# ASPECTS OF THE CHIRALITY OF WEAK b→c TRANSITIONS IN THE EXCLUSIVE SEMI-LEPTONIC DECAY $B(b)\to D^*(c)+\ell^-+\bar{\nu}_\ell$

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We consider the four-body angular decay distribution in the cascade decay $\overline{B^0}\to D^{*+}(\to D\pi)+\ell^-+\bar{\nu}_\ell$. We identify the parity violating (PV) angular contributions that are sensitive to the difference $H_+-H_-$ of the transverse current-induced $B\to D^*$ helicity amplitudes. The sign of these PV contributions directly reflects the chirality of the underlying weak b→c transition. We define two asymmetry observables that are sensitive to the PV difference $H_+-H_-$. The corresponding asymmetries are calculated in a specific spectator quark model.

In the standard model of electroweak interactions the weak b→c transition is left-chiral, i.e. the c-quark leaves the weak interaction vertex with predominant negative helicity. It is then an interesting and important question to ask whether and how this helicity information is transmitted to the final hadron or hadrons into which the c-quark hadronizes. It is clear that exclusive one-hadron decays are best suited for such an analysis since the helicity information is least degraded in the exclusive one-hadron hadronization.

In this paper we study some aspects of this interesting issue in the context of the exclusive semileptonic bottom meson decays $B\to D^*+\ell+\nu_\ell$. At the end of this paper we shall also briefly comment on some interesting consequences of the left-chirality of the b→c transitions in the exclusive non-leptonic bottom meson decays to baryon–antibaryon pairs where the left-chirality of the weak b→c transition leads to a number of interesting sum rules and selection rules.

To start with consider the full four-fold decay distribution in the cascade decay $\overline{B^0}\to D^{*+}(\to D\pi)+\ell^-+\bar{\nu}_\ell$. One has [1]

$$\frac{d\Gamma(\overline{B^0}\to D^{*+}(\to D\pi)+\ell^-+\bar{\nu}_\ell)}{dq^2\, d\cos\theta\, d\chi\, d\cos\theta^*} = B(D^{*+}\to D\pi)\,\frac{1}{2\pi}\bigg(\tfrac{3}{8}(1+\cos^2\theta)\tfrac{3}{4}\sin^2\theta^*\frac{d\Gamma_U}{dq^2}+\tfrac{3}{4}\sin^2\theta\,\tfrac{3}{2}\cos^2\theta^*\frac{d\Gamma_L}{dq^2}$$
$$-\tfrac{3}{4}\sin^2\theta\cos2\chi\,\tfrac{3}{4}\sin^2\theta^*\frac{d\Gamma_T}{dq^2}-\tfrac{9}{16}\sin2\theta\cos\chi\sin2\theta^*\frac{d\Gamma_I}{dq^2}+\tfrac{3}{4}\cos\theta\,\tfrac{3}{4}\sin^2\theta^*\frac{d\Gamma_P}{dq^2}$$
$$-\tfrac{9}{8}\sin\theta\cos\chi\sin2\theta^*\frac{d\Gamma_A}{dq^2}\bigg), \qquad (1)$$

where $q^2$ is the invariant momentum squared. $\theta$ is the polar angle of the lepton measured with respect to the D*-direction in the $(\ell^-\bar{\nu}_\ell)$ CM system (see fig. 1), $\chi$ is the azimuthal angle between the two decay planes spanned by $(D\pi)$ and $(\ell^-\bar{\nu}_\ell)$ (see fig. 1), and $\theta^*$ is the polar angle of the D relative to the D* in the D* rest frame[#1]. $B(D^{*+}\to D\pi)$ is the branching ratio $\Gamma_{D^{*+}\to D\pi}/\Gamma_{D^{*+}\to \mathrm{all}}$. We have used the zero lepton mass approxi-

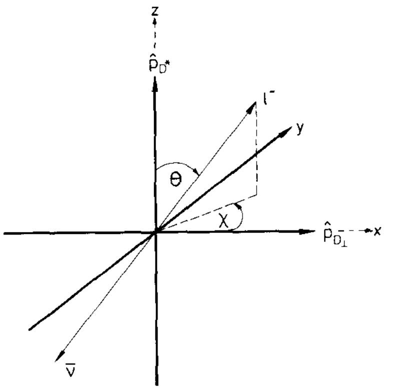


Fig. 1. Definition of the polar and azimuthal angles $\theta$ and $\chi$ of the lepton $\ell^-$ in the $(\ell^-\bar{\nu}_\ell)$ CM frame. $z$-axis is along $\boldsymbol{p}_{D^*}$ and $x$-axis in the $(\boldsymbol{p}_{D^*}, \boldsymbol{p}_D)$ plane with $p_{Dx} \geqslant 0$.

mation to derive the distribution (1). The general non-zero lepton mass case is discussed in ref. [1].

