Eur. Phys. J. C 16, 289–294 (2000)
Digital Object Identifier (DOI) 10.1007/s100520000425

THE EUROPEAN PHYSICAL JOURNAL C

CP violation in $\Lambda_b \to p\pi^-$ decay

R. Mohanta

Physics Department, Panjab University, Chandigarh–160014, India

Received: 15 February 2000 / Published online: 8 June 2000 – © Springer-Verlag 2000

Abstract. Using the next-to-leading order QCD corrected effective Hamiltonian, scale and scheme independent Wilson coefficients, we have estimated the branching ratio and CP violating asymmetries for the $\Lambda_b \to p\pi^-$ decay mode in the Standard Model using the framework of generalized factorization. The effects of nonfactorizable contributions are taken into account by treating the effective no. of colors (N_c^{eff}) as a free parameter. The form factors are evaluated in the nonrelativistic quark model. The estimated branching ratio is found to be $\mathcal{O}(10^{-6})$ which lies below the current experimental upper limit and the CP violating asymmetries are $a_{cp} \sim -8\%$ and $A(\alpha) \sim 2 \times 10^{-5}$.

1 Introduction

Despite many attempts CP violation still remains one of the most outstanding problem in particle physics [1]. So far it has been observed only within the neutral K meson system. Yet Kaon system by itself can not provide the whole picture of CP violation. Consequently it is essential to study CP violation outside this system which is important in order to understand whether the standard model provides a correct description of this phenomenon through the Cabibbo-Kobayashi-Maskawa (CKM) matrix. In this respect the B-meson system appears to be most promising, which is also reflected by the tremendous experimental efforts at the present and future B factories. Studies of CP violation in the B meson system [2] have suggested that large CP-violating asymmetries should be observed in the forthcoming experiments. The basic signal for CP violation is the partial rate difference between a B decay mode and its CP conjugate process. On the basis of CKM picture, either $B^0 - \bar{B}^0$ mixing or the absorptive part of penguin amplitudes [3] can give rise to significant CP violation in exclusive nonleptonic B decays. Since charged B mesons can not mix, a measurement of the CP violating observable in these decays would be a clear sign of direct CP violation. The existence of such CP asymmetries require the interference of two decay amplitudes with different weak and strong phase differences. The weak phase difference arises from the superposition of various penguin contributions and the usual tree diagrams while the strong phases are induced by final state interactions. It is worth emphasizing that the above decay modes are flavor self tagging processes which should be favored for experimental reconstuctions. It is also interesting to study CP violation in the bottom baryon system in order to find the physical channels which may have large CP asymmetries, even though the branching ratios for such processes are usually smaller than those for the corresponding processes of bottom mesons. The study of CP violation in the bottom system will be helpful for understanding the origin of CP violation and may provide useful information about the possible baryon asymmetry in our universe. Measurement of decay width of $\Lambda_b \to \Lambda J/\psi$ has been reported recently [5] and one expects more data in future in bottom baryon sector. In this paper we intend to study CP violation in the nonleptonic $\Lambda_b \to p\pi^-$ decay in the standard model. Like the search for CP violation in the $B^\pm$ decays, the study of beauty baryon decay does not need tagging processes of the associated beauty hadrons produced in the same event. CP violation in strange hyperon decays is extensively studied in [6] within the standard model and beyond it. The weak phase differences in these cases arise from the CKM matrix elements whereas the stong phase differences are evaluated considering the experimental data of $N\pi$ phase shifts. However for the $\Lambda_b \to p\pi^-$ decay mode the phase shifts are not known experimentally and have to be determined theoretically. The strong phases are generated by final state interactions (FSI). At the quark level, the strong phase differences arise through the absorptive parts of perturbative penguin diagrams (hard final state interactions) [3] and nonperturbatively (soft final state interactions) [4]. In the absence of an argument that parton-hadron duality should hold in exclusive processes, one can not exclude that the weak transition matrix elements receive phases originating from soft FSI. However the effects of soft FSI are extremely difficult to quantify. In the absence of a reliable theoretical calculation for soft FSI, we make the usual approximation of retaining the absorptive part from quark level calculation (hard FSI) for the strong phase difference in our analysis.

Here we use the standard theoretical framework to study the nonleptonic $\Lambda_b \to p\pi^-$ decay mode, which is based on the effective Hamiltonian approach in conjuction with the factorization hypothesis for hadronic matrix elements. The short distance QCD corrected Hamiltonian is calculated next to leading order. The renormalization scheme and scale problems with factorization approach for matrix elements can be circumvented by employing the scale and scheme independent effective Wilson coefficients. The form factors at maximum recoil have been calculated using the nonrelativistic quark model [7] and the nearest pole dominance has been used to extrapolate them to the required q^2 point.

