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Predictions for nonleptonic Λ_b and Θ_b decays

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Abstract

We study nonleptonic $\Lambda_b \to \Lambda_c \pi$, $\Sigma_c \pi$ and $\Sigma_c^* \pi$ decays in the limit $m_b, m_c, E_\pi \gg \Lambda_{\rm QCD}$ using the soft-collinear effective theory. Here $\Sigma_c = \Sigma_c(2455)$ and $\Sigma_c^* = \Sigma_c(2520)$. At leading order the $\Lambda_b \to \Sigma_c^{(*)} \pi$ rates vanish, while the $\Lambda_b \to \Lambda_c \pi$ rate is related to $\Lambda_b \to \Lambda_c \ell \bar{\nu}$, and is expected to be larger than $\Gamma(B \to D^{(*)} \pi)$. The dominant contributions to the $\Lambda_b \to \Sigma_c^{(*)} \pi$ rates are suppressed by $\Lambda_{\rm QCD}^2 / E_\pi^2$. We predict $\Gamma(\Lambda_b \to \Sigma_c^* \pi)/\Gamma(\Lambda_b \to \Sigma_c \pi) = 2 + \mathcal{O}[\Lambda_{\rm QCD}/m_Q, \alpha_s(m_Q)]$, and the same ratio for $\Lambda_b \to \Sigma_c^{(*)} \rho$ and for $\Lambda_b \to \Xi_c^{(\prime,*)} K$. "Bow tie" diagrams are shown to be suppressed. We comment on possible discovery channels for weakly decaying pentaquarks, $\Theta_{b,c}$ and their nearby heavy quark spin symmetry partners, $\Theta_{b,c}^*$.

Heavy baryon decays are interesting for many reasons. Heavy quark symmetry [1] is more predictive in semileptonic $\Lambda_b \to \Lambda_c \ell \bar{\nu}$ decay than in $B \to D^{(*)} \ell \bar{\nu}$, and may eventually give a determination of $|V_{cb}|$ competitive with meson decays [2]. In this Letter we concentrate on the more complicated case of nonleptonic $b \to c \bar{u} d$ baryon transitions, as shown in Table 1. These channels provide a testing ground for our understanding of QCD in nonleptonic decays. Our analysis is based on heavy quark symmetry and the soft-collinear effective theory (SCET) [4].

There is considerable experimental interest in these decays. Recently the CDF Collaboration measured [5]

$$\frac{f_{\Lambda_b}}{f_d} \frac{\mathcal{B}(\Lambda_b \to \Lambda_c^+ \pi^-)}{\mathcal{B}(\bar{B}^0 \to D^+ \pi^-)} = 0.66 \pm 0.11_{\rm (stat)} \pm 0.09_{\rm (syst)} \pm 0.18_{\rm (BR)}, \qquad (1)$$

where f_{Λ_b} and f_d are the fragmentation fractions of b quarks to Λ_b and $\bar{B}^0$, respectively. Using the input $f_{\rm baryon}/f_d = 0.304 \pm 0.053$, CDF quoted $\mathcal{B}(\Lambda_b \to \Lambda_c^+ \pi^-)/\mathcal{B}(\bar{B}^0 \to D^+ \pi^-) \simeq 2.2$ [5]. The measured lifetimes, $\tau(\Lambda_b) = 1.23$ ps and $\tau(B^0) = 1.54$ ps [6], then imply that $\Gamma(\Lambda_b \to \Lambda_c^+ \pi^-)/\Gamma(\bar{B}^0 \to D^+ \pi^-) \simeq 2.7$. More experimental results on semileptonic and other nonleptonic Λ_b decays are expected in the near future.

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Table 1
The decays considered in this Letter. Here s_l is the spin of the light degrees of freedom [3]. The mass shown for the $\Sigma_c^{(*)}$ is the average of the charge 0 and +1 states, and for the Ξ_c's the mass is the average in the doublet. The stability of the pentaquark states Θ_Q ($= \bar{Q}udud$) and their value of s_l are both conjectures

Notation	s_l	$I(J^P)$	Mass (MeV)	Decays considered
Λ_c, Λ_b	0	$0(\frac{1}{2}^+)$	2285, 5624	$\Lambda_b \to \Lambda_c^+ \pi^-$
$\Sigma_c = \Sigma_c(2455)$	1	$1(\frac{1}{2}^+)$	2452	$\Lambda_b \to \Sigma_c^+ \pi^-$, $\Sigma_c^0 \pi^0$, $\Sigma_c^0 \rho^0$
$\Sigma_c^* = \Sigma_c(2520)$	1	$1(\frac{3}{2}^+)$	2517	$\Lambda_b \to \Sigma_c^{*+} \pi^-$, $\Sigma_c^{*0} \pi^0$, $\Sigma_c^{*0} \rho^0$
Ξ_c, Ξ_c'	0, 1	$\frac{1}{2}(\frac{1}{2}^+)$	2469, 2576	$\Lambda_b \to \Xi_c'^0 K^0$
$\Xi_c^* = \Xi_c(2645)$	1	$\frac{1}{2}(\frac{3}{2}^+)$	2646	$\Lambda_b \to \Xi_c^{*0} K^0$
Θ_c, Θ_b	1	$0(\frac{1}{2}^+)$	m_{Θ_c}, m_{Θ_b}	$\Theta_b^+ \to \Theta_c^0 \pi^+$, $\Theta_c^0 \to \Theta^+ \pi^-$
Θ_c^*, Θ_b^*	1	$0(\frac{3}{2}^+)$	$\sim m_{\Theta_c} + 70$ $\sim m_{\Theta_b} + 22$	$\Theta_c^* \to \Theta_c \gamma$ or strongly $\Theta_b^* \to \Theta_b \gamma$

