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$\Lambda_b \rightarrow \Lambda_c + a_1$ decay in the heavy quark effective theory

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Abstract. Using the heavy quark approximation, we have studied the nonleptonic decay mode $\Lambda_b \rightarrow \Lambda_c a_1$. We have included nonfactorizable contributions as well as factorizable ones in our analysis. The estimated branching ratio for this process is (1.4 ± 0.1)% and the asymmetry parameter α found to be – 0.8.

Keywords. Nonfactorizable contributions; HQET; Isgur–Wise function.

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1. Introduction

In a recent letter one event of the decay mode $\Lambda_b \rightarrow \Lambda_c a_1$ has been reported by the DELPHI collaboration [1]. Though at present, there are not many experimental results available for heavy baryons, we expect in the near future that a large sample of Λ_b baryons will be available at the colliders. From the theoretical point of view considerable progress has been achieved in the last few years in understanding the weak decays of heavy hadrons due to the development of heavy quark effective theory (HQET) [2]. Contrary to the significant progress made in the studies of the meson decays, advancement in the arena of heavy baryons has been very slow. In particular the nonleptonic weak decays of heavy baryons have not been understood clearly till now. In fact some phenomenological approaches such as pole model, current algebra etc. have been employed to analyse these decay processes. The well-known factorization hypothesis [3] which has been applied successfully to the heavy meson decays can also be applied for heavy baryon cases [4]. However in a recent article it has been shown that the nonfactorization terms also contribute significantly to these decay processes [5]. In this paper we intend to study the decay mode $\Lambda_b \rightarrow \Lambda_c a_1$ in the heavy quark effective theory with $1/m_Q$ corrections. We have included the factorizable and nonfactorizable contributions in our analysis.

2. Methodology

The effective Hamiltonian including short distance QCD corrections for the decay process under consideration is given as

$$\mathcal{H}_{\rm eff} = \frac{G_{\rm F}}{\sqrt{2}} V_{ud}^* V_{cb} [C_1(m_b)(\bar{c}b)_\mu (\bar{d}u)^\mu + C_2(m_b)(\bar{c}u)_\mu (\bar{d}b)^\mu], \tag{1}$$

where $G_{\rm F}$ is the Fermi coupling constant and the quark current $(\bar{q}'q)_\mu$ is a short hand for the weak current $\bar{q}'_\alpha \gamma_\mu (1+\gamma_5) q_\alpha$; α is the color index. The values of the Wilson coefficients $C_{1,2}$ can be taken as [6]

$$C_1(m_b) = 1.128 \quad \text{and} \quad C_2(m_b) = -0.288. \tag{2}$$

Now applying Fierz transformation to the second term of eq. (1) we get

$$(\bar{c}u)_\mu (\bar{d}b)^\mu = \frac{1}{N_c} (\bar{c}b)_\mu (\bar{d}u)^\mu + \frac{1}{2} \sum_a (\bar{c}\lambda^a b)(\bar{d}\lambda^a u), \tag{3}$$

where the second term is a color singlet operator formed by the color octet currents, λ^a are the SU(3) Gell–Mann matrices and N_c is the number of colours. Hence the transition amplitude for the decay mode $\Lambda_b \to \Lambda_c a_1$ is given as

$$\mathcal{M}(\Lambda_b \to \Lambda_c a_1) = \frac{G_{\rm F}}{\sqrt{2}} V_{ud}^* V_{cb} \Bigg[\left(C_1(m_b) + \frac{C_2(m_b)}{N_c} \right) \langle \Lambda_c a_1 | (\bar{c}b)_\mu (\bar{d}u)^\mu | \Lambda_b \rangle + C_2(m_b) \langle \Lambda_c a_1 | \frac{1}{2} \sum_a (\bar{c}\lambda^a b)(\bar{d}\lambda^a u) | \Lambda_b \rangle \Bigg]. \tag{4}$$