The paper is organized as follows. The phenomenology of hyperon decays is presented in Sect. 2. In Sect. 3 we discuss the effective Hamiltonian together with the quark level matrix elements and the numerical value of the Wilson coefficients in the effective Hamiltonian approach. Using the factorization ansatz we evaluate the matrix elements in the nonrelativistic quark model [7]. In Sect. 4 the CP violating observables are discussed. Section 5 contains our results and discussions.

2 Phenomenology of hyperon decays

The study of CP violation in strange hyperon decays is extensively studied in [6], where the phenomenogy of hyperon decays are discussed in great detail. However for the sake of completeness we shall present here the basic features of their nonleptonic decays. The most general Lorentz-invariant amplitude for the decay $\Lambda_b \to p\pi^-$ can be written as

$$i\bar{u}_p(p_f)(a + b\gamma_5)u_{\Lambda_b}(p_i) \tag{1}$$

The corresponding matrix element for $\bar{\Lambda}_b \to \bar{p}\pi^+$ is then

$$i\bar{v}_{\bar{p}}(p_f)(-a^* + b^*\gamma_5)v_{\bar{\Lambda}_b}(p_i) \tag{2}$$

It is convenient to express the transition amplitude in terms of S-wave (parity violating) and P-wave (parity conserving) amplitudes S and P as

$$S + P\sigma \cdot \hat{\mathbf{q}} \tag{3}$$

where $\mathbf{q}$ is the proton momentum in the rest frame of Λ_b baryon and the amplitudes S and P are:

$$S = a\sqrt{\frac{\{(m_{\Lambda_b} + m_p)^2 - m_\pi^2\}}{16\pi m_{\Lambda_b}^2}}$$
$$P = b\sqrt{\frac{\{(m_{\Lambda_b} - m_p)^2 - m_\pi^2\}}{16\pi m_{\Lambda_b}^2}} \tag{4}$$

The experimental observables are the total decay rate Γ and the decay parameters α, β and γ which govern the decay-angular distribution and the polarization of the final baryon. The decay rate is given as

$$\Gamma = 2|\mathbf{q}|\{|S|^2 + |P|^2\} \tag{5}$$

and the asymmetry parameters are given as

$$\alpha = \frac{2Re(S^*P)}{\{|S|^2 + |P|^2\}}$$
$$\beta = \frac{2Im(S^*P)}{\{|S|^2 + |P|^2\}}$$
$$\alpha^2 + \beta^2 + \gamma^2 = 1 \tag{6}$$

The dominant term in the angular distribution is α hence we concentrate ourselves only to the parameter α only. Similar observables for the antihyperon decays are $\bar{\Gamma}$, $\bar{\alpha}$, $\bar{\beta}$ and $\bar{\gamma}$ where $\bar{\Gamma}$, $\bar{\alpha}$ are given as

$$\bar{\Gamma} = 2|\mathbf{q}|\{|\bar{S}|^2 + |\bar{P}|^2\}$$
$$\bar{\alpha} = \frac{2Re(\bar{S}^*\bar{P})}{\{|\bar{S}|^2 + |\bar{P}|^2\}} \tag{7}$$

3 Transition amplitude in the factorization approximation

The effective Hamiltonian $\mathcal{H}_{eff}$ for the decay process $\Lambda_b \to p\pi^-$ is given as

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}}\Big\{\lambda_u[c_1(\mu)O_1^u(\mu) + c_2(\mu)O_2^u(\mu)] + (\lambda_u + \lambda_c)\sum_{i=3}^{10} c_i(\mu)O_i(\mu)\Big\} + \text{h.c.}\,, \tag{8}$$

where $\lambda_u = V_{ub}V_{ud}^*$ and $\lambda_c = V_{cb}V_{cd}^*$ and $c_i(\mu)$ are the Wilson coefficients evaluated at the renormalization scale μ. The operators O_{1-10} are given as

$$\begin{aligned}
O_1^u &= (\bar{u}b)_{V-A}(\bar{d}u)_{V-A}\,,\\
O_2^u &= (\bar{u}_\alpha b_\beta)_{V-A}(\bar{d}_\beta u_\alpha)_{V-A}\,,\\
O_{3(5)} &= (\bar{d}b)_{V-A}\sum_{q'}(\bar{q}'q')_{V-A(V+A)}\,,\\
O_{4(6)} &= (\bar{d}_\alpha b_\beta)_{V-A}\sum_{q'}(\bar{q}'_\beta q'_\alpha)_{V-A(V+A)}\,,\\
O_{7(9)} &= \frac{3}{2}(\bar{d}b)_{V-A}\sum_{q'} e_{q'}(\bar{q}'q')_{V+A(V-A)}\,,\\
O_{8(10)} &= \frac{3}{2}(\bar{d}_\alpha b_\beta)_{V-A}\sum_{q'} e_{q'}(\bar{q}'_\beta q'_\alpha)_{V+A(V-A)}\,,
\end{aligned} \tag{9}$$

