Possibility of extracting the weak phase $\gamma$ from $\Lambda_b \to \Lambda D^0$ decays

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Abstract

We explore the possibility of extracting the weak phase $\gamma$ from pure tree decays $\Lambda_b \to \Lambda (D^0, \overline{D}^0, D^0_{CP})$ in a model independent way. The CP violating weak phase $\gamma$ can be determined cleanly, without any hadronic uncertainties, as these decay modes are free from the penguin pollutions. Furthermore, neither tagging nor time dependent studies are required to extract the angle $\gamma$ with these modes.
I. INTRODUCTION

CP violation still remains one of the unsolved problems till date in particle physics [1–3], even after almost four decades of its discovery in 1964 in neutral K meson system. Since then various attempts have been made by theorists and experimentalists to understand it but without much success. The Standard Model (SM) provides a simple description of this phenomenon through the complex CKM matrix [4]. Decays of B mesons provide a rich ground for investigating CP violation [4,5]. They allow stringent tests both for the SM and for studies of new sources of this effect. Within the SM, CP violation is often characterized by the so called unitarity triangle [7]. Detection of CP violation and the accurate determination of the unitarity triangle are the major goals of experimental B physics [8]. Although the exact cause of its origin is not yet understood completely but the recent results from SLAC and KEK B-factories, the first evidence of large CP violation in B-systems [9,10], have necessitated renewed interests in understanding the nature of CP violation within the framework of Standard Model (SM), instead of looking beyond it. Therefore, it is really high time in particle physics in general and B-physics in particular since CP violation is interlinked to many problems in particle physics. It can also give us a possible explanation for the baryon asymmetry of the universe. The general expectation is to check all possible ways to find out the root cause of it and explore as many decay modes as possible to arrive at a decisive conclusion. We now believe that the CKM explanation of SM (which explains the CP violation in K-systems), may also explain the mechanism of CP violation in B-systems. The knowledge gained in exploring various scenarios in B-systems within the SM will give us enough hints regarding the possible structure of CP violation in and beyond the SM, and may possibly also help us to narrow down our searches in that.

In Standard Model, the phenomenon of CP violation can be established if we can measure accurately the three angles ($\alpha \equiv \phi_2$, $\beta \equiv \phi_1$ and $\gamma \equiv \phi_3$) of the CKM unitarity triangle, which add up to 180°. The angle $\beta$ is the simplest one and can be measured in the gold plated mode $B \to J/\psi K_S$ without any uncertainties [11,12]. In fact the value of $\sin 2\beta$ has recently been reported [11,12]. The angle $\alpha$ can be measured from the decay mode $B \to \pi\pi$ [11,12]. Although the presence of penguin contribution introduces some uncertainties, but with isospin techniques [13] the penguin contamination can be disentangled so that $\alpha$ can be extracted without hadronic uncertainties. However, this analysis requires the measurement of $B^0 \to \pi^0\pi^0$, which is not feasible with first generation B-factory experiments. Alternatively, $\alpha$ can be determined by applying the isospin technique to $B \to \rho\pi$ decays [14]. The most difficult to measure is the angle $\gamma$. To this end, various interesting proposals have been made with a view to obtain $\gamma$ with lesser or no hadronic uncertainties and the search is going on to find out a gold plated mode from which one can extract $\gamma$ cleanly. It has been shown by Gronau, London and Wyler (GLW) [13] that $\gamma$ can be extracted from $B \to DK$ decay rates, employing amplitude triangle relations. Later, modification with different final states for these decay modes has been put forward by Atwood, Dunietz and Soni (ADS) [16] so that one can hope to find $\gamma$ without uncertainties from the above decay modes. In Ref. [17] the self tagging modes $B_s^0 \to D^0 K^{*0}$ have been considered for the extraction of $\gamma$. It has been discussed in Ref. [18] that $\gamma$ can also be determined cleanly from $B_s \to D_s^\pm K$ decays. The triangle approach was also extended to the $B_s$ system [13] to extract $\gamma$. It is shown that the $SU(3)$ relations [20] can be employed between $B \to K\pi, \pi\pi$ decay amplitudes for
the extraction of $\gamma$ and a nontrivial bound for $\gamma$ can be obtained. It is also argued that the phase $\gamma$ can be extracted from the $B_c \to DD_s$ decays \cite{21}, without any hadronic uncertainties, since these are pure tree decay modes and are free from penguin pollutions. We have recently studied another possibility to find the angle $\gamma$ cleanly from $B^0 \to \bar{D}^0\phi$ decay modes \cite{22}.