The first term in eq. (4) contains both factorizable and nonfactorizable contributions while the second term is purely nonfactorizable. Hence the amplitude becomes

$$\mathcal{M}(\Lambda_b \to \Lambda_c a_1) = \frac{G_{\rm F}}{\sqrt{2}} V_{ud}^* V_{cb} \Bigg[a_1 \Big\{ \langle \Lambda_c | (\bar{c}b)_\mu | \Lambda_b \rangle \langle a_1 | (\bar{d}u)^\mu | 0 \rangle + \langle \Lambda_c a_1 | (\bar{c}b)_\mu (\bar{d}u)^\mu | \Lambda_b \rangle^{\rm n.f} \Big\} + C_2(m_b) \langle \Lambda_c a_1 | \frac{1}{2} \sum_a (\bar{c}\lambda^a b)(\bar{d}\lambda^a u) | \Lambda_b \rangle^{\rm n.f} \Bigg], \tag{5}$$

where the coefficient $a_1 = C_1(m_b) + C_2(m_b)/N_c$.

Now let us define two parameters ϵ_1 and ϵ_8 that contain the nonfactorizable contributions as

$$\epsilon_1 = \frac{\langle \Lambda_c a_1 | (\bar{c}b)_\mu (\bar{d}u)^\mu | \Lambda_b \rangle^{\rm n.f.}}{\langle \Lambda_c | (\bar{c}b)_\mu | \Lambda_b \rangle \langle a_1 | (\bar{d}u)^\mu | 0 \rangle} \tag{6}$$

and

$$\epsilon_8 = \frac{\langle \Lambda_c a_1 | \frac{1}{2} \sum_a (\bar{c}\lambda^a b)(\bar{d}\lambda^a u) | \Lambda_b \rangle^{\rm n.f}}{\langle \Lambda_c | (\bar{c}b)_\mu | \Lambda_b \rangle \langle a_1 | (\bar{d}u)^\mu | 0 \rangle}. \tag{7}$$

With these two new parameters one can write eq. (5) as

$$\mathcal{M}(\Lambda_b \to \Lambda_c a_1) = \frac{G_{\rm F}}{\sqrt{2}} V_{ud}^* V_{cb} a_1^{\rm eff} \langle \Lambda_c | (\bar{c}b)_\mu | \Lambda_b \rangle \langle a_1 | (\bar{d}u)^\mu | 0 \rangle, \tag{8}$$

where $a_1^{\rm eff}$ is given as

$$a_1^{\rm eff} = a_1(1+\epsilon_1) + C_2\epsilon_8. \tag{9}$$

Thus we have found that the nonfactorizable effects amounts to a redefinition of the effective parameter a_1 which is assumed to be universal. So we still have a factorization scheme with the universal parameter $a_1^{\rm eff}$ to be determined from experiment. In a recent article, using the experimental data of nonleptonic B meson decays, Al-Shamali and Kamal [6] have obtained

$$a_1^{\rm eff} = 0.92 \pm 0.03. \tag{10}$$

Since in this analysis we are dealing with the same type of transition as for $B \to D\pi$ (i.e. $b \to c$), the value of $a_1^{\rm eff}$ for $\Lambda_b \to \Lambda_c a_1$ decay mode is assumed to be not very much different from the above one, we have taken $a_1^{\rm eff} = 0.92$ in our calculation.