where $O_{1,2}$ are the tree level current-current operators, O_{3-6} are the QCD and O_{7-10} are the EW penguin operators. $(\bar{q}_1 q_2)_{(V\pm A)}$ denote the usual $(V \pm A)$ currents. The sum over q' runs over the quark fields that are active at the scale $\mu = O(m_b)$ i.e. $(q' \in u, d, s, c, b)$. The Wilson coefficients depend (in general) in the renormalization scheme and the scale μ at which they are evaluated. In the next to leading order their values obtained in the naive dimensional regularization (NDR) scheme at $\mu = m_b(m_b)$ as [8]

$$\begin{aligned}
&c_1 = 1.082 \quad c_2 = -0.185 \quad c_3 = 0.014 \quad c_4 = -0.035\\
&c_5 = 0.009 \quad c_6 = -0.041 \quad c_7 = -0.002\ \alpha \quad c_8 = 0.054\ \alpha\\
&c_9 = -1.292\ \alpha \quad c_{10} = 0.263\ \alpha\,.
\end{aligned} \tag{10}$$

However the physical matrix elements $\langle p\pi|\mathcal{H}_{eff}|\Lambda_b\rangle$ are obviously independent of both scheme and the scale. Hence the dependence in the Wilson coefficients must be compensated by a comensurate calculation of the hadronic matrix elements in a nonperturbative framework, such as lattice QCD. Presently this is not a viable strategy as the calculation of the matrix elements $\langle \pi p|O_i|\Lambda_b\rangle$ is beyond the scope of the current lattice technology. However perturbation theory comes to (partial) rescue; with the help of which one-loop matrix elements can be rewritten in terms of the operators and the effective Wilson coefficients c_i^{eff} which are scheme and scale independent:

$$\langle d\bar{u}u|\mathcal{H}_{eff}|b\rangle = \sum_{i,j} c_i^{eff}(\mu)\langle d\bar{u}u|O_j|b\rangle^{tree} . \tag{11}$$

The effective Wilson coefficients $c_i^{eff}(\mu)$ may be expressed as [9]

$$\begin{aligned}
c_1^{eff}|_{\mu=m_b} &= c_1(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{1i} c_i(\mu) ,\\
c_2^{eff}|_{\mu=m_b} &= c_2(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{2i} c_i(\mu) ,\\
c_3^{eff}|_{\mu=m_b} &= c_3(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{3i} c_i(\mu)\\
&\quad -\frac{\alpha_s}{24\pi}(C_t + C_p + C_g) ,\\
c_4^{eff}|_{\mu=m_b} &= c_4(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{4i} c_i(\mu)\\
&\quad +\frac{\alpha_s}{8\pi}(C_t + C_p + C_g) ,\\
c_5^{eff}|_{\mu=m_b} &= c_5(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{5i} c_i(\mu)\\
&\quad -\frac{\alpha_s}{24\pi}(C_t + C_p + C_g) ,\\
c_6^{eff}|_{\mu=m_b} &= c_6(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{6i} c_i(\mu)\\
&\quad +\frac{\alpha_s}{8\pi}(C_t + C_p + C_g) ,\\
c_7^{eff}|_{\mu=m_b} &= c_7(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{7i} c_i(\mu)\\
&\quad +\frac{\alpha}{8\pi}C_e ,\\
c_8^{eff}|_{\mu=m_b} &= c_8(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{8i} c_i(\mu) ,\\
c_9^{eff}|_{\mu=m_b} &= c_9(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{9i} c_i(\mu)\\
&\quad +\frac{\alpha}{8\pi}C_e ,\\
c_{10}^{eff}|_{\mu=m_b} &= c_{10}(\mu) + \frac{\alpha_s}{4\pi}\left(\gamma^{(0)T}\ln\frac{m_b}{\mu} + \hat{r}^T\right)_{10i} c_i(\mu) .
\end{aligned} \tag{12}$$

where $\hat{r}^T$ and $\gamma^{(0)T}$ are the transpose of the matrices $\hat{r}$ and $\gamma^{(0)}$ arise from the vertex corrections to the operators $O_1 - O_{10}$ derived in [10], which are explicitly given in [11]