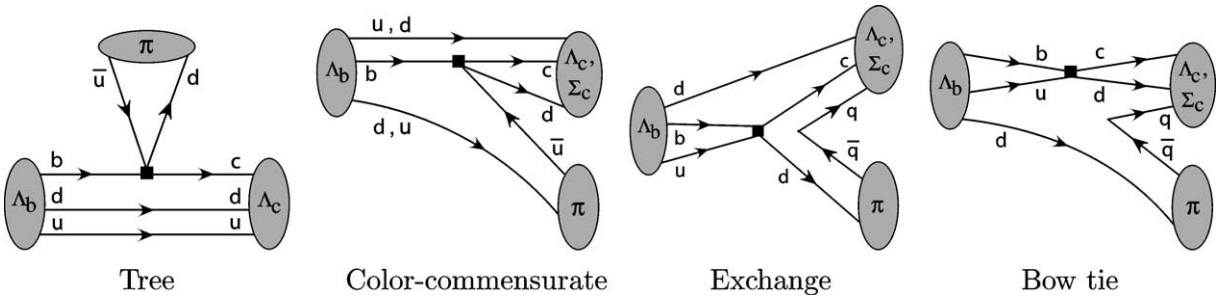

Fig. 1. Diagrams for Λ_b decays, giving amplitudes T, C, E, and B. Decay to Λ_c gets contributions from all four terms. Decays to $\Sigma_c^{(*)}$ and Ξ_c do not have T and T, C contributions, respectively.

The part of the weak Hamiltonian relevant for this Letter is

$$H_W = \frac{4G_F}{\sqrt{2}} V_{cb} V_{ud}^* \times \big[C_1(m_b) O_1(m_b) + C_2(m_b) O_2(m_b)\big], \quad (2)$$

where both the Wilson coefficients, C_i, and the four-quark operators

$$O_1 = (\bar{c}\gamma_\mu P_L b)\big(\bar{d}\gamma^\mu P_L u\big),$$
$$O_2 = (\bar{d}\gamma_\mu P_L b)\big(\bar{c}\gamma^\mu P_L u\big), \quad (3)$$

depend on the renormalization scale which we take to be m_b, and $P_L = (1 - \gamma_5)/2$. The combination $[C_1(m_b) + C_2(m_b)/3]|V_{ud}|$ is very close to unity.

Weak nonleptonic decays are sometimes characterized by diagrams corresponding to different Wick contractions. As shown in Fig. 1, there are more possibilities in baryon than in meson decays. In particular, a "Bow tie" contraction is unique to baryons. The color structure for baryons also differs from mesons: we find that the C diagram is of the same order in the large N_c limit as the T diagram.[1] Nonleptonic meson decay amplitudes are sometimes estimated using naive factorization, which would set $\langle \Lambda_c \pi | O_1 | \Lambda_b \rangle = \langle \Lambda_c | \bar{c}\gamma_\mu P_L b | \Lambda_b \rangle \langle \pi | \bar{d}\gamma_\mu P_L u | 0 \rangle$. In baryon decays the extra light quark implies that this procedure is ill-defined for all but the tree diagram. In naive factorization the $\Lambda_b \to \Sigma_c^{(*)} \pi$ decays are very suppressed, since the T contribution vanishes separately in the isospin and heavy quark limits [7] (just like the semileptonic $\Lambda_b \to \Sigma_c^{(*)} \ell \bar{\nu}$ decays), the C contribution vanishes after doing a Fiertz transformation on the four-quark operator, and the E and B amplitudes are

[1] If we treated the $N_c - 3$ additional quarks in the baryons as flavors that are sterile under the weak interaction then color-commensurate would become color-suppressed.

identically zero since the u and b fields are in different quark bilinears.

In this Letter we show that more rigorous techniques can still be applied to make reasonable predictions for all these decays. By expanding in m_b, m_c, $E_\pi \gg \Lambda_{\rm QCD}$ we show that for $\Lambda_b \to \Lambda_c^+ \pi^-$ the amplitudes corresponding to the diagrams in Fig. 1 satisfy $T \gg C \sim E \gg B$, and we find that the experimental result in Eq. (1) is consistent with theoretical expectations. Next we consider $\Lambda_b \to \Sigma_c^{(*)} \pi$ decays, and show the leading contributions to these nonleptonic rates are suppressed by $\Lambda_{\rm QCD}^2/E_\pi^2$, much like in $B^0 \to D^0 \pi^0$. Using heavy quark symmetry we derive a relation between the decay rate to Σ_c and Σ_c^* and comment on decays to Ξ_c. Finally we consider the detection of possible weakly decaying heavy pentaquarks, Θ_b and Θ_c, with nonleptonic decays.