To evaluate the transition matrix element between the two heavy baryons we employ HQET. One thus write the matrix element $\langle\Lambda_c|(\bar{c}b)_\mu|\Lambda_b\rangle$ in the HQET [2] as

$$\langle\Lambda_c^+(v',s')|\bar{c}\gamma_\mu(1+\gamma_5)b|\Lambda_b(v,s)\rangle = \eta(v\cdot v')\bar{u}_c(v',s')\gamma_\mu(1+\gamma_5)u_b(v,s). \tag{11}$$

The single form factor $\eta(v\cdot v')$ is given as

$$\eta(v\cdot v') = \left(\frac{\alpha_s(m_b)}{\alpha_s(m_c)}\right)^{a_I}\left(\frac{\alpha_s(m_c)}{\alpha_s(\mu)}\right)^{a_L}\eta_0(v\cdot v'), \tag{12}$$

where

$$a_I = -6/25 \quad \text{and} \quad a_L = 8/27[(v\cdot v')r(v\cdot v') - 1] \tag{13}$$

with

$$r(v\cdot v') = \frac{1}{\sqrt{(v\cdot v')^2-1}}\ln[v\cdot v' + \sqrt{(v\cdot v')^2-1}]. \tag{14}$$

$\alpha_s(m_b), \alpha_s(m_c)$ and $\alpha_s(\mu)$ are taken as 0.2, 0.32 and 0.26 respectively [2]. The function $\eta_0(v\cdot v')$ is the universal non-perturbative Isgur–Wise function which is given in the bound state soliton picture [7] (where the heavy baryons containing a single heavy charm or bottom quark are treated as the bound state of the chiral soliton with the heavy D or B meson) as

$$\eta_0(v\cdot v') = \left[\frac{4\sqrt{\mu_b\mu_c}}{(\sqrt{\mu}_b+\sqrt{\mu}_c)^2}\right]^{3/4}\exp\left(-\frac{(v\cdot v'-1)M_B^2}{\sqrt{\kappa}(\sqrt{\mu}_b+\sqrt{\mu}_c)}\right), \tag{15}$$

where M_B is the mass of the chiral soliton, i.e. the nucleon mass, μ_Q is the reduced mass of the soliton heavy meson bound system, κ is the spring constant and its value is taken to be [7] (440 MeV)^3. The product $(v\cdot v')$ is determined by considering the kinematics of the system as

$$v \cdot v' = \frac{M_{\Lambda_b}^2 + M_{\Lambda_c}^2 - M_{a_1}^2}{2M_{\Lambda_b} M_{\Lambda_c}}. \quad (16)$$

The matrix element $\langle a_1 | \bar{d}\gamma_\mu (1+\gamma_5) u | 0 \rangle$ is given as

$$\langle a_1 | \bar{d}\gamma^\mu (1+\gamma_5) u | 0 \rangle = f_{a_1} M_{a_1} \epsilon^\mu, \quad (17)$$

where ϵ^μ is the polarization of the axial vector meson a_1 and f_{a_1} is the decay constant.

Thus with eqs (11) and (17) the matrix element for the decay $\Lambda_b \to \Lambda_c a_1$ becomes

$$\mathcal{M}(\Lambda_b \to \Lambda_c a_1) = \frac{G_{\mathrm{F}}}{\sqrt{2}} V_{ud}^* V_{cb} a_1^{\mathrm{eff}} f_{a_1} M_{a_1} \eta(v \cdot v') \bar{u}_{\Lambda_c} \gamma_\mu (1+\gamma_5) u_{\Lambda_b} \epsilon_\mu. \quad (18)$$

The corresponding decay rate is given as

$$\Gamma(\Lambda_b \to \Lambda_c a_1) = \frac{|\vec{p}|}{8\pi M_{\Lambda_b}^2} G_{\mathrm{F}}^2 |V_{cb} V_{ud}|^2 (a_1^{\mathrm{eff}})^2 \eta^2 (v \cdot v') f_{a_1}^2 \times [(M_{\Lambda_b}^2 - M_{\Lambda_c}^2)^2 + M_{a_1}^2 (M_{\Lambda_b}^2 + M_{\Lambda_c}^2 - 2M_{a_1}^2)], \quad (19)$$