As emphasized above the determination of the angle $\gamma$ is most challenging, and therefore, it requires further scrutiny. It should be noted here that the precise value of $\gamma$ will play a crucial role for testing the validity of CKM model of CP violation in SM, since the standard technique of the unitarity triangle will largely depend on it. Hence, the goal is to check all possible decay modes and try to determine the angle $\gamma$ as cleanly as possible. In this paper we would like to show that the angle $\gamma$ can also be extracted from the pure tree decays of the $\Lambda_b$ baryon i.e, $\Lambda_b \to \Lambda\{D^0, D^{0*}, D_{CP}^0\}$. The advantage of these decay modes is that these are free from penguin pollutions and the amplitudes are of similar sizes. Furthermore, neither tagging nor time dependent studies are required for these decay modes, so $\gamma$ can be extracted cleanly without hadronic uncertainties. CP violation in the bottom baryon case has earlier been studied by Dunietz \cite{23}. One of us has recently studied the decay mode $\Lambda_b \to p\pi$ in and beyond SM \cite{24}. Here we would like to emphasize that the bottom baryon decay modes may serve as alternatives and/or may supplement the bottom meson decays for the study of CP violation. Although the branching ratios of $\Lambda_b$ decay modes are smaller in comparison to those of the $B$ counterpart modes, but there are certain advantages for the former case over the later. The usual problem with the bottom meson (for mixing induced CP violation) is the tagging and the time evolution of the decaying particle are not required for bottom baryon case. In Ref. \cite{23} it has been shown by Dunietz that the angle $\gamma$ can be determined from various baryonic decay modes. He has analyzed in detail the decay modes $\Lambda_b \to \Lambda\{D^0, D^{0*}, D_{CP}^0\}$ and the corresponding charge conjugate modes by considering three specific cases : i. the $p$-wave dominance, ii. the $s$-wave dominance iii. the $p$- to $s$-wave ratio to be constant, and has shown that $\gamma$ can be determined from the partial decay rates of the above six processes. In this paper we would like to explore the possibilities of extracting the angle $\gamma$, without assuming the dominance of either $s$-wave or $p$-wave. For this purpose we need only the experimental observables (i.e., the decay rate and angular distribution parameters) for the above six processes and the information on $\gamma$ can be extracted from these observables. The plan of the paper is as follows. In section II, we present the phenomenology of hyperon decays. The formalism of the extraction of the angle $\gamma$ from the decay modes $\Lambda_b \to \Lambda\{D^0, D^{0*}, D_{CP}^0\}$ is presented in Section III. Section IV contains our conclusion.

**II. PHENOMENOLOGY OF HYPERON DECAYS**

The study of CP violation in strange hyperon decays has been extensively studied in Ref. \cite{25}, where the phenomenology of hyperon decays has been discussed in great detail. However for the sake of completeness we shall present here the basic features of the nonleptonic hyperon decays. The most general Lorentz-invariant amplitude for the decay $\Lambda_b \to \Lambda D^0$ can be written as

$$\text{Amp}(\Lambda_b(p_i) \to \Lambda(p_f)D^0(q)) = i\bar{u}_\Lambda(p_f)(A + B\gamma_5)u_{\Lambda_b}(p_i), \quad (1)$$

3
where \( u_\Lambda \) and \( u_{\Lambda_b} \) are the Dirac spinors for \( \Lambda \) and \( \Lambda_b \) baryons respectively. The parameters \( A \) and \( B \) are complex in general. The matrix element for the corresponding CP conjugate process \( \bar{\Lambda}_b \rightarrow \bar{\Lambda} D^0 \) is given as

\[
\text{Amp}(\bar{\Lambda}_b(p_i) \rightarrow \bar{\Lambda}(p_f)D^0(q)) = i\bar{v}_\Lambda(p_f)(-A^* + B^*\gamma_5)v_{\Lambda_b}(p_i).
\] (2)

