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**Radiative Decays of Low-Lying Excited-State
Hyperons**

by

Simon John Taylor

A THESIS SUBMITTED
IN PARTIAL FULFILLMENT OF THE
REQUIREMENTS FOR THE DEGREE

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Abstract

Radiative Decays of Low-Lying Excited-State Hyperons

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The quark wave-functions of the lower-lying excited-state hyperons $\Lambda(1405)$, $\Sigma(1385)$, and $\Lambda(1520)$ are not well understood. For example, the $\Lambda(1405)$ may not be a regular three-quark state but a $\bar{K}N$ molecule. Several competing models have been proposed, but none have been convincingly eliminated. Measuring radiative decays provides a means of discriminating between the models. The radiative branching ratios are predicted to be small ($\sim 1\%$), but the radiative widths vary by factors of 2-10 from model to model.

The existing experimental data is sparse and inconsistent; moreover, the radiative decay of the $\Sigma(1385)$ has never been observed before (except for one event).

These lower-lying excited-state hyperons were produced in a tagged photon-beam experiment in the CLAS detector at TJNAF in the reaction $\gamma p \rightarrow K^+ Y^*$ for photon energies from threshold to 2.4 GeV. The radiative branching ratio for the $\Sigma^0(1385)$

relative to the $\Sigma^0(1385) \rightarrow \Lambda\pi^0$ channel was measured to be $0.021 \pm 0.008^{+0.004}_{-0.007}$.

corresponding to a partial width of $640 \pm 270^{+130}_{-220}$ keV.

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Chapter 1

Introduction

Although the lower-lying excited state hyperons $\Sigma(1385)$, $\Lambda(1405)$, and $\Lambda(1520)$ were discovered more than thirty years ago, their quark structures have proven to be elusive. In the non-strange sector the Isgur-Karl non-relativistic quark model (NRQM) [1] was remarkably successful in predicting the masses of N^* and Δ states but less successful in the strange sector. The lowest lying $\frac{1}{2}^-$ state was predicted to be mass-degenerate with the $\frac{3}{2}^-$ state, in direct contradiction with experiment: the $\frac{1}{2}^-$ state, identified with the $\Lambda(1405)$, and the $\frac{3}{2}^-$ state, identified with the $\Lambda(1520)$, differ in mass by about 113 MeV. Various bag model [2, 3, 4] calculations of the $\frac{1}{2}^-$ state have led theorists to conclude that the $\Lambda(1405)$ wave function may contain some admixture of $q^4\bar{q}$ structure along with the standard q^3 baryonic structure. In the extreme case, the $\Lambda(1405)$ may be a $\bar{K}N$ bound state and not a q^3 structure.

Measuring the radiative decays of these hyperons provides a means of discriminating between the various models. The radiative widths are sensitive to the particular hyperon wave-functions predicted by the models. Other observables such as the hyperon masses are less sensitive: most models get the mass of the $\Sigma(1385)$ more or less correct, for example. The predictions of the various models for the radiative widths can vary by factors of 2 up to an order of magnitude (see table 1.1), making

it experimentally tractable to rule out certain models. Despite this sensitivity, relatively few measurements of neutral hyperon radiative decays have been performed and the experimental results are either contradictory or model-dependent. The decay $\Sigma^0(1385) \rightarrow \gamma\Lambda(1116)$ has never been observed (except for one count reported by Meisner [5]), even though its radiative branching ratio is expected to be $\sim 1\%$.

As shown in figure 1.1, the decay hierarchy for the lower-lying excited states is quite complicated; however, the radiative widths for transitions between excited states are predicted to be small. In this thesis I will be focusing on the decays of the $\Lambda(1405)$, $\Lambda(1520)$, and $\Sigma(1385)$ to the ground-state Λ :

$$Y^* \longrightarrow \gamma\Lambda \quad (1.1)$$

Throughout this document the symbol Y^* will be used to denote unspecified excited-state hyperons.

The dominant decay modes of the lower-lying excited state hyperons are strong decays. The known decay channels for the $\Sigma(1385)$, $\Lambda(1405)$, and $\Lambda(1520)$ are

$$\begin{aligned} \Lambda(1405) &\longrightarrow \Sigma\pi \quad (BR \approx 100\%) \\ \Sigma(1385) &\longrightarrow \Lambda\pi \quad (BR = 88 \pm 2\%) \\ &\longrightarrow \Sigma\pi \quad (BR = 12 \pm 2\%) \\ \Lambda(1520) &\longrightarrow N\bar{K} \quad (BR = 45\%) \\ &\longrightarrow \Sigma\pi \quad (BR = 42\%) \\ &\longrightarrow \Lambda\pi\pi \quad (BR = 10\%) \end{aligned}$$

Figure 1.1 Radiative decay hierarchy of the lower-lying excited state hyperons. The dashed lines indicate transitions that are forbidden or very rare in most models. The J^P assignments for the states are on the outer boundaries of the diagram.

$$\longrightarrow \Sigma\pi\pi \quad (BR = 0.9\%)$$

$$\longrightarrow \Lambda\gamma \quad (BR = 0.8\%).$$

Figure 1.2 shows the decay hierarchy for the strong decays of the excited-state hyperons and the subsequent weak decays of the ground state hyperons.

The various competing models and their predictions for radiative widths will be discussed in more detail in section 1.1. All other branching ratios quoted in this document are from [6]. Previous experimental result will be reviewed in section 1.2. Finally, the details of the present experiment will be presented in section 1.3.

1.1 Theoretical overview

Since all baryons and mesons are color singlets, any model of a hadron must ensure that no free quarks or gluons exist outside the hadron. The two classes of model that I will discuss in some detail are constituent quark models that employ a confining potential and bag models that confine the quarks and gluons within a cavity known as a “bag”.

In the Isgur-Karl non-relativistic quark model (NRQM) [1], the confining force experienced by the quarks is linear in the separation distance. Consequently, the zeroth-order wave-functions are those of simple harmonic oscillators. Since the quarks carry spin, they can also interact with each other via the hyperfine interaction

$$H_{hyp}^{ij} = \frac{2\alpha_S}{3m_i m_j} \left[\frac{8\pi}{3} \delta^3(\vec{r}_{ij}) \vec{s}_i \cdot \vec{s}_j + \frac{1}{r_{ij}^3} (3\vec{s}_i \cdot \hat{r}_{ij} \vec{s}_j \cdot \hat{r}_{ij} - \vec{s}_i \cdot \vec{s}_j) \right], \quad (1.2)$$

Figure 1.2 Hierarchy for the decays of the $\Lambda(1405)$ and $\Sigma^0(1385)$ to the ground state hyperons. A) Decay chain involving $\Sigma^0(1193)$. B) Decay chain involving $\Sigma^+(1189)$. C) Decay chain involving $\Sigma^-(1197)$. Arrows indicating transitions are labeled by the non-baryonic decay particle and the probability for the decay to occur in parenthesis.

where $\vec{r}_{ij}$ is the separation between quarks i and j , m_i is the mass of quark i and $\vec{s}_i$ is the spin of quark i . α_S is treated as free parameter. The hyperfine interactions account for the mass splittings. Isgur and Karl ignore spin-orbit coupling, which they claim should be small.

The model gives the following wave-functions [7] for the hyperons:

$$|\Lambda(\frac{1}{2}^+)\rangle = 0.93 |\Lambda_8^2 S_S\rangle - 0.30 |\Lambda_8^2 S'_S\rangle - 0.20 |\Lambda_8^2 S_M\rangle - 0.05 |\Lambda_1^2 S_M\rangle, \quad (1.3)$$

$$|\Sigma(\frac{1}{2}^+)\rangle = 0.97 |\Sigma_8^2 S_S\rangle - 0.18 |\Sigma_8^2 S'_S\rangle - 0.16 |\Sigma_8^2 S_M\rangle - 0.02 |\Sigma_{10}^2 S_M\rangle \quad (1.4)$$

$$|\Sigma(\frac{3}{2}^+)\rangle = |\Sigma_{10}^4 S_S\rangle, \quad (1.5)$$

$$|\Lambda(\frac{1}{2}^-)\rangle = 0.90 |\Lambda_1^2 P_M\rangle + 0.43 |\Lambda_8^2 P_M\rangle - 0.06 |\Lambda_8^4 P_M\rangle, \quad (1.6)$$

$$|\Lambda(\frac{3}{2}^-)\rangle = 0.91 |\Lambda_1^2 P_M\rangle + 0.40 |\Lambda_8^2 P_M\rangle + 0.01 |\Lambda_8^4 P_M\rangle, \quad (1.7)$$

using the notation $|X_M^{2J+1} L_\sigma\rangle$, where M is the SU(3) multiplet, J is the total spin of the hadron, L is the orbital angular momentum, and σ is the SU(6) permutation symmetry. The S'_S notation indicates a symmetric spatial wave-function distinct from the S_S wave function. Isgur and Karl find that the $|\Lambda(\frac{1}{2}^-)\rangle$ and $|\Lambda(\frac{3}{2}^-)\rangle$ states are degenerate in mass (both 1490 MeV) [1].

Darewych, Horbatsch, and Koniuk [8] calculated the radiative widths and amplitudes within the framework of this model. They assumed that the photon is radiated through the de-excitation of a single quark. They approximated the ground state hyperon wave-functions $|\Lambda(\frac{1}{2}^+)\rangle$ and $|\Sigma(\frac{1}{2}^+)\rangle$ by their dominate components ($|\Lambda_8^2 S_S\rangle$ in both cases). Kaxiras *et al.* [9] also calculated the radiative widths using the Isgur-Karl wave-functions expressed in two bases: SU(6), in which the quark masses are

equal. and uds , in which the wave-function is symmetrized with respect to the exchange of u and d quarks and the strange quark s is allowed to have mass different from the u and d quarks.

In the non-relativistic limit the radiative width is given by [9]

$$\Gamma = 2\pi \frac{m_f}{m_i} \frac{k^2}{(2\pi)^3} \frac{4\pi}{2J_i+1} \sum_{M_i, \lambda} |A_{M_i, \lambda}|^2. \quad (1.8)$$

$$H_{int} = \left(\frac{2\pi}{k}\right)^{\frac{1}{2}} \left(-\frac{e_j}{m_j} \vec{p}_j \cdot \vec{\epsilon}_\lambda^*(\vec{k}) + i \frac{e_j}{m_j} \vec{s}_j \cdot \vec{k} \times \vec{\epsilon}_\lambda^*(\vec{k})\right) e^{-i\vec{k} \cdot \vec{r}_j}, \quad (1.9)$$

where e_j is the charge, m_j is the mass, $\vec{s}_j$ is the spin, and $\vec{p}_j$ is the momentum of the j th quark and where $\vec{\epsilon}_\lambda^*(\vec{k})$ is the photon polarization vector, $\vec{k} = k\hat{z}$ is the photon momentum, λ is the helicity of the photon, and M_i is the z-component of the angular momentum of the decaying baryon. $A_{M_i, \lambda}$ is the matrix element of H_{int} between the initial and final states.

