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Radiative annihilation of K^-p atoms and the $\Lambda(1405)$

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The radiative annihilations of K^-p atoms to $\Lambda\gamma$ and $\Sigma^0\gamma$ are shown to be dominated by the $\Lambda(1405)$ hyperon. These processes thus provide a sensitive probe of the properties of this controversial state. We show that present data on the $\Lambda\gamma$ branching ratio are in agreement with its interpretation as a normal quark-model state, and that further information on the $\Lambda\gamma$ and the comparably strong $\Sigma^0\gamma$ channel will provide compelling tests of the viability of this versus the $\bar{K}N$ bound-state interpretation.

I. INTRODUCTION

The $\Lambda(1405)$ with strangeness -1 and $J^P=\frac{1}{2}^-$ is one of the most controversial of the baryons. In the years immediately following its discovery it was widely discussed as a $\bar{K}N$ bound state produced by the dynamics of the meson-baryon interaction.[1] Since the advent of the quark model it has been argued that the $\Lambda(1405)$ has a more natural interpretation as an ordinary $3q$ state.[2–4] It is, however, extraordinarily difficult to decide between these two possibilities[5] with presently available strong-interaction data: since the $\Lambda(1405)$ is below the $\bar{K}N$ threshold it is seen only in $\Sigma\pi$ production and through its effect on the low-energy $\bar{K}N$ S wave.

This difficulty in establishing experimentally with any certainty the nature of the $\Lambda(1405)$ is exacerbated by the fact that it is not very well described in the quark model relative to other states. In the Isgur-Karl[3,4] model both the $\Lambda(1405)\frac{1}{2}^-$ and the $\Lambda(1520)\frac{3}{2}^-$ are dominantly states with $L=1$ in the relative coordinate between the s quark and the center of mass of the ud quark pair with spin $S_{ud}=0$: they just correspond to the two possible spin-orbit couplings. While the model correctly predicts the very low-lying position of the center of gravity of these two states (note that they lie below the lowest-lying negative-parity N^*'s), it fails to predict their splitting. This failure is often attributed to the model's neglect of all spin-orbit forces, but the fact that such forces seem to be small in the rest of the spectrum is at least a little discomforting to this point of view. On the side of overlooking this discrepancy is the fact that analyses[6,7] of the strong couplings of the low-lying negative-parity Λ's strongly support the simple predicted internal spin-space structure of the $\Lambda(1405)$ and $\Lambda(1520)$.

Partly in response to this flaw in the $3q$ interpretation, a number of recent calculations[8] have reexamined the $\bar{K}N$ bound-state picture of the $\Lambda(1405)$, finding reasonable interpretations of this type very similar in character to those of the earliest discussions of this state. Clearly, settling this issue requires either a much better understanding of quantum chromodynamics than is available today or some new experimental information.

Our main purpose here is to point out that the radiative annihilations of K^-p atoms[9] into $\Lambda\gamma$ and $\Sigma^0\gamma$ provide additional very clean information on the $\Lambda(1405)$ which should help to decide this issue. Note that $\Sigma\pi$, $\Lambda\gamma$, and $\Sigma^0\gamma$ are the only open two-body channels for this state, so that a measurement of these two additional couplings[10,11] triples our information on its internal structure, and completes the set of available observables.