The matrix element for the corresponding CP conjugate process \( \bar{\Lambda}_b \rightarrow \bar{\Lambda} D^0 \) is given as

\[
\text{Amp}(\bar{\Lambda}_b(p_i) \rightarrow \bar{\Lambda}(p_f)D^0(q)) = \chi^\dagger_{\Lambda_b}(S + P\sigma \cdot \hat{q})\chi_{\Lambda},
\] (3)

where \( \chi_{\Lambda_b} \) and \( \chi_{\Lambda} \) are the Pauli spinors, \( \sigma \) are the Pauli spin matrices and \( q \) is the c.o.m. momentum of the final particles in the rest frame of \( \Lambda_b \) baryon. The amplitudes \( S \) and \( P \) are given as

\[
S = A\sqrt{\frac{\{(m_{\Lambda_b} + m_\Lambda)^2 - m_D^2\}}{16\pi m_{\Lambda_b}^2}}
\]

\[
P = B\sqrt{\frac{\{(m_{\Lambda_b} - m_\Lambda)^2 - m_D^2\}}{16\pi m_{\Lambda_b}^2}}.
\] (4)

The experimental observables are the total decay rate \( \Gamma \) and the asymmetry parameters \( \alpha' \), \( \beta' \) and \( \gamma' \) which govern the decay-angular distribution and the polarization of the final baryon. The decay rate is given as

\[
\Gamma = 2|q|\{|S|^2 + |P|^2\}
\] (5)

and the asymmetry parameters are given as

\[
\alpha' = \frac{2\text{Re}(S^*P)}{|S|^2 + |P|^2}
\]

\[
\beta' = \frac{2\text{Im}(S^*P)}{|S|^2 + |P|^2}
\]

\[
\gamma' = \frac{|S|^2 - |P|^2}{|S|^2 + |P|^2}.
\] (6)

However, these three angular distribution parameters are not independent, and are related as

\[
\alpha'^2 + \beta'^2 + \gamma'^2 = 1.
\] (7)

Thus there are three independent observables (\( \Gamma \) and any two of \( \alpha' \), \( \beta' \) and \( \gamma' \)). Likewise, the observables for the antihyperon decays (\( \bar{\Gamma}, \bar{\alpha}', \bar{\beta}' \) and \( \bar{\gamma}' \)) can be written similar to Eqs. (4) and (5) with \( S \) and \( P \) replaced by \( \bar{S} \) and \( \bar{P} \) respectively.

With Eqs. (4) and (5), we can alternatively obtain the following three observables for nonleptonic hyperon decays.
\[ |S|^2 = \frac{\Gamma}{4|q|}(1 + \gamma') \]
\[ |P|^2 = \frac{\Gamma}{4|q|}(1 - \gamma') \]
\[
\tan \Delta = \frac{\beta'}{\alpha'} \tag{8}
\]
where \( \Delta \) is the relative strong phase between the \( S \) and \( P \) wave amplitudes. After knowing the phenomenology of hyperon decays, we now proceed to extract the weak phase \( \gamma \) from the decay modes \( \Lambda_b \to \Lambda \{D^0, \bar{D}^0, D^0_{CP}\} \).

### III. EXTRACTION OF THE ANGLE \( \gamma \)

With the advent of hadronic \( b \) facilities it becomes possible to produce \( \Lambda_b \) baryons in large numbers. In this paper we show that the baryonic counterpart of the \( B_d^0 \to D^0 K_s \) decay modes could be well suited to extract the CKM angle \( \gamma \). The corresponding \( \Lambda_b \) decays are \( \Lambda_b \to \Lambda \{D^0, \bar{D}^0, D^0_{CP}\} \), where \( D^0_{CP} \) is the CP eigenstate of neutral \( D \) meson. Let us now write the decay amplitudes for the above processes. Both these amplitudes proceed via the colour suppressed tree diagrams only and hence are free from penguin pollutions. The amplitude for \( \Lambda_b \to \Lambda D^0 \) arises from the quark level transition \( b \to e\bar{u}s \) and hence has no weak phase in the Wolfenstein parametrization, while the amplitude \( \Lambda_b \to \Lambda \bar{D}^0 \) arises from \( b \to u\bar{c}s \) transition and carries the weak phase \( e^{-i\gamma} \). The amplitudes also have the strong phases \( e^{i\delta_1} \) and \( e^{i\delta_2} \), where \( i = S \) or \( P \). Thus we can write the decay amplitude for \( \Lambda_b \to \Lambda D^0 \) process as (here and henceforth we will not write explicitly the spinors \( \chi_{\Lambda_b} \) and \( \chi_{\Lambda} \) in the decay amplitudes)