Warns, Pfeil, and Rollnik [10] calculated the radiative transitions within the framework of a relativized constituent quark model (CQM). In this model the hyperon consists of three point-like but massive quarks. The electromagnetic Hamiltonian $\mathcal{H}_{EM}$ contains the non-relativistic Hamiltonian (c.f. equation 1.9) corrected up to $\mathcal{O}(1/c^3)$ and has the structure

$$\mathcal{H}_{EM} = H_{SQ}^{nr} + H_{SQ}^{(2)} + H_{SQ}^{(3)} + H_I^{(2)} + H_I^{(3)} + H_{CM}^{(2)} + H_{CM}^{(3)} + \mathcal{O}\left(\frac{1}{c^4}\right), \quad (1.10)$$

where H_{SQ}^{nr} is the non-relativistic contribution due to the coupling between the photon field and a single quark (summed over the three quarks in the problem), $H_{SQ}^{(2)}$ and $H_{SQ}^{(3)}$ are second- and third-order corrections (in $1/c$) to the non-relativistic single

quark Hamiltonian. $H_I^{(2)}$ and $H_I^{(3)}$ are second- and third-order corrections arising from the interaction potential between the quarks, and $H_{CM}^{(2)}$ and $H_{CM}^{(3)}$ are second- and third-order corrections arising from the center-of-mass motion of the three quark system.

Warns *et al.* used the Isgur-Karl wave-functions (equations 1.3-1.7). The u and d quark masses were taken to be 350 MeV each and the strange quark mass was 580 MeV. The harmonic oscillator strength α_ρ was 0.3 GeV. The radiative widths are computed from the matrix elements of $\mathcal{H}_{EM}$ using equation 1.8. They obtain larger radiative widths than their non-relativized model predicts. They obtain a smaller $\Lambda(1405) \rightarrow \Sigma^0 \gamma / \Lambda(1405) \rightarrow \Lambda \gamma$ ratio than seen in the data [11, 12]. These observations lead them to conclude that a non-three quark model for the $\Lambda(1405)$ may be more appropriate.

In the MIT bag model [13], the fields describing the dynamical behavior of the quarks and gluons comprising the hadron are confined to the hadron interior by the inclusion of the term $-g_{\mu\nu}\theta_s B$ to the energy-momentum tensor. The step function θ_s ensures that the fields are confined because it is only nonzero within the volume of the hadron. B is a parameter of the theory with the units of pressure. Inside the bag the quarks are quasi-free, subject to interactions with each other via gluons and the boundary conditions imposed by the bag. Besides B , there are three more parameters of the theory. The strength of the quark-gluon interaction is characterized by the strong coupling constant α_s . The u and d quarks are assumed to have the same mass

(usually taken to be zero), but the strange quark is allowed to have a non-zero mass, m_s . Finally, the parameter Z_0 characterizes the magnitude of the zero-point energy of the quantum modes within the bag. Gluon interactions are treated to first order in α_s (this is the so-called One Gluon Exchange approximation, or OGE). Hadron masses are usually calculated using the fixed-radius static-cavity approximation in which the radius of the cavity is R . Under this assumption, the effective Hamiltonian determining the hadron mass M is given by

$$M(R) = \frac{4}{3}\pi R^3 B - \frac{Z_0}{R} + H_{KE} + H_g, \quad (1.11)$$

where H_{KE} is the bag energy (which includes the rest masses and the kinetic energy of the quarks in the bag) and H_g is the perturbative gluon Hamiltonian. The actual radius R_0 of the bag is determined by requiring $\partial M(R)/\partial R = 0$ and the actual hadron mass is given by $M(R_0)$. In the treatment of DeGrand *et al.*, the strange quark mass m_s is chosen such that the Ω^- mass is correct. The other parameters (B , Z_0 , and α_s) are determined by fitting to the known masses of the non-strange particles p , Δ , and ω . For the case where the u and d quarks are taken to be massless, the values of the parameters are $B^{1/4} = 145$ MeV, $Z_0=1.84$, $\alpha_s=0.55$, and $m_s=279$ MeV.

A non-zero mass for the strange quark breaks $SU(3)$ symmetry and allows for mixing between flavor-singlet $|\Lambda_1^{-2}\rangle$ and flavor-octet $|\Lambda_8^{-2}\rangle$ states (using the notation $|\chi_M^{-2J+1}\rangle$, as before). The positive parity hyperon states ($\Lambda(\frac{1}{2}^+, 1116)$, $\Sigma(\frac{1}{2}^+, 1193)$, $\Sigma^*(\frac{3}{2}^+, 1385)$) are considered to consist of three quarks (u , d , and s) each in the $S_{\frac{1}{2}}$ orbit. These hyperon states are “pure” in the sense that there is no mixing between

the multiplets: $\Sigma^*(\frac{3}{2}^+, 1385)$ belongs to the ground-state SU(3) flavor-decuplet and $\Lambda(\frac{1}{2}^+, 1116)$ and $\Sigma(\frac{1}{2}^+, 1193)$ are part of the ground-state SU(3) flavor-octet. For the negative parity states, bag model calculations predict substantial mixing between the multiplets. The two lowest-lying $\frac{1}{2}^-$ states in this model are:

$$\left| \Lambda \left(\frac{1}{2}^- \right) \right\rangle_1 = 0.39 \left| \Lambda_1, \frac{1}{2}^- \right\rangle + 0.42 \left| \Lambda_8, \frac{1}{2}^- \right\rangle_s + 0.46 \left| \Lambda_8, \frac{1}{2}^- \right\rangle_m + 0.67 \left| \Lambda_8, \frac{3}{2}^- \right\rangle. \quad (1.12)$$

$$\left| \Lambda \left(\frac{1}{2}^- \right) \right\rangle_2 = 0.93 \left| \Lambda_1, \frac{1}{2}^- \right\rangle - 0.11 \left| \Lambda_8, \frac{1}{2}^- \right\rangle_s - 0.21 \left| \Lambda_8, \frac{1}{2}^- \right\rangle_m - 0.26 \left| \Lambda_8, \frac{3}{2}^- \right\rangle. \quad (1.13)$$

using the notation $\left| X_M, S^P \right\rangle$, where M is the flavor-SU(3) multiplet, S is the spin of the excited quark, P is the parity of the excited quark, and the subscript $s(m)$ indicates totally (mixed) symmetric wave-functions in pseudo-spin for the octet states. The masses corresponding to $\left| \Lambda(\frac{1}{2}^-) \right\rangle_1$ and $\left| \Lambda(\frac{1}{2}^-) \right\rangle_2$ were determined to be 1364 MeV and 1446 MeV, respectively. Neither mass agrees with the mass of the $\Lambda(1405)$. Fitting the $\Lambda(1520)$ presents a similar problem. The lowest-lying $\frac{3}{2}^-$ state in the model is

$$\begin{aligned} \left| \Lambda \left(\frac{3}{2}^- \right) \right\rangle_1 = & 0.86 \left| \Lambda_1, \frac{3}{2}^- \right\rangle + 0.35 \left| \Lambda_8, \frac{3}{2}^-, S=0 \right\rangle - 0.35 \left| \Lambda_8, \frac{3}{2}^-, S=1 \right\rangle \\ & - 0.14 \left| \Lambda_8, \frac{1}{2}^- \right\rangle. \end{aligned} \quad (1.14)$$

where the notation is the same as for equations 1.12 and 1.13 with the added notation of $S=0$ ($S=1$) indicating the total spin of the two $S_{\frac{1}{2}}$ quarks. The state $\left| \Lambda(\frac{3}{2}^-) \right\rangle_1$ has a mass of 1257 MeV, clearly too low.

Kaxiras *et al.* [9] calculate the radiative widths within the MIT bag model framework. The radiative width is determined by

$$\Gamma_{J_f J_i} = 2k \frac{1}{2J_i + 1} \sum_{M_i, M_f} \sum_{\lambda=\pm 1} \left| \langle J_f M_f | \int d^3 r \vec{\epsilon}_\lambda^*(\vec{k}) \cdot \vec{j}(\vec{r}) e^{-i\vec{k} \cdot \vec{r}} | J_i M_i \rangle \right|^2. \quad (1.15)$$

$$\vec{j}(\vec{r}) = e_f \psi^\dagger(\vec{r}) \vec{\alpha} \psi(\vec{r}), \vec{\alpha} = \begin{bmatrix} 0 & \vec{\sigma} \\ \vec{\sigma} & 0 \end{bmatrix}. \quad (1.16)$$

Here J_i, M_i (J_f, M_f) are the total angular momentum and the z-projection of the angular momentum for the initial (final) baryon states; $\vec{\sigma}$ are the Pauli spin matrices and $\psi(\vec{r})$ is the Dirac spinor for the quark undergoing the radiative transition. The three-momentum of the photon is $\vec{k} \equiv k \hat{z}$, with $k = m_i - m_f$, where m_i (m_f) is the physical mass of the initial (final) baryon. The radiative decay is assumed to occur via the de-excitation of a single quark whose charge is e_f . The photon polarization vector is denoted $\vec{\epsilon}_\lambda^*(\vec{k})$. The sum is performed over the two helicity states ($\lambda = \pm 1$) of the photon.

Because of the lack of agreement between the mass of the $\Lambda(1405)$ and the masses corresponding to $|\Lambda(\frac{1}{2}^-)\rangle_1$ and $|\Lambda(\frac{1}{2}^-)\rangle_2$, Kaxiras *et al.* calculated the width using both wave-functions. They calculated the radiative width for the $|\Lambda(\frac{3}{2}^-)\rangle$ state by using 1.14 and by assuming that the state is a flavor-singlet ($|\Lambda_1, \frac{3}{2}\rangle$). Inclusion of the octet contribution noticeably changes the calculated radiative widths. In their calculation, the $\Lambda(1520) \rightarrow \gamma \Sigma$ width is 100 keV for the flavor-singlet state, compared to 17 keV for the mixed state. The $\Lambda(1520) \rightarrow \gamma \Lambda$ width is 27 keV for the flavor-singlet state, compared to 46 keV for the mixed state.

In the chiral bag model of Umino and Myhrer[14, 2], the bag is dressed by a mesonic cloud. The meson cloud contribution to the Hamiltonian is considered perturbatively. The meson fields are excluded from the interior of the bag and continuity of axial current is required at the bag surface. Both kaon and pion fields are included in the calculation. To first order in the meson fields, the Lagrangian in their model is given by

$$\begin{aligned} \mathcal{L} = & \left(i\bar{q}(x)\gamma^\mu\partial_\mu q(x) - \frac{1}{4}\vec{F}^{\mu\nu}(x) \cdot \vec{F}_{\mu\nu}(x) - B \right) \theta_S \\ & - \frac{1}{2}\bar{q}(x) \left(1 + \frac{i}{f}\vec{\lambda} \cdot \vec{\phi}(x)\gamma_5 \right) q(x)\delta_S + \frac{1}{2} \left(\partial_\mu \vec{\phi}(x) \right)^2 \overline{\theta}_S, \end{aligned} \quad (1.17)$$

where $q(x)$ are the quark fields, $\vec{\phi}(x)$ are the meson fields, $\vec{\lambda}$ are the Gell-Mann matrices, B is the bag constant, and f is the meson decay constant (either f_π or $f_K = 1.09f_\pi$). Confinement is accomplished via the step function

$$\theta_S = \begin{cases} 1 & \text{if } x \text{ is inside the bag} \\ 0 & \text{otherwise} \end{cases} \quad (1.18)$$

$\overline{\theta}_S$ is the inverse of θ_S and $\delta_S = \delta(\vec{r} - \vec{R}_0)$ is the Dirac delta-function at the surface of the bag with radius R_0 . The third term in 1.17 is the Klein-Gordon Lagrangian for a free scalar field. They performed the calculations for the negative parity hyperons only under the assumption that two of the quarks are in the $S_{\frac{1}{2}}$ state and one is in either the $P_{\frac{1}{2}}$ or the $P_{\frac{3}{2}}$ state. They assumed a static spherical bag with the bag parameters $\alpha_s = 1.5$, $B^{1/4} = 145$ MeV, $m_s = 250$ MeV, and $Z_0=0.25$ [2].