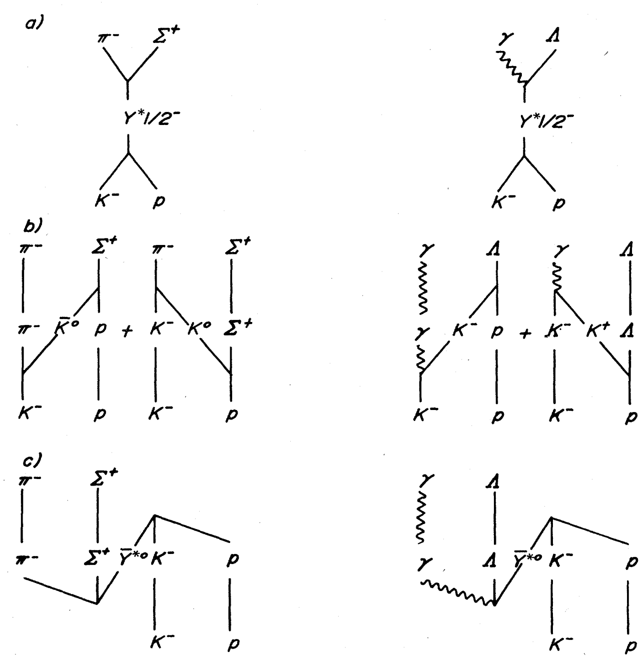

FIG. 1. Some time-ordered perturbation-theory diagrams. (a) An s-channel resonance contributing to $K^-p\to\pi^-\Sigma^+$ and $K^-p\to\gamma\Lambda$. (b) Some t-channel meson exchanges contributing to $K^-p\to\pi^-\Sigma^+$ and $K^-p\to\gamma\Lambda$. (c) Z graphs contributing to $K^-p\to\pi^-\Sigma^+$ and $K^-p\to\gamma\Lambda$.

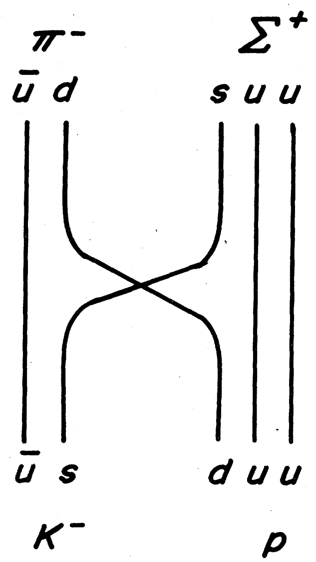

FIG. 2. A quark-exchange contribution to $K^-p \to \pi^-\Sigma^+$.

II. CALCULATION OF THE ANNIHILATION RATES

When a K^-p atom annihilates, it can produce any one of the final states $\Sigma^+\pi^-$, $\Sigma^-\pi^+$, $\Sigma^0\pi^0$, $\Lambda\pi^0$, $\Lambda\gamma$, or $\Sigma^0\gamma$. Experimentally one can only measure straightforwardly the branching ratios of these channels so we will need to study them all theoretically to extract information on the radiative width of the $\Lambda(1405)$. If we denote these channels generically by AB then the transition matrix element for radiative annihilation of an orbital-angular-momentum-zero, spin-s state $|0,s\rangle$ of the kaonic atom to AB will be

$$\langle AB \,|\, T \,|\, 0,s\rangle = \int d^3p \langle AB \,|\, T \,|\, K^-(\mathbf{p})p(-\mathbf{p}s)\rangle \phi_0(\mathbf{p}) \,, \tag{1}$$

TABLE I. Parameters of the low-lying $\Lambda\frac{1}{2}^-$ and $\Sigma\frac{1}{2}^-$ hyperons.

Y^*	Width (on resonance) (MeV)	^{21}	Mixing angles[a] ^{28}	^{48}	^{210}
$\Lambda(1405)$	40	$+0.90$	$+0.43$	-0.06	
$\Lambda(1670)$	35	-0.39	$+0.75$	-0.58	
$\Lambda(1800)$	300	-0.18	$+0.50$	-0.85	
$\Sigma(1620)$	50		-0.82	$+0.54$	-0.17
$\Sigma(1750)$	50		$+0.46$	-0.81	$+0.35$
$\Sigma(1810)$	100		-0.33	-0.21	$+0.92$

[a]Conventions are those of Ref. 7.