The proof of factorization at leading order for $\Lambda_b \to \Lambda_c \pi$ decay follows closely that for $\bar{B}^0 \to D^{(*)+} \pi^-$ [8], so we do not review it here. In this case the nonperturbative expansion parameter for SCET is $\lambda = \Lambda_{\rm QCD}/E_\pi$ [9]. Since E_π is set by the bottom and charm quark masses, we take this to be of the same order as the expansion parameter for the heavy quark effective theory (HQET), i.e., $\lambda \sim \Lambda_{\rm QCD}/m_Q$ ($Q = c, b$). Working at leading order in λ and $\alpha_s(m_b)$ and neglecting the pion mass, the $\Lambda_b \to \Lambda_c \pi$ matrix element factorizes in the standard way,

$$\begin{aligned}
&\langle \Lambda_c(v',s')\pi \,|\, H_W \,|\, \Lambda_b(v,s)\rangle \\
&\quad = \sqrt{2} G_F \left(C_1 + \frac{C_2}{3}\right) V_{cb} V_{ud}^* f_\pi E_\pi \\
&\qquad \times \langle \Lambda_c(v',s') | \bar{c} \not{n} P_L b \,|\, \Lambda_b(v,s)\rangle, \qquad (4)
\end{aligned}$$

where $f_\pi = 131$ MeV is the pion decay constant, n is a light-like four-vector along the direction of the pion's four-momentum, $p_\pi^\mu = E_\pi n^\mu$, and the four-velocities of the Λ_b and Λ_c are v and v', respectively. Perturbative corrections induce a multiplicative factor in Eq. (4), $\langle T(x)\rangle_\pi = \int_0^1 {\rm d}x\, T(x)\phi_\pi(x)$, where $T(x)$ is computable and ϕ_π is the nonperturbative light-cone pion distribution function [10,11], and a term proportional to the matrix element of $\bar{c} \not{n} P_R b$. At leading order in $\alpha_s(m_Q)$, we can set $\langle T(x)\rangle_\pi = 1$ and the term involving $\bar{c} \not{n} P_R b$ to 0. This implies that the nonleptonic rate is related to the semileptonic differential decay rate at maximal recoil,

$$\begin{aligned}
&\Gamma(\Lambda_b \to \Lambda_c \pi) \\
&\quad = \frac{3\pi^2 (C_1 + C_2/3)^2 |V_{ud}|^2 f_\pi^2}{m_{\Lambda_b}^2 r_\Lambda} \\
&\qquad \times \left(\frac{{\rm d}\Gamma(\Lambda_b \to \Lambda_c \ell \bar{\nu})}{{\rm d}w} \right)_{w_{\rm max}}, \qquad (5)
\end{aligned}$$

where $r_\Lambda = m_{\Lambda_c}/m_{\Lambda_b}$, $w = v \cdot v' = (m_{\Lambda_b}^2 + m_{\Lambda_c}^2 - q^2)/(2 m_{\Lambda_b} m_{\Lambda_c})$, and $w_{\rm max}$ corresponds to $q^2 = m_\pi^2$ ($\simeq 0$).

The semileptonic $\Lambda_b \to \Lambda_c \ell \bar{\nu}$ form factors are

$$\begin{aligned}
&\langle \Lambda_c(p',s') \,|\, V_\mu \,|\, \Lambda_b(p,s)\rangle \\
&\quad = \bar{u}(p',s')[f_1 \gamma_\mu + f_2 v_\mu + f_3 v'_\mu] u(p,s), \\
&\langle \Lambda_c(p',s') \,|\, A_\mu \,|\, \Lambda_b(p,s)\rangle \\
&\quad = \bar{u}(p',s')[g_1 \gamma_\mu + g_2 v_\mu + g_3 v'_\mu] \gamma_5 u(p,s), \qquad (6)
\end{aligned}$$

where f_i and g_i are functions of w, and the relevant currents are $V_\nu = \bar{c}\gamma_\nu b$ and $A_\nu = \bar{c}\gamma_\nu\gamma_5 b$. The spinors are normalized to $\bar{u}(p,s)\gamma^\mu u(p,s) = 2p^\mu$. In the heavy quark limit,

$$\begin{aligned}
&\zeta(w) = f_1(w) = g_1(w), \\
&0 = f_2(w) = f_3(w) = g_2(w) = g_3(w), \qquad (7)
\end{aligned}$$

where $\zeta(w)$ is the Isgur–Wise function for ground state baryons. The differential decay rate is given by

$$\begin{aligned}
&\frac{{\rm d}\Gamma(\Lambda_b \to \Lambda_c \ell \bar{\nu})}{{\rm d}w} \\
&\quad = \frac{G_F^2 m_{\Lambda_b}^5 |V_{cb}|^2}{24\pi^3} r_\Lambda^3 \sqrt{w^2-1} \\
&\qquad \times \left[6w + 6w r_\Lambda^2 - 4r_\Lambda - 8 r_\Lambda w^2\right] \mathcal{F}_\Lambda^2(w), \qquad (8)
\end{aligned}$$

where in the $m_Q \gg \Lambda_{\rm QCD}$ limit $\mathcal{F}_\Lambda(w)$ is equal to the Isgur–Wise function, $\zeta(w)$, and in particular $\mathcal{F}_\Lambda(1) = 1$. In terms of the original form factors

$$\begin{aligned}
&\mathcal{F}_\Lambda(w)^2 \\
&\quad = \left[6w + 6w r_\Lambda^2 - 4r_\Lambda - 8 r_\Lambda w^2\right]^{-1} \\
&\qquad \times \Big\{ (w-1)\left[(1+r_\Lambda) f_1 + (w+1)(r_\Lambda f_2 + f_3)\right]^2 \\
&\qquad\quad + (w+1) \\
&\qquad\quad \times \left[(1-r_\Lambda) g_1 - (w-1)(r_\Lambda g_2 + g_3)\right]^2 \\
&\qquad\quad + 2\left(1 - 2 r_\Lambda w + r_\Lambda^2\right) \\
&\qquad\quad \times \left[(w-1) f_1^2 + (w+1) g_1^2\right] \Big\}. \qquad (9)
\end{aligned}$$

