Two-body hadronic decays $\Lambda_b(\frac{1}{2})^+ \to B^*(\frac{3}{2}) + P$ in a quark model

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The framework under which decays $\Lambda_b(\frac{1}{2}^+) \to B(\frac{1}{2}^+) + M$ are analyzed is not applicable for the decays $\Lambda_b(\frac{1}{2}^+) \to B^*(\frac{3}{2}^+) + M$. These decays occur through a baryon pole Σ_c^0 which is generated by the W-exchange diagram in the process $b + u \xrightarrow{W} c + d$. The effective Hamiltonian which arises from the W-exchange diagram is expressed in the nonrelativistic limit. Since Σ_c^0 belongs to representation 6 of SU(3), it contributes to two sets of decays: $\Lambda_b \to \Delta^0 D^0, \Delta^- D^+, \Sigma^{*-} D_s^+$ and $\Lambda_b \to \Sigma_c^{*-}\pi^+, \Sigma_c^{*0}\pi^0, \Xi_c^{*0} K^0$. The branching ratios for these decays are evaluated and can be compared with their experimental values when the data become available. Another prediction of the model is that asymmetry parameter $\alpha = 0$ since the baryon pole contributes to the parity conserving (p-wave) amplitude and does not contribute to the parity violating (d-wave) amplitude.

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I. INTRODUCTION

In the standard model, two-body hadronic decays of heavy flavor mesons and baryons are analyzed in terms of the effective Lagrangian or Hamiltonian [1,2,3],

$$H_{\mathrm{eff}} \equiv V_{cb}V^*_{qq'}\left[\left(C_1 + \frac{1}{3}C_2\right)(\bar{q}'^{\beta} q_\beta)_{V-A}(\bar{c}^\alpha b_\alpha)_{V-A} + \left(C_2 + \frac{1}{3}C_1\right)(\bar{c}^\beta q_\beta)_{V-A}(\bar{q}'^\alpha b_\alpha)_{V-A}\right], \quad (1)$$

where $q = c, q' = s$ or $q = u, q' = d$. The above Hamiltonian corresponds to decays which are not Cabbibo suppressed.

The effective Hamiltonian arises from the transition $b \to c + s + \bar{c}$ or $b \to c + d + \bar{u}$. The short distance QCD effects are incorporated in the Wilson coefficients C_1 and C_2. In the factorization ansatz, long-distance strong interaction effects are shifted to the evaluation of the baryon form factors in some models.

Finally, the effective Hamiltonian is written in the form

$$H_{\mathrm{eff}} = V_{cb}V^*_{cs}[a_1(\bar{s}c)_{V-A}(\bar{c}b)_{V-A} + a_2(\bar{c}c)_{V-A}(\bar{s}b)_{V-A}] \quad (2)$$

$$H_{\mathrm{eff}} = V_{cb}V^*_{ud}[a_1(\bar{d}u)_{V-A}(\bar{c}b)_{V-A} + a_2(\bar{c}u)_{V-A}(\bar{d}b)_{V-A}], \quad (3)$$

where

$$a_1 = C_1 + \frac{1}{3}C_2\text{: tree diagram}, \quad (4)$$

$$a_2 = C_2 + \frac{1}{3}C_1\text{: color suppressed tree diagram}. \quad (5)$$

In Ref. [4], these decays were analyzed. The form factors which are functions of $s = q^2 = (p - p')^2$, where $p = p' + q$ was evaluated at the $s = (\text{mass of } M)^2$ in the quark model, using heavy quark spin symmetry. In particular, using $a_2 \approx 0.10$, for $C_1(m_b) \approx 1.121$ and $C_2(m_b) \approx -0.275$ [3] and recent experimental values of other parameters, one finds $\mathrm{BR}(\Lambda_b \to \Lambda J/\psi) \simeq 1.21 \times 10^{-4}$ and the asymmetry $\alpha \simeq -0.18$ to be compared with the experimental values [3]: $((6.3 \pm 0.8) \times 10^{-4})$, $\alpha = 0.05 \pm 0.18$. However, replacing $C_2 + \frac{1}{3}C_1$ with the more general value $C_2 + \zeta C_1$ and taking $\zeta = 0.46$, we get $\mathrm{BR} = 6.0 \times 10^{-4}$ [4].

