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Weak Radiative Decays of Beauty Baryons

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Abstract

Weak radiative decays of beauty baryons into strange baryons, induced by the electroweak penguin, are estimated by using a quark model approach. Relations between formfactors in the semileptonic and in the weak radiative decays are derived within the heavy quark effective theory. The partial decay widths are found to be of the order of 10^{-15}MeV for $\Lambda_b \to \Lambda\gamma$ and $\Xi_b \to \Xi\gamma$ and of the oder of 10^{-13}MeV for $\Omega_b \to \Omega\gamma$. The Ω_b radiative decay is thus expected at the sizable branching ratio of approximately 10^{-4}.

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The investigation of the electroweak penguin transition $b \rightarrow s\gamma$ is of prime importance, both as a test of the standard model[1] and as a possible window of new physics[2]. The recent observation of the exclusive process $B \rightarrow K^*\gamma$ with a branching ratio of $(4.5 \pm 1.5 \pm 0.9) \times 10^{-5}$[3] and of the inclusive process $B \rightarrow X_s\gamma$ with a branching ratio of $(2.32 \pm 0.57 \pm 0.35) \times 10^{-4}$[4] constitute solid evidence in the support of the interpretation of these decays in terms of the short-distance $b \rightarrow s\gamma$ transition.

The theoretical treatment of this basic loop process using renormalization group equations to treat perturbative QCD corrections has been improved during recent years[5], following the original assertion[6] on the role of the QCD corrections in increasing the rate of this process to bring it into the realm of observability. The full leading order calculation of the process has been completed[7] and the next-to-leading order calculation has also been partially performed[8]. At present, the theoretical uncertainty of the standard model calculations is of the order of 30%[9] and the experimental errors are even slightly higher[3, 4].

Along with the improvement in the theoretical calculations by the completion of the next-to-leading order calculation which should reduce the theoretical uncertainty, and the expected increased accuracies of future measurements, it is of obvious interest to investigate additional physical processes which are driven by the $b \rightarrow s\gamma$ transition. The weak radiative decays of beauty baryons are natural candidates for this task. Indeed, preliminary estimates on such decays were undertaken recently by Cheng *et al*[10] and by Cheng and Tseng[11] and an overview of this topic is given in Ref. [12].

In the present work, the weak radiative decays of the beauty baryons are investigated by the use of the quark model employed by Hussain *et al*[13] and by Körner and Krämer[14] to treat the semileptonic and the nonleptonic decays of heavy baryons. In their model, spin interactions between the spectator and active quarks are ignored and the q^2-dependence of all the formfactors are taken as pole-like. This approach is consistent with the heavy quark effective theory(HQET), as applied to baryons[15].

The lowest lying baryons containing one heavy beauty quark can be classified into $\bar{3}$, 6 and 6^* under the $SU(3)$ flavor symmetry of the light quarks[15, 16]. From among these states, only the spin 1/2 baryons Λ_b^0 , $\Xi_b^{-,0}$ of the antitriplet-baryons and Ω_b of the sextet-baryons are expected to decay weakly. We consider here the weak radiative exclusive processes which can be induced by the short-distance $b \to s\gamma$ transition; these are $\Lambda_b^0 \to \Lambda^0\gamma$, $\Xi_b^{-,0} \to \Xi^{-,0}\gamma$[10-12] and $\Omega_b^- \to \Omega^-\gamma$. Additional radiative decays like $\Lambda_b \to \Sigma_c\gamma$, $\Xi_b \to \Xi_c\gamma$, which are caused by weak bremsstrahlung quark processes, were found to be very rare[10] and will not be of our concern here.

The basic mechanism for the decays considered here is assumed to be the quark level transition $b \to s\gamma$, given by the following amplitude[5-9]

$$\mathcal{A} = \frac{G_F}{\sqrt{2}}\frac{e}{8\pi^2}c_7^{eff}(\mu)V_{tb}V_{ts}^*F_{\mu\nu}\bar{s}\sigma_{\mu\nu}[m_b(1+\gamma_5)+m_s(1-\gamma_5)]b, \tag{1}$$

where V_{tb} and V_{ts} are Cabibbo-Kobayashi-Maskawa matrix elements and $F_{\mu\nu}$ is the field strength tensor of the photon. The coefficient $c_7^{eff}(\mu)$, which is the combination of several Wilson coefficients running from $\mu \sim m_t$ to $\mu \sim m_b$, has been calculated[7-9] to be

$$c_7^{eff}(m_b = 4.5\text{GeV}) = 0.32, \tag{2}$$

when one uses $m_t = 174$GeV and $\Lambda_{QCD} = 200MeV$. A different choice for the renormalization point μ introduces the large uncertainty we mentioned.