Umino and Myhrer found the following wave-functions for the lowest-lying negative parity hyperons for the case where they only include pion fields [14]:

$$\left| \Lambda(\tfrac{1}{2}^-) \right\rangle = -0.77 \left| {}^2 1_{1/2} \right\rangle + 0.003 \left| {}^4 8_{1/2} \right\rangle + 0.64 \left| {}^2 8_{1/2} \right\rangle. \quad (1.19)$$

$$\left| \Lambda(\tfrac{3}{2}^-) \right\rangle = -0.95 \left| {}^2 1_{3/2} \right\rangle - 0.09 \left| {}^4 8_{3/2} \right\rangle + 0.29 \left| {}^2 8_{3/2} \right\rangle. \quad (1.20)$$

where we introduce yet another notation $\left| {}^{2S+1} M_J \right\rangle$, in which M is the SU(3) multiplet, J is the total angular momentum of the state, and $2S+1$ indicates the spin multiplet of the three-quark system. The masses of the $\left| \Lambda(\tfrac{1}{2}^-) \right\rangle$ and $\left| \Lambda(\tfrac{3}{2}^-) \right\rangle$ states were 1539 MeV and 1534 MeV, respectively. In other words, similar to the Isgur-Karl model, the $\Lambda(1405)$ and the $\Lambda(1520)$ are nearly mass degenerate in the model. Inclusion of the kaon fields did not affect the wave-functions, but lowered the mass of the $\left| \Lambda(\tfrac{3}{2}^-) \right\rangle$ state from 1534 MeV to 1528 MeV while raising the mass of the $\left| \Lambda(\tfrac{1}{2}^-) \right\rangle$ state from 1539 MeV to 1557 MeV compared to the case where only pion fields were considered[14, 2], thus alleviating the mass degeneracy but pulling the prediction for the $\left| \Lambda(\tfrac{1}{2}^-) \right\rangle$ state further away from the observed mass of the $\Lambda(1405)$.

The radiative width in the chiral bag model differs from 1.15 by the inclusion of a meson current term:

$$\begin{aligned} \Gamma_{J_f J_i} &= 2k \frac{1}{2J_i+1} \sum_{M_i, M_f} \sum_{\lambda=\pm 1} \left| \langle J_f M_f | \vec{\epsilon}_\lambda^*(\vec{k}) \cdot (\vec{I}_q + \vec{I}_m) | J_i M_i \rangle \right|^2; \\ \vec{I}_q &\equiv \int_0^R d^3 r \vec{J}_q(\vec{r}) e^{-i\vec{k} \cdot \vec{r}}; \\ \vec{I}_m &\equiv \int_R^\infty d^3 r \vec{J}_m(\vec{r}) e^{-i\vec{k} \cdot \vec{r}}, \end{aligned} \quad (1.21)$$

where R is the bag radius. Here $\vec{J}_q(\vec{r})$ is the quark electromagnetic current operator, given by 1.16, and $\vec{J}_m(\vec{r})$ is a two-body operator of the form $\sum_{i,j} \hat{O}(i)\hat{O}(j)$ for the case where quark i emits a meson that is absorbed by quark j . The meson cloud can only contribute to the radiative width via the coupling of the photon with the mesons in flight. The dominant contribution to the radiative width is from the quark core via the de-excitation of single quarks.

The results of all of the calculations described above are tabulated in table 1.1 and displayed graphically in figure 1.3. Figure 1.3 demonstrates that the widths predicted by bag models are generally lower than those predicted by constituent quark models.

The flavor part of the $\Sigma^0(1385)$ wave function is the same in both the Non-relativistic Quark Model and the MIT Bag Model. The full wave function, however, is a product of four components:

$$\psi = \psi(color)\psi(space)\psi(flavor)\psi(spin). \quad (1.22)$$

Here $\psi(color)$ is an anti-symmetric singlet state: it is the same for all baryons. Where the two models can differ is in the spatial wave function. This presumably accounts for some of the difference in magnitude of the radiative width predictions in the two models.

1.1.1 The $\Lambda(1405)$ problem

Purely q^3 wave-functions do not give satisfactory results for the $\Lambda(1405)$. As mentioned earlier, the $\Lambda(1405)$ wave-function might have some admixture of $q^4\bar{q}$;

Transition	NRQM [9, 8]	MIT BAG[9]	Chiral BAG[2]	RELQM [10]	Experiment K^-p
$\Lambda(1520) \rightarrow \Lambda\gamma$	156	46	32	215	134 ± 23 [23] 30 ± 11 [24]
$\Lambda(1405) \rightarrow \Lambda\gamma$	200	$60^\dagger, 17^\ddagger$	75	118	27 ± 8 [12, 11]
$\Sigma(1385) \rightarrow \Lambda\gamma$	273	152		267	
$\Sigma(1193) \rightarrow \Lambda\gamma$	8.5	4.6		4.1	8.6 ± 1.4 [15]
$\Lambda(1520) \rightarrow \Sigma\gamma$	55	17	51	293	47 ± 17 [24]
$\Lambda(1405) \rightarrow \Sigma\gamma$	72	$18^\dagger, 2.7^\ddagger$	1.9	46	10 ± 4 [12, 11] 23 ± 7 [12, 11]
$\Sigma(1385) \rightarrow \Sigma\gamma$	22	15		23	

Table 1.1 Comparison of theoretical predictions for the radiative widths (in keV) for the transitions $Y^* \rightarrow \Lambda\gamma$ and $Y^* \rightarrow \Sigma\gamma$ to experimental results. In the case of the predictions of the MIT bag model for the $\Lambda(1405)$, $\dagger$ marks the value corresponding to the $\left| \Lambda \left(\frac{1}{2}^- \right) \right\rangle_1$ state (equation 1.12) and $\ddagger$ marks the value corresponding to the $\left| \Lambda \left(\frac{1}{2}^- \right) \right\rangle_2$ state (equation 1.13). [§] The result of Mast *et al.* has been recomputed using a full width of 15.6 ± 1.0 MeV.

Figure 1.3 Graphical representation of the radiative widths from table 1.1. In the case of the MIT bag model [9], the $\Lambda(1405)$ widths correspond to the wave-function $|\Lambda(\frac{1}{2}^-)\rangle_1$. The predictions of the various models are indicated by bars and the experimental results by symbols with error bars. Solid bars: NRQM[9, 8]; open bars: relativized CQM[10]; forward slashes: chiral bag[2]; backward slashes: MIT bag [9]; squares: Mast *et al.* [23]; triangles: Bertini[24]; circles: Burkhardt and Lowe [11].

indeed, the $q^4\bar{q}$ component may dominate the wave-function. Mulders, for example, proposed the approximate wave-function [16]

$$\left| \Lambda\left(\frac{1}{2}^{-}\right) \right\rangle \approx 0.83 \left| \Lambda\left(\frac{1}{2}^{-}, q^4\bar{q}\right) \right\rangle \pm 0.56 \left| \Lambda\left(\frac{1}{2}^{-}, q^3\right) \right\rangle. \quad (1.23)$$

Goldman *et al.* [17] quote estimates of 20 keV in the constructive case and 6 keV in the destructive case for the $\Lambda(1405) \rightarrow \gamma\Lambda$ radiative width. They quote a rather large width (325 keV) for the $\Lambda(1405) \rightarrow \gamma\Sigma^0$ channel. Veit *et al.* use the Cloudy-Bag model, which allows the meson fields to couple to the baryon fields throughout the volume of the bag instead of just at the bag surface. They find that the $\Lambda(1405)$ is predominately a $\bar{K}N$ bound state with some admixture of a three-quark state. They claim that the lowest “pure” excited three-quark state is the $\Lambda(1520)$ [3, 4].

Other models for the $\Lambda(1405)$ abandon the three-quark picture. In the Skyrme model [18, 19], the $\Lambda(1405)$ is treated as a negative parity kaon-soliton bound state, clearly not a q^3 structure at all. In the model of Nacher *et al.*, the $\Lambda(1405)$ is created dynamically from a set of ten meson-baryon coupled channels, namely K^-p , $\bar{K}^0N$, $\pi^0\Lambda$, $\pi^0\Sigma^0$, $\eta\Lambda$, $\eta\Sigma^0$, $\pi^+\Sigma^-$, $\pi^-\Sigma^+$, $K^+\Xi^-$, and $K^0\Xi^0$. In the restricted kinematical regime they consider, only the $\pi^+\Sigma^-$, $\pi^-\Sigma^+$, and $\pi^0\Sigma^0$ channels contribute significantly to the resonance shape, although the $\pi^0\Lambda$ channel adds a small background underneath the resonance and the $\bar{K}N$ channel opens up above the resonance. The most significant point about the model is that the shapes of the $\pi^+\Sigma^-$, $\pi^-\Sigma^+$, and $\pi^0\Sigma^0$ cross-sections are asymmetric and that they peak at different masses due to non-negligible $I=1$, $I=0$ cross terms, as shown in figure 1.4, taken from [20].

Figure 1.4 Meson-baryon cross sections in a dynamical resonance model of the $\Lambda(1405)$ [20].

1.2 Previous Experimental results

The largest sample of $\Lambda(1405)$ events was obtained by Hemingway using a mono-energetic (4.2 GeV/c) K^- beam impinging on a hydrogen bubble chamber. The decay of the $\Sigma(1660)$ into $\Sigma\pi\pi$ provided a clean sample of $\Lambda(1405)$'s because most of these decays proceed via $\Sigma(1660) \rightarrow \Lambda(1405)\pi \rightarrow (\Sigma\pi)^0\pi$. He observes more $\Sigma^+\pi^-\pi^+$ events satisfying his cuts than $\Sigma^-\pi^+\pi^+$ events (1106 events compared to 766 events). Simple isospin arguments assign equal probability to both channels. This discrepancy was attributed to interference in the $\Sigma^-\pi^+\pi^+$ channel. The $\Sigma^+\pi^-$ invariant mass showed a significant asymmetric peak, identified as the $\Lambda(1405)$, at about 1.4 GeV. The $\Lambda(1405)$ line shape could not be fit well using a Breit-Wigner parameterization. [21]

Thomas *et al.* found that neither the $\Lambda(1405)$ nor the $\Sigma(1385)$ are produced polarized in π^-p interactions [22].

The radiative decay of the $\Lambda(1405)$ has never been measured directly. Radiative widths for $\Lambda(1405) \rightarrow \Lambda\gamma$ and $\Lambda(1405) \rightarrow \Sigma^0\gamma$ have been inferred from studies of the decays of the K^-p system. Whitehouse *et al.* [12] obtained a branching ratio of $(0.86 \pm 0.12) \times 10^{-3}$ for the $\Lambda\gamma$ channel and $(1.44 \pm 0.23) \times 10^{-3}$ for the $\Sigma^0\gamma$ channel. They measured the photon spectrum using a NaI(Tl) detector from a mono-energetic (680 MeV/c) K^- beam stopping in a liquid hydrogen target. They observed 499 $\Lambda\gamma$ events and 850 $\Sigma^0\gamma$ events.