It is clear that $(\bar{c}b)_{V-A}$ and $(\bar{s}b)_{V-A}$ in Eq. (2) belong to the singlet and triplet representations of $SU(3)$, respectively. Thus, in the factorization ansatz, the only possible decay modes for B_b belonging to representation $\bar{3}$ are $\Lambda_b \to \Lambda_c^+ D_s^-$, $\Xi_b^0 \to \Xi_c^+ D_s^-$, $\Xi_b^- \to \Xi_c^0 D_s^-$ for the first term. For the second term, since $\bar{3} \times 3 = 8 + 1$, possible decay modes are $\Lambda_b \to \Lambda J/\psi$, $\Xi_b^0 \to \Xi^0 J/\psi$, $\Xi_b^- \to \Xi^- J/\psi$, where Λ, Ξ^0 and Ξ^- are members of the octet representation of $SU(3)$.

Hence, for the decays $B_b(\frac{1}{2}^+) \to B_b^*(\frac{3}{2}^+)P$, where B_b belongs to representation $\bar{3}$ of $SU(3)$, the above framework is not applicable as $B_b^*(\frac{3}{2}^+)$ either belongs to the decuplet or sextet representation of $SU(3)$. In this paper, the framework needed for the above decays is formulated.

II. EFFECTIVE HAMILTONIAN FOR DECAYS $\Lambda_b(p) \to B^*(p') + P(q)$

For the decays of type $\Lambda_b(p) \to B^*(p') + P(q)$, the Lorentz structure of the T matrix is given by [5]

$$T = \frac{1}{(2\pi)^{\frac{9}{2}}}\sqrt{\frac{mm^*}{2p_0p_0'q_0}}\frac{q^\lambda}{f_P}\bar{u}_\lambda(p')[C - D\gamma_5]u(p), \quad (6)$$

where u_λ is the Raita-Schwinger spinor, u is the Dirac spinor, and $q = p - p'$. It is clear that C is the parity conserving (p-wave) amplitude and D is the parity-violating (d-wave) amplitude. f_P is the pseudo scalar meson decay constant, which is introduced here to make the amplitudes C and D dimensionless. For the Rarita-Schwinger spinor $u_\lambda(p')$

$$\sum_{\text{spin}} u_\lambda(p')\bar{u}_\mu(p') = \frac{\gamma.p' + m'}{2m'}\left[\eta_{\lambda\mu} - \frac{1}{3}\gamma_\lambda\gamma_\mu + \frac{i}{3m'}(\gamma_\lambda p'_\mu - \gamma_\mu p'_\lambda) - \frac{2}{3m'^2}p'_\lambda p'_\mu\right]$$

$$\sum_{\text{spin}} u(p)\bar{u}(p) = \frac{1}{2}\frac{\gamma.p + m}{2m} \tag{7}$$

Using above equations and taking traces, we get the decay rate [5]

$$\Gamma = \frac{1}{6\pi}\frac{m|q|^3}{m'^2 f_P^2}[(E' + m')|C|^2 + (E' - m')|D|^2] \tag{8}$$

and for asymmetry α is given by

$$\alpha = \frac{2|\mathbf{q}|ReCD^*}{(E' + m')|C|^2 + (E' - m')|D|^2} \tag{9}$$

In order to determine the amplitudes C and D, one needs basic H_{eff}. In the nonleptonic decays of hyperons, there is important contribution viz the baryon-pole contribution (Born term) to the parity conserving (p-wave) decay amplitude for which the W exchange is relevant [6]. Such a contribution for Λ_b decays arises from the W exchange $b + u \overset{W}{\to} c + d$.