The partial helicity rates $d\Gamma_i/dq^2$ can be expressed in terms of reduced hadron tensor components $\hat{H}_i$, viz.

$$\frac{d\Gamma_i}{dq^2} = \frac{G^2}{(2\pi)^3} |V_{bc}|^2 \frac{pq^2}{12M_1^2} \hat{H}_i , \tag{2}$$

where $\hat{H}_i$ (i=U, L, T, I, P, A) are bilinear expressions of the three helicity amplitudes $H_+$, $H_-$ and $H_0$ describing the current induced transitions $B \to D^*$, see eq. (4). One has [1]

$$\begin{aligned}
&\hat{H}_U = |H_+|^2 + |H_-|^2, && \text{unpolarized transverse,}\\
&\hat{H}_L = |H_0|^2, && \text{longitudinal,}\\
&\hat{H}_T = \mathrm{Re}(H_+ H_-^*), && \text{transverse interference,}\\
&\hat{H}_I = \tfrac{1}{2}\mathrm{Re}(H_+ H_0^* + H_- H_0^*) && \text{transverse–longitudinal interference}\\
&\hat{H}_P = |H_+|^2 - |H_-|^2, && \text{parity-odd,}\\
&\hat{H}_A = \tfrac{1}{2}\mathrm{Re}(H_+ H_0^* - H_- H_0^*), && \text{parity-asymmetric}.
\end{aligned} \tag{3}$$

The labeling of the reduced hadron tensor components $\hat{H}_i$ in terms of the polarization components of the gauge boson $W_{\overline{\text{off-shell}}}$ follows the conventional notation of one-gauge-boson exchange physics.

Further symbols used in eq. (2) are: $M_1$ denotes the mass of the $\overline{B}^0$, $p$ is the momentum of the $D^{*+}$ in the $\overline{B}^0$ rest frame, $G$ is the fermi coupling constant $G \cong 1.02\, m_p^{-2} \times 10^{-5}$, and $V_{bc}$ is the $b \to c$ KM matrix element. Note that we have dropped angular terms in eq. (1) that are multiplied by $\mathrm{Im}(H_i H_j^*)$, $i \neq j$, assuming that the three helicity amplitudes are relatively real.

Turning now to the dynamics of the process one qualitatively expects $|H_-| > |H_+|$, i.e. dominance of the transverse negative helicity of the $D^*(c)$ over the transverse positive helicity. The c-quark emerges from the left-chiral weak interaction vertex with dominant negative helicity. After recombining with the spectator quark (which has equal probability of both helicities) to form the $D^*(c)$, one concludes that the $D^*(c)$ must have dominant transverse negative helicity irrespective of the details of the underlying quark model dynamics that is used to describe the hadronization phase.

[#1] The polar angle $\theta^*$ distribution of $D^* \to D\pi$ has also been considered in ref. [2]. The relevant $\theta^*$ decay distribution can be obtained from (1) by $\theta$ and $\chi$ integration.

Two of the angular contributions in the distribution eq. (1) are sensitive to this chirality effect. They are the parity-odd and parity-asymmetric contributions that multiply the two parity-violating (PV) bilinear terms $\hat{H}_P$ and $\hat{H}_A$ that are antisymmetric under chirality exchange $H_+ \leftrightarrow H_-$. A measurement of the sign of the coefficients of these two angular terms would clearly establish the "chirality" of the D* and thereby the chirality of the underlying b→c transition.