The extent to which the $\Lambda(1405)$ contributes to these decays is an open model-dependent question. Burkhardt and Lowe [11] observe that the coupling constants for the vertices of the diagrams used in an isobar-model calculation are shared between K^-p pionic and radiative decays because many of the same vertices are involved. They calculated the widths for the pionic channels in the isobar model and deduced the magnitudes of the products and ratios of the coupling constants through comparison with the experimentally well known rates for $K^-p \rightarrow \Lambda\pi^0$, $\Sigma^+\pi^-$, $\Sigma^0\pi^0$, and $\Sigma^-\pi^+$. Burkhardt and Lowe obtain the following radiative widths:

$$\Gamma_{\Lambda(1405) \rightarrow \Lambda\gamma} = 27 \pm 8 \text{ keV},$$

$$\Gamma_{\Lambda(1405) \rightarrow \Sigma^0\gamma} = 10 \pm 4 \text{ or } 23 \pm 7 \text{ keV}.$$

They also obtained a second solution of ~ 0.1 keV for the $\Lambda\gamma$ channel, which they dismissed because it differs greatly from previously published calculations.

The only published measurement of the $\Lambda(1520) \rightarrow \Lambda\gamma$ decay was performed by Mast *et al.* [23]. They observed 258 $K^-p \rightarrow \Lambda\gamma$ events in a sample of some 51000 $K^-p \rightarrow \Lambda + \text{neutrals}$ events in a hydrogen bubble chamber experiment. Fifteen percent of these events were assumed to be $\Lambda(1520) \rightarrow \Sigma^0\gamma$ reactions. This background subtraction is model-dependent because they assumed that the $\Lambda(1520)$ is an SU(3) singlet. By contrast, the models described in section 1.1 predict a substantial amount of mixing among SU(3) multiplets. Mast *et al.* concluded that the branching ratio for the radiative decay of the $\Lambda(1520)$ was 0.0086 ± 0.0014 , from which they determined a partial width of 0.15 ± 0.03 MeV relative to a full width of 17.4 MeV. If the nominal

value of 15.6 ± 1.0 [6] is used, the partial width becomes 0.134 ± 0.023 MeV. The angular distribution of the photons from the decay was consistent with an electric dipole transition. A later CERN experiment reported a radiative width of 33 ± 11 keV for the $\gamma\Lambda$ channel and 47 ± 17 keV for the $\gamma\Sigma^0$ channel [24], but the details of the experiment have not been subsequently published. The two results for the $\gamma\Lambda$ channel differ by a factor of four.

Numerous measurements of the ratio $\Gamma(\Sigma\pi)/\Gamma(\Lambda\pi)$ for the $\Sigma(1385)$ exist [6], but the agreement is poor: values range from 0.08 to 0.21 with large error bars. The Particle Data Group reports an average of 0.135 ± 0.011 .

Attempts to measure the radiative decay of the $\Sigma(1385)$ have not met with much success. Meisner [5] observed one event consistent with $\bar{K}^0 p \rightarrow \Sigma(1385)\pi^+ \rightarrow \Lambda\gamma\pi^+$, from which he determined a branching ratio of 0.17 ± 0.17 . He identified the photon candidate via the missing mass technique and applied a kinematic fit to the event. Colas *et al.* [25] determined the following limits on the $\Sigma(1385)$ radiative decay:

$$\frac{\Gamma(\Sigma(1385) \rightarrow \Lambda\gamma)}{\Gamma(\Sigma(1385) \rightarrow \Lambda\pi^0)} < 0.06 \quad (1.24)$$

$$\frac{\Gamma(\Sigma(1385) \rightarrow \Sigma^0\gamma)}{\Gamma(\Sigma(1385) \rightarrow \Lambda\pi^0)} < 0.05. \quad (1.25)$$

Their results do not discriminate between any of the models, whose predictions are unilaterally smaller than these limits.

1.3 The experiment

Few photo-production experiments have been done for the lower-lying excite-state hyperons. The $\Sigma^0(1385)$ production cross section was measured in the late 1960's using bubble chambers [26]. The dependence on the incident photon energy E_γ was determined to be

$$\sigma(\Sigma^0(1385)) = \frac{1.5 \pm 0.9}{E_\gamma^2} \mu b \quad (1.26)$$

for E_γ (in units of GeV) between 2 and 5.8 GeV. Boyarski *et al.* [27] found that the sum of the cross sections for the $\Lambda(1405)$ and the $\Sigma^0(1385)$ had the dependence

$$\sigma(\Sigma^0(1385)) + \sigma(\Lambda(1405)) = \frac{3.7 \pm 1.3}{E_\gamma^2} \mu b. \quad (1.27)$$

using an 11 GeV Bremsstrahlung beam. One experiment [28] studied the production of the $\Lambda(1520)$ and angular distributions for the decay of the $\Lambda(1520)$ to K^-p , but the researchers did not attempt to look for radiative decays. Experiments studying the $\Lambda(1405)$ and the $\Sigma(1385)$ generally used kaon or pion beams. Experiments involving kaon beams suffer from the drawback that both the $\Lambda(1405)$ and the $\Sigma(1385)$ are below the $\bar{K}N$ threshold. Direct production experiments ($K^-p \rightarrow Y^*$) can only access the high mass tails of these states. On the other hand, photo-production experiments do not distinguish between the $I=0$ $\Lambda(1405)$ state and the $I=1$ $\Sigma(1385)$ state in the production channel because the electromagnetic interaction does not conserve isospin.

In this thesis I will be describing the results of an experiment in which the hyperons are photo-produced in the reaction

$$\gamma p \longrightarrow K^+ Y^*, \quad (1.28)$$

where Y^* denotes the excited state hyperon. The data were taken during the first G1 (real photon) running period (May-June 1998) in CLAS at Jefferson Lab. The accelerator, the CLAS detector, and details of the trigger will be described in section 2. The primary electron beam energy E_0 was about 2.5 GeV. The photon energy range was between 20% and 95% of E_0 . The photon beam was not polarized. About 250 million analyzable triggers were taken, of which about 50-60% had tracks.

Chapter 2

Experimental Apparatus

The Thomas Jefferson National Accelerator Facility (TJNAF) is the site of a recirculating linear electron accelerator capable of delivering beam to three experimental halls simultaneously. A beam of 45 MeV electrons is injected into one of two parallel linear accelerators connected on both ends by a set of five recirculation arcs, as shown in figure 2.1. Each time the electron beam passes through one of the linacs the electrons' energies are boosted by ΔE , where ΔE is determined by the linac settings. Typically ΔE is about 0.4 GeV. At the end of the second linac, the beam can be extracted from any of the five passes. At the typical linac setting the beam energies available to the halls range from 800 MeV to 4.0 GeV in multiples of 800 MeV.

The electron beam circulates in bunches separated by about two-thirds ns. The extraction separates the beam into three beams each with 499 MHz bunch structure. In other words, each bunch is separated by 2.004 ns when the electron beam reaches one of the experimental halls. The accelerator can deliver beam currents ranging from ~ 100 pA to ~ 100 μ A.

The three experimental end stations are known as Halls A, B, and C. Hall B is the middle and smallest end station. A view of the hall is shown in figure 2.2. Both electron and photon beam experiments have been run in this hall. In the latter

Figure 2.1 Sketch of the accelerator facility.

case, the electron beam is converted to a photon beam by the Hall B Bremsstrahlung tagging system. The photon then interacts in a target enclosed by the CEBAF Large acceptance spectrometer (CLAS).

CLAS is an almost 4π toroidal spectrometer composed of many layers of detectors. Figure 2.3 shows a plan view with the beam direction pointing from left to right; figure 2.4 shows a head-on view looking in the direction of the beam. The spectrometer is divided into six sectors by a set of superconducting magnets that produce a roughly toroidal magnetic field. Each sector is instrumented with the detector components listed below. Three drift chamber regions enclose the target to track particles through CLAS and measure their momenta. For electron-beam running, cherenkov counters are used to identify electrons in the forward angles (below 45°). The drift chambers

Figure 2.2 Hall B

Figure 2.3 Side cut-away view of the CEBAF Large Acceptance spectrometer

Figure 2.4 View looking downstream of the CEBAF Large Acceptance Spectrometer.

and cherenkov counters are enclosed by a shell of scintillators intended for charged particle identification. Behind the scintillators downstream of the target lie the electromagnetic calorimeters used for pion and electron identification and detection of neutrons and photons.

2.1 Detector components

2.1.1 The tagging spectrometer (tagger)

At the entrance to Hall B, a thin 10^{-4} ($646 \mu\text{g}/\text{cm}^2$) radiation-length gold foil radiator converts the incident electron beam into a photon beam with a characteristic bremsstrahlung energy spectrum. Ignoring the negligible recoil energy of the target nucleus, the energy E_γ of the photon is determined by

$$E_\gamma = E_0 - E_e, \quad (2.1)$$

where E_0 is the energy of the incident electron beam and E_e is the energy of the scattered electron. The tagger magnet bends the scattered electron's trajectory away from the incident direction toward the magnet's focal plane covered by 61 two cm thick scintillators ("T-counters") for good timing and 384 smaller scintillators ("E-counters") for finer energy resolution. Figure 2.5 shows the tagger magnet and the E- and T-counter planes. The E-counters overlap each other to form 767 energy bins, thus providing $\sim 1 \times 10^{-3}$ momentum resolution. The tagging range is between 20% and 95% of the incident electron energy. Figure 2.6 shows a representative energy

spectrum for 2.543 GeV incident electron beam. The spiky behavior arises from differing sizes of the bins. For more details see [29].

Each T-counter is read out from two sides and the signals are discriminated by Phillips 715 constant fraction discriminators. The resulting pulses are then digitized by LeCroy 1872A FASTBUS TDCs and counted by four CAEN V 560 E 16-channel VME scalers. The E-counters are read out with LeCroy 1877 multi-hit common-stop FASTBUS TDCs.

2.1.2 Collimators

Two collimators separated by two sweeping magnets were placed in the beam-line between the radiator and the target. These collimators were used to cut the tails of the photon beam. The bore diameter of the first collimator was 8.61 mm; the diameter of the second was 17.3 mm. The purpose of the sweeper magnets was to bend the charged debris produced by the tails of the photon beam interacting with the primary collimator material away from the beam axis. The second cleanup collimator helped remove any remaining background.

2.1.3 The target

The CLAS target was a 185 mm long $\times$ 59.68 mm diameter liquid hydrogen cell. The cell walls were 170 μm thick Mylar. The volume of the target was 447.88 cm^3 . The target is housed in a vacuum chamber whose walls are formed from a 1.5 mm

Figure 2.5 The Hall B tagging system.

thick carbon-fiber tube whose density is about 1.6 g/cm^3 . A schematic diagram of the target and assembly is shown in figure 2.7.

2.1.4 The start counter

The start counter is a set of three independent 3 mm thick BC408 scintillators, each arranged to cover two adjacent sectors. The scintillators formed a hexagonal cone (nose region) and a hexagonal tube (leg region). A sketch of one of the three detectors is shown in figure 2.8. The distance of closest approach between the leg regions and the beam line is approximately 11 cm. The scintillators are bent in three places such that the “nose” region covers down to the 7° angle with respect to the beam line from the center of CLAS. These counters are intended to provide a signal

Figure 2.6 Energy spectrum of the incident photons. This histogram was used as the photon energy source for the event generator.

