Photon decays of baryons with strangeness

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The photon-decay amplitudes and widths of the low-lying strange baryons are calculated within the framework of the nonrelativistic quark model of Isgur and Karl. A comparison with experiment is made.

INTRODUCTION

The nonrelativistic quark model of Isgur and Karl[1,2] has been very successful in correlating a large amount of experimental data in hadron physics. Diverse physical quantities such as baryon magnetic moments, mass spectra, decay widths, and charge radii have been calculated and found to be in good agreement with experiment.

In the near future the radiative decay widths of the low-lying strange baryons will be measured,[3] thus providing another testing ground for quark models. At present there exists only one calculation of these widths, which is within the framework of the MIT bag model.[4] Therefore a calculation within the nonrelativistic quark model framework is of interest.

It was found previously[5,6] that the compositions given by the Isgur-Karl model, when coupled with a simple model for baryon decay, predicted both $N\gamma$ and baryon-meson decay amplitudes with the following interesting features: (i) The "missing-resonance problem" is resolved; many of the states anticipated by the naive nonrelativistic quark model which are not seen experimentally are predicted to decouple from elastic channels. (ii) The coupled states correspond to observed states in both their masses and decay amplitudes. (iii) The observed violations of SU(6) selection rules are accounted for; for example, the amplitude of the SU(6)-forbidden process $D15\to p\gamma$ is predicted correctly in both magnitude and sign.

The present calculation is an extension of the work in Ref. 6. The spectrum and decay scheme of the hyperon states of interest are given in Fig. 1.

THE DECAY MODEL

The Isgur-Karl model for baryon structure can be summarized as follows. Quantum chromodynamics (QCD) is phenomenologically incorporated by subjecting quarks to confining two-body potentials perturbed by the strong color spin-spin forces due to one-gluon exchange. Diagonalization of the full Hamiltonian leads to the following baryon compositions[2]:

$$\begin{aligned}
&\Lambda(1116)\tfrac{1}{2}^+:\ |{}^28S_s\tfrac{1}{2}^+\rangle\ ,\\
&\Sigma(1193)\tfrac{1}{2}^+:\ |{}^28S_s\tfrac{1}{2}^+\rangle\ ,\\
&\Sigma(1385)\tfrac{3}{2}^+:\ |{}^410S_s\tfrac{3}{2}^+\rangle\ , \qquad (1)\\
&\Lambda(1405)\tfrac{1}{2}^-:\ +0.90\,|{}^21P_M\tfrac{1}{2}^-\rangle+0.43\,|{}^28P_M\tfrac{1}{2}^-\rangle\\
&\qquad -0.06\,|{}^48P_M\tfrac{1}{2}^-\rangle\ ,\\
&\Lambda(1520)\tfrac{3}{2}^-:\ +0.91\,|{}^21P_M\tfrac{3}{2}^-\rangle+0.40\,|{}^28P_M\tfrac{3}{2}^-\rangle\\
&\qquad +0.01\,|{}^48P_M\tfrac{3}{2}^-\rangle\ ,
\end{aligned}$$

where we have used the notation $|{}^{2S+1}XL_\sigma J^P\rangle$. Here X is the SU(3) multiplicity, S, L, P, and J are the total spin, total orbital angular momentum, parity, and total angular momentum, and σ is the permutational symmetry of the SU(6) state. A complete discussion of wave functions and conventions can be found in Ref. 6.

We assume that photoemission occurs via the deexcitation of a single quark.[7] A nonrelativistic reduction of the pointlike γ^μ quark-photon interaction, when sandwiched between initial and final baryon states, leads to the T-matrix element:

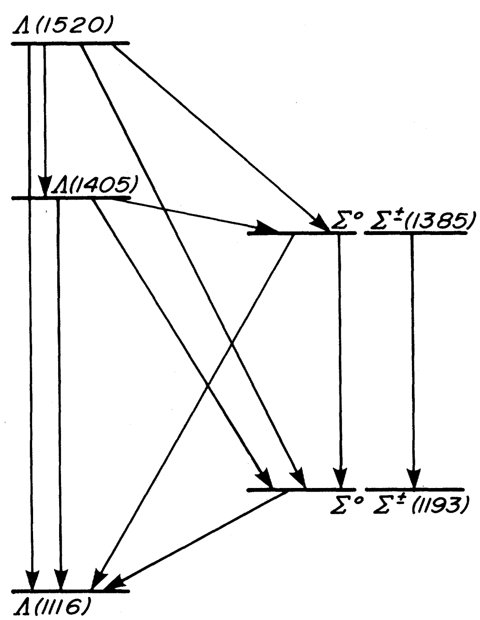

FIG. 1. Hyperon decay scheme.