In order to be more quantitative let us turn to the spectator quark model approach to exclusive semileptonic bottom meson decays developed in ref. [3][#2]. Let us briefly recall the salient features of the KS model [3]. One matches the particle helicity amplitudes to the free quark decay helicity amplitudes at $q^2=0$, assuming that the spectator quark is spin-inert. Thus one has

$$H_0 = \langle D^{*+}(\text{long.}) | J_0 | \overline{B^0} \rangle \simeq \tfrac{1}{2} I \langle c\downarrow | J_0 | b\downarrow \rangle ,$$

$$H_{-(+)} = \langle D^{*+}\downarrow(\uparrow) | J_{-(+)} | \overline{B^0} \rangle \simeq +(-) \frac{1}{\sqrt{2}} I \langle c\downarrow(\uparrow) | J_{-(+)} | b\uparrow(\downarrow) \rangle , \tag{4}$$

where the factors $\frac{1}{2}$ $(1/\sqrt{2})$ are spin projection factors and $I$ is to be interpreted as the B→D* wave function overlap.

We have specified eq. (4) to the particular decay case $\overline{B^0}(b) \to D^{*+}(c) + \ell^- + \bar{\nu}_\ell$ involving the b→c transition. The case $B(\bar{b}) \to D^*(\bar{c}) + \ell^+ + \nu_\ell$ can be discussed along similar lines [1].

The above matching procedure leads to the $q^2=0$ ratios [3]

$$\sqrt{q^2}\, H_0(0) : H_-(0) : H_+(0) = (M_1^2 - M_2^2) : 2M_1 : 2M_2 , \tag{5a}$$

where $M_1$ and $M_2$ are the $\overline{B^0}$ and $D^{*+}$ masses, respectively. $q^2$ is the 4-momentum transfer squared $q^2 = (p_1 - p_2)^2$. Equivalently, in terms of a standard set of invariant form factors (see e.g. refs. [1,3]) one has

$$F_1^A(0) : F_2^A(0) : F^V(0) = (M_1 + M_2) : \frac{-2}{M_1 + M_2} : \frac{-2}{M_1 + M_2} \tag{5b}$$

Similar ratios are found in the approach of ref. [4].

The invariant form factors $F_1^A$, $F_2^A$ and $F^V$ are then continued to $q^2 \neq 0$ by using pole type form factors with a power behaviour given by the QCD power counting rules [3].

All the hadron dynamics of the SL decay process is contained in the six $q^2$-dependent helicity rate functions $d\Gamma_i/dq^2$. In fig. 2 their $q^2$-dependence is plotted as given in the KS model [3] with a wave function overlap value of $I=0.7$ as in ref. [4].

The unpolarized transverse rate function $d\Gamma_U/dq^2$ dominates for $q^2 \gtrsim 4.5\ \text{GeV}^2$ but is suppressed towards smaller $q^2$-values due to a spin-kinematical $q^2$-factor (see refs. [1,3]) where the longitudinal contribution dominates. The interference rates $d\Gamma_i/dq^2$ ($i=$T, I, P, A) are smaller than $d\Gamma_U/dq^2$ and $d\Gamma_L/dq^2$ over most of the available $q^2$-range but are nevertheless non-negligible. For the integrated total helicity rates $\Gamma_i$ one obtains

$$(\Gamma_U, \Gamma_L, \Gamma_T, \Gamma_I, \Gamma_P, \Gamma_A) = (12.7,\ 13.1,\ 5.3,\ 8.2,\ -6.9,\ -2.6) \times |V_{bc}|^2 \times 10^{12}\ \text{s}^{-1} . \tag{6}$$

The transition rate given by $\Gamma_U + \Gamma_L$ and the polarization of the D* (which is sensitive to the ratio $\Gamma_L/\Gamma_U$) and their dependence on the form factors have extensively been studied in the literature [2]. E.g. the ratio $\Gamma_L/\Gamma_U$ was measured by the ARGUS Collaboration in the polar angle distribution of D*→Dπ to be $0.85 \pm 0.45$ [5] which is consistent with the theoretical value $\Gamma_L/\Gamma_U = 1.03$ calculated from (6) and with the models of refs. [2,4].

In this paper we mainly concentrate on the two PV rate functions $d\Gamma_P/dq^2$ and $d\Gamma_A/dq^2$. An inspection of the angular decay distribution (1) shows that their contributions can be projected out by defining the following asymmetry ratios:

[#2] For brevity's sake we shall refer to the model of ref. [3] as the KS model.