\[
A_1 = \text{Amp}(\Lambda_b \to \Lambda D^0) = S_1 e^{i\delta_S^0} + P_1 e^{i\delta_P^0} \sigma \cdot \hat{q} = e^{i\delta_S^0} \left( S_1 + P_1 e^{i\Delta_1} \sigma \cdot \hat{q} \right), \tag{9}
\]

where \( S_1 \) and \( P_1 \) are magnitudes of the \( S \) and \( P \) wave amplitudes and \( \Delta_1 = \delta_P^0 - \delta_S^0 \), is the relative strong phase between them. From the decay mode \( \Lambda_b \to \Lambda D^0 \), one can extract the three observables \( S_1, P_1 \) and \( \Delta \) using Eq. \( \text{(8)} \). Similarly, one can write the decay amplitude for the process \( \Lambda_b \to \Lambda \bar{D}^0 \) as

\[
A_2 = \text{Amp}(\Lambda_b \to \Lambda \bar{D}^0) = e^{-i\gamma} \left( S_2 e^{i\delta_S^0} + P_2 e^{i\delta_P^0} \sigma \cdot \hat{q} \right) = e^{-i\gamma} e^{i\delta_S^0} \left( S_2 + P_2 e^{i\Delta_2} \sigma \cdot \hat{q} \right). \tag{10}
\]

Thus from this decay mode we can extract another set of three observables \( S_2, P_2 \) and \( \Delta_2 \). Now, the amplitudes for the corresponding CP conjugate processes are given as

\[
\bar{A}_1 = \text{Amp}(\bar{\Lambda}_b \to \bar{\Lambda} \bar{D}^0) = \bar{S}_1 e^{i\delta_S^0} + \bar{P}_1 e^{i\delta_P^0} \sigma \cdot \hat{q} = e^{i\delta_S^0} \left( \bar{S}_1 + \bar{P}_1 e^{i\Delta_1} \sigma \cdot \hat{q} \right),
\]
\[
\bar{A}_2 = \text{Amp}(\bar{\Lambda}_b \to \bar{\Lambda} D^0) = e^{i\gamma} \left( \bar{S}_2 e^{i\delta_S^0} + \bar{P}_2 e^{i\delta_P^0} \sigma \cdot \hat{q} \right) = e^{i\gamma} e^{i\delta_S^0} \left( \bar{S}_2 + \bar{P}_2 e^{i\Delta_2} \sigma \cdot \hat{q} \right), \tag{11}
\]

where \( \bar{S}_{1,2} = -S_{1,2} \) and \( \bar{P}_{1,2} = P_{1,2} \). Considering these two above decay modes, the observables \( \bar{S}_{1,2}, \bar{P}_{1,2} \) and \( \Delta_{1,2} \) can be determined.

We now consider the decay modes \( \Lambda_b \to \Lambda D^0 \), where \( D^0 \) denote the neutral \( D \) meson even/odd \( CP \) states, defined as \( D^0_\pm = (D^0 \pm \bar{D}^0)/\sqrt{2} \). The \( CP \) even state \( D^0_+ \) can be
identified by the CP even decay products such as $\pi^+\pi^-$ and $K^+K^-$, whereas the CP odd state $D_0^-$ can be identified by the CP odd products such as $K_S\pi^0$, $K_S\rho^0$, $K_S\omega$ and $K_S\phi$. One can use either of these two CP eigenstates for the extraction of $\gamma$. Here we are considering the CP even eigenstate ($D_0^+$), however, the same argument will also hold for the CP odd state ($D_0^-$).