where $\phi_0(\mathbf{p})$ is the normalized momentum-space wave function of the kaonic atom. Assuming $\langle AB \,|\, T \,|\, K^-(\mathbf{p})p(-\mathbf{p}s)\rangle$ is a slowly varying function of $\mathbf{p}$, and noting that $\mathbf{p} \simeq 0$, this becomes

$$\langle AB \,|\, T \,|\, 0,s\rangle \simeq (2\pi)^{3/2}\psi_0(\mathbf{0})\langle AB \,|\, T \,|\, K^-(0)p(0s)\rangle \,, \tag{2}$$

where $\psi_0(\mathbf{r})$ is the wave function of the relative coordinate $\mathbf{r}=\mathbf{r}_K-\mathbf{r}_p$ and $\langle AB \,|\, T \,|\, K^-(0)p(0s)\rangle$ is the T-matrix element for the zero-momentum annihilation process. [Branching ratios are independent of $\psi_0(0)$, which is fortunate since strong interaction effects dominant at short distances can produce large departures from hydrogenic behavior.] Figure 1 shows some (old-fashioned) *time-ordered perturbation-theory* diagrams for various contributions to the matrix elements for $K^-p \to \Sigma^+\pi^-$ and $K^-p \to \Lambda\gamma$. Figure 2 shows an additional process which can contribute to $K^-p \to \Sigma^+\pi^-$ because the hadrons are composite objects: quark exchange. Note that this latter process is physically distinct from t-channel meson exchange which involves a $q^5\bar{q}^2$ intermediate state.

Since the hadrons are extended objects, the usual Feynman-diagram expansion, designed for pointlike ob-

TABLE II. Decay amplitudes for the states of Table I. Note the following. (i) $K=q^2/6\alpha^2$ where q is the boson momentum and $x=m_u/m_s$. (ii) The full strong amplitudes are obtained by multiplying entries in the table by $i\tilde{S}\alpha$ [see note (iv)]. (iii) The full electromagnetic amplitudes are obtained by multiplying entries in the table by $i\sqrt{4\pi}\mu_p\alpha e^{-K}$ [see note (iv)]. (iv) Parameters have been fixed by the previous analyses of Refs. 3 and 7: their values are $\tilde{S}=-7$ GeV^{-1}, $\alpha=0.41$ GeV, $\mu_p=0.13$ GeV^{-1}, and $x=0.6$. (v) Conventions are those of Ref. 7.

$^{2s+1}X_M$	K^-p	$\Sigma^+\pi^-$	$\Sigma^0\pi^0$	$\Sigma^-\pi^+$	$\Lambda\pi^0$	$\Lambda\gamma$	$\Sigma\gamma$
$^2\Lambda_1$	$\frac{1}{\sqrt{6}}$	$\frac{1}{\sqrt{6}}$	$-\frac{1}{\sqrt{6}}$	$\frac{1}{\sqrt{6}}$	0	$-\frac{1}{3}(1+3K)\left[\frac{2x+1}{3}\right]$	$-\frac{1}{\sqrt{3}}(1+3K)$
$^2\Lambda_8$	$\frac{1}{\sqrt{6}}$	$-\frac{1}{6}\sqrt{2/3}$	$\frac{1}{6}\sqrt{2/3}$	$-\frac{1}{6}\sqrt{2/3}$	0	$-\frac{1}{3}\left[\frac{2x+1}{3}+(2x-1)K\right]$	$\frac{1}{\sqrt{3}}(1+K)$
$^4\Lambda_8$	0	$-\frac{1}{3}\sqrt{2/3}$	$\frac{1}{3}\sqrt{2/3}$	$-\frac{1}{3}\sqrt{2/3}$	0	$\frac{1}{3}K$	$\frac{-1}{\sqrt{3}}K$
$^2\Sigma_8$	$\frac{\sqrt{2}}{18}$	$\frac{5\sqrt{2}}{18}$	0	$-\frac{5\sqrt{2}}{18}$	$\frac{\sqrt{6}}{18}$	$\frac{1}{\sqrt{3}}(1+\frac{1}{6}K)$	$\frac{1}{12}\left[2\left[1+\frac{4x-1}{3}\right]+K\left[1-\frac{4x-1}{9}\right]\right]$
$^4\Sigma_8$	$-\frac{2\sqrt{2}}{9}$	$-\frac{\sqrt{2}}{9}$	0	$\frac{\sqrt{2}}{9}$	$\frac{\sqrt{6}}{9}$	$-\frac{1}{6\sqrt{3}}K$	$-\frac{1}{18}\frac{4x-1}{3}K$
$^2\Sigma_{10}$	$\frac{\sqrt{2}}{18}$	$-\frac{\sqrt{2}}{18}$	0	$\frac{\sqrt{2}}{18}$	$\frac{\sqrt{6}}{18}$	$-\frac{1}{\sqrt{3}}(1-\frac{1}{6}K)$	$\frac{1}{3}(1-\frac{1}{6}K)\left[\frac{1+2x}{3}\right]$

TABLE III. Quark-exchange amplitudes.