TABLE I. Full photon amplitudes in standard pion photoproduction conventions are obtained by multiplying entries in this table by $\mu_p \alpha e^{-K}(2\pi/q)^{1/2}$ where α is the harmonic-oscillator strength parameter, $K=q^2/6\alpha^2$, $x=m_u/m_s$, and μ_p is the proton magnetic moment.

Initial state \ Final state	$\vert\Lambda\,^21P_M\frac{1}{2}^-\rangle$ $A_{3/2}$	$\vert\Lambda\,^21P_M\frac{1}{2}^-\rangle$ $A_{1/2}$	$\vert\Lambda\,^28P_M\frac{1}{2}^-\rangle$ $A_{3/2}$	$\vert\Lambda\,^28P_M\frac{1}{2}^-\rangle$ $A_{1/2}$
$\vert\Lambda\,^21P_M\frac{3}{2}^-\rangle$	$\frac{\sqrt{3}}{27}\frac{q}{\alpha}(1-x)$	$\frac{q}{27\alpha}(1-6K)(1-x)$	$-\frac{\sqrt{3}}{18}\frac{q}{\alpha}\left[\frac{2x+1}{3}\right]$	$-\frac{q}{6\alpha}(1-2K)\left[\frac{2x+1}{3}\right]$
$\vert\Lambda\,^28P_M\frac{3}{2}^-\rangle$	$-\frac{\sqrt{3}}{6}\frac{q}{\alpha}\left[\frac{2x+1}{3}\right]$	$-\frac{q}{6\alpha}(1-2K)\left[\frac{2x+1}{3}\right]$	$\frac{\sqrt{3}}{27}\frac{q}{\alpha}(1-x)$	$\frac{q}{27\alpha}(1-x+6Kx)$
$\vert\Lambda\,^48P_M\frac{3}{2}^-\rangle$				
$\vert\Lambda\,^21P_M\frac{1}{2}^-\rangle$				
$\vert\Lambda\,^28P_M\frac{1}{2}^-\rangle$				
$\vert\Lambda\,^48P_M\frac{1}{2}^-\rangle$				
$\vert\Sigma^0\,^410S_s\frac{3}{2}^+\rangle$				
$\vert\Sigma^0\,^28S_s\frac{1}{2}^+\rangle$				

Initial state \ Final state	$\vert\Lambda\,^28S_s\frac{1}{2}^+\rangle$ $A_{3/2}$	$\vert\Lambda\,^28S_s\frac{1}{2}^+\rangle$ $A_{1/2}$	$\vert\Sigma^0\,^410S_s\frac{3}{2}^+\rangle$ $A_{3/2}$	$\vert\Sigma^0\,^410S_s\frac{3}{2}^+\rangle$ $A_{1/2}$	$\vert\Sigma^0\,^410S_s\frac{3}{2}^+\rangle$ $A_{-1/2}$
$\vert\Lambda\,^21P_M\frac{3}{2}^-\rangle$	$-\frac{i}{\sqrt{6}}\left[\frac{2x+1}{3}\right]$	$-\frac{i\sqrt{2}}{6}\left[\frac{2x+1}{3}\right](1-6K)$	0	0	0
$\vert\Lambda\,^28P_M\frac{3}{2}^-\rangle$	$-\frac{i}{\sqrt{6}}\left[\frac{2x+1}{3}\right]$	$-\frac{i\sqrt{2}}{6}\left[\frac{2x+1}{3}-(2x-1)2K\right]$	0	$\frac{i2}{\sqrt{3}}K$	$i2K$
$\vert\Lambda\,^48P_M\frac{3}{2}^-\rangle$	$\frac{i\sqrt{15}}{5}K$	$\frac{i}{3\sqrt{5}}K$	$-\frac{i\sqrt{10}}{5}(1+3K)$	$-\frac{i2\sqrt{30}}{15}(1+K)$	$-\frac{i\sqrt{10}}{5}(1-K)$
$\vert\Lambda\,^21P_M\frac{1}{2}^-\rangle$		$-\frac{i}{3}(1+3K)\left[\frac{2x+1}{3}\right]$		0	0
$\vert\Lambda\,^28P_M\frac{1}{2}^-\rangle$		$-\frac{i}{3}\left[\frac{2x+1}{3}+(2x-1)K\right]$		$-\frac{i\sqrt{6}}{3}K$	$i\sqrt{2}K$
$\vert\Lambda\,^48P_M\frac{1}{2}^-\rangle$		$\frac{i}{3}K$		$-\frac{i}{\sqrt{6}}(1+4K)$	$-\frac{i}{\sqrt{2}}(1+2K)$
$\vert\Sigma^0\,^410S_s\frac{3}{2}^+\rangle$	$-\frac{q}{\alpha}$	$-\frac{q}{\sqrt{3}\alpha}$			
$\vert\Sigma^0\,^28S_s\frac{1}{2}^+\rangle$		$\frac{\sqrt{6}}{3}\frac{q}{\alpha}$			