The quantities C_t, C_p and C_g are arising from the peguin type diagrams of the opearators $O_{1,2}$, the peguin type diagrams of the opearators $O_3 - O_6$ and the tree level diagrams of the dipole operator O_g respectively which are given in the NDR scheme (after $\overline{\text{MS}}$ renormalization) by

$$\begin{aligned}
C_t &= -\left(\frac{\lambda_u}{\lambda_t}\tilde{G}(m_u) + \frac{\lambda_c}{\lambda_t}\tilde{G}(m_c)\right) c_1\\
C_p &= [\tilde{G}(m_d) + \tilde{G}(m_b)]c_3 + \sum_{i=u,d,s,c,b}\tilde{G}(m_i)(c_4 + c_6)\\
C_g &= -\frac{2m_b}{\sqrt{\langle k^2\rangle}}c_g^{eff} , \qquad c_g^{eff} = -1.043\\
C_e &= -\frac{8}{9}\left(\frac{\lambda_u}{\lambda_t}\tilde{G}(m_u) + \frac{\lambda_c}{\lambda_t}\tilde{G}(m_c)\right)(c_1 + 3c_2)\\
\tilde{G}(m_q) &= \frac{2}{3} - G(m_q, k, \mu)
\end{aligned} \tag{13}$$

$$G(m,k,\mu) = -4\int_0^1 dx\; x(1-x)\ln\left(\frac{m^2 - k^2 x(1-x)}{\mu^2}\right) , \tag{14}$$

It should be noted that the quantities C_t, C_p and C_g depend on the CKM matrix elements, the quark masses, the scale μ and k^2, the momentum transferred by the virtual particles apearing in the penguin diagrams. In the factorization approximation there is no model independent way to keep track of the k^2 dependence; the actual value of k^2 is model dependent. From simple kinematics of charmless nonleptonic B decays [12] one expects k^2 to be typically in the range

$$\frac{m_b^2}{4} \le k^2 \le \frac{m_b^2}{2} . \tag{15}$$

Assuming that in the rest frame of the Λ_b baryon, the spectator diquarks both in the initial and final baryon have negligible momentum and the momentum shared equally by between the two quarks of the pion, we have found $k^2 \approx m_b^2/2$. For numerical calculation we have taken the CKM matrix elements expressed in terms of the Wolfenstein parameters with values $A = 0.815$, $\lambda = \sin\theta_c$=0.2205, $\rho = 0.175$ and $\eta = 0.37$ [9]. Using the mass renormalization equations with three loop β function, the values of the current quark masses are evaluated at various energy scales in [13]. Since the energy released in the decay mode $\Lambda_b \to p\pi^-$ is of the order of m_b, we take the current quark mass values at scale $\mu \sim m_b$ from [13] as: $m_u(m_b) = 3.2$ MeV, $m_d(m_b) = 6.4$ MeV, $m_s(m_b) = 90$ MeV, $m_c(m_b)$= 0.95 GeV and $m_b(m_b)$= 4.34 GeV. Thus we obtain the values of the effective renormalization scheme and scale independent Wilson coefficients for $b \to d$ transitions as:

$$\begin{aligned}
&c_1^{eff} = 1.168 \quad c_2^{eff} = -0.365 \quad c_3^{eff} = 0.0224 + i0.0038\\
&c_4^{eff} = -(0.0455 + i0.0115) \quad c_5^{eff} = 0.0131 + i0.0038\\
&c_6^{eff} = -(0.0475 + i0.0115)\\
&c_7^{eff}/\alpha = -(0.0294 + i0.0329) \quad c_8^{eff}/\alpha = 0.055\ \alpha\\
&c_9^{eff}/\alpha = -(1.426 + i0.0329) \quad c_{10}^{eff}/\alpha = 0.48 .
\end{aligned} \tag{16}$$

After obtaining the effective Wilson coefficients now we want to calculate the matrix element $\langle \pi p|O_i|\Lambda_b\rangle$ where O_i are the four quark current operators listed in (9), using the factorization approximation. In this approximation, the hadronic matrix elements of the four quark operators $(\bar{d}b)_{(V-A)}(\bar{u}d)_{(V-A)}$ split into the product of two matrix elements, $\langle \pi|(\bar{d}u)_{(V-A)}|0\rangle$ and $\langle p|(\bar{u}b)_{(V-A)}|\Lambda_b\rangle$ where

Fierz transformation has been used so that flavor quantum numbers of the currents match with those of the hadrons. Since Fierzing yield operators which are in the color singlet-singlet and octet-octet forms, this procedure results in general the matrix elements which have the right flavor quantum numbers but involve both singlet-singlet and octet-octet current operators. However there is no experimental information available for the octet-octet part. So in the factorization approximation, one discards the color octet-octet piece and compensates this by treating N_c, the numbers of colors as a free parameter, and its value is extracted from the data of two body nonleptonic decays.