In order to calculate the baryonic transitions induced by $b \to s\gamma$, our basic tool is the heavy quark effective theory which permits to relate[15, 17] in the heavy quark limit the formfactors of the magnetic transition (1) to those of the semileptonic decays. Moreover, we treat the decays under consideration here as heavy-to-light ($b \to s$)

quark transitions. Throughout the present paper, we neglect corrections in the mass parameter of the order $1/m_b$. Needless to say, a more accurate approach should consider the mass correction as well, especially when the above mentioned uncertainty in (2) will be reduced.

We turn now to derive the formfactors required here. For a baryonic transition induced by a $V-A$ current between spin 1/2 baryons, one generally has six formfactors. However, in the limit of the heavy quark mass $m_Q \to \infty$, using HQET one can express the relevant matrix element in terms of two independent formfactors only[13, 15]. Thus one has for $\Lambda_b \to \Lambda$ (and a similar expression for $\Xi_b \to \Xi$),

$$< \Lambda|\bar{s}\gamma_\mu(1-\gamma_5)b|\Lambda_b >= \bar{u}_\Lambda(P_2)[F_1(q^2) + F_2(q^2)\frac{\not{P}_1}{m_{\Lambda_b}}]\gamma_\mu(1-\gamma_5)u_{\Lambda_b}(P_1). \tag{3}$$

For the transition $\Omega_b \to \Omega$, where we have a heavy spin 1/2 baryon belonging to the 6-representation decaying into a spin 3/2 baryon belonging to the decuplet, the matrix element involves six independent formfactors[17], which for the $V-A$ current is

$$\begin{aligned} &< \Omega|\bar{s}\gamma_\mu(1-\gamma_5)b|\Omega_b > \\ &= \bar{u}_\Omega(P_2)_\alpha[g^{\alpha\beta}(C_1 + C_2\frac{\not{P}_1}{m_{\Omega_b}}) + \frac{{P_1}^\alpha}{m_{\Omega_b}}\frac{{P_1}^\beta}{m_{\Omega_b}}(C_3 + C_4\frac{\not{P}_1}{m_{\Omega_b}}) \\ &+ \frac{{P_1}^\alpha}{m_{\Omega_b}}\gamma^\beta(C_5 + C_6\frac{\not{P}_1}{m_{\Omega_b}})]\gamma_\mu(1-\gamma_5)\sqrt{\frac{1}{3}}(\gamma_\beta + \frac{P_{1\beta}}{m_{\Omega_b}})u_{\Omega_b}(P_1). \end{aligned} \tag{4}$$

Relating now by HQET[15] the formfactors in (3)(4) to the matrix elements for the radiative decays, and contracting these with the electromagnetic tensor $F_{\mu\nu}$, one obtains

$$F^{\mu\nu} < \Lambda|\bar{s}\sigma_{\mu\nu}(1+\gamma_5)b|\Lambda_b >= F^{\mu\nu}\bar{u}_\Lambda(P_2)[F_1(q^2) + F_2(q^2)\frac{m_\Lambda}{m_{\Lambda_b}}]\sigma_{\mu\nu}(1+\gamma_5)u_{\Lambda_b}(P_1) \tag{5}$$

for the weak radiative decay $\Lambda_b \to \Lambda\gamma$ (or $\Xi_b \to \Xi\gamma$), and likewise

$$F^{\mu\nu} < \Omega|\bar{s}\sigma_{\mu\nu}(1+\gamma_5)b|\Omega_b > \ = \ F^{\mu\nu}\bar{u}_\Omega(P_2)_\alpha[g^{\alpha\beta}(C_1 + C_2\frac{m_\Omega}{m_{\Omega_b}}) + \frac{P_1{}^\alpha}{m_{\Omega_b}}\frac{P_1{}^\beta}{m_{\Omega_b}}(C_3 + C_4\frac{m_\Omega}{m_{\Omega_b}})$$
$$+\frac{P_1{}^\alpha}{m_{\Omega_b}}\gamma^\beta(C_5 + C_6\frac{m_\Omega}{m_{\Omega_b}})]\sigma_{\mu\nu}(1+\gamma_5)\sqrt{\frac{1}{3}}(\gamma_\beta + \frac{P_{1\beta}}{m_{\Omega_b}})u_{\Omega_b}(P_1) \tag{6}$$

for $\Omega_b \to \Omega\gamma$. A possible alternative approach which was used before[10, 11, 18] is to establish firstly the relations between formfactors in the rest frame of the initial heavy hadron and then to boost them to a general Lorentz frame. However, certain difficulties occur[11] when carrying out this procedure for the $\Lambda_b \to \Lambda\gamma$, $\Xi_b \to \Xi\gamma$ decays.