TABLE I. (*Continued.*)

Initial state \ Final state	$\lvert\Sigma^0\,{}^28S_s\frac{1}{2}^+\rangle$ $A_{3/2}$	$\lvert\Sigma^0\,{}^28S_s\frac{1}{2}^+\rangle$ $A_{1/2}$	$\lvert\Sigma^\pm\,{}^28S_s\frac{1}{2}^+\rangle$ $A_{3/2}$	$\lvert\Sigma^\pm\,{}^28S_s\frac{1}{2}^+\rangle$ $A_{1/2}$
$\lvert\Lambda\,{}^21P_M\frac{3}{2}^-\rangle$	$-\frac{i}{\sqrt{2}}$	$-\frac{i}{\sqrt{6}}(1-6K)$		
$\lvert\Lambda\,{}^28P_M\frac{3}{2}^-\rangle$	$\frac{i}{\sqrt{2}}$	$\frac{i}{\sqrt{6}}(1-2K)$		
$\lvert\Lambda\,{}^48P_M\frac{3}{2}^-\rangle$	$-\frac{i3}{\sqrt{5}}K$	$-\frac{i}{\sqrt{15}}K$		
$\lvert\Lambda\,{}^21P_M\frac{1}{2}^-\rangle$		$-\frac{i}{\sqrt{3}}(1+3K)$		
$\lvert\Lambda\,{}^28P_M\frac{1}{2}^-\rangle$		$\frac{i}{\sqrt{3}}(1+K)$		
$\lvert\Lambda\,{}^48P_M\frac{1}{2}^-\rangle$		$-\frac{i}{\sqrt{3}}K$		
$\lvert\Sigma^0\,{}^410S_s\frac{3}{2}^+\rangle$	$+\frac{q}{\sqrt{3}\alpha}\left[\frac{2x+1}{3}\right]$	$\frac{q}{3\alpha}\left[\frac{2x+1}{3}\right]$		
$\lvert\Sigma^0\,{}^28S_s\frac{1}{2}^+\rangle$				
$\lvert\Sigma^+\,{}^410S_s\frac{3}{2}^+\rangle$			$\frac{2q}{\sqrt{3}\alpha}\left[\frac{2+x}{3}\right]$	$\frac{2q}{3\alpha}\left[\frac{2+x}{3}\right]$
$\lvert\Sigma^-\,{}^410S_s\frac{3}{2}^+\rangle$			$\frac{2\sqrt{3}q}{9\alpha}(x-1)$	$\frac{2q}{9\alpha}(x-1)$

$$\langle B'(p',s')\gamma(K,\lambda)\mid T\mid B(p,s)\rangle$$

$$=\frac{-3ie}{(2\pi)^{3/2}}\left\langle B'(p',s')\right|\frac{e_3}{e}\left[\vec{\sigma}_3\cdot\frac{(\vec{K}\times\vec{\epsilon}^{\,*})}{2m_3}+i\frac{\vec{p}_3'\cdot\vec{\epsilon}^{\,*}}{m_3}\right]\times e^{-i\vec{K}\cdot\vec{r}_3}\left|B(p,s)\right\rangle \qquad (2)$$

for the process $B\rightarrow B'\gamma$. Here $\vec{\epsilon}(K,\lambda)$ is the photon polarization vector, e_3, $\frac{1}{2}\vec{\sigma}_3$, $\vec{r}_3$, and m_3 are the charge, spin, position, and mass of the third quark, and $\vec{p}_3'$ is the momentum of the third quark in B'. Overall permutational symmetry of the wave functions ensures that the full amplitude is obtained by calculating for emission from the third quark and multiplying by 3. At most three independent helicity amplitudes are obtained when all possible photon polarizations and initial- and final-baryon spin states are considered for the present systems. We name these amplitudes in the conventional way: $A_{3/2}$, $A_{1/2}$, and $A_{-1/2}$ where the subscript denotes the J_z value of the initial baryon which emits a photon of positive helicity in the $\hat{z}$ direction:

$$A_{J_z}=\langle B'(J',J_z'=J_z-1)\gamma(K,+)\mid T\mid B(J,J_z)\rangle\ . \qquad (3)$$

The radiative widths are given by

$$\Gamma_\gamma=\frac{1}{(2J+1)}\frac{M_{B'}}{M_B}\frac{q^2}{2\pi}4\sum_{J_z}|A_{J_z}|^2\ , \qquad (4)$$

where J and M_B are the angular momentum and mass of the decaying hyperon.