Combining the above results for $\Lambda_b \to \Lambda_c^+ \pi^-$ decay with the analogous ones for $\bar{B}^0 \to D^{(*)+}\pi^-$ we find that

$$\frac{\Gamma(\Lambda_b \to \Lambda_c \pi^-)}{\Gamma(\bar{B}^0 \to D^{(*)+}\pi^-)} = \frac{8 m_{\Lambda_b}^3 (1 - r_\Lambda^2)^3 r_{D^{(*)}}}{m_B^3 (1 - r_{D^{(*)}}^2)^3 (1 + r_{D^{(*)}})^2} \left(\frac{\zeta(w_{\max}^\Lambda)}{\xi(w_{\max}^{D^{(*)}})} \right)^2, \quad (10)$$

where ξ is the Isgur–Wise function for $B \to D^{(*)}$ semileptonic decay, and $r_{D^{(*)}} = m_{D^{(*)}}/m_B$. When the $\Lambda_b \to \Lambda_c^+ \ell^- \bar{\nu}$ rate is measured, one can directly test factorization using Eq. (5) or Eq. (10). In the absence of this data, we have to resort to using model predictions for the baryon Isgur–Wise function. If the ratio of Isgur–Wise functions in Eq. (10) is unity then the prefactor in Eq. (10) implies that $\Gamma(\Lambda_b \to \Lambda_c \pi^-)/\Gamma(\bar{B}^0 \to D^{(*)+}\pi^-) = 1.6(1.8)$. This enhancement is in rough agreement with the data in Eq. (1). A similar result also follows from the small velocity limit ($m_Q \gg m_b - m_c \gg \Lambda_{\rm QCD}$), in which the nonleptonic rates satisfy $\Gamma(\Lambda_b \to \Lambda_c \pi) : \Gamma(B \to D^* \pi) : \Gamma(B \to D\pi) = 2:1:1$, while for the semileptonic rates $\Gamma(\Lambda_b \to \Lambda_c \ell \bar{\nu}) : \Gamma(B \to D^* \ell \bar{\nu}) : \Gamma(B \to D \ell \bar{\nu}) = 4:3:1$.

The large N_c limit provides some support for the ratio of baryon to meson Isgur–Wise functions being close to unity at maximal recoil. In the large N_c limit the heavy baryons can be treated as bound states of chiral solitons and mesons containing a heavy quark. In this picture, the baryon Isgur–Wise function, $\zeta(w)$, is predicted to be $\zeta(w) = 0.99 e^{-1.3(w-1)}$ [12].[2] This gives $\zeta(w_{\max}^\Lambda = 1.4) = 0.57$, which is indeed close to $\xi(w_{\max}^{D^*} = 1.5) \simeq 0.55$ [13]. Using this model for $\zeta(w)$, $|V_{cb}| = 0.04$, $\tau_{\Lambda_b} = 1.23$ ps, and Eqs. (5) and (8) yield the prediction that $\mathcal{B}(\Lambda_b \to \Lambda_c^+ \pi^-) = 4.6 \times 10^{-3}$. As expected, this is larger than $\mathcal{B}(\bar{B}^0 \to D^{(*)+}\pi^-) \simeq 2.7 \times 10^{-3}$. However, the uncertainty in this prediction is quite large, particularly given that large N_c strictly only applies for w near 1. The same large N_c inputs predict $\mathcal{B}(\Lambda_b \to \Lambda_c^+ \ell^- \bar{\nu}) \approx 6\%$, i.e., this channel is expected to make up a large part of the inclusive $\Lambda_b \to X_c \ell^- \bar{\nu}$ rate, with the $s_l^P = 1^-$ excited Λ_c states making up a significant fraction of the remainder [14].

Order $\Lambda_{\rm QCD}/m_Q$ corrections to these predictions may be significant. The $\Lambda_b \to \Lambda_c^+ \pi^-$ amplitude receives contributions from the T, C, E, and B classes of diagrams in Fig. 1. In SCET, $|E/T|$ and $|C/T|$ are of order $\Lambda_{\rm QCD}/m_Q$ [15], and we will show later that $|B/T|$ is further suppressed. In $B \to D\pi$ decay we know from $\mathcal{B}(B^- \to D^0 \pi^-)/\mathcal{B}(\bar{B}^0 \to D^+ \pi^-) \simeq 1.8$ [6] that $\Lambda_{\rm QCD}/m_Q$ corrections affect the amplitudes at the 15–30% level. In particular $|A(\bar{B}^0 \to D^+ \pi^-)| = |T + E| = (5.9 \pm 0.3) \times 10^{-7}$ GeV and $|A(B^- \to D^0 \pi^-)| = |T + C| = (7.7 \pm 0.3) \times 10^{-7}$ GeV. The ratio of these amplitudes can be reproduced by a power correction involving a hadronic parameter $|s_{\rm eff}| \simeq 430$ MeV, which is of natural size [15]. Since $B_s \to D_s^- \pi^+$ only has a T contribution, accurate measurement of this rate will improve our understanding of the size of E and C. CDF recently measured $[f_s \mathcal{B}(B_s \to D_s^- \pi^+)]/[f_d \mathcal{B}(B^0 \to D^- \pi^+)] = 0.35 \pm 0.05_{\rm (stat)} \pm 0.04_{\rm (syst)} \pm 0.09_{\rm (BR)}$ [16], and using $f_s/f_d = 0.26 \pm 0.03$ yields $\mathcal{B}(B_s \to D_s^- \pi^+)/\mathcal{B}(B^0 \to D^- \pi^+) \simeq 1.35$. Neglecting $SU(3)$ breaking[3] this implies $|A(B_s \to D_s^- \pi^+)| = |T| = (7.3 \pm 1.5) \times 10^{-7}$ GeV and that $|C|$ and $|E|$ may be comparable. The errors are still too large to draw any definite conclusions.