In our approach to the weak radiative decays, we use eqs. (5)(6) with formfactors from the quark model as determined previously for the $(V-A)$ current induced transitions. To calculate the decay rates for the various processes, we proceed as follows. For the $\frac{1}{2} \to \frac{1}{2}\gamma$ transitions, we use (5) with the formfactors of Ref. [11] with monopole behavior and pole masses $M_V = 5.42$GeV and $M_A = 5.85$GeV, which gives

$$F_1(q^2 = 0) = 0.059(0.11), F_2(q^2 = 0) = -0.025(-0.019), \tag{7}$$

where the two given values are for the vector (axial) currents. Using now (7) and (1) with $V_{tb} = 1$, $V_{ts} = 0.04$, this leads to

$$\Gamma(\Lambda_b \to \Lambda\gamma) = 1.45 \times 10^{-15}\text{MeV}; \quad \Gamma(\Xi_b \to \Xi\gamma) = 2.18 \times 10^{-15}\text{MeV}. \tag{8}$$

If we use a lifetime of $\tau(\Lambda_b) = 1.04$ps[19], which has an experimental uncertainty of about 20%, we arrive at a predicted branching ratio for $\Lambda_b \to \Lambda\gamma$

$$BR(\Lambda_b \to \Lambda\gamma) \simeq 2.3 \times 10^{-6}. \tag{9}$$

A similar figure would be obtained for $\Xi_b \to \Xi\gamma$ decays, except that no measurement exists at present for the Ξ_b lifetime. Our result in (9) is smaller than the 10^{-5} order of magnitude obtained in [10], though quite close to the recent result of Ref. [11].

A similar approach is used for the decay $\Omega_b \to \Omega\gamma$. In this case, using the quark model of Ref.[13], we obtain

$$\begin{aligned}
C_1(q^2) &= \frac{(m_{\Omega_b}+m_\Omega)^2-q^2}{4m_{\Omega_b}m_\Omega}H(q^2),\\
C_3(q^2) &= -\frac{1}{2}H(q^2),\\
C_2(q^2) &= C_4(q^2)=C_5(q^2)=C_6(q^2)=0,\\
H(q^2) &= \sqrt{3}\frac{1-(m_{\Omega_b}-m_\Omega)^2/m_{pole}^2}{1-q^2/m_{pole}^2},
\end{aligned} \tag{10}$$

where again a monopole behavior, but with same pole mass of 5.42GeV has been used for all the formfactors. Taking a mass for Ω_b of 6.08GeV, we get

$$\Gamma(\Omega_b \to \Omega\gamma) = 1.63\times 10^{-13}\mathrm{MeV}. \tag{11}$$

It is interesting to remark that transitions involving $\Omega_c \to \Omega$ are similarly larger than those for $\Lambda_c \to \Lambda$ and $\Xi_c \to \Xi$ by one or two orders of magnitude in the decay rates[13, 14, 20]. These results are apparently a result of the strong ovelaps in the light flavour wavefunctions between the initial and final states in $\Omega_{b,c} \to \Omega$ transitions. Additional support for this picture is the fact that the lifetime of Ω_c, recently measured to be 0.055×10^{-12}sec[21], is considerably shorter than those of other weakly decaying charmed hadrons, which can also be understood qualitatively in this way[22]. The branching ratio for $\Omega_b \to \Omega\gamma$ is then

$$\mathrm{Br}(\Omega_b \to \Omega\gamma) = 1.3\times 10^{-4}[\frac{\tau(\Omega_b)}{0.5\times 10^{-12}\mathrm{sec}}], \tag{12}$$

where we have normalized the Ω_b lifetime to 0.5×10^{-12}sec. The branching ratio in (12) is larger by two orders of magnitude than that of $\Lambda_b \to \Lambda\gamma$ in (10). Thus, from the present calculation, the $\Omega_b \to \Omega\gamma$ decay mode is appearing as the most attractive baryonic exclusive channel for an alternative investigation of the $b \to s\gamma$ electroweak penguin.

We conclude by a discussion on additional advantages of studying the weak radiative decay of the beauty baryons. Among these processes, $\Xi_b^- \to \Xi^-\gamma$ and $\Omega_b \to \Omega\gamma$ are the cleanest on the theoretical side. Neither of these processes involves the W-exchange or the W-annihilation diagrams, which give additional contributions in the weak radiative decays of other b-hadrons. Also, the long distance contributions in this b-sector are small and under control[23]. Thus , the measurements of $\Xi_b^- \to \Xi^-\gamma$ and of $\Omega_b \to \Omega\gamma$ in future experiments, together with their improved theoretical estimations, will be helpful to obtain the needed insight on long distance contribution in weak radiative decays of the other beauty hadrons, as well as on the question of W-exchange contributions in those decays.

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