All the parameters used in the present calculation, m_u, m_d, m_s, and α (the harmonic-oscillator strength parameter), are fixed by previous analyses.[8] The amplitudes are presented in Table I and the numerical widths in Table II.

RESULTS AND CONCLUSIONS

The existing experimental information in this sector is rather limited. There exist measurements of the $\Sigma^0(1193)\rightarrow\Lambda(1116)\gamma$ (Ref. 9) and $\Lambda(1520)\rightarrow\Lambda(1116)\gamma$ (Ref. 10) and large upper bounds on the radiative decays of the $\Sigma^0(1385)$.[11]

First, our good prediction of 8.6 keV compared to the experimental value 11.5±2.6 keV for the width $\Sigma^0(1193)\rightarrow\Lambda(1116)\gamma$ reflects the "classical" SU(3) prediction.[12]

Second, our prediction of 96 keV also compares favorably with the measured value of 150±30 keV for the $\Lambda(1520)\rightarrow\Lambda(1116)\gamma$ width. One should note that the experimental figure should probably be revised slightly upward as a 15% background subtraction was performed on the basis of the fact that the $\Lambda(1520)$ is a pure SU(3) singlet. The necessary subtraction is reduced if one assumes the composition quoted in Eq. (1).

Finally we are certainly within the rather weak measured upper bounds of ~2000 keV for the radiative widths of the $\Sigma(1385)$.

We would like to close by pointing out two interesting

TABLE II. Decay widths in keV. Entries in this table are obtained by folding baryon compositions [Eq. (1)] against amplitudes in Table I and substituting into Eq. (4). We take standard numerical values $\mu_p=0.13\ \text{GeV}^{-1}$, $x=m_u/m_s \sim 0.6$, and $\alpha=0.41$ GeV (see Ref. 8).

Initial state \ Final state	$\Lambda(1405)$	$\Lambda(1116)$	$\Sigma^0(1385)$	$\Sigma^0(1193)$	$\Sigma^+(1193)$	$\Sigma^-(1193)$
$\Lambda(1520)$	0.2	96	≈ 0	74		
		150 ± 30[a]				
$\Lambda(1405)$		143	0.3	91		
$\Sigma^0(1385)$		232		19		
$\Sigma^+(1385)$					104	
$\Sigma^-(1385)$						2.5
$\Sigma^0(1193)$		8.6				
		11.5 ± 2.6[a]				

[a]Experimental measurement.

features of our results.

If one recomputes the widths assuming the baryon compositions which result from turning off strong hyperfine forces, the $\Lambda(1520),\Lambda(1405)\to\Sigma(1193)\gamma,\Sigma(1385)\gamma$ widths become nearly zero (~ 1 keV). This is due to the fact that in the absence of hyperfine forces the lowest-lying excited Λ's are states in which the strange quark oscillates against the nonstrange pair.[1] (It costs less energy to excite this mode because of the larger strange-quark mass.) Thus the spatial wave function is symmetric with respect to interchange of the light quarks. The Λ isospin wave function is antisymmetric and thus to maintain overall symmetry (as we have necessarily a color singlet) the spin wave function is symmetric with respect to light quark interchange. The ground-state Σ's are of course symmetric with respect to light quark interchange in their isospin and spin wave functions. Thus the decay can only proceed through a spin flip which is suppressed by q^2/α^2. This reflects the selection rule observed in bag-model calculations.[4] The observation of a large decay width for the $\Lambda(1405)$ and $\Lambda(1520)$ to the ground-state Σ's would be further indication that color hyperfine forces were present in these systems.

We also mention the amusing result that the $\Sigma^-(1385)\to\Sigma^-(1193)\gamma$ decay width is zero in the SU(3) limit; i.e., the width is proportional to the strange-light-quark mass difference. This is due to the fact that all the quarks in Σ^- have the same charge and thus as far as the Σ^- is concerned the interaction Hamiltonian is an SU(3) singlet (in this limit) and cannot induce transitions between the octet and decuplet. A similar selection rule was pointed out previously by Lipkin.[13]

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