Process	Amplitude[a] (in units of $\mathrm{Re}[K^-p\to\Lambda(1405)\to\Sigma^+\pi^-]$)
$K^-p\to\Sigma^+\pi^-$	$+2\epsilon$
$K^-p\to\Sigma^0\pi^0$	$-\epsilon$
$K^-p\to\Sigma^-\pi^+$	0
$K^-p\to\Lambda\pi^0$	$-\sqrt{3}\epsilon$

[a]From spin-independent interquark interactions.

jects, can be very misleading. Figure 1(c) is perhaps the best illustration of this fact: the Feynman diagrams with an s-channel $\Lambda\frac{1}{2}^+$ or $\Sigma\frac{1}{2}^+$ propagator are really representing the time-ordered processes which it displays (a time-ordered $\frac{1}{2}^+$ pole would not conserve parity). In the form of time-ordered diagrams we can see that these Z-graph processes will be very strongly suppressed relative to their pointlike values since (1) the basic vertex (e.g., the $Y^*\bar{K}N$ vertex) is being used far from the $\mathbf{q}\simeq 0$ decay point where it is "measured": compare to the $pp\gamma$ vertex at $\mathbf{q}=0$ and in $\bar{p}p$ pair production, and (2) the graphs require that vacuum fluctuations create (or destroy) three extended hadrons right on top of another pair of extended hadrons. Thus time-ordered perturbation theory clearly leads us to ignore these graphs.

A similar analysis indicates that one should neglect the diagrams of Fig. 1(b) and related t-channel exchanges. Consider, for purposes of illustration, an exchange so heavy that the resulting interaction can be approximated by a δ function. In this case these diagrams contribute only when the interacting hadrons are right on top of one another; but in this extreme the hadrons become hopeless approximations to the physical states: ordinary hadrons represent only the asymptotic confined colorless states of QCD, and cannot be expected to have any significant amplitude at small interhadron separations.

We therefore conclude that, in a sum over hadronic states, the diagrams of Fig. 1(a) will dominate. Fortunately, these are precisely the amplitudes that can be calculated with the most reliability. We have

$$\langle AB\,|\,T\,|\,K^-(0)p(0s)\rangle=(2\pi)^3\sum_{Y^*=\Lambda\frac{1}{2}^-,\Sigma\frac{1}{2}^-}\frac{\langle AB\,|\,H_{ABY^*}(0)\,|\,Y^*(0s)\rangle\langle Y^*(0s)\,|\,H_{Y^*K^-p}(0)\,|\,K^-(0)p(0s)\rangle}{M_{K^-}+M_p-M_{Y^*}+i\Gamma_{Y^*}/2}\,, \tag{3}$$

where the matrix elements are taken at the indicated off-resonance kinematic points and where Γ_{Y^*} is the width at the off-resonance $E=M_{K^-}+M_p$ relevant to the Y^* propagator. The required Y^* masses and couplings (and their dependences on E at such a low energy) are all predicted without any free parameters by the model of Refs. 3, 7, and 10. These relevant resonances and their properties are detailed in Tables I and II, where a few previously unpublished radiative decay matrix elements are shown for the first time.

As already mentioned, we would expect, in addition to these $Y^*\frac{1}{2}^-$ contributions, some modest effects in the strong-annihilation amplitudes from the quark exchange diagrams of Fig. 2. They will exist for precisely the same reasons that Figs. 1(b) and 1(c) are suppressed: because the hadrons are extended objects. We have not attempted a direct calculation of these amplitudes (though we believe such a calculation, while difficult,[12] is feasible) since they have little effect on our predictions for the radiative branching ratios. We have instead parametrized them by a strength parameter ϵ which we find to be $\simeq +0.27$; the four strong channels of interest then have the exchange amplitudes given in Table III.

