Heavy Baryons and electromagnetic decays

I. Scimemi a, *

aDepartament de Física Teòrica, IFIC, Universitat de València – CSIC
E-46100 Burjassot (València), Spain

In this talk I review the theory of electromagnetic decays of the ground state baryon multiplets with one heavy quark calculated using Heavy Hadron Chiral Perturbation Theory [1]. The M1 and E2 amplitudes for $S^* \rightarrow S\gamma$, $S^* \rightarrow T\gamma$ and $S \rightarrow T\gamma$ are separately analyzed. All M1 transitions are calculated up to $\mathcal{O}(1/\Lambda^2_\chi)$. The E2 amplitudes contribute at the same order for $S^* \rightarrow S\gamma$, while for $S^* \rightarrow T\gamma$ they first appear at $\mathcal{O}(1/(m_Q\Lambda^2_\chi))$ and for $S \rightarrow T\gamma$ are completely negligible. Once the loop contributions is considered, relations among different decay amplitudes are derived. In ref. [1] it is shown that the coupling of the photon to light mesons is responsible of a sizable enhancement of these decay widths. Furthermore, one can obtain an absolute prediction for $\Gamma(\Xi^*_c \rightarrow \Xi\gamma)$.

1. Introduction

In Heavy Hadron Chiral Perturbation Theory (HHCPT) one constructs an effective Lagrangian whose basic fields are heavy hadrons and light mesons $\bar{q}q$ and $\bar{q}q'$. In ref. [1], the formalism is extended to include also electromagnetism. In this talk I describe how, using this formalism, one can calculate the electromagnetic decay width of some baryons containing a $c$ or a $b$ quark. The details of this computation are reported in ref. [1] and here I limit myself to trace its guidelines. In order to classify these baryons one observe that the light degrees of freedom in the ground state of a baryon with one heavy quark can be either in a $s_l = 0$ or in a $s_l = 1$ configuration. The first one corresponds to $J^P = 1/2^+$ baryons, which are annihilated by $T_i(v)$ fields which transform as a $3$ under the chiral $SU(3)_{L+R}$ and as a doublet under the HQET $SU(2)_v$. In the second case, $s_l = 1$, the spin of the heavy quark and the light degrees of freedom combine together to form $J^P = 3/2^+$ and $J = 1/2^+$ baryons which are degenerate in mass in the $m_Q \rightarrow \infty$ limit. The spin-$4/3$ ones are annihilated by the Rarita-Schwinger field $S^{(v)}_{\chi}(v)$ while the spin-$1/2$ baryons are destroyed by the Dirac field $S\gamma(v)$. They transform as a $6$ under $SU(3)_{L+R}$ and as a doublet under $SU(2)_v$, and are symmetric in the $i$, $j$ indices. I consider the decays $S^* \rightarrow S\gamma$ and $S^{(*)} \rightarrow T\gamma$. For most of these decays the available phase space is small, so that the emission of a pion is suppressed or even forbidden and the electromagnetic process becomes relevant. Moreover these kinds of decays are getting measured [1]. In the case of $S^* \rightarrow S\gamma$ all contributions up to order $\mathcal{O}(1/\Lambda^2_\chi)$ are calculated for M1 and E2 transitions. All divergences and scale dependence can be absorbed in the redefinition of one $\mathcal{O}(1/\Lambda_\chi)$ coupling for each type of process (M1, E2). Eliminating the unknown constants it is possible to find relations among the amplitudes which are valid up to the considered order.

An analogous calculation can be performed for $S^* \rightarrow T\gamma$. In this case, the E2 contribution has to be computed up to order $\mathcal{O}(1/m_Q\Lambda^2_\chi)$, implying the intervention of two new constants. Finally for $S \rightarrow T\gamma$ the M1 amplitude is calculated up to order $\mathcal{O}(1/\Lambda^2_\chi)$, while the E2 contribution is found to be extremely suppressed. In the case $S^{(*)} \rightarrow T\gamma$ it exists a process which do not receive any contribution from local terms in the Lagrangian.
grangian and therefore its width is described by a finite chiral loop calculation: \( \Gamma(\Xi_{c}^{0}(\tau) \to \Xi_{c}^{0}(\gamma)) \) (and analogously \( \Gamma(\Xi_{c}^{0}(\tau) \to \Xi_{c}^{0}(\gamma)) \)). In the following I comment these results and I refer to ref. \([1]\) for the formalism and for a more complete comparison with other results existing in literature. A similar formalism can be applied to the study of the magnetic moments of the same baryons \([3]\).

### 2. Results for \( S^{*} \to S \gamma \) decays

The decay amplitudes are decomposed by

\[
A(B^{*} \to B \gamma) = A_{M1} O_{M1} + A_{E2} O_{E2},
\]

where the corresponding M1 and E2 operators are defined by

\[
O_{M1} = e B \gamma_{\mu} \gamma_{5} B^{\mu}, \\
O_{E2} = i e B \gamma_{\mu} \gamma_{5} B^{\mu} v_{a} (\partial^{\mu} F^{a\nu} + \partial^{\nu} F^{a\mu}),
\]

The leading contributions to M1 transitions come from the light– and heavy–quark magnetic interactions which are of \( O(1/A_{\chi}) \) and \( O(1/m_{Q}) \), respectively. We have computed the next-to-leading chiral corrections of \( O(1/A_{\chi}^{2}) \), which originate from the loop diagrams shown in fig. \([3]\).