The effective Hamiltonian for the W-exchange diagram $b + u \to c + d$ is given by

$$H_{\text{eff}} = V_{cb}V_{ud}^*\left[\left(C_1 + \frac{1}{3}C_2\right)(\bar{d}^\beta\gamma^\mu(1 - \gamma^5)u_\beta) \times (\bar{c}^\alpha\gamma_\mu(1 - \gamma^5)b_\alpha) + \left(C_2 + \frac{1}{3}C_1\right) \times (\bar{c}^\alpha\gamma^\mu(1 - \gamma^5)u_d)(\bar{d}^\beta\gamma_\mu(1 - \gamma^5)b_\beta)\right] \tag{10}$$

after taking into account the QCD correction. For the case considered in this paper, the first term is relevant. Corresponding to the first term, the M matrix for the W-exchange diagram is given by

$$M \approx V_{cb}V_{ud}^*\frac{G_F}{\sqrt{2}}\left(C_1 + \frac{1}{3}C_2\right) \times [\bar{u}(p'_i)\gamma^\mu(1 - \gamma_5)\alpha_i^+ u(p_i)] \times [\bar{u}(p'_j)\gamma_\mu(1 - \gamma_5)\gamma_j^+ u(p_j)], \tag{11}$$

where $q = p_i - p'_i = p'_j - p_j$ and $q^2 \ll m_W^2$. In the Pauli representation of the γ matrices, the four-component Dirac spinor can be written as $u = \begin{pmatrix} u_A \\ u_B \end{pmatrix}$, where each u_A and u_B has two components. In the nonrelativistic limit u_B is of order v/c compared to u_A. Thus, only the bilinears

$$\bar{u}\gamma^0 u \approx u_A^\dagger u_A + O(v^2/c^2)$$
$$\bar{u}\gamma^i\gamma_5 u \approx u_A^\dagger\sigma^i u_A + O(v^2/c^2)$$

are large (see for example [7]). Using above results, after writing M in terms of two component spinors and then taking the Fourier transform, one gets the effective Hamiltonian in the leading nonrelativistic limit for the W exchange [5,6],

$$H_W^{pc} = \frac{G_F}{\sqrt{2}}V_{cb}V_{ud}^*\left(C_1 + \frac{1}{3}C_2\right) \times \sum_{i\neq j}\alpha_i^+\gamma_j^-(1 - \sigma_{\mathbf{i}}.\sigma_{\mathbf{j}})\delta^3(x)$$
$$H_W^{pv} = 0, \tag{12}$$

where α_i^+ and γ_j^- are operators which convert b-quark to c-quark and u-quark to d-quark.

$$\alpha_i^+|b\rangle = |c\rangle$$
$$\gamma_j^-|u\rangle = |d\rangle.$$

Following comments are in order. H_W^{pc} in the leading nonrelativistic limit was first derived in Ref. [6] for the parity conserving nonleptonic decays of baryons. The result obtained insure $\Delta I = 1/2$ rule (or octet dominance) in agreement with experiment. Other results obtained were also in agreement with the experimental values.

In Ref. [5], the decays $\Lambda_c \to \Delta^{++}K^-, \Sigma^{*0}\pi^+, \Xi^{*0}K^+$, were analysed in the same framework. The branching ratio: $\text{BR}(\Lambda_c \to \Delta^{++}K^-) = 9.0 \times 10^{-3}$ obtained is in good agreement with the experimental value $(8.6 \pm 1.0) \times 10^{-3}$. For the case, analyzed in this paper, the W-exchange diagram $b + u \to c + d$, involves two heavy quarks, hence the leading nonrelativistic approximation valid upto $O(v^2/c^2)$ is viable to apply. We note that our approach is somewhat analogous with that considered in Ref. [8]: Finally, the QCD correction has also been incorporated, it gives a factor $(C_1 + \frac{1}{3}C_2)^2 \approx 1.06$

Before we proceed further we note that [9]

$$|\Lambda^0_{b,c}\rangle = \frac{1}{\sqrt{2}}|(ud-du)b, c > \chi_{MA} \tag{13}$$