(in conjunction with the time-of-flight counters) of a hadronic event in CLAS and to link this event with a hit in the tagger. For more information, see [30].

2.1.5 The main torus

The magnetic field for CLAS is provided by six superconducting magnets constructed by Oxford instruments. The magnets are placed at 60° intervals about the beam line and arranged so that the field is approximately toroidal. The maximum intensity of the magnetic field is about 2 Tesla. Each magnet is about 2.5 m high and 5 m wide; the shape is shown in figure 2.9. The polarity is such that negative particles bend towards the beam axis and positive particles bend away from the beam axis. The magnetic field does not appreciably change the azimuthal angles of particle trajectories [31].

2.1.6 Drift Chambers

The drift chambers are composed of eighteen chambers divided into three regions (R1, R2, and R3) and consisting of 35,148 cells. Each region is divided into two superlayers consisting of six layers of hexagonal drift cells (figure 2.10). In each region, the wires in one superlayer (the “axial” layer) are arranged axially with respect to the magnetic field and the wires in the second region (the “stereo” layer) are oriented at a 6° stereo angle with respect to the radial direction. The region 2 and 3 drift chambers each consist of 12 sense-wire layers; due to space constraints, region 1 drift chambers consist of 10 layers. Each cell consists of six “field” wires surrounding a

Figure 2.8 Sketch of one of the three start counter scintillators showing the support structures, light guides and phototubes.

Figure 2.9 Schematic view of the six superconducting toroidal magnets showing the approximate shape of the coils.

single “sense” wire. The sense wires were 20 μm diameter gold-plated tungsten; the field wires were 140 μm diameter gold-plated aluminum. The field wires were held at high voltage.

The chambers were filled with a gas mixture of 90% Ar, 10% CO_2 . Charged particles traversing CLAS ionize the gas as they pass through. The liberated electrons drift towards the sense wires. The sense wires nearest to a particle’s trajectory register “hits” along the trajectory. The time it takes for the electrons to drift to the sense wires can be used to determine the drift distance between the sense wires and the particle’s trajectory by using a drift time-to-drift distance function. Particles passing through region 2 experience the largest magnetic field. The fringe fields in regions 1 and 3 are very low. Consequently, most of the information about the charge and the momentum of a charged particle traversing CLAS comes from hits in Region 2.

The drift chamber signals are capacitively decoupled from the high voltage and then amplified and discriminated. The resulting pulses from pairs of wires are multiplexed together. The multiplexed drift chamber signals are read out with LeCroy 1877 common stop multi-hit FASTBUS TDCs.

The design goals for 1 GeV/c particles were $\delta p/p \leq 0.5\%$ and $\delta\theta, \delta\phi \leq 2$ mrad. The average angular resolution determined from studies of electron-proton elastic scattering was ± 10 mrad; the momentum resolution $\delta p/p$ was about 1-2.5%. Since both the electron and the proton tracks contribute to the error, for a single track

Figure 2.10 Representative view of part of region 3 showing a track passing through the drift chambers.

the angular resolution and percent momentum error are expected to be closer to the design goal[31].

2.1.7 Time-of-flight Scintillators

The Time-of-flight system comprises a shell of 342 BC-408 scintillators, 57 per sector, enclosing the drift chambers and cherenkov counters. The scintillators cover a range in theta between 8° downstream of the target to 142° upstream of the target. Each scintillator is 5.08 cm thick; they vary in length between 32 cm in the extreme downstream direction to 450 cm for larger angles. In the forward angle (up to about 45°) the scintillators are 15 cm wide; for the larger angles the widths are 22 cm, except for the very back angle counters, which are again 15 cm wide. The width of each scintillator subtends about 1.5° in the lab angle and the length covers the azimuthal range from -29° to 29° relative to the center of each sector. As a cost-cutting measure the last eighteen counters in the back angles were passively ganged together into nine pairs. Because of this pairing, each sector comprises 48 electronic channels per side. The analog signals are read out by LeCroy 1881M FASTBUS ADCs to provide pulse height measurements. The signals for each channel are also discriminated using LeCroy 2313 CAMAC leading-edge discriminators to provide timing information that is digitized by LeCroy 1872A-mod-100 FASTBUS TDCs. The average time resolution for electrons is 163 ps; for pions from electron-pion coincidences, the resolution ranges between 150 ps in the forward angle to ~ 300 -350 ps in the back angles [32].

2.1.8 Electromagnetic Calorimeters

Each sector of CLAS is covered by a wedged-shaped sampling calorimeter consisting of 39 lead-scintillator layers comprising sixteen radiation lengths of material. Each calorimeter subtends an angular range between 8° and 45° in θ_{LAB} . The scintillators in each layer are 10 cm thick: the interleaving lead sheets are 2.2 mm thick. Each scintillator layer is divided into 36 paddles arranged parallel to a side of the wedge in one of three orientations (“views”). Each layer’s orientation after the first is rotated 120° with respect to the previous layer. The three views form 13 layers: the first five of these layers comprise the “inner” stack and the remaining eight comprise the “outer” stack. For each stack, the light from all of the 36 scintillators in each view are read out by a single photomultiplier tube. The photomultiplier tubes were a combination of EMI 9954 and Phillips XP2262 PMTs[33].

Limited large-angle coverage is provided by the Large Angle Calorimeters in sectors one and two. The calorimeters subtend lab angles between 45° and 70° in θ_{lab} . Similar in design to the forward calorimeters, the large angle calorimeters comprised 33 layers of 1.5 cm thick NE110A plastic scintillator interleaved with 2 mm thick lead. Each layer is orientated at right angles with respect to the previous layer. The scintillators are separated into two (inner and outer) layers. In each layer, all the scintillators with the same orientation but at different distances from the beam line are grouped together into “stacks”; the light from each scintillator in a stack is transported via light guides to a single EMI 9945A photomultiplier tube.

As for the Time-of-flight counters, pulse height information is provided by LeCroy 1881M FASTBUS ADCs and timing information is obtained from using LeCroy 2313 CAMAC discriminators feeding LeCroy 1872A-mod-100 FASTBUS TDCs.

2.1.9 Downstream devices

The photon flux is measured by a set of three detectors: the Pair Spectrometer (PS), the Pair Counter (PC), and the Total Absorption Shower Counter (TASC), in order of position along the beam direction. A schematic diagram of these devices is shown in figure 2.11.

The pair spectrometer consists of a converter, a dipole magnet, and a set of four scintillator paddles on each side of the beam line. A small fraction of the number of photons that reach the downstream devices are converted into e^+e^- pairs by a converter. The converters' thicknesses were 1% or 2% of a radiation length. The 1% converter was generally used. The electrons and positrons are swept in opposite directions by the magnet and deposit energy in the scintillators.

The pair counter is a photon-beam monitoring device consisting of a veto scintillator followed by a thin lead converter, another thin scintillator, and a set of four partially overlapping scintillators. Electron-positron pairs are created by the photon beam striking the lead sheet with an efficiency of about 1.5%. Coincidences are measured between the scintillator nearest to the converter and the four overlapping scintillators behind it.

Figure 2.11 Sketch of the downstream devices: not drawn to scale.

The total absorption counter (TASC) is a set of four lead-glass blocks arranged in a two-by-two array to which four photo-multipliers are attached. A Phillips XP4312B phototube is attached to one end of each block. A LeCroy 1881M ADC is used to measure the energy for each block. Unlike the pair counter and the pair spectrometer, the efficiency of the TASC for detecting photons is almost 100%, but it can only be operated in low intensity beams ($\sim 100pA$). Special normalization runs are required to calibrate the pair spectrometer and the pair counter with respect to the total absorption counter so that the photon flux may be measured during production running at higher beam intensity (typically ~ 8 nA).

Four CAEN V 560 E 16-channel VME scalers were used to count coincidences between the downstream devices and the 61 T-counters. The first scaler scaled co-

incidences between the even T-counters and the pair counter. The second scaler scaled coincidences between the odd T-counters and the pair spectrometer. During normalization runs, the third and fourth scalers scaled coincidences between the total absorption counter and all 61 of the T-counters. During production runs, the third scaler counted accidental coincidences between the odd T-counters and the pair spectrometer signal and the fourth counted accidental coincidences between the even T-counters and the pair counter signal. For production runs, pair spectrometer/T-counter coincidences are only counted for half the tagger focal plane. The distribution is assumed to be smoothly varying so that the unmeasured coincidences can be interpolated from the measured ones.

2.2 Trigger electronics

For photon-beam running, the trigger for reading out the FASTBUS crates is provided by a coincidence between the Start Counter, the Tagger Master OR, and the Time-of-flight counters.

2.2.1 Tagger Master-OR

For each T-counter, a coincidence between the left and right signals and one of the set of corresponding E-counters is required. The resulting signals from all 61 T-counters are then OR'd together. Figure 2.12 is a schematic diagram of the logic.

Figure 2.12 Schematic diagram of the tagger trigger and readout electronics.

2.2.2 Prescaling

In the energy range of the tagged photon beam (0.5-2.4 GeV) used in the G1 experiment, the γp total cross section is not flat: the larger cross sections are for the lower beam energies (figure 2.13). Furthermore, the bremsstrahlung spectrum itself has a strong $1/E_\gamma$ dependence (figure 2.14). Most of the physics of interest comes from the higher energy photons. Due to limitations in the data acquisition system, the lower energy part of the spectrum was prescaled in order to make the event rate more even across the energy range. The photon energy spectrum was divided into four energy regions. Table 2.1 summarizes the energy regions and the prescale factors applied to the trigger signals.

2.2.3 Start Counter trigger

The signals from each side of a paddle are actively summed using a Phillips Model 740 Fan-In/Out. The result is discriminated by a Phillips 730 Tri-mode discriminator. Logic signals from each pair are OR'ed together to provide the start counter trigger signal. A schematic diagram of the circuit is shown in figure 2.15.

Energy range	T-counter range	Prescale factor
0.80-0.95	1-16	1
0.55-0.80	17-29	2
0.33-0.55	30-45	4
0.2-0.33	46-61	8

Table 2.1 Tagger trigger prescale ranges.

Figure 2.13 The γp total cross section. The shaded region is the region covered by the tagger for 2.5 GeV incident electron beam. Figure courtesy of Reinhard Schumacher.

Figure 2.14 The bremsstrahlung spectrum $d\sigma/dE_\gamma$ as a function of the photon energy E_γ .

Figure 2.15 Schematic diagram of the start counter trigger electronics.

2.2.4 TOF pretrigger boards

The time-of-flight pretrigger boards provide inputs to the Level 1 Trigger electronics. The pretrigger boards discriminate the analogue sum of signals from adjacent pairs of counters. The resulting signal is discriminated to produce a so-called “double” logic signal. “Double” signals from adjacent counters are OR’ed together to produce “triple” logic signals. For each sector, the input to the trigger consists of the sixteen “double” pulses from the first sixteen counters and sixteen “triples” pulses from the remaining counters (17-48). The width of the signals produced by the TOF pretrigger boards was 100 ns. The thresholds were set to 200 mV. A schematic diagram of the pretrigger electronics is shown in figure 2.16.