Now we turn to $\Lambda_b \to \Sigma_c \pi$ decays. As shown in Table 1, there are two Σ_c states with different spin which we refer to as Σ_c and Σ_c^*. They form a heavy quark spin symmetry doublet with the spin and parity of the light degrees of freedom, $s_l^{\pi_l} = 1^+$. Under isospin, the Λ_b is $I = 0$, the $\Sigma_c^{(*)}$ is $I = 1$, and the Hamiltonian is $I = 1$, so the $\Sigma_c^{(*)}\pi$ final state must be $I = 1$ (it cannot be $I = 0$ or 2). Therefore the rates to the two different charge channels are equal,

$$\Gamma\left(\Lambda_b \to \Sigma_c^{(*)0} \pi^0\right) = \Gamma\left(\Lambda_b \to \Sigma_c^{(*)+} \pi^-\right). \quad (11)$$

Based on B decay data and the SCET power counting, we expect $\Gamma(\Lambda_b \to \Sigma_c^{(*)}\pi)$ to be up to about an order of magnitude smaller than $\Gamma(\Lambda_b \to \Lambda_c \pi)$, since the leading contributions to $\Lambda_b \to \Sigma_c^{(*)}\pi$ are power suppressed.

[2] Updating the parameters by fitting the mass splitting to give $\kappa = (0.411\ {\rm GeV})^3$, and using $m_N = \bar{\Lambda} = 0.8$ GeV (instead of M_N) for the mass of the light degrees of freedom leaves the exponent essentially unchanged.

[3] In the heavy quark limit of the T amplitude $SU(3)$ will be tested by the measurement of $B_s \to D_s \ell \bar{\nu}$.

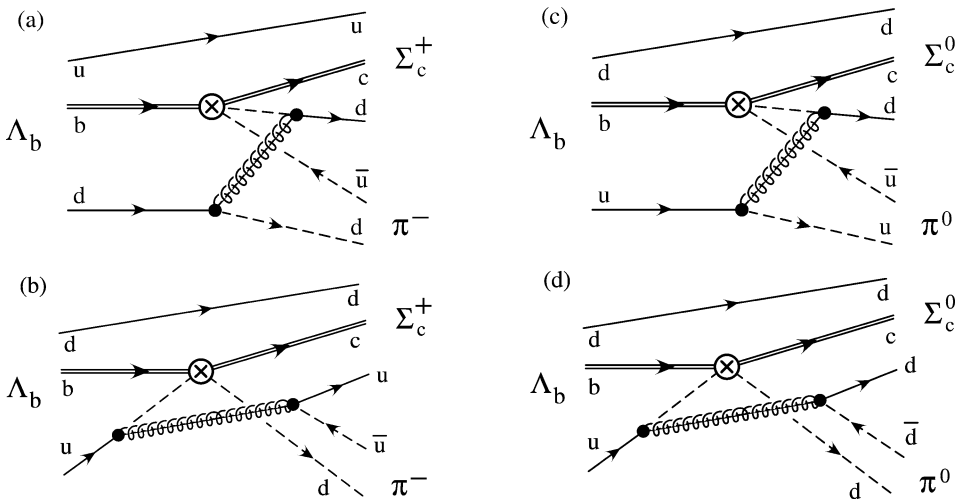

Fig. 2. Contributions in $\mathrm{SCET_I}$ to $\Lambda_b \to \Sigma_c^+ \pi^-$ [(a) and (b)], and to $\Sigma_c^0 \pi^0$ [(c) and (d)]. Solid dots denote insertions of the suppressed usoft-collinear Lagrangian, $\mathcal{L}_{\xi q}^{(1)}$, the double lines are heavy quarks, the dashed lines are collinear quarks, the solid lines are usoft quarks, and the "looped lines" are collinear gluons. The nonleptonic weak vertex is denoted by $\otimes$.

Again, we use SCET to expand in $\Lambda_{\rm QCD}/m_Q$, $\Lambda_{\rm QCD}/E_\pi$, and $\alpha_s(m_Q)$, keeping only the leading terms that cause the $\Lambda_b \to \Sigma_c^{(*)}\pi$ transitions. These come from the C and E diagrams in Fig. 1 and their contributions can be studied following the analysis of $\bar{B}^0 \to D^{(*)0}\pi^0$ in Ref. [15]. The leading diagrams in $\mathrm{SCET_I}$ that determine the matching onto power suppressed operators are shown in Fig. 2. To match the C and E diagrams, two insertions of the mixed usoft-collinear Lagrangian, $\mathcal{L}_{\xi q}^{(1)}$ [17], is required, each yielding a suppression of $\sqrt{\Lambda_{\rm QCD}/E_\pi}$. This yields the power counting that $|C/T|$ and $|E/T|$ are $\mathcal{O}(\Lambda_{\rm QCD}/E_\pi)$. In contrast, matching the B diagram in Fig. 1 requires four insertions of $\mathcal{L}_{\xi q}^{(1)}$ (or other higher-dimensional terms in the Lagrangian), and B is therefore power suppressed compared to C and E by at least an additional $\Lambda_{\rm QCD}/E_\pi$.