To summarize: we have argued that, aside from small quark rearrangement terms, atomic K^-p annihilation rates will be dominated by s-channel $\frac{1}{2}^-$ resonances; other processes that are usually considered will almost certainly be strongly suppressed and will certainly be incorrectly estimated if the Feynman-diagram expansion is used. In fact we will see in the next section that these rates are dominated by the $\Lambda(1405)$ resonance itself and so provide important new sources of information on its structure. Note that, as a result of our conclusion regarding processes like those of Figs. 1(b) and 1(c), we differ significantly from the last of Refs. 9 on the extent of $\Lambda(1405)$ dominance in these processes.

III. RESULTS AND DISCUSSION

Table IV gives the results of our calculation. We should emphasize at this point that aside from ϵ this calculation has no free parameters whatsoever: all baryon compositions and couplings are taken directly from previ-

TABLE IV. Branching fractions. The results of our calculation are given in the Theory column. Column I is a calculation with all resonances omitted except the $\Lambda(1405)$, but including quark exchange. Column II is a calculation with the $\Lambda(1405)$ resonance only.

Channel	Theory	Experiment[a]	I	II
$\Sigma^+\pi^-$	0.20	0.20 ± 0.01	0.23	0.34
$\Sigma^0\pi^0$	0.29	0.27 ± 0.01	0.31	0.34
$\Sigma^-\pi^+$	0.43	0.46 ± 0.01	0.40	0.32
$\Lambda\pi^0$	0.07	0.07 ± 0.01	0.06	0
$\Lambda\gamma$	3.4×10^{-3}	$(2.8\pm0.8)\times10^{-3}$	3.6×10^{-3}	2.9×10^{-3}
$\Sigma\gamma$	2.6×10^{-3}		2.3×10^{-3}	1.8×10^{-3}

[a]Reference 13.

ous work. Since the model for baryon structure and decay is quite successful in describing other baryons, and since our predicted radiative branching ratios are (as demonstrated in the table) quite insensitive to everything but the $\Lambda(1405)$, we believe these predictions constitute a significant test of the model's ability to describe the $\Lambda(1405)$. If this state does not behave as predicted here, then we would have serious doubts about its interpretation as an ordinary three-quark system.

Given the present experimental information, there seems to be no immediate reason to raise such doubts. Clearly, however, a better determination of the $\Lambda\gamma$ rate and a first measurement of the comparable $\Sigma^0\gamma$ rate are crucial to making the case for the quark-model interpretation of the old and problematical $\Lambda(1405)$.

Note added in proof. Kim Maltman has pointed out to us that spin-dependent interactions, which we neglected, could be important in the quark-exchange amplitudes of Table III; his explicit calculations of the full amplitudes (using the model of Ref. 12) are, however, within a few percent of the $-2{:}1{:}0{:}-\sqrt{3}$ ratios given there. The predicted magnitude and sign of ϵ is under study (K. Maltman, private communication).

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[9]These processes have been discussed theoretically before. For early work in the context of meson-exchange models see G. Ya. Korenman and S. Popov, Phys. Lett. **40B**, 628 (1972); J. O. Eeg and H. Pilkuhn, Nuovo Cimento **32A**, 44 (1976). More recent work which is similar to ours in spirit, but very different in detail, has been done by H. Burkhardt, J. Lowe, and A. S. Rosenthal [Nucl. Phys. **A440**, 653 (1985)], who also show the important role of the $\Lambda(1405)$ in these processes.

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[12]For a related calculation, see K. Maltman and N. Isgur, Phys. Rev. D **29**, 952 (1984), where quark-exchange contributions to *NN* scattering are computed.

[13]J. Lowe *et al.* [Nucl. Phys. **B209**, 16 (1982)] have made the most recent experimental study of $K^-p \to \Lambda\gamma$. For the strong channels see the review by D. J. Miller, K. J. Nowak, and T. Tymieniecka, in *Low and Intermediate Energy Kaon-Nuclear Physics,* edited by E. Ferrari and G. Violini (Reidel, Holland, 1980), p. 251.