The amplitude for $\Lambda_b \to \Lambda D_0^+$ is thus given as

$$A_+ = \text{Amp}(\Lambda_b \to \Lambda D_0^+) = S_+ e^{i\delta_+^S} + P_+ e^{i\delta_+^P} \sigma \cdot \hat{q} = e^{i\delta_+^S} \left( S_+ + P_+ e^{i\Delta_+} \sigma \cdot \hat{q} \right),$$

(12)

where $S_+$, $P_+$ are the magnitudes of the $S$ and $P$ wave amplitudes with phases $e^{i\delta_+^S}$ and $e^{i\delta_+^P}$. It should be noted that these phases contain both strong and weak components. Thus from this decay mode we can extract the observables $S_+$, $P_+$ and $\Delta_+ = \delta_+^P - \delta_+^S$. We can write the amplitude for this decay mode in another form, i.e.,

$$A_+ = \frac{1}{\sqrt{2}} \left[ \text{Amp}(\Lambda_b \to \Lambda D_0^0) + \text{Amp}(\Lambda_b \to \Lambda D_0^0) \right]$$

$$= \frac{1}{\sqrt{2}} \left[ e^{i\delta_+^S} \left( S_1 + S_2 e^{i(\sigma_+^S - \gamma)} \right) + e^{i\delta_+^P} \left( P_1 + P_2 e^{i(\sigma_+^P - \gamma)} \right) \sigma \cdot \hat{q} \right],$$

(13)

where $\sigma_+^{S,P} = \delta_+^{S,P} - \delta_+^S$. Thus comparing Eqs. (12) and (13) we obtain

$$S_+ e^{i\delta_+^S} = \frac{1}{\sqrt{2}} e^{i\delta_+^P} \left( S_1 + S_2 e^{i(\sigma_+^S - \gamma)} \right),$$

$$P_+ e^{i\delta_+^P} = \frac{1}{\sqrt{2}} e^{i\delta_+^S} \left( P_1 + P_2 e^{i(\sigma_+^P - \gamma)} \right).$$

(14)

Now the amplitude for the corresponding CP conjugate process, i.e., $\Lambda_b \to \Lambda D_0^0$ is given as

$$\bar{A}_+ = \text{Amp}(\Lambda_b \to \Lambda D_0^0) = \frac{1}{\sqrt{2}} \left[ \bar{A}_1 + \bar{A}_2 \right]$$

$$= \bar{S}_+ e^{i\delta_+^S} + \bar{P}_+ e^{i\delta_+^P} \sigma \cdot \hat{q} = e^{i\delta_+^S} \left( \bar{S}_+ + \bar{P}_+ e^{i\Delta_+} \sigma \cdot \hat{q} \right),$$

(15)

where $\bar{S}_+$, $\bar{P}_+$ are the magnitudes of the $S$ and $P$ wave amplitudes with phases $e^{i\delta_+^S}$ and $e^{i\delta_+^P}$. The observables obtained from this decay modes are $\bar{S}_+$, $\bar{P}_+$ and $\bar{\Delta}_+ = \delta_+^P - \delta_+^S$. Substituting the values of $A_1$ and $A_2$ from Eq. (12) in Eq. (15) we obtain the relations similar to Eq. (14) as

$$\bar{S}_+ e^{i\delta_+^S} = \frac{1}{\sqrt{2}} e^{i\delta_+^P} \left( \bar{S}_1 + \bar{S}_2 e^{i(\sigma_+^S + \gamma)} \right),$$

$$\bar{P}_+ e^{i\delta_+^P} = \frac{1}{\sqrt{2}} e^{i\delta_+^S} \left( \bar{P}_1 + \bar{P}_2 e^{i(\sigma_+^P + \gamma)} \right).$$

(16)

We now use Eqs. (14) and (16) to obtain the weak phase $\gamma$. To derive the analytic expression for $\gamma$, we define the combinations of observables

$$X = \frac{2S_2^2 - S_1^2 - S_2^2}{2 S_1 S_2},$$

$$\bar{X} = \frac{2\bar{S}_2^2 - \bar{S}_1^2 - \bar{S}_2^2}{2 \bar{S}_1 \bar{S}_2}$$

(17)
Here we have considered only the $S$ wave components, but similar combinations can also be derived from the $P$ wave observables and one can use either set, for the extraction of $\gamma$. Thus, in this method we need to know only the magnitudes of $S$ and $P$ waves but not the phase difference between them.