In addition there is a further matching onto $\mathrm{SCET_{II}}$. The resulting matrix element involves soft and collinear operators which factor [15].[4] The matrix element of the weak Hamiltonian, $\langle \Sigma_c^0(v',s')\pi^0|H_W|\Lambda_b(v,s)\rangle$, can be written (neglecting $\alpha_s(m_Q)$ corrections) as a convolution integral of a jet function, $J(x,k_1^+,k_2^+)$, with the matrix element involving the collinear fields, $\langle \pi|O_c(x)|0\rangle$ which gives $\phi_\pi(x)$, and that involving the soft fields, $\langle \Sigma_c(v',s')|O_s(k_j^+)|\Lambda_b(v,s)\rangle$. In what follows we only need the form of the soft operator [15]

$$O_s\big(k_j^+\big) = \big[\big(\bar{h}_{v'}^{(c)} S\big)\not{n} P_L\big(S^\dagger h_v^{(b)}\big)\big]\big[(\bar{d}S)_{k_1^+}\not{n} P_L\big(S^\dagger u\big)_{k_2^+}\big], \tag{12}$$

where $h_v^{(Q)}$ is an HQET heavy quark field, S is a soft Wilson line, and the subscripts denote the momentum carried by the fields. For our purposes the most important aspect of the analysis is that O_s only involves combination $\bar{h}_{v'}^{(c)}\not{n} P_L h_v^{(b)}$. Thus, by heavy quark symmetry

$$\begin{aligned}
&\big\langle \Sigma_c(v',s')\big|O_s\big|\Lambda_b(v,s)\big\rangle \\
&\quad = \frac{1}{\sqrt{3}}\bar{u}_{\Sigma_c}(v',s')\big(\gamma^\mu - v'^\mu\big)\gamma_5\not{n} P_L u_{\Lambda_b}(v,s) X_\mu, \\
&\big\langle \Sigma_c^*(v',s')\big|O_s\big|\Lambda_b(v,s)\big\rangle \\
&\quad = \bar{u}^\mu_{\Sigma_c^*}(v',s')\not{n} P_L u_{\Lambda_b}(v,s) X_\mu,
\end{aligned} \tag{13}$$

where v and v' are the four-velocities of the Λ_b and $\Sigma_c^{(*)}$, respectively. The spinor field normalizations are $\bar{u}(v,s)u(v,s) = 1$ for the Λ_b and Σ_c, and $\bar{u}_\alpha(v,s)u^\alpha(v,s) = -1$ for the Σ_c^*. X_μ is the most gen-

[4] Since the messenger modes from Ref. [18] do not spoil factorization for cases with a product of color singlet soft and collinear operators, we can neglect them in our analysis.

eral vector compatible with the symmetries of QCD,

$$X_\mu = a n_\mu + b v_\mu + c v'_\mu. \tag{14}$$

Note that in Eq. (12) the part of O_s involving the light quark fields is parity violating, so we need not worry about the fact that $\Lambda_b \to \Sigma_c$ is an "unnatural" transition. Using $m_{\Lambda_b} v = m_{\Sigma_c} v' + E_\pi n$ to eliminate the term proportional to v_μ in Eq. (14), it is easy to see that any term in X_μ proportional to v'_μ does not contribute, so only n_μ remains. Hence the ratio of the rates for $\Lambda_b \to \Sigma_c \pi$ and $\Lambda_b \to \Sigma_c^* \pi$ are determined model independently at leading order in $\Lambda_{\rm QCD}/m_Q$ and $\alpha_s(m_Q)$, similar to the $\bar{B}^0 \to D^{(*)0}\pi^0$ case. We find

$$\frac{\Gamma(\Lambda_b \to \Sigma_c^* \pi)}{\Gamma(\Lambda_b \to \Sigma_c \pi)} = 2 + \mathcal{O}\big[\Lambda_{\rm QCD}/m_Q, \alpha_s(m_Q)\big]. \tag{15}$$

To evaluate the square of the matrix element in Eq. (13), we used the spin sums from Ref. [14] for the various spin $\Sigma_c^{(*)}$ states. The explicit calculation shows that the rate to Σ_c^* with $|s'| = 3/2$ vanishes, as required by angular momentum conservation.

A practical complication in testing this prediction is that the $\Sigma_c^{(*)}$ states decay to $\Lambda_c \pi$, and so both decay channels $\Lambda_b \to \Sigma_c^{(*)0}\pi^0 \to \Lambda_c \pi^- \pi^0$ and $\Lambda_b \to \Sigma_c^{(*)+}\pi^- \to \Lambda_c \pi^0 \pi^-$ contain a π^0 that makes the reconstruction hard at hadron colliders. This can be circumvented by studying $\Lambda_b \to \Sigma_c^{(*)0}\rho^0$ decays. In this case the final states are $\Lambda_c \pi^- \pi^+ \pi^-$. Decays to a vector meson are potentially more complicated due to the fact that "long-distance" contributions can induce transverse polarizations at the same order in $\Lambda_{\rm QCD}/E_\pi$. However, at leading order in $\alpha_s(m_Q)$ these long-distance contributions vanish for the ρ^0 final state [15] and we obtain

$$\frac{\Gamma(\Lambda_b \to \Sigma_c^{*0}\rho^0)}{\Gamma(\Lambda_b \to \Sigma_c^0 \rho^0)} = 2 + \mathcal{O}\big[\Lambda_{\rm QCD}/m_Q, \alpha_s(m_Q)\big]. \tag{16}$$