![Figure 1. Meson loops contributing to \( S^{*} \to S \gamma \).](image)

The resulting M1 amplitudes can be written as:

\[
A_{M1}(B^{*}) = \frac{1}{\sqrt{3}} \left( \frac{-Q_{Q}}{m_{Q}} - \frac{2c_{a}}{3A_{\chi}} a_{\chi}(B^{*}) \right) + \frac{g_{2}^{2}}{4(4\pi f_{+})^{2}} a_{g_{2}}(B^{*})
\]

| c quark | \( a_{\chi} \) | \( a_{g_{2}} \) |
|--------|----------------|----------------|
| \( \Sigma_{c}^{++} \to \Sigma_{c}^{++} \gamma \) | 2 | \( I_{\pi} + I_{K} \) |
| \( \Sigma_{c}^{+} \to \Sigma_{c}^{+} \gamma \) | 1/2 | \( I_{K}/2 \) |
| \( \Sigma_{c}^{0*} \to \Sigma_{c}^{0*} \gamma \) | -1 | \(-I_{\pi}\) |
| \( \Xi_{c}^{0} \to \Xi_{c}^{0} \gamma \) | -1 | \(-I_{K}\) |
| \( \Xi_{c}^{0*} \to \Xi_{c}^{0*} \gamma \) | 1/2 | \( I_{\pi}/2 \) |
| \( \Omega_{c}^{0} \to \Omega_{c}^{0} \gamma \) | -1 | \(-I_{K}\) |

Table 1 Contributions to M1 amplitudes for \( S^{*} \to S \gamma \). The values of \( a_{g_{2}} \) can be deduced from the ones of \( a_{g_{2}} \) with the substitution \( I_{\pi} \to m_{i}/m_{K} \) (\( i = \pi, K \)).

\[
+ \frac{g_{2}^{2}}{4(4\pi f_{+})^{2}} a_{g_{2}}(B^{*})
\]

\[
\text{In Table 1 we show the values of the coefficients } a_{\chi}(B^{*}) \text{ for the decays of baryons containing one charm or bottom quark. In the table, }
\]

\[
I_{i} \equiv I(\Delta_{ST}, m_{i}) = 2 \left( -2 + \log \frac{m_{i}^{2}}{\mu^{2}} \right) + 2 \frac{\sqrt{\Delta_{ST}^{2} - m_{i}^{2}}}{\Delta_{ST}} \times \log \left( \frac{\Delta_{ST} + \sqrt{\Delta_{ST}^{2} - m_{i}^{2}}}{\Delta_{ST} - \sqrt{\Delta_{ST}^{2} - m_{i}^{2}}} \right)
\]

where \( \Delta_{ST} \) is the mass difference between \( S \) and \( T \)–baryons. Due to flavor symmetry, all contributions are equal for charm and bottom baryons, with the only exception of the term proportional to the heavy quark electric charge \( Q_{c} = +2/3 \), \( Q_{b} = -1/3 \). The main things to be observed are the following:

- the corrections proportional to \( g_{2}^{2} \) are obtained performing a one–loop integral (fig. \([2]\) with an \( S \) baryon running in the loop) that has to be renormalized. It can be demonstrated \([1]\) that the scale \( \mu \) dependence of the loop integrals is exactly canceled by the corresponding dependence of the coefficient \( c_{S}(\mu) \);
- the contribution proportional to \( g_{2}^{2} \) involves a loop integral with a baryon of the \( T \) mul-
amplitude for done in the previous paragraph, we write the M1 defined as in Eq. (2). Similarly to what we have
constants are eliminated can be easily found. A complete list of them is reported in ref. [1].
Looking at table 1 one sees that relations among the decay amplitudes in which all unknown constant
A factor (E, /Λc) 2 ∼ 5%. In principle, it should be possible to determine experimentally the ratio
The E2 amplitudes come at higher chiral order which break spin symmetry, O
The E2 amplitudes come at higher chiral order to note that at this order it appears an operator,
which gives rise to divergent loop diagrams. Since at O(1/Λc 2 ) this decay does not get any contribution from local terms,
looking at table 1 one sees that relations among the decay amplitudes in which all unknown constants
A factor (E, /Λc) 2 ∼ 5%. In principle, it should be possible to determine experimentally the ratio
The E2 amplitudes come at higher chiral order which break spin symmetry, O
The E2 amplitudes come at higher chiral order which gives rise to divergent loop diagrams. Since at O(1/Λc 2 ) this decay does not get any contribution from local terms,
looking at table 1 one sees that relations among the decay amplitudes in which all unknown constants
A factor (E, /Λc) 2 ∼ 5%. In principle, it should be possible to determine experimentally the ratio