It is now very simple to see that one can obtain an expression for $\gamma$ from Eq. (17) as

$$2\gamma = \arccos \bar{X} - \arccos X,$$

with some discrete ambiguities. One can also obtain the value of $\sin^2 \gamma$ from Eq. (17) via the relation

$$\sin^2 \gamma = \frac{1}{2} \left[1 - X \bar{X} \pm \sqrt{(1 - X^2)(1 - \bar{X}^2)}\right].$$

One of the sign in Eq. (19) will give the correct value of $\sin^2 \gamma$, while the other will give the value of the strong phase $\sin^2 \sigma_S^+$. One can use both $S$ and $P$ wave observables so that some of the ambiguities will be reduced and $\gamma$ can be determined cleanly.

Now let us find out the values of the branching ratios for the above decay modes. The decay processes $\Lambda_b \to \Lambda D^0$ and $\Lambda_b \to \Lambda \bar{D}^0$ are governed by the quark level transitions $b \to c \bar{u}s$ and $b \to u \bar{c}s$ respectively. The effective Hamiltonians for such transitions are given as

$$H_{\text{eff}}(b \to c \bar{u}s) = \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* \left[C_1(m_b)(\bar{c}b)(\bar{s}u) + C_2(m_b)(\bar{c}u)(\bar{s}b)\right]$$

and

$$H_{\text{eff}}(b \to u \bar{c}s) = \frac{G_F}{\sqrt{2}} V_{ub} V_{cs}^* \left[C_1(m_b)(\bar{u}b)(\bar{s}c) + C_2(m_b)(\bar{u}c)(\bar{s}b)\right]$$

respectively, where $C_1(m_b) = 1.13$ and $C_2(m_b) = -0.29$ are the Wilson coefficients evaluated at the renormalization scale $m_b$. $(\bar{c}b) = \bar{c}\gamma^\mu(1 - \gamma_5)b$, etc. are the usual left handed color singlet quark currents.

The hadronic matrix elements of the four-fermion current operators present in the effective Hamiltonian (20) and (21) are very difficult to evaluate from the first principle of QCD. So to evaluate the matrix elements of the effective Hamiltonian, we use the generalized factorization approximation. In this approximation, the hadronic matrix elements of the four quark operators split into the product of two matrix elements, one describing the transition of the initial baryon to the final baryon state and the other describing the formation of the final meson from vacuum. Furthermore, the nonfactorizable contributions, which play an important role in colour suppressed decays, are incorporated in a phenomenological way: they are lumped into the coefficients $a_1 = C_1 + C_2/N_c$ and $a_2 = C_2 + C_1/N_c$, so that the effective coefficients $a_1^{\text{eff}}$ and $a_2^{\text{eff}}$ are treated as free parameters and their values can be extracted from the experimental data. In this paper we shall denote $a_2^{\text{eff}}$, which describes the color suppressed decays, by simply $a_2$. Thus the factorized amplitude for the decay process $\Lambda_b \to \Lambda D^0$ is given as

$$\text{Amp}(\Lambda_b \to \Lambda D^0) = \frac{G_F}{\sqrt{2}} V_{cb} V_{us}^* a_2 |D^0\rangle (\bar{c}u) |0\rangle \langle \Lambda | (\bar{s}b) |\Lambda_b\rangle.$$
The general expression for the baryon transition is given as

\[
\langle \Lambda(p_f)|V_\mu - A_\mu|\Lambda_b(p_i)\rangle = \bar{u}_\Lambda(p_f)\left\{f_1(q^2)\gamma_\mu + if_2(q^2)\sigma_{\mu\nu}q^\nu + f_3(q^2)q_\mu
\right.