2.2.5 CAMAC logic

Coincidences between the tagger Master OR signal and the Start Counter or the downstream devices are remote-controllable via a set of CAMAC units. The start counter and Master-OR signals are plugged into the test inputs of two LeCroy 4413 leading-edge discriminators. This serves two purposes: (1) it regenerates the signals, which have been degraded over the cable length needed to reach the CAMAC logic; (2) the widths of the logic pulses can now be controlled independently from the electronics that produced the original signals. The start counter trigger signal width was set to 15 ns; the MOR trigger signal width was set to ~ 7 ns. Diagrams of the CAMAC logic are shown in figures 2.17 and 2.18.

Figure 2.16 Schematic diagram of the time-of-flight pretrigger electronics.

Figure 2.17 Schematic diagram of the CAMAC logic for the CLAS trigger: (A) Discriminators

Figure 2.18 Schematic diagram of the CAMAC logic for the CLAS trigger: (B) Logic unit

2.2.6 Level 1 trigger

Signals from the time-of-flight pretrigger boards provide inputs to the Level 1 trigger [34]. The Level 1 hardware compares the input pattern to the pattern loaded into the Level 1 memory lookup tables and latches the trigger signal if a match is found. The memory lookup occurs in three stages. The first stage maps the pretrigger signals in each sector to four intermediate bits. This allows for the trigger to include sub-ranges of the scintillator array in each sector. The intermediate bits are sent on the second stage memory-lookup, which checks combinations of the intermediate bits called “sector events”. Finally, “sector events” are correlated in the third stage. The tagger-start counter coincidence signal enters Level 1 asynchronously at the third stage. The trigger signal is then passed on to the Trigger Supervisor, which generates the gates and interrupts for the VME and FASTBUS hardware.

2.3 Data Acquisition

A schematic diagram of the data acquisition system is shown in figure 2.19. Data words are read out in the FASTBUS crates for each ADC and TDC, translated into BOS format at the front end, and transported via fast Ethernet to a Sun SMP, where the fragments of each event are assembled and written out in BOS format to a RAID array of hard drives. See chapter 3.1 for a brief description of BOS. Finally, the data is shipped across dedicated lines to the Jefferson Lab computer center and stored on tape in a tape silo.

Figure 2.19 Schematic diagram of the data acquisition system (diagram courtesy of Konstantin Loukachine).

Chapter 3

Reconstruction Software and Data Reduction

3.1 Reconstruction

The CLAS reconstruction code consists of thousands of lines of code in both Fortran and C. Its function is to convert the raw data (hits registered in ADCs and TDCs) into a form more useful for physics analysis in a process known as “cooking”. Each detector component has a corresponding reconstruction package.

The tracking code is the most complicated part of the reconstruction. It is divided into two steps: hit-based and time-based tracking. During the hit-based phase, the digitized hits on the wires are grouped into segments within single superlayers. The clusters in different superlayers within the same sector are then linked together to form tracks. Time-based tracking uses the timing information from the time-of-flight scintillators to improve the quality of the tracks by correcting the measured drift times. The drift times are converted to distances and each track is refitted using the corrected distances between the track and the wires[31].

The outputs of the packages are reconstructed result banks stored in the BOS format. BOS is a Fortran-77 based dynamic memory management system [35]. A “bank” is a collection of data words attached to a “header” containing information

about the number of rows and columns of data in the bank, the bank identifier, and the location of the next bank in memory. It is approximately analogous to a C structure in that the data words can be a mix of floating point numbers, characters, and integers. Data is transferred between the detector packages via these result banks. Details specific to photon beam running will be elaborated in the following sections.

3.1.1 Tagger reconstruction

The tagger reconstruction package [36] reports the time of the photon at the nominal center of the target for each “good” hit in the focal plane. First, overflow and underflow TDC hits in the T-counters are removed and the TDC channels are converted into nanoseconds. The E-counter times are required to be within 700-1200 TDC channels. Then the E and T counter hits are matched geometrically according to a table and required to be in time with each other within a window of 20 ns. Figure 3.1 shows the relationship between T- and E-counter hits. The true E-T coincidences show up in the solid band. Several E-counters can match geometrically with a single T-counter. The T-counter times are corrected to the nearest RF bucket (accelerator clock). The energy of the photon for each E-counter that survives the cuts is obtained from another table as a fraction of the incident electron beam energy.

3.1.2 Start counter reconstruction

Because of the nonlinear shape of each segment of the start counter, the speed of light propagation in the material depends on whether the light goes directly to the

Figure 3.1 E-counter/T-counter coincidence.

tube or must go around the bend to reach the tube. Consequently, the average of the times at the two tubes is not independent of position. This is further complicated by the possibility that two particles could hit a single segment. For one charged particle striking one of the three pairs of scintillators, the best time is determined by the pulse-height weighted average of the two tubes' times (see [30] for more information). For the case where two tracks point to the same start counter pair, tracking is used to determine the position of the hit and the nearest single-tube time is used, corrected to the hit position.

3.1.3 Time-of-flight reconstruction

The Time-of-flight reconstruction code reports the time and energy deposition for each hit in the scintillator array. The raw ADC and TDC information for each struck scintillator in each event is stored in the SC bank. Each "hit" consists of a number identifying the scintillator and the ADC and TDC information for both the left side and the right side of the paddle. The reconstruction algorithm for the time-of-flight comprises four main steps.

Initialization

At the beginning of the run the calibration constants are read from the Map database and stored in global memory. The map file is SC_CALIBRATIONS.map and is located in \$CLAS_PARMS/Maps.

Construction of the SC1 banks

The SC1 bank contains those hits in the SC bank that have valid timing information. Those hits that have no useful TDC values on either side of the paddle are discarded.

The data remains separated between left and right tubes for each hit but the TDC values are converted to time units (ns) and the ADC values are converted to energy units (MeV). The conversion for the ADC values is

$$E = \frac{10(A - P)}{A_{NMIP}} \text{MeV}, \quad (3.1)$$

where A is the ADC value of the hit in channels, P is the pedestal, and A_{NMIP} is the pedestal-subtracted ADC value corresponding to a normally incident minimum ionizing particle (NMIP), which deposits 10 MeV in the scintillator material. The TDC values are converted to nanoseconds and corrected for slewing according to

$$T = c_0 + c_1 t + c_2 t^2 + t_w(A, P), \quad (3.2)$$

where t is the TDC value in channels, c_0 is a constant that relates the timing of each of the time-of-flight channels to each other, c_1 is the conversion factor between TDC channels and nanoseconds, c_2 is a small constant that accounts for the nonlinearity of the TDC and t_w is a correction dependent on the pulse height:

$$t_w = f_w\left(\frac{600}{V_T}\right) - f_w\left(\frac{A - P}{V_T}\right); \quad (3.3)$$

$$f_w(x) = \begin{cases} \frac{w_2}{x^{w_3}} & \text{if } x < w_o \\ \frac{w_2}{w_o^{w_3}} (1 + w_3) - \frac{w_2 w_3}{w_o^{w_3+1}} & \text{if } x > w_o \end{cases}$$

V_T is the discriminator threshold in ADC channels. The time-walk correction is set to zero for ADC zeroes or overflows or TDC zeroes or overflows. The time T is set to the error value of 100000.0 for TDC overflows or -100000.0 for TDC underflows.

Construction of the SCR banks

During the construction of the SCR bank, information from both sides of each paddle with an entry in the SC1 bank is combined to form the best timing, energy and position information for the hit. If good timing values for both sides of a given struck scintillator are found, the time of the hit is determined by

$$T = \frac{T_L + T_R}{2} - y \frac{v_R - v_L}{2v_L v_R}, \quad (3.4)$$

where T_L (T_R) is the walk-corrected time in ns for the left (right) tube and v_L (v_R) is the speed of light propagation toward the left (right) tube. The position of the hit along the scintillator relative to the center of the paddle is given by

$$y = \frac{v_L v_R (T_L - T_R)}{v_L + v_R}. \quad (3.5)$$

The propagation speeds v_L and v_R could in principle be different, in which case the time of the hit would depend on the position of the hit along the bar, but in practice a single average value is measured and the position dependence cancels. If both left and right ADCs are present and within range for a given paddle, the energy deposition is determined by

$$E = \sqrt{E_L E_R} e^{y \frac{\lambda_R - \lambda_L}{\lambda_L \lambda_R}}, \quad (3.6)$$

where E_L and E_R are the pedestal-corrected ADC values in units of MeV from the SC1 bank. y is the position along the length of the paddle relative to the center of the paddle, and λ_R (λ_L) is the attenuation length for light propagation to the right (left) tube. In the case where only the left ADC is present and valid, the energy deposition is given by

$$E = E_L e^{-y/\lambda_L}. \quad (3.7)$$

In the case where only the left TDC is valid, the time is given by

$$T = T_L - \frac{y}{v_R}. \quad (3.8)$$

Analogous equations hold for the case where only the right ADC or TDC is valid. By default, the code tries to use tracking information to find y . If this fails, and both ADCs are present for the given hit, then the position can be calculated rather crudely according to:

$$y = \frac{\lambda_L \lambda_R}{\lambda_L + \lambda_R} \ln \left(\frac{E_L}{E_R} \right). \quad (3.9)$$

Errors are computed for the time, the energy deposition, and the y -position of each hit.

A status word is set according to table 3.1 depending on the quality of the hit, 15 being best, 1 and 4 being worst.

Construction of the SCRC banks

The SCRC bank takes individual hits in the SCR bank and looks for clusters in which adjacent hits arise from a single particle depositing energy in two scintillators.

Value	Meaning
1	Left TDC present: no right tube data.
2	Left ADC present: hit discarded.
3	Left ADC and TDC present: no right tube data.
4	Right TDC present: no left tube data.
5	Both TDCs present, no ADC information
6	Left ADC, right TDC present
7	Left ADC and TDC, right TDC present but no right ADC
8	Right ADC present; hit discarded.
9	Left TDC, right ADC present
10	Both ADCs present; no TDC information: hit discarded.
11	Both ADCs present, left TDC present, no right TDC
12	Right TDC and ADC only: no left tube data.
13	Both TDCs present: right ADC present; left ADC missing.
14	Both ADCs present, right TDC present; left TDC missing.
15	Complete set: left TDC, ADC; right TDC, ADC

Table 3.1 SCR status word definition.

The clusters are chosen based on overlap in space and time. If the difference in times or positions is less than three times the uncertainties in the differences, then the two hits in the adjacent counters are considered to be one cluster. If three consecutive counters fire in a given event, then the adjacent pair that are better matched in space and time are clusterized; the other scintillator information enters the SCRC bank as a single hit. If a cluster is found, the energies of the two hits are added together to form the "cluster energy" and the time is computed from an energy-weighted average of the two hits forming the cluster. The x- and z-positions are simple averages of the individual hit positions. The y-position (along the length of the bar) is an energy-weighted average. If no cluster is found, the SCRC status word is the same as the

SCR status word and the i.d. is the same as the SCR i.d.; otherwise, the SCRC status word is the sum of the status word for the first hit in the cluster and 100 times the status word for the second hit in the cluster. The i.d. is the SCR i.d. for the first hit in the cluster.

3.1.4 Determining the start time of the event

To eliminate out-of-time accidentals, a coincidence between the tagger and the start counter is required in software. For each track that hits a time-of-flight scintillator, the particle's speed is determined from the difference between the times at the start counter plane and the TOF plane and the track-length difference between the start counter plane and the TOF plane. The track is then propagated back to the "vertex" in order to determine a "vertex time". For each track, the vertex time is compared to each entry in the tagger result bank. The tagger code reports the time at the nominal center of the target. This time is corrected for the difference (in ns) between the center and the projection of each track to an arbitrary plane known as the "vertex" plane. The best tagger hit is the one for which the difference between the tagger time at the target and the start counter vertex time is the smallest less than 5 ns. Figure 3.2 shows the relative timing between the tagger and the start counter. Accidentals between the start counter and the tagger show up as a diagonal band in figure 3.2A. The true coincidences fall in the large peak centered on zero in figure 3.2B.