Thus Λ_b and Λ_c belong to triplet rep. $\bar{3}$ of SU(3). The spin $\frac{1}{2}^+$ and spin $\frac{3}{2}^+$ baryons which belong to representation 6 of SU(3), are

$$S_{ij}(S^*_{ij}) = \frac{1}{2}|(q_iq_j + q_jq_i)c > \chi_{MS}(\chi_S) \tag{14}$$

where $q = u$, d, s. The spin wave functions χ_{MA}, χ_{MS}, and χ_S are [9]

$$\chi_{MA} = \frac{1}{\sqrt{2}}|(\uparrow\downarrow - \downarrow\uparrow)\uparrow\rangle, \tag{15}$$

$$\chi_{MS} = \frac{1}{\sqrt{6}}|-(\uparrow\downarrow + \downarrow\uparrow)\uparrow + 2\uparrow\uparrow\downarrow\rangle, \tag{16}$$

$$\chi_S = \frac{1}{\sqrt{3}}|\uparrow\uparrow\downarrow + (\uparrow\downarrow + \downarrow\uparrow)\uparrow\rangle. \tag{17}$$

In particular, we note that

$$S_{22} = \sqrt{2}ddc\chi_{MS} = \sqrt{2}\Sigma^0_c. \tag{18}$$

It is clear that relevant operators are

$$\alpha^+_3\gamma^-_1(1-\sigma_1.\sigma_3) + \alpha^+_3\gamma^-_2(1-\sigma_2.\sigma_3). \tag{19}$$

Hence, we get

$$\sum_{i\neq j}\alpha^+_i\gamma^-_j(1-\sigma_i.\sigma_j)|\Lambda_b\rangle = \sqrt{6}|\Sigma^0_c\rangle. \tag{20}$$

We note

$$\bar{3}\times 10 = 6 + 24, \qquad 8\times 6 = \bar{3} + 6 + 15 + 24.$$

Thus, only possible decays through the Σ^0_c pole are

Set I (II)

$$\Lambda_b \to \Sigma^0_c \to \Delta^0 D^0, \qquad \Delta^- D^+,$$
$$\Sigma^{*-}D^+_s(\Sigma^{*-}_c\pi^+, \Sigma^{*0}_c\pi^0, \Xi^*_c K^0).$$

Hence, in the SU(3) limit, the p-wave (parity conserving) amplitude C for the two sets of decays is given by

$$C = F\left(C_1 + \frac{1}{3}C_2\right)\frac{\langle\Sigma^0_c|H^{pc}_W|\Lambda_b\rangle}{m_{\Lambda_b} - m_{\Sigma_c}}, \tag{21}$$

where F is $(2, 2\sqrt{3}, 2)g$ and $(\sqrt{2}, -\sqrt{2}, \sqrt{2})g_c$ for sets I and II, respectively. The weak matrix elements $\langle\Sigma^0_c|H^{pc}_w|\Lambda_b\rangle$ in using Eq. (12) and Eq. (20) are given by

$$\langle\Sigma^0_c|H^{pc}_W|\Lambda_b\rangle = \left[\frac{G_F}{\sqrt{2}}V_{cb}V^*_{ud}\right]\sqrt{6}d', \tag{22}$$

where [8]

$$d' = \langle\psi_0|\delta^3(x)|\psi_0\rangle = \frac{3(m_\Delta - m_{\Sigma_c})}{8\pi\alpha_s}\bar{m}^2 \tag{23}$$

for set I and, for set II,

$$d' = \frac{3(m_{\Sigma^*_c - m_{\Sigma_c}})}{8\pi\alpha_s}\bar{m}^2_c. \tag{24}$$

Since $H^{p.v.}_W = 0$, it follows that the baryon pole cannot generate the d-wave amplitude; hence, $D = 0$ and, thus, the asymmetry $\alpha = 0$. This is in accordance with the two-particle nonleptonic decay of Ω^- for which the experimental value of $\alpha = 0$. This is the first prediction of the framework used without detailed analysis.