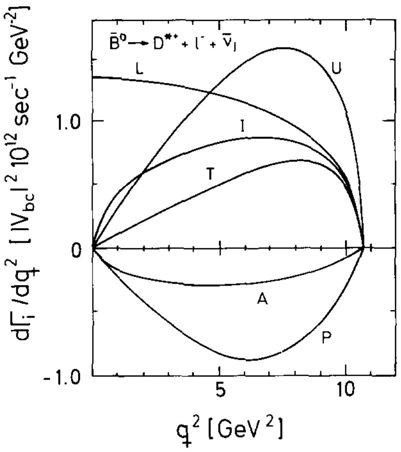


Fig. 2. Helicity rates $d\Gamma_i/dq^2$ for $W_{\text{off-shell}}$ polarization components $i$=U, L, T, I, P, A in the KS model [3].

$$\text{P:}\quad A_{FB}=\frac{d\Gamma(\theta)-d\Gamma(\pi-\theta)}{d\Gamma(\theta)+d\Gamma(\pi-\theta)},\quad \pi/2\leqslant\theta\leqslant\pi\,, \tag{7a}$$

$$\text{A:}\quad A_A=\frac{d\Gamma(\theta^*,\chi)-d\Gamma(\theta^*,\pi+\chi)-d\Gamma(\pi-\theta^*,\chi)+d\Gamma(\pi-\theta^*,\pi+\chi)}{d\Gamma(\theta^*,\chi)+d\Gamma(\theta^*,\pi+\chi)+d\Gamma(\pi-\theta^*,\chi)+d\Gamma(\pi-\theta^*,\pi+\chi)},\quad 0\leqslant\theta^*\leqslant\pi/2,\ -\pi/2\leqslant\chi\leqslant\pi/2. \tag{7b}$$

We have used a notation in eq. (7) where the angles that do *not* appear in the arguments of the differential rate $d\Gamma$ in (7) have been integrated out over their physical ranges ($0\leqslant\theta$, $\theta^*\leqslant\pi$, $0\leqslant\chi\leqslant 2\pi$). Integrating over the remaining variables (numerator and denominator separately!) we finally obtain the following values for the asymmetry ratios in the KS model:

$$A_{FB}=-\tfrac{3}{4}\Gamma_P/\Gamma_{U+L}=0.20\,, \tag{8}$$

$$A_A=-\tfrac{3}{2}\Gamma_A/\Gamma_{U+L}=0.15\,. \tag{9}$$

Present $e^+e^-$ experiments running on the $\Upsilon$(4S) produce bottom mesons which are practically at rest. Since leptons can only be detected and measured above a certain threshold energy $E_\ell^{\text{cut}}$ in these experiments (which is typically 0.5 GeV for electrons and 1.0 GeV for muons) this excludes the extreme backward region $\cos\theta\to 1$. The experimentally accessible angular range is then

$$-1\leqslant\cos\theta\leqslant\text{Min}(\cos\theta(q^2,E_\ell^{\text{cut}});1)\,, \tag{10}$$

where

$$\cos\theta(q^2,E_\ell^{\text{cut}})=-\frac{4M_1E_\ell^{\text{cut}}-M_1^2-q^2+M_2^2}{2M_1 p}\,. \tag{11}$$

It is interesting to consider asymmetries subject to the experimental constraint (10). One needs, however, to symmetrize the angular range (10) for the asymmetry definition. We shall accordingly define the asymmetries $A_{FB}$ and $A_A$ also in the symmetrized restricted angular range

$$-\text{Min}(\cos\theta(q^2,E_\ell^{\text{cut}});1)\leqslant\cos\theta\leqslant\text{Min}(\cos\theta(q^2,E_\ell^{\text{cut}});1)\,. \tag{12}$$

This is illustrated in fig. 3 where we show a plot of the double decay distribution $d\Gamma_p/dq^2\,d\cos\theta$ as a function of $q^2$ and $\cos\theta$. The forward–backward asymmetry is clearly evident in fig. 3a. Figs. 3c and 3d show the same