where $|\vec{p}|$ is the magnitude of the three momentum of the emitted particle in the rest frame of the Λ_b baryon. Using the particle masses from ref. [8] as $M_{\Lambda_b} = 5624$ MeV, $M_{\Lambda_c} = 2284.9$ MeV, $M_{a_1} = 1230$ MeV and the value of the decay constant [6] $f_{a_1} = (228 \pm 10)$ MeV we obtain the branching ratio for the process

$$\mathrm{Br}(\Lambda_b \to \Lambda_c a_1) = (1.1 \pm 0.1)\%. \quad (20)$$

To obtain the asymmetry parameter, we compare the transition amplitude to the general form in the rest frame of Λ_b,

$$\chi_{\Lambda_c}^\dagger [S\vec{\sigma} + P_1 \hat{p} + iP_2 (\hat{p} \times \vec{\sigma}) + D(\vec{\sigma} \cdot \hat{p})\hat{p}] \cdot \epsilon \chi_{\Lambda_b}, \quad (21)$$

where $\hat{p}$ is the unit vector in the direction of outgoing baryon. We can obtain the values for S, P_1, P_2 and D as [9]

$$S = \frac{G_{\mathrm{F}}}{\sqrt{2}} V_{ud}^* V_{cb} a_1^{\mathrm{eff}} f_{a_1} \eta(v \cdot v') M_{a_1}, \quad (22)$$

$$\frac{P_1}{S} = -\left(\frac{M_{\Lambda_b} + M_{\Lambda_c}}{E_{\Lambda_c} + M_{\Lambda_c}}\right) \frac{|\vec{p}|}{E_{a_1}} = -1.268, \quad (23)$$

$$\frac{P_2}{S} = \frac{|\vec{p}|}{(E_{\Lambda_c} + M_{\Lambda_c})} = 0.397, \quad (24)$$

$$\frac{D}{S} = \frac{|\vec{p}|^2}{E_{a_1}(E_{\Lambda_c} + M_{\Lambda_c})} = 0.345. \quad (25)$$

From these expressions we can obtain the asymmetry parameter α of the final Λ_c with respect to Λ_b polarization as

$$\alpha = 2\mathrm{Re}\,\frac{[(1+D/S)^* P_1/S + 2(P_2^*/S)M_{a_1}^2/E_{a_1}^2]}{K}, \tag{26}$$

where

$$K = [|1+D/S|^2 + |P_1/S|^2 + 2(1+|P_2/S|^2)M_{a_1}^2/E_{a_1}^2]. \tag{27}$$

Thus the asymmetry parameter α is obtained as $\alpha = -0.8$, where the negative sign reflects the $V-A$ structure of the current.

Now for an improvement over the result for the decay rate, one should include the $1/m_Q$ corrections to the decay widths. These corrections arise from the modification of the current operator between the full theory of the QCD and the heavy quark effective theory as well as from the modification of hadronic states.

We first evaluate the corrections arising from the current modifications. Up to order $(1/m_Q)$, the heavy quark current operator $\bar{c}\gamma_\mu(1+\gamma_5)b$ in the full theory can be expanded in terms of the current operator in the effective theory [10] as

$$\begin{aligned}\bar{c}\gamma_\mu(1+\gamma_5)b = {} & \bar{h}_{v'}^{(c)}\gamma_\mu(1+\gamma_5)h_v^{(b)} + \frac{1}{2m_b}\bar{h}_{v'}^{(c)}\gamma_\mu(1+\gamma_5)(i\overrightarrow{\not{D}})\,h_v^{(b)} \\ & + \frac{1}{2m_c}\bar{h}_{v'}^{(c)}(-i\overleftarrow{\not{D}})\,\gamma_\mu(1+\gamma_5)h_v^{(b)}.\end{aligned} \tag{28}$$

