\mathscr{H}}_2$. We will denote this relative phase by δ and treat it as a free parameter. The strong decay amplitudes are then expressed as

$$\langle p\pi^0|\hat{\mathscr{H}}_1|N^+\rangle=ig_1\overline{u}_p\gamma_5u_N, \tag{8}$$

and

$$\langle\pi^0\pi^-|\hat{\mathscr{H}}_2|\rho^-\rangle=e^{i\delta}g_2(p_{\pi^-}-p_{\pi^0})\cdot\epsilon_{\rho^-}, \tag{9}$$

respectively, where the effective coupling constants g_1 and g_2 can be expressed as

$$g_1^2=\frac{8\pi m_N^2\Gamma_{N^+\to N\pi}}{3\lambda_N(m_N^2+m_\rho^2-2m_Nm_p-m_\pi^2)}, \tag{10}$$

$$g_2^2=\frac{6\pi m_\rho^2\Gamma_{\rho^-\to\pi^0\pi^-}}{\lambda_\rho^3}, \tag{11}$$

with m_p being the mass of proton, $\Gamma_{N^+\to N\pi}$ and $\Gamma_{\rho^-\to\pi^0\pi^-}$ being the partial decay widths for $N^+(1440)\to N(939)\pi$ and $\rho^-(770)\to\pi^0\pi^-$, respectively, and

$$\lambda_N=\frac{1}{2m_N}\sqrt{\left[m_N^2-(m_p+m_\pi)^2\right]\cdot\left[m_N^2-(m_p-m_\pi)^2\right]}, \tag{12}$$

$$\lambda_\rho=\frac{1}{2}\sqrt{m_\rho^2-4m_\pi^2}. \tag{13}$$

The differential CP asymmetry is then defined as

$$A_{CP}=\frac{\overline{|\mathscr{M}|^2}-\overline{|\bar{\mathscr{M}}|^2}}{\overline{|\mathscr{M}|^2}+\overline{|\bar{\mathscr{M}}|^2}}, \tag{14}$$

where $\bar{\mathscr{M}}$ is the decay amplitude of the CP conjugate process, $\overline{\Lambda_b^0}\to\overline{p}\pi^+\pi^0$, and the overlines above $|\mathscr{M}|^2$ and $\left|\bar{\mathscr{M}}\right|^2$ represent averaging and summing over the spin states of the initial and final particles, respectively. After some algebra, one has

$$\begin{aligned}
\overline{|\mathscr{M}|^2}=\Big\{&|\lambda_1|^2\left[(m_{\Lambda_b}^2-s_-)(s_0-m_p^2)-m_\pi^2(m_{\Lambda_b}-m_p)^2+m_\pi^4\right]\\
&+|\lambda_2|^2\left[(m_{\Lambda_b}^2-s_0)(s_--m_p^2)-m_\pi^2(m_{\Lambda_b}-m_p)^2+m_\pi^4\right]\\
&+2\mathscr{R}\left(\lambda_1\lambda_2^*\right)\Big[s_0s_-+m_{\Lambda_b}m_p(m_{\Lambda_b}^2-m_{\Lambda_b}m_p\\
&+m_p^2-s_0-s_-)-m_\pi^4\Big]\Big\}\\
&+\left\{m_{\Lambda_b}\to-m_{\Lambda_b}\right\},
\end{aligned}\tag{15}$$

where s_- is the invariant mass squared of the system $p\pi^-$,

$$\lambda_1=\frac{m_{\Lambda_b}^2-s_0}{m_{\Lambda_b}-m_p}\frac{g_1}{s_N}\eta^N+e^{i\delta}\frac{g_2}{s_\rho}m_\rho\eta^p, \tag{16}$$

$$\lambda_2=\frac{m_{\Lambda_b}(m_p-m_N)+m_pm_N-s_0}{m_{\Lambda_b}-m_p}\frac{g_1}{s_N}\eta^N-e^{i\delta}\frac{g_2}{s_\rho}m_\rho\eta^p, \tag{17}$$

and $s_N=s_0-m_N^2+im_N\Gamma_N$, $s_\rho=s-m_\rho^2+im_\rho\Gamma_\rho$. In order to obtain the expression for $\overline{\left|\bar{\mathscr{M}}\right|^2}$, all one needs to do is to replace the CKM matrix elements in Eq. (15) with their complex conjugates.

In order to see where the CP asymmetry arises, let's first display the weak and strong phases in λ_1 and λ_2 explicitly. For a fixed point in the phase space, λ_1 and λ_2 can be expressed as

$$\lambda_i=\lambda_i^{\rm Tree}e^{i(\phi_i^{\rm Tree}+\alpha_i^{\rm Tree})}+\lambda_i^{\rm Penguin}e^{i(\phi_i^{\rm Penguin}+\alpha_i^{\rm Penguin})}, \tag{18}$$

where $i=1,2$, $\lambda_i^{\rm Tree}$ and $\lambda_i^{\rm Penguin}$ are the tree and penguin parts of λ_i, respectively, $\phi_i^{\rm Tree}$ and $\phi_i^{\rm Penguin}$ are the corresponding weak phases, which take the values

$$\phi_i^{\rm Tree}={\rm Arg}\left(V_{ub}V_{ud}^*\right), \tag{19}$$