It is also worth noting that

$$\begin{aligned}\frac{\Gamma(\bar{B}^0 \to D^0\pi^0)}{\Gamma(\bar{B}^0 \to D^0\rho^0)} &= \frac{\Gamma(\Lambda_b \to \Sigma_c^0 \pi^0)}{\Gamma(\Lambda_b \to \Sigma_c^0 \rho^0)} \\ &\quad + \mathcal{O}\big[\Lambda_{\rm QCD}/m_Q, \alpha_s(\sqrt{m_Q \Lambda_{\rm QCD}})\big],\end{aligned} \tag{17}$$

where in contrast to Eqs. (15) and (16) this prediction requires a perturbative expansion at the intermediate scale $\sqrt{\Lambda_{\rm QCD} m_Q}$.

The decays $\Lambda_b \to \Xi_c K$ decays are also Cabibbo-allowed. (These decays involve "$\bar{s}s$ popping" so only $\Xi_c^0 K^0$ is allowed, not $\Xi_c^+ K^-$.) They are similar to $\Lambda_b \to \Sigma_c^{(*)}\pi$ in the sense that the leading contribution in the heavy quark limit vanishes. As shown in Table 1 there are three Ξ_c "ground states", Ξ_c, Ξ_c', and Ξ_c^*. The Ξ_c and Ξ_c' can mix, but the former is expected to be mostly the state that transforms as $\bar{\mathbf{3}}$ under flavor $SU(3)$, while the latter is mostly a $\mathbf{6}$. The Ξ_c^* also transforms as a $\mathbf{6}$, and forms a heavy quark spin symmetry doublet with the Ξ_c'. Thus, a relation similar to Eq. (15) also holds in this case, i.e., $\Gamma(\Lambda_b \to \Xi_c^* K)/\Gamma(\Lambda_b \to \Xi_c' K) = 2 + \mathcal{O}[\Lambda_{\rm QCD}/m_Q, \alpha_s(m_Q)]$. This prediction may be hard to test since Ξ_c' decays to $\Xi_c \gamma$. One can also consider Cabibbo-suppressed Λ_b decays, e.g., $\Lambda_b \to \Xi_c \pi$, and the weak decays of other baryons containing a heavy bottom quark.

Perhaps the most exciting possibility is the existence of heavy baryonic pentaquark states. Recently several experiments claimed to observe a baryon $\Theta^+(1540)$ with the quantum numbers of $K^+ n$. A possible explanation is to consider the Θ^+ as a bound state of two spin-zero ud diquarks in a P-wave with an $\bar{s}$ antiquark [19]. If diquarks play an important role in these exotic states then the analogous heavy flavor states, $\Theta_c = \bar{c}[ud]^2$ and $\Theta_b = \bar{b}[ud]^2$, may be below threshold for strong decays by $\Delta E \simeq -100$ MeV and $\Delta E \simeq -160$ MeV, respectively [19].[5] Since the spin of the light degrees of freedom is $s_l = 1$, we expect from heavy quark symmetry that Θ_Q come with a doublet partner of similar mass, Θ_Q^*, as shown in Table 1, with a mass splitting of order $\Lambda_{\rm QCD}^2/m_Q$. From the mass splittings for the Σ_c and Ξ_c we expect $m_{\Theta_c^*} - m_{\Theta_c} \sim 70$ MeV and $m_{\Theta_b^*} - m_{\Theta_b} \sim 22$ MeV. In this case the Θ_Q^* may also be stable with respect to the strong interactions and decay to $\Theta_Q \gamma$. Since the splitting for $\Theta_c^{(*)}$ is larger, it is possible that the

[5] It is possible that the Θ_Q are above the strong decay thresholds [20]. The assumptions in our analysis are that (i) Θ_Q decay weakly; and (ii) the spin of the light degrees of freedom is $s_l = 1$, as suggested by [19]. If (i) is correct but (ii) is not, it would be easy to modify our predictions, including Eq. (19).

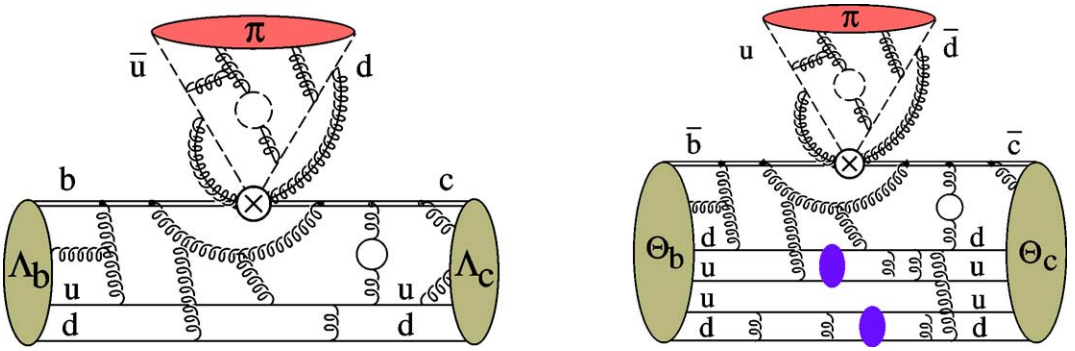

Fig. 3. Comparison of the weak nonleptonic decays of the Λ_b and Θ_b.