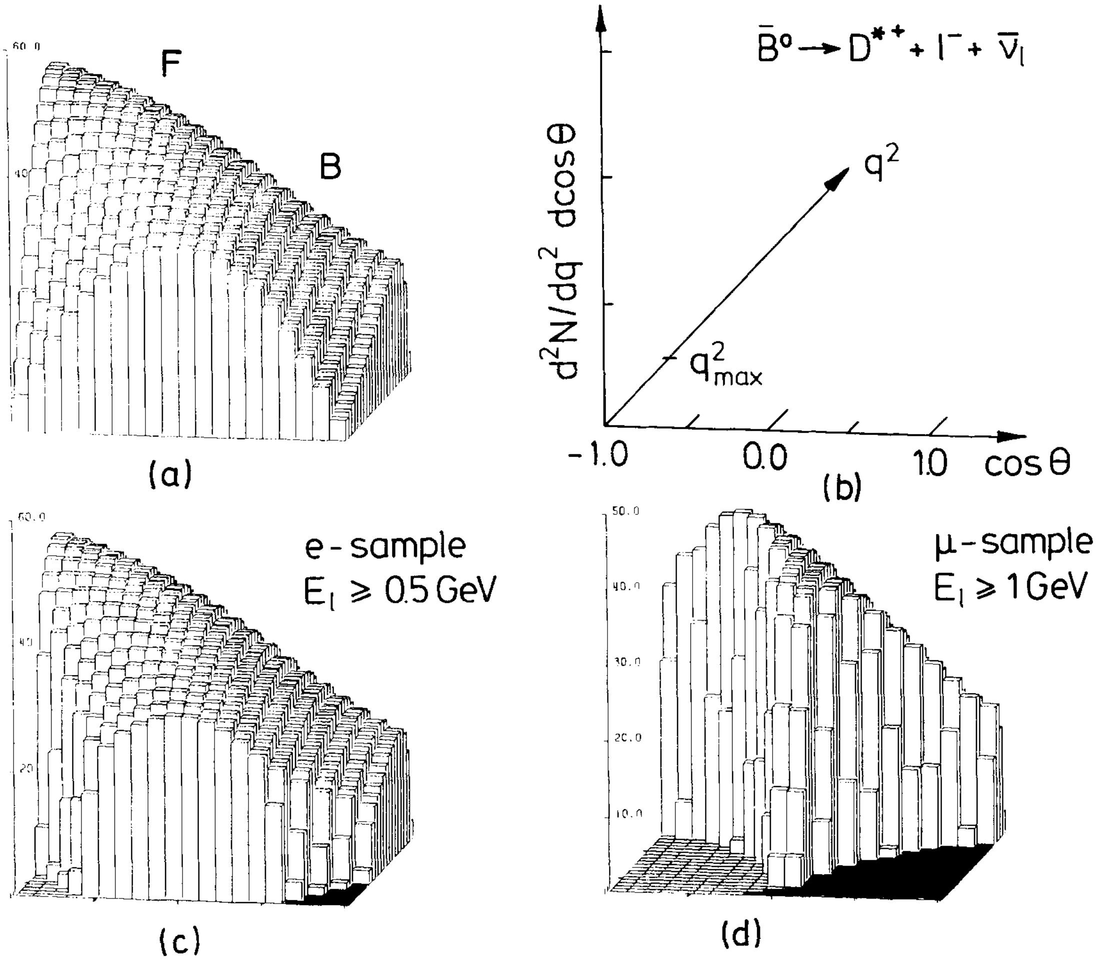


Fig. 3. Unnormalized double differential rate $d\Gamma/dq^2\, d\cos\theta$ for $\bar{B}^0 \to D^{*+} + \ell^- + \bar{\nu}_\ell$ in the KS model [3]. (a) Complete phase space, (b) coordinate scales, (c) and (d) symmetrized restricted phase space domains for (c) e-sample $E_\ell^{cut} = 0.5$ GeV, (d) μ-sample $E_\ell^{cut} = 1.0$ GeV. F: forward hemisphere $-1 \leqslant \cos\theta \leqslant 0$, B: backward hemisphere $0 \leqslant \cos\theta \leqslant 1$.

rate distribution but within the restricted symmetric Dalitz plot domain according to eq. (12). The asymmetric domain eq. (10) is also indicated in the plots by blacking out the relevant backward phase space region.

In fig. 4 we show a plot of the two asymmetry ratios $A_{FB}$ and $A_A$ as a function of $E_\ell^{cut}$. $E_\ell^{cut}$ defines the symmetrized restricted phase space domain (12) through eq. (11). The asymmetries rise from zero at $E_\ell^{cut} = (M_1 - M_2)/2 = 1{,}63$ GeV (zero phase space) to their largest values $A_{FB} = 0.20$ and $A_A = 0.15$, see eqs. (8), (9), at $E_\ell^{cut} = 0$ when the whole phase space domain is accessible for the asymmetry measurement.