III. DETAILED ANALYSIS OF $\Lambda_b \to \Delta D$ AND $\Lambda_b \to \Sigma^*_c\pi$

We first discuss the decay $\Lambda_b \to \Delta D$ decay. In the rest frame of Λ_b, $|\mathbf{q}| = 2.327$ GeV for the final state $\Delta^0 D^0, \Delta^- D^+$ and $|\mathbf{q}| = 2.242$ GeV for $\Sigma^{*-}D^+_s$. Using experimental values for $m = m_{\Lambda_b}$ and $m^* = m_\Delta$ or m_{Σ^*} $f_D = 207$ MeV and $f_{D_s} = 257$ MeV [10], we get from Eq. (8)

$$\Gamma = 2.17\times 10^2|C|^2 \text{ GeV} \tag{25}$$

$$= 1.13\times 10^2|C|^2 \text{ GeV} \tag{26}$$

for $\Delta^0 D^0, \Delta^- D^+$ and $\Sigma^{*-}D^+_s$ First we note that the constant g in Eq. (24) can be estimated by using partial conservation of the axial vector current (PCAC) and the nonrelativistic quark model (NQM):

$$g = g_{\Sigma^0_c D^0\Delta^0} = g_{\Sigma^0_c D^+\Delta^-} = \frac{m_{\Sigma^0_c + m_\Delta}}{f_D}g_A. \tag{27}$$

Now, in NQM [11], $g_A = -\frac{2\sqrt{2}}{3\sqrt{3}}$. Thus,

$$g = 9.686, \tag{28}$$

and

$$g = g_{\Sigma^0_c D^+_s}\textstyle\sum^{*-} = \frac{m_{\Sigma^0_c} + m_{\Sigma^+}}{f_{D_s}}g_A = 8.024. \tag{29}$$

In Eq. (24), in order to take into account large momentum transfer in heavy flavor baryon decays, we take

$$\bar{m}^2 = \frac{m^4_d}{(m_b + m_c)^2} \approx 3.18\times 10^{-4}. \tag{30}$$

Using constituent quark masses, $m_d \approx 0.334$ GeV, $m_c \approx 1.4$ GeV, $m_b \approx 4.85$ GeV, $\alpha_s \approx 0.32$ and experimental values

$$m_{\Sigma_c} - m_\Delta \approx 1.22 \text{ GeV}, m_{\Sigma_c} - m_{\Sigma^*} \approx 1.07 \text{ GeV}, \quad (31)$$

we get from Eqs. (22) and (23) and Eqs. (28)–(30)

$$C \approx 7.75 \times 10^{-10}, \qquad 5.62 \times 10^{-10}. \quad (32)$$

Hence, from Eqs. (25) and (26),

$$\Gamma(\Lambda_b \to \Delta^0 D^0) \approx 1.23 \times 10^{-16} \text{ GeV}$$

$$\Gamma(\Lambda_b \to \Delta^- D^+) = 3\Gamma(\Lambda_b \to \Delta^0 D^0) \approx 3.69 \times 10^{-16} \text{ GeV} \quad (33)$$

$$\Gamma(\Lambda_b \to \Sigma^{*-} D_s^+) \approx 3.38 \times 10^{-17} \text{ GeV}. \quad (34)$$

Using $\tau_{\Lambda_b} \approx 1.451 \times 10^{-12}$ s, we get

$$\mathrm{BR}(\Lambda_b \to \Delta^0 D^0) \approx 2.70 \times 10^{-4},$$
$$\mathrm{B.R}(\Lambda_b \to \Delta^- D^+) \approx 8.11 \times 10^{-4}, \quad (35)$$

$$\mathrm{BR}(\Lambda_b \to \Sigma^{*-} D_s^+) \approx 7.04 \times 10^{-5}, \quad (36)$$

Set II

$$\Lambda_b \to \Sigma_c^{*-}\pi^+, \qquad \Sigma_c^{*0}\pi^0, \qquad \Xi_c^{*0}K^0$$
$$|\vec{q}| = 2.244 \text{ GeV}, \qquad 2.240 \text{ GeV}$$

From Eq. (8), using $f_\pi \approx 131$ MeV and $f_K \approx 160$ MeV, we get

$$\Gamma(\Lambda_b \to \Sigma_c^{*-}\pi^+) \approx 1.90 \times 10^2 |C|^2 \text{ GeV} \quad (37)$$

$$\Gamma(\Lambda_b \to \Xi_c^{*0}K^0) \approx 1.21 \times 10^2 |C|^2 \text{ GeV} \quad (38)$$