The matrix elements of the $(V-A)(V+A)$ operators i.e. (O_6 & O_8) can be calculated as follows. After Fierz ordering and factorization they contribute as [14]

$$\langle p\pi|O_6|\Lambda_b\rangle = -2\sum_q \langle\pi|\bar{d}(1+\gamma_5)q|0\rangle\langle p|\bar{q}(1-\gamma_5)b|\Lambda_b\rangle \tag{17}$$

Using the Dirac equation the matrix element can be rewritten as

$$\langle p\pi|O_6|\Lambda_b\rangle = \Big[R_1\langle p|V_\mu|\Lambda_b\rangle - R_2\langle p|A_\mu|\Lambda_b\rangle\Big]\langle\pi|A_\mu|0\rangle\ , \tag{18}$$

with

$$R_1 = \frac{2m_\pi^2}{(m_b-m_u)(m_d+m_u)}\ ,$$
$$R_2 = \frac{2m_\pi^2}{(m_b+m_u)(m_d+m_u)}\ , \tag{19}$$

where the quark masses are the current quark masses. The same relation works for O_8.

Thus under the factorization approximation the baryon decay amplitude is governed by a decay constant and baryonic transition form factors. The general expression for the baryon transition is given as

$$\begin{aligned}\langle p(p_f)|V_\mu - A_\mu|\Lambda_b(p_i)\rangle &= \bar{u}_p(p_f)\Big\{f_1(q^2)\gamma_\mu + if_2(q^2)\sigma_{\mu\nu}q^\nu + f_3(q^2)q_\mu \\ &-[g_1(q^2)\gamma_\mu + ig_2(q^2)\sigma_{\mu\nu}q^\nu + g_3(q^2)q_\mu]\gamma_5\Big\}u_{\Lambda_b}(p_i)\ ,\end{aligned} \tag{20}$$

where $q = p_i - p_f$. In order to evaluate the form factors at maximum momentum transfer, we have employed nonrelativistic quark model [7], where they are given as:

$$\begin{aligned}f_1(q_m^2)/N_{fi} &= 1 - \frac{\Delta m}{2m_i} + \frac{\Delta m}{4m_i m_q}\left(1-\frac{\Lambda_b}{2m_f}\right) \\ &\quad\times(m_i+m_f-\eta\Delta m) \\ &\quad-\frac{\Delta m}{8m_i m_f}\frac{\bar{\Lambda}}{m_Q}(m_i+m_f-\eta\Delta m) \\ f_3(q_m^2)/N_{fi} &= \frac{1}{2m_i} - \frac{1}{4m_i m_f}(m_i+m_f-\eta\Delta m) \\ &\quad-\frac{\bar{\Lambda}}{8m_i m_f m_Q}[(m_i+m_f)\eta+\Delta m] \\ g_1(q_m^2)/N_{fi} &= \eta + \frac{\Delta m\bar{\Lambda}}{4}\left(\frac{1}{m_i m_q} - \frac{1}{m_f m_Q}\right)\eta \\ g_3(q_m^2)/N_{fi} &= -\frac{\bar{\Lambda}}{4}\left(\frac{1}{m_i m_q} - \frac{1}{m_f m_Q}\right)\eta\end{aligned} \tag{21}$$

where $\bar{\Lambda} = m_f - m_q$, $\Delta m = m_i - m_f$ (m_i and m_f are the initial and final baryon masses), $q_m^2 = \Delta m^2$, $\eta = 1$, m_Q and m_q are the constituent quark masses of the interacting quarks of initial and final baryons with values m_u=338 MeV and m_b=5 GeV. N_{fi} is the flavour factor:

$$N_{fi} =_{\text{flavor spin}} \langle p|b_u^\dagger b_b|\Lambda_b\rangle_{\text{flavor spin}} = \frac{1}{\sqrt{2}} \tag{22}$$

Since the calculation of q^2 dependence of form facors is beyond the scope of the nonrelativistic quark model we will follow the conventional practice to assume a pole dominance for the form factor q^2 behaviour as

$$f(q^2) = \frac{f(0)}{(1-q^2/m_V^2)^2} \qquad g(q^2) = \frac{g(0)}{(1-q^2/m_A^2)^2} \tag{23}$$

where $m_V(m_A)$ is the pole mass of the vector (axial vector) meson with the same quantum number as the current under consideration. The pole masses are taken as $m_V = 5.32$ GeV and $m_A = 5.71$ GeV. Assuming a dipole q^2 behaviour for form factors, and taking the masses of the baryons and pion from [15] we found

$$\begin{aligned}f_1(m_\pi^2) &= 0.043 \qquad & m_i f_3(m_\pi^2) &= -0.009 \\ g_1(m_\pi^2) &= 0.092 \qquad & m_i g_3(m_\pi^2) &= -0.047\end{aligned} \tag{24}$$