- \left[g_1(q^2)\gamma_\mu + ig_2(q^2)\sigma_{\mu\nu}q^\nu + g_3(q^2)q_\mu\gamma_5\right\}u_{\Lambda_b}(p_i),
\]

where \(q = p_i - p_f\). In order to evaluate the form factors at maximum momentum transfer, we have employed nonrelativistic quark model [23], where they are given as:

\[
\begin{align*}
\frac{f_1(q^2)}{N_{f_i}} &= 1 - \frac{\Delta m}{2m_{\Lambda_b}} + \frac{\Delta m}{4m_{\Lambda_b}m_s} \left(1 - \frac{L}{2m_A}\right)(m_{\Lambda_b} + m_\Lambda - \Delta m) \\
&\quad - \frac{\Delta m}{8m_{\Lambda_b}m_A m_b}(m_{\Lambda_b} + m_\Lambda - \Delta m) \\
\frac{f_3(q^2)}{N_{f_i}} &= \frac{1}{2m_{\Lambda_b}} - \frac{1}{4m_{\Lambda_b}m_s} \left(1 - \frac{L}{2m_A}\right)(m_{\Lambda_b} + m_\Lambda - \Delta m) \\
&\quad - \frac{L}{8m_{\Lambda_b}m_A m_b}[(m_{\Lambda_b} + m_\Lambda) + \Delta m] \\
\frac{g_1(q^2)}{N_{f_i}} &= 1 + \frac{\Delta mL}{4} \left(\frac{1}{m_{\Lambda_b}m_s} - \frac{1}{m_A m_b}\right) \\
\frac{g_3(q^2)}{N_{f_i}} &= -\frac{L}{4} \left(\frac{1}{m_{\Lambda_b}m_s} + \frac{1}{m_A m_b}\right)
\end{align*}
\]

where \(L = m_\Lambda - m_s\), \(\Delta m = m_{\Lambda_b} - m_\Lambda\), \(q^2 = \Delta m^2\), \(m_b\) and \(m_s\) are the constituent quark masses of the interacting quarks of initial and final baryons with values \(m_s = 510\) MeV and \(m_b = 5\) GeV. \(N_{f_i}\) is the flavour factor:

\[
N_{f_i} = \text{flavor spin} \langle \Lambda|b_s^\dagger b_b|\Lambda_b\rangle_{\text{flavor spin}} = \frac{1}{\sqrt{3}}
\]

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

\[
\begin{align*}
f(q^2) &= \frac{f(0)}{(1 - q^2/m_V^2)^2} & g(q^2) &= \frac{g(0)}{(1 - q^2/m_A^2)^2}
\end{align*}
\]

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 \(m_V = 5.42\) GeV and \(m_A = 5.86\) GeV. Assuming a dipole \(q^2\) behavior for form factors, and taking the masses of the baryons and \(D^0\) meson from Ref. [27] we found

\[
\begin{align*}
f_1(m_D^2) &= 0.08, & m_{\Lambda_b}f_3(m_D^2) &= -0.01, & g_1(m_D^2) &= 0.13 & m_{\Lambda_b}g_3(m_D^2) &= -0.05.
\end{align*}
\]

The matrix element \(\langle D^0|A_\mu|0\rangle\) is related to the \(D^0\) meson decay constant \(f_D\) as

\[
\langle D^0(q)|A_\mu|0\rangle = -if_D q_\mu
\]
Hence we obtain the transition amplitude for $\Lambda_b \to \Lambda D^0$ as

$$A(\Lambda_b \to \Lambda D^0) = \frac{G_F}{\sqrt{2}} V_{cb} V_{ub}^* i f_D a_2 \bar{u}_\Lambda(p_f) \left[ (f_1(m_D^2)(m_{\Lambda_b} - m_\Lambda) + f_3(m_D^2)m_D^2) + \left( g_1(m_D^2)(m_{\Lambda_b} + m_\Lambda) - g_3(m_D^2)m_D^2 \right) \gamma_5 \right] u_{\Lambda_b}(p_i) .$$

Comparing the Eqs. (29) and (30) we obtain the values of the parameters $A$ and $B$ as

$$A = \frac{G_F}{\sqrt{2}} V_{cb} V_{ub}^* a_2 f_D \left[ f_1(m_D^2)(m_{\Lambda_b} - m_\Lambda) + f_3(m_D^2)m_D^2 \right]$$
$$B = \frac{G_F}{\sqrt{2}} V_{cb} V_{ub}^* a_2 f_D \left[ g_1(m_D^2)(m_{\Lambda_b} + m_\Lambda) - g_3(m_D^2)m_D^2 \right]$$