Now PCAC gives

$$g_c = g_{\Sigma_c^{*-}\Sigma_c^0\pi^+} = \frac{m_{\Sigma_c^*} + m_{\Sigma_c}}{f_\pi} g_{Ac} \approx 25.30$$
$$g_{\Xi_c^{*0}\Sigma_c^0 K^0} = \frac{m_{\Xi^0} + m_{\Sigma_c}}{f_K} g_{Ac} \approx 21.24$$

on using the NQM value $g_{Ac} = 2/3$. Hence, from Eqs. (22) and (24), using

$$\bar{m}_c^2 = \left(\frac{m_d m_c}{m_b + m_c}\right)^2 \approx 5.60 \times 10^{-3}$$

and

$$(m_{\Sigma_c^*} - m_{\Sigma_c}) \approx 0.064 \text{ GeV}, \qquad (\Xi_c^* - m_{\Sigma_c}) \approx 0.191 \text{ GeV}, \quad (39)$$

we get

$$C \approx 1.31 \times 10^{-9} \quad \text{and} \quad C \approx 3.30 \times 10^{-9} \quad (40)$$

for $\Sigma_c^{*-}\pi^+$ and $\Xi_c^{*0}K^0$.

Hence, for the decay rates and the branching ratios, we get from Eqs. (37) and (38),

$$\Gamma(\Lambda_b \to \Sigma_c^{*-}\pi^+) = \Gamma(\Lambda_b \to \Sigma_c^{*0}\pi^0) \approx 2.92 \times 10^{-16} \text{ GeV} \quad (41)$$

$$\Gamma(\Lambda_b \to \Xi_c^{*0}K^0) \approx 12.46 \times 10^{-16} \text{ GeV} \quad (42)$$

$$\mathrm{BR}(\Lambda_b \to \Sigma_c^{*-}\pi^+) = B.R(\Lambda_b \to \Sigma_c^{*0}\pi^0) \approx 6.47 \times 10^{-4} \quad (43)$$

$$\mathrm{BR}(\Lambda_b \to \Xi_c^{*0}K^0) \approx 2.74 \times 10^{-3}. \quad (44)$$

To conclude: No experimental data for the branching ratios of two sets of decay channels,

$$\Lambda_b \to \Delta^0 D^0, \qquad \Delta^- D^+, \qquad \Sigma_c^{*-} D_s^+$$

and

$$\Lambda_b \to \Sigma_c^{*-}\pi^+, \qquad \Sigma_c^{*0}\pi^0, \qquad \Xi_c^{*0}K^0,$$

are available to test the branching ratios given in Eq. (35) and (36) and Eqs. (43) and (44). One notes that relative branching ratios viz

$$\frac{\Gamma(\Lambda_b \to \Delta^- D^+)}{\Gamma(\Lambda_b \to \Delta^0 D^0)} \approx 3, \qquad \frac{\Gamma(\Lambda_b \to \Sigma^{*-} D_s^+)}{\Gamma(\Lambda_b \to \Delta^0 D^0)} \approx 0.28 \quad (45)$$

and

$$\frac{\Gamma(\Lambda_b \to \Sigma_c^{*0}\pi^0)}{\Gamma(\Lambda_b \to \Sigma_c^{*-}\pi^+)} \approx 1, \qquad \frac{\Gamma(\Lambda_b \to \Xi_c^{*0}K^0)}{\Gamma(\Lambda_b \to \Sigma_c^{*-}\pi^+)} \approx 4.28 \quad (46)$$

are independent of the parameters $\bar{m}^2, \bar{m}_c^2$ and the axial vector coupling constants g_A and g_{Ac}. Thus, Eqs. (45) and (46), together with the prediction that asymmetry parameter $\alpha = 0$, will test the general framework used in the analysis of decays $\Lambda_b(\frac{1}{2}^+) \to B^*(\frac{3}{2}^+) + P$.