The matrix element $\langle\pi|A_\mu|0\rangle$ is related to the pion decay constant f_π as

$$\langle\pi(q)|A_\mu|0\rangle = -if_\pi q_\mu \tag{25}$$

Hence one obtains the transition amplitude for $\Lambda_b \to p\pi^-$ as (where the factor $G_F/\sqrt{2}$ is suppressed)

$$\begin{aligned}&A(\Lambda_b \to p\pi^-) \\ &= if_\pi\bar{u}_p(p_f)\Big[\Big\{\lambda_u\,(a_1+a_4+a_{10}+(a_6+a_8)R_1) \\ &\quad+\lambda_c\,(a_4+a_{10}+(a_6+a_8)R_1)\Big\}\left(f_1(m_\pi^2)(m_i-m_f)\right. \\ &\quad\left.+\,f_3(m_\pi^2)m_\pi^2\right) + \Big\{\lambda_u\,(a_1+a_4+a_{10}+(a_6+a_8)R_2) \\ &\quad+\lambda_c\,(a_4+a_{10}+(a_6+a_8)R_2)\Big\}\left(g_1(m_\pi^2)(m_i+m_f)\right. \\ &\quad\left.-\,g_3(m_\pi^2)m_\pi^2\right)\gamma_5\Big]u_{\Lambda_b}(p_i)\ .\end{aligned} \tag{26}$$

The coefficients $a_1,\ a_2\cdots a_{10}$ are combinations of the effective Wilson coefficients given as

$$\begin{aligned}a_{2i-1} &= c_{2i-1}^{eff} + \frac{1}{(N_c^{eff})_{2i-1}}c_{2i}^{eff} \\ a_{2i} &= c_{2i}^{eff} + \frac{1}{(N_c^{eff})_{2i}}c_{2i-1}^{eff} \quad i=1,2\cdots5\ ,\end{aligned} \tag{27}$$

$$a_{cp} = \frac{2Im(\lambda_u\lambda_c^*)\,(Im(S_uS_c^*)+Im(P_uP_c^*))}{|\lambda_u|^2(|S_u|^2+|P_u|^2)+|\lambda_c|^2(|S_c|^2+|P_c|^2)+2Re(\lambda_u\lambda_c^*)\,(Re(S_uS_c^*)+Re(P_uP_c^*))}$$
$$= \frac{-2\sin\gamma\left(|S_uS_c|\sin(\delta_u^S-\delta_c^S)+|P_uP_c|\sin(\delta_u^P-\delta_c^P)\right)}{|\frac{\lambda_u}{\lambda_c}|(|S_u|^2+|P_u|^2)+|\frac{\lambda_c}{\lambda_u}|(|S_c|^2+|P_c|^2)+2\cos\gamma\,(|S_uS_c|\cos(\delta_u^S-\delta_c^S)+|P_uP_c|\cos(\delta_u^P-\delta_c^P))}\,, \quad (35)$$

where N_c^{eff} is the effective no. of colors treated as free parameter in order to model the nonfactorizable contributions to the matrix elements and its value can be extracted from the two body nonleptonic B decays. Naive factorization implies $N_c = 3$. A recent analysis of $B \to D\pi$ data gives $N_c^{eff} \sim 2$ [16]. On the other hand Mannel et al [17] have used $N_c^{eff} = \infty$ to study the nonleptonic decays of Λ_b baryon. So here we have taken three sets of values i.e., 2, 3 and ∞ for the effective no. of colors.

Using (4) and (20)-(26) we obtain the parity violating (S) and parity conserving (P) amplitudes (in the unit of $f_\pi G_F/\sqrt{2}\times 10^{-3}$) for three different sets of effective no. of colors as

$$\begin{aligned} S &= \lambda_u(30.215-0.624i)-\lambda_c(2.453+0.624i)\\ P &= \lambda_u(64.662+1.334i)-\lambda_c(5.245-1.334i)\,,\\ &\quad (\text{for } N_c^{eff}=2) \end{aligned} \quad (28)$$

$$\begin{aligned} S &= \lambda_u(32.1-0.66i)-\lambda_c(2.581+0.66i)\\ P &= \lambda_u(68.7-1.411i)-\lambda_c(5.52+1.411i)\\ &\quad (\text{for } N_c^{eff}=3) \end{aligned} \quad (29)$$

$$\begin{aligned} S &= \lambda_u(35.871-0.738i)-\lambda_c(2.845+0.738i)\\ P &= \lambda_u(76.769-1.578i)-\lambda_c(6.084+1.578i)\,,\\ &\quad (\text{for } N_c^{eff}=\infty) \end{aligned} \quad (30)$$