(A) Tagger time vs start counter/tagger difference

(B) Projection onto (start counter - tagger) axis

Figure 3.2 Start counter/ tagger relative timing showing the prominent “true” peak and several accidental bumps. The main peak has a resolution of 434 ps.

3.1.5 Photon flux

In order to calculate an absolute cross section, the incident flux of photons on the target must be known. The shape of the bremsstrahlung spectrum is determined from the low intensity normalization runs. Coincidences between hits in the tagger focal plane and the total absorption counter are required in software and the number of photons ($N_{\gamma, norm}$) satisfying this coincidence requirement are accumulated in 767 E-counter bins. The scaler counts for T-counter hits in coincidence with the total absorption counter and the pair spectrometer are also integrated over each normalization run. The number of photons for a production run is scaled up from the number of photons from a reference normalization run by the ratio of pair-spectrometer scaler counts. For each production run, the ratio is given by

$$r_{PS}(T_i) = \frac{N_{prod}^{in\ time}(T_i) - \alpha N_{prod}^{delayed}(T_i)}{N_{norm}^{in\ time}(T_i)}, \quad (3.10)$$

where $\alpha \approx 0.9$ is an empirical constant that accounts for a small difference in effective gate widths between the scalers measuring the “in time” coincidences and the scalers measuring accidental coincidences between the tagger counters and the pair spectrometer. $N_{prod}^{in\ time}$ ($N_{norm}^{in\ time}$) is the accumulated number of “in time” coincidences between T-counter T_i and the pair spectrometer for the production (normalization) run, and $N_{prod}^{delayed}$ is the accumulated number of accidental coincidences between T-counter T_i and the pair spectrometer for the production run. The number of photons for each

E-counter bin E_j for the production run is given by

$$N_{\gamma,prod}(E_j) = \frac{P_{norm}}{P_{prod}} R_{PS}(E_j) N_{\gamma,norm}(E_j) \frac{L_{prod}}{L_{norm}}. \quad (3.11)$$

$$R_{PS}(E_j) = \frac{\sum_{T_i=T_i^{min}(E_j)}^{T_i^{max}(E_j)} r_{PS}(T_i)}{T_i^{max}(E_j) - T_i^{min}(E_j) + 1}. \quad (3.12)$$

where P_{norm} (P_{prod}) is the DAQ prescale factor for the normalization (production) run. L_{norm} (L_{prod}) is the livetime for the normalization (production) run. and $T_i^{max}(E_j)$ ($T_i^{min}(E_j)$) is the maximum (minimum) T-bin corresponding to E-bin E_j . The normalization procedure is described in more detail in [37].

3.2 Data Reduction

Not all of the events written to disk by the data acquisition system contain tracks. Furthermore, due to the open trigger used during the G1A running period, only a small fraction of the “good” events pertain to the particular reactions of interest. In order to reduce the volume of data to a manageable size, we employed the following data reduction technique.

During the process of “cooking” the alc program converts the data from raw BOS banks to reconstructed BOS banks. Only those events that contain tracks at the hit-based level are written to tape. The cooked files are then filtered for events with kaon candidates, which are written out to a separate file. The kaon filter checks for at least one track whose time-of-flight mass is between $0.3 \text{ GeV}/c^2$ and $0.7 \text{ GeV}/c^2$. The details of this technique will be described in section 4. This filter reduces the size of the data set by about two orders of magnitude. A second filter is applied to

the filtered files to separate out events with both a loosely defined kaon and a loosely defined proton. The proton cut was $0.7\text{-}1.3\text{ GeV}/c^2$.

Chapter 4

Particle Identification

As show in figure 4.1, pions, protons, and kaons form well separated bands in their speed β as a function of momentum up to about 2 GeV/c. The kaon band is enhanced in this figure by the kaon filtering process. The speed β depends on the flight time and the path length l from the target to the scintillator according to the equation

$$\beta = \frac{l}{c\Delta t}. \quad (4.1)$$

Kaons are identified by calculating the mass from the time-of-flight Δt and the momentum p . The mass is given by

$$M = \frac{p}{\gamma\beta}, \quad (4.2)$$

where $\gamma = (1 - \beta^2)^{-1/2}$. Figure 4.2 shows a representative mass spectrum focusing on the kaon peak. A significant pion and to a lesser extent proton background exists underneath the kaon peak. As the momentum increases the resolution on the momentum gets worse. As shown in figure 4.3A, this causes the mass of the kaon peak to broaden as a function of increasing momentum. Consequently, using a fixed range in mass to identify kaons will allow more background into the cut relative to the signal for low kaon momenta than for high momenta. To address this problem I fit the mass

Figure 4.1 β as a function of the momentum for the filtered data sample. The sharp edges of the kaon band reflect the selection cuts.

as a function of momentum in the lab with a Gaussian plus a polynomial background for slices of momentum. Figure 4.3B shows that the standard deviation of the fit generally increases with momentum. Assuming Gaussian statistics, a $\pm 3\sigma$ cut should contain 99.7% of the signal. The kaon identification cut is defined as follows:

$$\begin{aligned}
 0.35 < M < 0.65 \text{ GeV}/c^2; \\
 |M - M_K| < 3\sigma; \\
 \sigma = -0.000873 + 0.0199641p + 0.0033829p^2
 \end{aligned} \tag{4.3}$$

where p is the momentum of the kaon in the lab frame. Figure 4.3A also illustrates the problem of missing the correct RF bucket. The diagonal bands correspond to pions whose start time is off by some multiple of 2.004 ns.

Because the problem of contamination is much smaller for protons than for kaons, protons are identified with a simple mass cut in the range 0.85-1.2 GeV/c^2 .

The particle identification cuts for the pions was slightly more complicated. Since the mass as calculated in equation 4.2 blows up as β approaches 1, a cut on β was applied for the upper bound:

$$\beta < \frac{p}{\sqrt{p^2 + 0.08^2}} + 0.02. \tag{4.4}$$

The lower bound was a simple mass cut of 0.3 GeV. The β cut has the desirable effect of eliminating much of the electron contamination. Figure 4.4 shows the curves for the upper and lower bounds superimposed on the beta versus momentum plot for pion candidates.

Figure 4.2 Hadronic mass for positive particles as calculated from the time-of-flight and the momentum.

Figure 4.3 Dependence of the mass from time-of-flight on momentum p . (A) Two dimensional histogram of mass vs. momentum; (B) Resolution of the kaon peak as a function of momentum. The dashed curve is a second-order polynomial fit to the data.

Generally, before time-based tracking can be applied to a hit-based track candidate, a match between the track and one of the scintillators is required so that the speed of the particle can be calculated; however, since the polarity of the field was such that negative particles were bent towards the beam pipe, soft pi-minuses often did not hit any of the scintillators. For these cases, negative particles with no beta assigned to them were considered to be pi minuses; beta was calculated from a mass assumption. For reactions involving the $\Lambda(1405)$ and the $\Sigma(1385)$, this assumption is safe because the threshold for producing a K minus in the reaction corresponds to a hyperon mass of $M_Y = M_p + M_{K^-} = 1432$ MeV.

The hyperons are identified by calculating the missing mass for the reaction $\gamma p \rightarrow K^+ X$, where X represents the “missing” particle. In terms of the four-momenta P_N , k , and P_K of the proton target, the incident photon, and the final-state kaon, respectively, the mass M_Y of the hyperon is determined by

$$\begin{aligned} M_Y^2 &= (P_N + k - P_K) \cdot (P_N + k - P_K) \\ &= s + M_K^2 - 2P_K \cdot (P_N + k), \end{aligned} \quad (4.5)$$

where $s = E_{CM}^2$ is the invariant Mandelstam quantity. In the lab frame, $k = (E_\gamma^{LAB}, 0, 0, E_\gamma^{LAB})$ and $P_N = (M_p, 0, 0, 0)$, so that

$$M_Y^2 = s + M_K^2 - 2E_K^{LAB}(E_\gamma^{LAB} + M_p) + 2p_K^{LAB} E_\gamma^{LAB} \cos \theta_K^{LAB}, \quad (4.6)$$

Figure 4.4 β versus momentum for pion candidates. The curves indicate the boundaries of the pion identification cut.

Figure 4.5 Missing mass off the kaon candidates.

where p_K^{LAB} is the magnitude of the kaon's three-momentum and θ_K^{LAB} is the angle between the kaon's direction and the beam direction in the lab frame, and

$$s = M_p^2 + 2M_p E_\gamma^{LAB}. \quad (4.7)$$

An uncorrected missing mass spectrum is shown in figure 4.5. Besides the prominent $\Lambda(1116)$ and $\Sigma^0(1193)$ peaks, the $\Lambda(1520)$ and the $\Lambda(1405)/\Sigma^0(1385)$ excited states can be seen clearly above the background. The primary source of the background is contamination from pions mis-identified as kaons.

I will be focusing on the $\Sigma(1385)/\Lambda(1405)$ region. Since the masses of the three charged stages of the ground-state Σ are slightly different, the thresholds for the channels to open up are different:

$$\Sigma^+\pi^- : M_{\Sigma^+} + M_{\pi^-} = 1328.94 \text{ MeV};$$

$$\Sigma^0\pi^0 : M_{\Sigma^0} + M_{\pi^0} = 1327.53 \text{ MeV};$$

$$\Sigma^-\pi^+ : M_{\Sigma^-} + M_{\pi^+} = 1337.01 \text{ MeV}.$$

The $\bar{K}N$ channel opens up at $M_K + M_p = 1.432 \text{ GeV}$. The $\Sigma^-\pi^+$ and K^-p thresholds define a natural range of 1.34-1.43 GeV for identifying the $\Lambda(1405)$ and $\Sigma^0(1385)$.

4.1 Empirical corrections

4.1.1 Correcting for energy loss

Charged particles lose energy as they pass through the target material, the target wall, the carbon epoxy beam pipe, the start counter, and to lesser extent the drift

chambers. Consequently, the momentum as measured by the drift chambers is less than the initial momentum just after the reaction occurs. In order to properly correct for this energy loss, the interaction point must be determined.

Determination of the primary vertex

Because of their relatively long lifetimes, the Λ , Σ^+ , and Σ^- particles resulting from the decays of the excited state hyperons decay some distance from the primary vertex determined by the kaon and the excited-state hyperon. In the case of $Y^* \rightarrow \Sigma^+\pi^-$, the K^+ and the π^- determine the primary vertex. In the case of $Y^* \rightarrow \Sigma^-\pi^+$, the K^+ and the π^+ determine the primary vertex. The determination of the primary vertex for neutral hyperon decays involving the Λ is less obvious for several reasons. The only charged particle coming from the primary vertex is the K^+ . The lifetime of the Λ is the longest of the three ground state hyperons discussed ($c\tau = 7.89$ cm). A significant number of these particles decay outside the target. Consequently, using the charged particles from the decay of the Λ may give skewed results. Furthermore, the beam cannot be used as a constraint for the vertex position because the size of the photon beam spot at the target is about 1 cm for 2.5 GeV primary electron beam. The length of the target (18.5 cm) does not allow the z-position to be constrained very well. Finally, the current vertex reconstruction code is limited to using charged tracks. The solution to this quandary comes from the observation that the proton takes most of the momentum in the decay $\Lambda \rightarrow p\pi^-$. For this case, the primary vertex

is approximately determined from the intersection between the proton and the kaon tracks.