Finally, decays with three particles in the final state through resonances

$$\begin{aligned}\Lambda_b &\to \Delta^0 D^0 \to p\pi^- D^0\\ &\to \Delta^- D^+ \to n\pi^- D^+\\ &\to \Sigma^{*-} D_s^+ \to \Lambda\pi^- D_s^+\\ &\to \Sigma^0 \pi^- D_s^+\end{aligned}$$

are of considerable interest. For the decuplet $\Delta, m_\Delta = 1232$ MeV, $\Gamma_\Delta \approx 117$ MeV, $m_{\Sigma^*} = 1385$ MeV, SU (3) gives [1]

$$\Delta^0 \to p\pi^-:\ \sqrt{2}g^*$$
$$\to n\pi^0 = 2g^*$$
$$\Delta^- \to n\pi^-:\ \sqrt{6}g^*$$
$$\Sigma_c^{*-} \to \Sigma^-\pi^0:\ g^*$$
$$\to \Sigma^0\pi^-:\ g^*$$
$$\to \Lambda\pi^-:\ -\sqrt{3}g^*.$$

Using physical masses and the phase space factor,

$$\Gamma(\Delta^0 \to p\pi^-) \simeq 39\ \text{MeV},$$
$$\Gamma(\Delta^0 \to n\pi^0) \simeq 78\ \text{MeV}$$
$$\Gamma(\Delta^- \to n\pi^-) \simeq 117\ \text{MeV}$$
$$\Gamma(\Sigma^{*-} \to \Lambda\pi^-) \simeq 32.7\ \text{MeV}$$
$$\Gamma(\Sigma^{*-} \to \Sigma^0\pi^-) = \Gamma(\Sigma^{*-} \to \Sigma^*\pi^0) \simeq 2.8\ \text{MeV}$$

Thus, $\Gamma_{\Sigma^{*-}} \simeq 38.3$ MeV, to be compared with the experimental value $\Gamma_{\Sigma^{*-}} \approx (39.1 \pm 2.1)$ MeV [3]. Finally, we get

$$\frac{\Gamma(\Lambda_b \to \Delta^- D^+ \to n\pi^- D^+)}{\Gamma(\Lambda_b \to \Delta^0 D^0 \to p\pi\Sigma^- D^0)} \approx 9 \quad (47)$$

$$\frac{\Gamma(\Lambda_b \to \Sigma^{*-} D_s^+ \to \Lambda\pi^- D_s^+)}{\Gamma(\Lambda_b \to \Delta^0 D^0 \to p\pi^- D^0)} \approx (0.28)\frac{(32.7)}{39} \approx 0.23. \quad (48)$$

For the second set of decays, we note that $\Xi_c^0, \Xi_c^+, \Lambda_c$ belong to representation $\bar{3}$ and $\Sigma_c^{*0}, \Xi_c^{*0}, \Xi_c^{*+}$ belong to representation 6 of SU(3). Now $\bar{3} \times 8 \to 6 + \bar{3} + 15$; thus, SU(3) gives

$$\Sigma_c^{*0} \to \Lambda_c^+\pi^-:\ -\sqrt{2}g_c^*$$
$$\Xi_c^{*0,+} \to \Xi_c^{+,0}\pi^{-,+}:\ \pm g_c^*$$
$$\to \Xi_c^{0,+}\pi^0:\ \mp \frac{1}{\sqrt{2}}g_c^*.$$

The first prediction of the above analysis taking into account the phase space is that the total decay width of Ξ_c^{*0}:

$$\Gamma_{\Xi_c^{*0}} \approx \frac{3}{2}(0.11)\Gamma_{\Sigma_c^{*0}} = \frac{3}{2}(0.11)(14.5 \pm 1.5)\ \text{MeV}$$
$$= (2.39 \pm 0.25)\ \text{MeV} = \Gamma_{\Xi_c^{*+}}$$

on using the experimental value

$$\Gamma_{\Sigma_c^{*0}} = \Gamma(\Sigma_c^{*0} \to \Lambda_c^+\pi^-) = (14.5 \pm 1.5)\ \text{MeV}.$$

The experimental limits on decay width $\Gamma_{\Xi_c^{*0}} < 5.5$ MeV, $\Gamma_{\Xi_c^{*+}} < 3.1$ MeV, [3].