4 CP violating asymmetry

For $\Lambda_b \to p\pi^-$ decay the CP violating rate asymmetry in partial decay rate and α (Asymmetry parameter) are defined as follows,

$$a_{cp} = \frac{\Gamma(\Lambda_b\to p\pi^-)-\Gamma(\bar{\Lambda}_b\to \bar{p}\pi^+)}{\Gamma(\Lambda_b\to p\pi^-)+\Gamma(\bar{\Lambda}_b\to \bar{p}\pi^+)}\,, \quad (31)$$

$$A(\alpha) = \frac{\alpha+\bar{\alpha}}{\alpha-\bar{\alpha}}\,. \quad (32)$$

As these decays are all self tagging the measurement of these CP violating asymmetry is essentially a counting experiment in well defined final states. Their rate asymmetries require both weak and strong phase differences in interfereing amplitudes. The weak phase difference arises from the superposition of amplitudes from various tree (current-current) and penguin diagrams. The strong phase which are needed to obtain nonzero values for a_{cp} are generated by final state interactions.

Without loss of generality, we can write the parity violating/conserving transition amplitudes for the hyperon decay as

$$\begin{aligned} S &= \lambda_u|S_u|e^{i\delta_{su}}+\lambda_c|S_c|e^{i\delta_{sc}}\\ P &= \lambda_u|P_u|e^{i\delta_{pu}}+\lambda_c|P_c|e^{i\delta_{pc}} \end{aligned} \quad (33)$$

Table 1. Branching ratio and CP violating parameters for the decay $\Lambda_b \to p\pi^-$ decay mode

N_c^{eff}	BR (Theory)	BR Expt. [15]	a_{cp}	$A(\alpha)$
2	0.83×10^{-6}		-8.3%	-2.4×10^{-5}
3	0.93×10^{-6}	$<5\times10^{-5}$	-8.3%	-2.3×10^{-5}
∞	1.16×10^{-6}		-8.3%	-1.5×10^{-5}

where $\lambda_q = V_{qb}V_{qd}^*$, $(S/P)_u$ and $(S/P)_c$ denote the contribution from hadronic matrix elements proportional to the product of CKM matrix elements λ_u and λ_c respectively for (S/P) waves. The corresponding strong phases are denoted by $\delta_u^{(S/P)}$ and $\delta_c^{(S/P)}$ respectively.The corresponding quantities for the antihyperon decay are given as

$$\begin{aligned} \bar{S} &= -\left(\lambda_u^*|S_u|e^{i\delta_{su}}+\lambda_c^*|S_c|e^{i\delta_{sc}}\right)\\ \bar{P} &= \lambda_u^*|P_u|e^{i\delta_{pu}}+\lambda_c^*|P_c|e^{i\delta_{pc}} \end{aligned} \quad (34)$$

Thus the CP violating rate asymmetry is given as, (see (35) on top of the page) where the weak phases entering in the $b \to d$ transition is equal to $-\gamma$, as we are using Wolfenstein approximation in which λ_c has no weak phase and the phase of λ_u is $-\gamma$ which is obtained from the relation $\tan\gamma = \left(\frac{\eta}{\rho}\right)$. The strong phases $(\delta_u^S - \delta_c^S)$ are obtained from

$$\begin{aligned} \cos(\delta_u^S-\delta_c^S) &= \frac{1}{|S_uS_c|}(ReS_u\ ReS_c+ImS_u\ ImS_c)\\ \sin(\delta_u^S-\delta_c^S) &= \frac{1}{|S_uS_c|}(ImS_u\ ReS_c-ImS_c\ ReS_u) \end{aligned} \quad (36)$$

Similar relations hold for $(\delta_u^P - \delta_c^P)$ with the amplitude S is replaced by P.

5 Results and discussions

In this section we have estimated the branching ratio and the CP violating asymmetries for the decay mode $\Lambda_b \to p\pi^-$. The magnitude of the parity-conserving (P wave) and parity violating (S-wave) amplitudes are given in (28)–(30) for three different sets of effective no. of colors. Using the pion decay constant f_π to be 130.7 MeV and G_F from [15] we have obtained the branching ratio with (5) as given in Table 1. It is seen that the branching ratio is maximum for $N_c^{eff} = \infty$, but its value for all three sets of N_c^{eff} lies below the present experimental upper limt $BR(\Lambda_b \to p\pi^-) < 5\times10^{-5}$ [15]. The CP asymmetry a_{cp} is found to be $\sim -8\%$ in all the three cases and the

CP violating parameter $A(\alpha)$ obtained using (6), (7) and (32) increases with the effective no. of colors as seen from Table 1.