In the effective theory the matrix elements of the operators on the right hand side of (28) between the Λ_b and Λ_c baryon states are given as

$$\begin{aligned}\langle\Lambda_c(v',s')|\bar{h}_{v'}^{(c)}\gamma_\mu(1+\gamma_5)h_v^{(b)}|\Lambda_b(v,s)\rangle = {} & \eta(v\cdot v')\bar{u}_c(v',s') \\ & \times\gamma_\mu(1+\gamma_5)u_b(v,s),\end{aligned} \tag{29}$$

$$\begin{aligned}\langle\Lambda_c(v',s')|\bar{h}_{v'}^{(c)}\gamma_\mu(1+\gamma_5)(i\overrightarrow{\not{D}})\,h_v^{(b)}|\Lambda_b(v,s)\rangle = {} & \eta_\alpha(v,v')\bar{u}_c(v',s') \\ & \times\gamma_\mu(1+\gamma_5)\gamma_\alpha u_b(v,s),\end{aligned} \tag{30}$$

$$\begin{aligned}\langle\Lambda_c(v',s')|\bar{h}_{v'}^{(c)}(-i\overleftarrow{\not{D}})\,\gamma_\mu(1+\gamma_5)h_v^{(b)}|\Lambda_b(v,s)\rangle = {} & \eta_\alpha(v',v)\bar{u}_c(v',s') \\ & \gamma_\alpha\gamma_\mu(1+\gamma_5)u_b(v,s).\end{aligned} \tag{31}$$

The most general form of $\eta_\alpha(v,v')$ is given as

$$\eta_\alpha(v,v') = \eta_+(w)(v+v')_\alpha + \eta_-(w)(v-v')_\alpha, \tag{32}$$

where $w = v\cdot v', \eta_\pm(w)$ are given as [11]

$$\eta_+(w) = \frac{\bar{\Lambda}}{2}\,\frac{w-1}{w+1}\eta(w), \tag{33}$$

$$\eta_-(w) = \frac{\bar{\Lambda}}{2}\,\eta(w). \tag{34}$$

The parameter $\bar{\Lambda}$ is the scale of light degrees of freedom given as $\bar{\Lambda} = M_{\Lambda_Q} - m_Q \approx 700\,\text{MeV}$.

As for the corrections arising from the hadronic states we include $1/m_Q$ corrections to the effective Lagrangian i.e.,

$$\mathcal{L}_{\text{eff}} = \sum_Q \left[h_v^{(Q)} iv \cdot D h_v^{(Q)} + \frac{\mathcal{L}_I}{2m_Q} \right], \tag{35}$$

where

$$\mathcal{L}_I = \bar{h}_v^Q \left[(iD)^2 + \frac{g}{2}\sigma_{\mu\nu}G^{\mu\nu} \right] h_v^Q = \mathcal{O}_1 + \mathcal{O}_2. \tag{36}$$

Now inserting the higher order term of the effective Lagrangian $(\mathcal{L}_I)$ into the matrix elements of the lowest order current $J = \bar{h}'\Gamma h$, we obtain [2,11]

$$\langle \Lambda(v',s')|iT \int \mathrm{d}x \{J(0), \mathcal{O}_1(x)\}|\Lambda(v,s)\rangle = \bar{\Lambda}\chi(w)\bar{u}_\Lambda(v',s')\Gamma u_{\Lambda_b}(v,s). \tag{37}$$

Insertion of the kinetic operator $\mathcal{O}_1$ preserves the Dirac structure of the current and transforms trivially under the spin symmetry, contributing to the form factors in the same proportion as the leading order term (11) and hence effectively correcting the Isgur–Wise function [2,11]. The contributions from the chromomagnetic operator $\mathcal{O}_2$ vanishes by Lorentz invariance. For the subleading form factor, vector current conservation implies

$$\chi(1) = 0. \tag{38}$$

Since $\chi(w)$ only corrects the Isgur–Wise function, we assume a simple form being consistent with the constraint (38) as done in ref. [2]

$$\chi(w) = \frac{w-1}{w+1}\eta(w). \tag{39}$$

Now with eqs (28)–(39) we obtain the decay rate for the $\Lambda_b \to \Lambda_c a_1$ process including $1/m_Q$ corrections as