Thus with Eqs. (30), (4) and (5) the branching ratio for the process $\Lambda_b \to \Lambda D^0$ is given as

$$BR(\Lambda_b \to \Lambda D^0) = 4.56 \times 10^{-6} ,$$

where we have used $a_2 = 0.3$, $f_D = 300$ MeV [27] and the values of the CKM matrix elements from Ref. [27]. The branching ratio for the decay mode is $\Lambda_b \to \Lambda D^0$ is also found with Eqs. (30), (4) and (5) and by substituting the appropriate CKM matrix elements

$$BR(\Lambda_b \to \Lambda D^0) = 8.29 \times 10^{-7} .$$

Although the branching ratios for these decay modes are quite small, they could be observed in the future $B$ experiments.

Now let us study the experimental feasibility of such decay modes. The observables are the partial decay rate ($\Gamma$) and the angular distribution parameters $\alpha'$, $\beta'$ and $\gamma'$, which characterize the strength of $S$ and $P$ waves. Measuring the partial rate requires no polarization while the measurement of $\alpha'$, $\beta'$ and $\gamma'$ do. In processes where the final baryon subsequently decays as $\Lambda_b \to \Lambda(\to p\pi^-)D^0$, $\alpha'$ is determined from the decay distribution of the final baryon $\Lambda \to p\pi^-$. While the observables $\beta'$ and $\gamma'$ require an initial polarization. Here, we would like to point out that because only a few percent of the $D^0$ modes are $CP$ eigenstates, it might be difficult to extract the angular distribution parameters for $\Lambda_b \to \Lambda D^0_\pm$ modes.

The BTeV experiment, with luminosity $2 \times 10^{32} \text{cm}^{-2}\text{s}^{-1}$, will produce $2 \times 10^{11}$ $b\bar{b}$ hadrons per $10^7$ sec of running [28]. If we assume the production fraction as

$$\mathcal{B}_d : B^- : \mathcal{B}_s : \Lambda_b = 0.375 : 0.375 : 0.15 : 10$$

we get around $2 \times 10^{10}$ numbers of $\Lambda_b$ baryon per year of running at BTeV. If we take the branching ratios as : $BR(\Lambda_b \to \Lambda D^0) \sim 4.5 \times 10^{-6}$, $BR(D^0 \to K^-\pi^+)$ and $K^-\pi^+\pi^-\pi^+$ = 0.12, the reconstruction efficiency as 0.05 and the trigger efficiency level as 0.9, we expect to get 486 $\Lambda D^0$ events per year.

### IV. CONCLUSION

In this paper, we have analyzed the possibility of extracting the weak phase $\gamma$ from the decay modes $\Lambda_b \to \Lambda \{D^0, D^0, D^0_{CP}\}$. The transition amplitudes and branching ratios of
these decay processes are calculated using the effective Hamiltonian and the generalized factorization approximation. The advantage of these modes are that, first, they are described by the color suppressed tree diagrams only and are free from penguin contaminations. Second, neither tagging nor time dependent studies are required for these decay modes and hence $\gamma$ can be determined cleanly without any hadronic uncertainties. The price one has to pay is that the branching ratios for these processes are to be very small, of the order $(10^{-6} - 10^{-7})$, which are one or two order smaller than those of the corresponding $B$ decays. Nevertheless, the absence of usual difficulties, as in the case of bottom mesons, make such modes worthy of careful study. To do so, we have to wait for the future high statistics $B$ experiments, where large samples of $\Lambda_b$ events are expected to be available.

To summarize, in this paper we have shown that the decay modes $\Lambda_b \rightarrow \Lambda \{D^0, \bar{D}^0, D^{0}_{CP}\}$ appear to be ideally suited for the clean determination of the angle $\gamma$. Here we have considered only the Standard Model contributions to the decay processes. So if the extracted value of $\gamma$, from these modes differs from its value constrained by SM, then this would be a clear indication of the possible existence of new physics. Therefore, outside the $B$ meson system, these decay modes may possibly give us valuable information regarding the nature of CP violation and guide us to know physics beyond Standard Model. The strategies presented here are particularly interesting for future $B$ experiments such as BTeV, LHCb and beyond.

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