The forward–backward asymmetry $A_{FB}$ drops to 0.19 and 0.15 from its maximal value 0.20 at the two cut values $E_\ell^{cut} \cong 0.5$ GeV and $E_\ell^{cut} \cong 1.0$ GeV for the e- and μ-samples, respectively. The asymmetry $A_A$ remains practically flat at its maximal value of 0.15 up to cut values of $E_\ell^{cut} \simeq 1.1$ GeV which includes the two above $E_\ell^{cut}$ cut values relevant to the e- and μ-samples.

Taking the KS model as a measure, only ≅5% and ≅40%, respectively, of the total SL rate $B(b) \to D^*(c)(\to D\pi) + \ell^- + \bar{\nu}_\ell$ have been forfeited at these two cut values. The two suggested asymmetry mea-

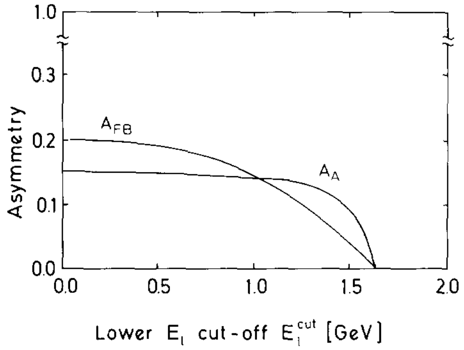


Fig. 4. Forward–backward asymmetry $A_{FB}$ and asymmetry $A$ as a function of lower energy cut $E_\ell^{cut}$ in the KS model [3].

surements with their expected asymmetry values in excess of 0.15 for realistic experimental cuts should be experimentally feasible in the near future.

A measurement of the sign of the asymmetries alone would be a significant experimental achievement as it would allow one to check on the chirality of the underlying b→c transition irrespective of the details of the underlying quark model that is used to describe the exclusive hadronization phase. Measuring in addition the magnitude of the asymmetries would provide a probe of the details of the bound state quark dynamics involved in the description of the current-induced $B\to D^*$ transition.

We further remark that a similar analysis of the exclusive SL decays of bottom mesons into the light vector mesons ω and ρ, and charmed mesons into the vector mesons $K^*$ (ρ, ω) would allow one to similarly conclude for the chirality of the fundamental b→u and c→s (c→d) current-induced quark transitions [1].

We would finally like to point out that there are also some interesting consequences of the left-chirality of the b→c (and b→u) transitions in the exclusive non-leptonic (NL) bottom meson decays into baryon–antibaryon pairs. In refs. [6,7] it was shown that the W-exchange and penguin contributions to the NL baryonic decays are small and therefore negligible. Thus the W-decay contributions dominate in the baryonic bottom meson decays.

For the b→u transitions the Cabibbo favoured mode involves the decay $b\to(ud)+\bar{u}$ where the diquark (ud) couples to the baryon in the final state. Assuming that the b→u transitions are left chiral one has an effective $(V-A)(V-A)$ current–current structure for the transition $b\to(ud)+\bar{u}$. In this situation the Körner–Pati–Woo (KPW) theorem applies which states that the (ud) diquark in the baryon is in an antisymmetric flavour state and thus has $I=0$ [8,9]. One then immediately concludes that the baryonic decay modes involving the $b\to(ud)+\bar{u}$ transition must satisfy a $\Delta I=1/2$ rule, and also that the decay into the ground state decuplet is forbidden, i.e. $B(b)\not\to 3/2^+ + X$, but $B(b)\to\overline{3/2^+}+X$ [6,7,10].

For the baryonic decay modes involving the quark decay $b\to(cd)+\bar{u}$ it is not immediately clear whether the KPW theorem applies to the (cd) diquark in the baryon because of possible flavour symmetry breaking effects. However, in ref. [6] it was shown that the flavour breaking inherent in the basic weak interaction process is compensated for by the flavour breaking of the baryon wave function with the consequence that the (cd) diquark is flavour antisymmetric. Thus, the KPW theorem also applies in this flavour breaking situation. One again obtains a number of interesting sum rules connecting different baryonic decay modes involving charmed baryons [6] that are too numerous to be listed here. Also one has the selection rule $B(b)\not\to 3/2^+ + X$, but $B(b)\to\overline{3/2^+}+X$, for the ground state $3/2^+$ charmed baryons [6].

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