$$\begin{aligned}
\Gamma(\Lambda_b(v) \to \Lambda_c(v')a_1(p)) &= \frac{G_{\mathrm{F}}^2}{8\pi M_{\Lambda_b}^2}|V_{ud}^* V_{cb}|^2 (a_1^{\text{eff}})^2 f_{a_1}^2 \eta^2(v\cdot v')|\vec{p}| \\
&\times \Bigg[A \left\{ 1 + \frac{\bar{\Lambda}}{w+1}(2w-1)\left(\frac{1}{m_b} + \frac{1}{m_c}\right) \right. \\
&\left. + \left(\frac{\bar{\Lambda}}{w+1}\right)^2 w(w-1)\left(\frac{1}{m_b}+\frac{1}{m_c}\right)^2 \right\} \\
&- \frac{\bar{\Lambda}}{w+1}\left(\frac{B}{m_c} + \frac{C}{m_b}\right) + \left(\frac{\bar{\Lambda}}{w+1}\right)^2 \left(\frac{D}{m_c^2} + \frac{E}{m_b^2} + \frac{F}{m_c m_b}\right) \Bigg],
\end{aligned} \tag{40}$$

where

$$\begin{aligned}
A &= (M_{\Lambda_b}^2 - M_{\Lambda_c}^2)^2 + M_{a_1}^2(M_{\Lambda_b}^2 + M_{\Lambda_c}^2 - 2M_{a_1}^2),\\
B &= 2M_{\Lambda_b}M_{\Lambda_c}[2(M_{\Lambda_b} - M_{\Lambda_c}w)^2 + M_{a_1}^2],\\
C &= 2M_{\Lambda_b}M_{\Lambda_c}[2(M_{\Lambda_b}w - M_{\Lambda_c})^2 + M_{a_1}^2],\\
D &= M_{\Lambda_b}M_{\Lambda_c}(1-w)[2(M_{\Lambda_b} - M_{\Lambda_c}w)^2 + M_{a_1}^2],\\
E &= M_{\Lambda_b}M_{\Lambda_c}(1-w)[2(M_{\Lambda_b}w - M_{\Lambda_c})^2 + M_{a_1}^2],\\
F &= 2M_{\Lambda_b}M_{\Lambda_c}(1-w)[(1+w)^2(M_{\Lambda_b} - M_{\Lambda_c})^2\\
&\quad + w(w-1)(M_{\Lambda_b} + M_{\Lambda_c})^2 + M_{a_1}^2].
\end{aligned} \tag{41}$$

Using the quark masses as $m_c = 1.5$ GeV and $m_b = 5$ GeV we obtain the branching ratio as

$$\mathrm{Br}(\Lambda_b \to \Lambda_c a_1) = (1.4 \pm 0.1)\%. \tag{42}$$

3. Conclusion

In this paper we have presented estimates of the branching ratio and asymmetry parameter for the decay mode $\Lambda_b \to \Lambda_c a_1$. Heavy quark effective theory with $1/m_Q$ corrections to both hadronic states and current operator has been employed. We have also included the nonfactorizable contributions as well as factorizable ones in our analysis. Since HQET does not predict the shape of the universal form factor i.e. the Isgur–Wise function, besides the normalization condition $\eta_0(w = 1) = 1$, evaluation in a suitable model is necessary for the entire kinematic region. Here we have evaluated the IW function in the bound state soliton picture at the particular kinematic point of interest. Data in the baryonic b-sector have started to pour in and in particular for the decay $\Lambda_b \to \Lambda_c a_1$ one event has already been observed by the DELPHI Collaboration [1]. We expect more data to come in the near future (b-sector) from the hadron colliders to test our result and to enrich our understanding in the baryonic hadronic transition and above all the standard model of particle interactions.

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