3. Results for S∗ → Tγ decays

The M1 and E2 operators for these decays are defined as in Eq. (2). Similarly to what we have done in the previous paragraph, we write the M1 amplitude for S∗ → Tγ decays as

\[ A_{M1}(B^*) = -\sqrt{2} \frac{c_{ST}}{\Lambda_c} a_2(B^*) \]

where the dominant error come from the uncertainty on g2,3.
The E2 amplitude in S∗ → Tγ is suppressed by an extra power of 1/mQ. The first non-zero contributions comes at O(1/mQΛc 2 ). It is important to note that at this order it appears an operator, which break spin symmetry,

\[ L' = i \frac{g'}{m_Q} \chi \left[ \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} (\xi_1)^j S^k_l + \epsilon_{ijk} S^k_l \sigma_{\mu\nu} (\xi_1)^j T_i \right] , \]

where the dominant error come from the uncertainty on g2,3.

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]

Both the contributions coming from eq. 8 which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

\[ -i \frac{g'}{m_Q} \epsilon_{ijk} T^{ij}_1 \sigma_{\mu\nu} Q_1 S^k_l \partial_\nu \bar{F}^{\mu\nu} . \]
is also given by a finite loop calculation. Unfortunately, since the coupling \( g' \) is not known, there is no absolute prediction in this case. An experimental measurement of these E2 amplitudes would provide a direct estimate of \( g' \).

4. Results for \( S \to T\gamma \)

The calculation of the M1 amplitude for \( S \to T\gamma \) decays is analogous to that of the previous section. Now the M1 operator is defined as

\[
O_{M1} = i e \bar{B}_T \sigma_{\mu\nu} B_S F^{\mu\nu}
\]

and the corresponding amplitude can be written in the form

\[
A_{M1}(B) = \frac{1}{\sqrt{6}} \frac{c_{ST}}{\Lambda_\chi} a_\chi(B) - g_2 g_3 \frac{\Delta_{ST} a_g(B)}{4\sqrt{6}(4\pi f_p)^2},
\]

where the coefficients satisfy

\[ a_\chi(B) = a_\chi(B^*), \quad a_g(B) = a_g(B^*) \quad (14) \]

Therefore, the relation (13) is also valid in this case. The widths of the decays \( \Xi^{0'}_c \to \Xi^{0}_c\gamma \) and \( \Xi^{-'}_b \to \Xi^{-}_b\gamma \) can be predicted through a finite loop calculation. From

\[
\Gamma(S \to T\gamma) = 16\alpha_{em} \frac{E_3^2 M_T}{M_S} |A_{M1}|^2,
\]

we find

\[
\Gamma(\Xi^{0'}_c) = (1.2 \pm 0.7) \text{ KeV},
\]

\[
\Gamma(\Xi^{-'}_b) = (3.1 \pm 1.8) \text{ KeV}. \quad (16)
\]

Again the dominant error in Eq. (16) is given by the uncertainty of \( g_{2,3} \).

For these decays the E2 amplitude is further suppressed than in the previous cases. The lowest–order contribution appears at \( O(1/m_Q^3 \Lambda_\chi^2) \) and, therefore, can be neglected.

REFERENCES

1. M.C. Banuls, A. Pich and I. Scimemi, Phys. Rev. D61 (2000) 094009.
2. M. Wise, Phys. Rev. D45 (1992) R2188.
3. G. Burdman and J. Donoghue, Phys. Lett. B280 (1992) 287.
4. T.M. Yan et al., Phys. Rev. D46 (1992) 1148; Erratum Phys. Rev. D55 (1997) 5851.
5. P. Cho, Phys. Lett. B285 (1992) 145.
6. P. Cho and H. Georgi, Phys. Lett. B296 (1992) 408; Phys. Lett. B300 (1993) 410 (E).
7. C.P. Jessop et al., CLEO Collaboration, Phys. Rev. Lett. 82 (1999) 492.
8. M.C. Banuls et al., Phys. Rev. D61 (2000) 074007.
9. M.J. Savage, Phys. Lett. B345 (1995) 61.
10. H.J. Rose and D.M. Brink, Rev. Mod. Phys. 39 (1967) 306.
11. M.N. Butler et al., Phys. Lett. B304 (1993) 353; B314 (1993) 122.
12. E-791 Collaboration (E.M. Aitala et al.), Phys. Lett. B379 (1996) 292.
13. Z. Guralnik, M. Luke and A.V. Manohar, Nucl. Phys. B390 (1993) 474.
14. A. Grozin and O.I. Yakovlev, Eur. Phys. J. C2 (1998) 721.
15. CLEO Collaboration, G. Brandenburg et al., Phys. Rev. Lett. 78 (1997) 2304.
16. H.-Y. Cheng, Phys. Lett. B399 (1997) 281.
17. M. Lu, M.J. Savage and J. Walden, Phys. Lett. B369 (1996) 337.