In this work we have studied the direct CP violation in $\Lambda_b \to p\pi^-$ decay mode. Using the next-to-leading order QCD corrected effective Hamiltonian, we have obtained the branching ratio and CP asymmetries within the framework of generalized factorization. The nonfactorizable contributions are parametrized in terms of the effective no. of colors N_c^{eff}. So in addition to the naive factorization approach ($N_c^{eff} = 3$), here we have taken two more values for N_c^{eff} i.e., $N_c^{eff} = 2$ and ∞. The baryonic form factors at maximum momentum transfer (q_m^2) are evaluated using the nonrelativistic quark model and the extrapolation of the form factors from q_m^2 to the required q^2 value is done by assuming the pole dominance. The weak phases in our analysis are obtained from the CKM matrix elements whereas the strong phases are obtained from the absorptive part of the penguin diagrams. The obtained branching ratio lies within the present experimental upper limit and the CP violating observables are found to be $a_{cp} \sim -8\%$ and $A(\alpha) \simeq 2 \times 10^{-5}$. In the future there will be more data on the heavy Λ_b baryon from different experimental groups, hence it will be very interesting to look for such CP violating asymmetries in the experiments in order to get a deeper understanding of the mechanism of CP violation. Furthermore, the study of CP violation in Λ_b decays may provide insight into the baryon asymmetry phenomena required for baryogenesis.

Acknowledgements. The author would like to thank Professor M. P. Khanna and Dr. A. K. Giri for many useful discussions and also to CSIR, Govt. of India, for financial support.

References

1. I. I. Bigi, V. A. Khoze, N. G. Uraltsev, A. I. Sanda, in CP Violation, ed. C. Jarlskog (World Scientific, Singapore, 1989), p.175; I. Dunietz, Ann. Phys. **184**, 350 (1988) and references there in
2. For a recent review, see: R. Fleischer, Int. J. Mod. Phys. A **12**, 2459 (1997) and references there in
3. M. Bander, D. Silverman, A. Soni, Phys. Rev. Lett. **43**, 242 (1979); J. M. Gérard, W. S. Hou, Phys. Rev. D **43**, 2909 (1991)
4. J. F. Donoghue, E. Golowich, A. A. Petrov, J. M. Soares, Phys. Rev. Lett. **77**, 2178 (1996)
5. CDF Collaboration, F. Abe et al, Phys. Rev. D **55**, 1142 (1997): UA1 Collaboration, C. Albarjar et al., Phys. Lett. B **273**, 540 (1991): CDF Collaboration, F. Abe et al., Phys. Rev. D **47**, 2639 (1993); P. Abreu et al., Phys. Lett. B **374**, 351 (1996)
6. J. F. Donoghue, S. Pakvasa, Phys. Rev. Lett **55**, 162 (1985); J. F. Donoghue, X-G. He, S. Pakvasa, Phys. Rev. D **34**, 833 (1986); D. Chang, X-G. He, S. Pakvasa, Phys. Rev. Lett. **74**, 3927 (1995); X-G. He, H. Steger, G. Valencia, Phys. Lett. B **272**, 411 (1991); X-G. He, G. Valencia, Phys. Rev. D **52**, 5257 (1995); J. Tandean, G. Valencia, Phys. Lett. **451**, 382 (1999)
7. H. Y. Cheng, B. Tseng, Phys. Rev. D **53**, 1457 (1996)
8. G. Buchalla, A.J. Buras, M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996)
9. Y. H. Chen, H. Y. Cheng, B. Tseng, K. C. Yang, Phys. Rev. D **60**, 094014 (1999); A. Ali, G. Kramer, C. D. Lü, Phys. Rev. D **58**, 094009 (1998); A. Ali, C. Greub, Phys. Rev. D, **57**, 2996 (1998)
10. A. J. Buras et al, Nucl. Phys. B **370**, 69 (1992); M. Ciuchini et al, Zeit. Phys. C **68**, 239 (1995)
11. Y. H. Chen, H. Y. Cheng, B. Tseng, K. C. Yang, Phys. Rev. D **60**, 094014 (1999)
12. N. G. Deshpande, J. Trampetic, Phys. Rev. D **41**, 2926 (1990)
13. H. Fusaoku, Y. Koide, Phys. Rev. D **57**, 3986 (1998)
14. G. Kramer, W. F. Palmer, H. Simma, Zeit. Phys. C **66**, 429 (1995)
15. C. Caso et al, Euro. Phys. J. C **1**, 1 (1998)
16. H. Y. Cheng, K. C. Yang, Phys. Rev. D **59**, 092004, (1999)
17. T. Mannel, W. Roberts, Z. Phys. **59**, 179 (1993)