Finally, the branching ratios for three particle states $\Lambda_c^+\pi^-\pi^0$ and $\Xi_c^+\pi^-K^0$ through resonances Σ_c^{*0} and Ξ_c^{*0} are given by

$$\frac{\Gamma(\Lambda_b \to \Xi_c^{*0}K^0 \to \Xi_c^+\pi^-K^0)}{\Gamma(\Lambda_b \to \Sigma_c^{*0}\pi^0 \to \Lambda_c^+\pi^-\pi^0)} = (4.24)(0.11) \approx 0.47.$$

To summarize, we have analyzed two sets of decays of Λ_b:

$$\Lambda_b \to \Sigma_c^0 \to \Delta^0 D^0, \qquad \Delta^- D^+, \qquad \Sigma^{*-}D_s^+$$
$$\Lambda_b \to \Sigma_c^0 \to \Sigma^{*+}\pi^-, \qquad \Sigma_c^{*0}\pi^0, \qquad \Xi_c^{*0}K^0.$$

We note that the other two members of triplet $\bar{3}$ are $\Xi_b^0 = \frac{1}{\sqrt{2}}(us - su)b\chi_{MA}$ and $\Xi_b^- = \frac{1}{\sqrt{2}}(ds - sd)b\chi_{MA}$. The W exchange cannot generate a baryon pole for Ξ_b^-; thus, the decays $\Xi_b^- \to B^*(\frac{3}{2}^+) + P$ are not possible. This is another prediction of our formalism. However, for Ξ_b^0, the W exchange gives

$$H_W^{p.c}|\Xi_b^0\rangle = -\frac{1}{2}[-2(d^\uparrow s^\uparrow + s^\uparrow d^\uparrow)c^\downarrow + (d^\uparrow s^\downarrow + d^\downarrow s^\uparrow)c^\uparrow$$
$$+ (s^\uparrow d^\uparrow + s^\downarrow d^\uparrow)c^\uparrow$$
$$= \sqrt{6}|\Xi_c'^0\rangle.$$

Thus, $\langle \Xi_c'^0|H_W^{p.c.}|\Xi_b^0\rangle = \sqrt{6}$.

Hence, the pole diagram gives two set of decays:

$$\Xi_b^0 \to \Xi_c'^0 \to \Sigma^{*0}D^0, \qquad \Sigma^{*-}D^+, \qquad \Xi^{*-}D_s^+$$
$$\Xi_b^0 \to \Xi_c'^0 \to \Xi_c^{*+}\pi^-, \qquad \Xi_c^{*0}\pi^0, \qquad \Omega_c^{*0}K^0.$$

Hence, in the SU(3) limit, for two sets of decays, the p-wave (parity conserving) amplitude C is given by set I:

$$C = \sqrt{2}(1,1,1)g_c\frac{\langle \Xi_c'^0|H_W^{pc}|\Xi_b^0\rangle}{m_{\Xi_b^0} - m_{\Xi_c'^0}}$$

set II:

$$C = \frac{1}{\sqrt{2}}(\sqrt{2}, -1, 2)g_c\frac{\langle \Xi_c'^0|H_W^{pc}|\Xi_b^0\rangle}{m_{\Xi_b^0} - m_{\Xi_c'^0}}.$$

Following exactly the same procedure as for the Λ_b decays, one can calculate the branching ratios $\Xi_b^0 \to B(\frac{3}{2}^+) + P$ decays. At present, no experimental data for $\Lambda_b \to B^*(\frac{3}{2}^+) + P$ and $\Xi_b^0 \to B^*(\frac{3}{2}^+) + P$ are available to test the prediction of our model. In the future, it is expected that more data for heavy flavor hadron decays will be coming from the LHCb, including the decays considered in this paper.

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[1] For a review, see, for example, Fayyazuddin and Riazuddin, *A Modern Introduction to Particle Physics*, 3rd ed. (World Scientific, Singapore, 2011), Ch. 15, Sec. 15.3.4. See also Refs. [2] and [3].

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