Θ_c^* is just above the strong decay threshold, making the spectroscopy even more interesting (like in D^* decays).

It may be possible to discover the $\Theta_{b,c}$ via the decay chains

$$\Theta_b^+ \to \Theta_c^0 \pi^+,$$
$$\Theta_c^0 \to \Theta^+ \pi^- \to K_S p \pi^- \to \pi^+ \pi^- p \pi^- \tag{18}$$

that are Cabibbo-allowed and lead to all charged final states. The most interesting aspect of the $\Theta_b^+ \to \Theta_c^0$ decay is that in the diquark picture the correlation is maintained, as shown in Fig. 3, and so no additional suppression factor is expected. In weak Θ_b decays to ordinary baryons this would not be the case. While we do not know the Θ_Q production rates, we can estimate the branching ratios in Eq. (18). The lifetime of a weakly decaying $\Theta_{b,c}$ is expected to be comparable with other weakly decaying hadrons that contain a charm or a bottom quark. The $\Theta_b^+ \to \Theta_c^0 \pi^+$ amplitude factorizes, and is related to $\Theta_b^+ \to \Theta_c^0 \ell \bar{\nu}$ via a formula identical to Eq. (5). For the nonleptonic rate we obtain in the heavy quark limit

$$\begin{aligned}
&\frac{\Gamma(\Theta_b^+ \to \Theta_c^0 \pi^+)}{\Gamma(\Lambda_b \to \Lambda_c \pi^-)} \\
&\quad = \frac{m_{\Theta_b}^3 (1 - r_\Theta^2)^3}{m_{\Lambda_b}^3 (1 - r_\Lambda^2)^3} \frac{1}{\zeta(w_{\max}^\Lambda)^2} \frac{1}{144 r_\Theta^4} \\
&\qquad \times \Big\{ 4\big[\eta_1(w_{\max}^\Theta)\big]^2 r_\Theta^2 \big(1 + 18 r_\Theta^2 + r_\Theta^4\big) \\
&\qquad\quad - 4\eta_1(w_{\max}^\Theta)\eta_2(w_{\max}^\Theta) r_\Theta \big(1 - r_\Theta^2\big)^2 \big(1 + r_\Theta^2\big) \\
&\qquad\quad + \big[\eta_2(w_{\max}^\Theta)\big]^2 \big(1 - r_\Theta^2\big)^4 \Big\},
\end{aligned} \tag{19}$$

where $r_\Theta = m_{\Theta_c}/m_{\Theta_b}$, and $\eta_{1,2}$ are the two Isgur–Wise functions that parameterize the weak $\Theta_b \to \Theta_c^{(*)}$ matrix elements where $\eta_1(1) = 1$. In particular

$$\begin{aligned}
&\big\langle \bar{\Theta}_c(v', s') \big| \bar{c} \Gamma b \big| \bar{\Theta}_b(v, s) \big\rangle \\
&\quad = \frac{1}{3}\big[g^{\alpha\beta}\eta_1(w) - v^\alpha v'^\beta \eta_2(w)\big] \bar{u}(v', s') \\
&\qquad \times \gamma_5 (\gamma_\alpha + v'_\alpha) \Gamma (\gamma_\beta + v_\beta) \gamma_5 u(v, s).
\end{aligned} \tag{20}$$

Thus, $\mathcal{B}(\Theta_b^+ \to \Theta_c^0 \pi^+)$ is expected to be similar to $\mathcal{B}(\Lambda_b \to \Lambda_c \pi)$. If the Θ_Q states exist then an analysis of the $\Lambda_{\rm QCD}/m_Q$ corrections would be warranted, as the mass of the light degrees of freedom is sizable. We expect $\mathcal{B}(\Theta_c^0 \to \Theta^+ \pi^-)$ to be at the few percent level, while the other branching ratios in Eq. (18) may be of order unity.

In summary, we studied nonleptonic Λ_b decays to $\Lambda_c \pi$, $\Sigma_c \pi$ and $\Sigma_c^* \pi$. Eqs. (10), (15), (16), and (19) are our main results. In the $m_Q \gg \Lambda_{\rm QCD}$ limit the $\Lambda_b \to \Lambda_c \pi$ rate is related to $\Lambda_b \to \Lambda_c \ell \bar{\nu}$, and we found that $\Gamma(\Lambda_b \to \Lambda_c \pi)$ is expected to be larger than $\Gamma(B \to D^{(*)} \pi)$, as observed by CDF. At leading order in $\Lambda_{\rm QCD}/m_Q$ the $\Lambda_b \to \Sigma_c^{(*)} \pi$ rates vanish, but an analysis of the leading contributions suppressed by $\Lambda_{\rm QCD}/m_Q$ was still possible. We predict $\Gamma(\Lambda_b \to \Sigma_c^* \pi)/\Gamma(\Lambda_b \to \Sigma_c \pi) = \Gamma(\Lambda_b \to \Sigma_c^* \rho)/\Gamma(\Lambda_b \to \Sigma_c \rho) = 2 + \mathcal{O}[\Lambda_{\rm QCD}/m_Q, \alpha_s(m_Q)]$. We also discussed properties of pentaquarks with a $\bar{b}$ or $\bar{c}$, including a possible discovery channel if they decay weakly.

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