Since the directions in the lab of the proton and the Λ are not necessarily the same, using the K^+ and the proton to determine the primary vertex could introduce systematic errors. In order to quantify this error, $K^+\Lambda X$ events were generated with a Gaussian distribution in x and y ($\sigma=0.6$ cm) and a flat distribution in z (from -9 cm to +9 cm) for the primary vertex. The reconstructed vertices were then compared to the generated vertices. Figure 4.6 shows the results. Although the differences in x and y center on zero, the difference in z between the reconstructed vertices and the generated vertices is offset about 0.25 cm on average. The reconstructed vertices are slightly systematically downstream of the true vertices.

The primary vertex distributions for the real data are shown in figure 4.7. The x and y distributions are approximately Gaussian but with long non-Gaussian tails probably arising from poorly reconstructed tracks. The tails of the z-distribution are similarly long. In the light of these two observations I applied a rough cylindrical cut on the primary vertex from -10.5 cm to +9.5 cm in z and 3.0 cm radially.

Once the primary vertex is determined, all charged particles coming from the primary interaction point are corrected for energy losses in the target material, the target walls, the carbon epoxy beam pipe, and the start counter. The energy losses in the drift chamber material and the air are ignored.

Figure 4.6 Comparison between generated primary vertices and reconstructed vertices using the proton and kaon tracks.

Figure 4.7 K^+p vertex distributions. A) Beam profile in x and y; B) Projection on the y-axis; C) projection onto x-axis; D) z-projection of the K^+p vertices showing the extent of the target. The dotted lines in plots B) and C) indicate the position of the target walls. The σ 's of the x and y distributions were 0.74 cm and 0.76 cm, respectively.

For the case where $\Lambda \rightarrow p\pi^-$, the secondary vertex is well-defined by the intersection of the proton and the π^- tracks. The momenta of these particles are also corrected for energy losses.

4.1.2 Correcting for angular dependencies

Since M_Y is an invariant quantity, it should not depend on the kaon lab angles: however, figures 4.8 and 4.10A show dependence of M_Λ and M_{Σ^0} on θ_K^{LAB} and ϕ as measured by the drift chambers. For small lab angles ($< 25^\circ$), the measured hyperon mass decreases rapidly with angle. To correct for this deviation I adjust the magnitude of the momentum according to the formula:

$$\frac{p'}{p} = 1 + \frac{A}{1 - \cos \theta} + \frac{B}{(1 - \cos \theta)^2}, \quad (4.8)$$

where $A \approx -2.0 \times 10^{-5}$ and $B \approx -4.7 \times 10^{-6}$. The correction differs slightly from sector to sector. Each sector needs to be adjusted with respect to one another: some of the widths of missing mass peaks come from poor intercalibration between the sectors. I correct the ϕ dependence with a second-order polynomial form. This dependence also varies from sector to sector. The same corrections were applied to the momenta of all positive particles. No correction was applied to the momenta of negative particles. These are purely empirical corrections intended to account for some systematic errors in the momentum reconstruction that most likely arise from inadequate understanding of the magnetic field. These problems are still under investigation by the tracking and drift chamber experts.

Figure 4.8 Ground-state hyperon mass dependence on $\cos \theta_{lab}^K$ before corrections are applied.

Figure 4.9 Ground-state hyperon mass dependence on $\cos\theta_{lab}^K$ after corrections are applied.

Figure 4.10 Ground-state hyperon mass dependence on ϕ^K (A) before the ϕ dependent corrections are applied and (B) after the corrections are applied.

Figures 4.9 and 4.10B show the dependencies after the corrections have been applied. Since the $\Lambda(1116)$ and the $\Sigma^0(1193)$ are so narrow (the lifetimes are $(2.632 \pm 0.020) \times 10^{-10} s$ and $(7.4 \pm 0.7) \times 10^{-20} s$, respectively), the widths of the peaks in figure 4.5 are purely instrumental. The standard deviations are 7.31 ± 0.04 MeV and 6.24 ± 0.03 MeV, respectively. Before corrections, they were 8.67 ± 0.05 MeV and 7.20 ± 0.04 MeV, respectively.

4.1.3 Correction of the photon beam energy

Even after correcting for energy loss and distortions of the momenta as a function of angle, the Λ and Σ^0 peaks are lower than the nominal values for the masses by about 3.8 MeV and 2.8 MeV, respectively. I correct the mass peaks by multiplying the photon energy reported by the tagger result bank by 1.006. This is within the errors of the primary beam and tagger magnet calibrations.

4.2 Fiducial cuts

Not all physical regions of the CLAS volume are active due to the shadow of the six superconducting coils and the holes in the upstream and downstream directions relative to the target. Poorly reconstructed tracks can point toward these dead regions. Furthermore, at the edges of the active volume of CLAS the acceptance is poor and can change rapidly. In order to minimize these problems, I analyzed only those events whose tracks fell within the fiducial volume described below and shown

in figure 4.11. Tracks pointing to the coils were eliminated, as were very forward-pointing tracks. The shadow of the coils increases in width as particle's trajectory points closer to the forward hole. The cuts were determined empirically. Given the lab quantities p , θ , and ϕ , and the charge q , the cut values are given by

$$\left. \begin{aligned} \theta_{min}^+ &= \theta_0^+ + \theta_s^+ \left(1 - \frac{pI_{max}}{5I}\right)^{24} \\ \Delta\phi^+ &= \phi_0^+ (\cos(\theta - \theta_{cut}))^{\alpha^+} \\ \alpha^+ &= \left(\frac{pI_{max}}{5I}\right)^{1/8} \end{aligned} \right\} \text{ for } q > 0: \quad (4.9)$$

$$\left. \begin{aligned} \theta_{min}^- &= \theta_0^- + \theta_s^- \frac{I}{0.05I + pI_{max}} \\ \Delta\phi^- &= \phi_0^- (\sin(\theta - \theta_{min}))^{\alpha^-} \\ \alpha^- &= 0.1 \left(\frac{pI_{max}}{I}\right)^{1/2} \end{aligned} \right\} \text{ for } q < 0. \quad (4.10)$$

with the general conditions that $\Delta\phi < 24^\circ$ and $\theta > 10^\circ$ and where I is the magnitude of the current (1920 A) in the superconducting coils used during the G1A experiment and I_{max} is the maximum current (3860 A) allowed. The constants θ_0^+ , θ_s^+ , θ_{cut} , θ_0^- and θ_s^- are 10° , 25° , 75° , 10° , and 15° , respectively.

The minimum photon energy required in the software was 1.275 GeV. The minimum momentum for the kaon candidates required in the software was 0.35 GeV/c. The maximum kaon momentum in software was 1.7 GeV/c.

4.3 Minimizing K^+ misidentification

For events that contain identified proton and π^- tracks, the most significant sources of background due to misidentified kaons are the production and subsequent

Figure 4.11 Fiducial volume.

decays of $\omega(782)$ and $\rho(770)$ in the reactions

$$\gamma p \rightarrow \begin{cases} p\rho: & \rho \rightarrow \pi^+\pi^- \\ p\omega: & \omega \rightarrow \pi^+\pi^-\pi^0 \end{cases} \quad (4.11)$$

In these cases the mis-identified kaons were actually π^+ 's. Assuming that the kaon candidates were pions, I computed the missing mass from the reaction $\gamma p \rightarrow p\pi^+\pi^-X$. The result is shown in figure 4.12A. A large spike at zero missing mass squared can be seen next to a smaller peak at the mass squared of the π^0 . The large hump of counts above the π^0 peak is due to the true kaons.

Figures 4.12C and 4.12B show the narrow ω peak and the broader ρ peak selected by cutting on the zero missing mass squared peak and the π^0 peak, respectively. To remove this contamination I cut above the π^0 peak in figure 4.12A at 0.0275 GeV^2 .

4.4 Identifying the $\Lambda(1116)$

The dominant decay mode $\Lambda \rightarrow p\pi^-$ occurs 63.9% of the time; $\Lambda \rightarrow n\pi^0$ occurs 35.8% of the time. With this in mind we look in the analysis for exactly three charged particles (proton, kaon, and pi-minus) in the final state. Once all the corrections to the momenta are applied, the four-momentum of the $\Lambda(1116)$ is reconstructed from the four-momenta of the proton and pi minus with the condition that the invariant mass of the $p\pi^-$ be consistent with a Λ . Figure 4.13 shows the invariant mass of the proton-pi minus pair for events in the excited-state region of hyperon mass (1.25-1.8 GeV). It peaks at the $\Lambda(1116)$ mass with a long tail to larger invariant masses. To

(A) Missing mass squared from $\gamma p \rightarrow p \pi^+ \pi^- X$ (GeV^2)

(B) $\pi^+ \pi^-$ mass (GeV)

(C) $\pi^+ \pi^- \pi^0$ mass (GeV)

Figure 4.12 K^+ misidentification: A) Missing mass squared from $\gamma p \rightarrow p \pi^+ \pi^- X$ with the assumption that the kaon candidates are π^+ 's. The shaded region is dominated by true kaonic events. B) $\pi^+ \pi^-$ invariant mass showing ρ peak. C) $\pi^+ \pi^- \pi^0$ invariant mass showing ω peak.

Figure 4.13 Invariant mass of proton-pi minus pairs in events that contain kaon candidates and excited-state hyperon masses between 1.25-1.8 GeV. The dashed lines show the limits of the cut used to identify the Λ .

identify the ground-state Λ I put a cut between 1.1118 GeV and 1.1218 GeV about the peak. Figure 4.14 shows the exponential fall-off for the traversal of the Λ from the primary to the secondary vertex in units of $ct = x/\gamma\beta$. An exponential fit to this distribution gives $c\tau = 7.8 \pm 0.3$ cm for the decay constant, in good agreement with the Particle Data Group's value of 7.89 cm.

Figure 4.15 shows the missing mass squared from the reaction $Y^* \rightarrow \Lambda X$ for hyperon masses in the range 1.25-1.8 GeV. Cutting on the π^0 peak in the range 0.007-0.0275 GeV^2 in the missing mass squared plot for hyperon masses in the range 1.25-1.8 GeV (figure 4.15) reveals a clean $\Sigma^0(1385)$ peak in the missing mass (figure 4.16A); similarly, the missing mass off the kaon shows the shape of the $\Lambda(1405)$ when the missing mass squared is cut above the π^0 peak (figure 4.16B). The $\Lambda(1520)$ also shows up, as expected, because of the large branch (42%) to $\Sigma\pi$. Finally, figure 4.16C shows the missing mass for the counts in the photon peak between -0.007 GeV^2 and 0.0035 GeV^2 .

Figure 4.14 Lifetime of the $\Lambda(1116)$ for excited-state hyperons in the 1.25-1.8 GeV mass region. The x-axis is the proper time in units of cm. The curve is an exponential fit to the data.

Figure 4.15 Missing mass squared from $Y^* \rightarrow \Lambda X$ for hyperon masses in the range 1.25-1.8 GeV.

(A) $K^+\Lambda(\pi^0)$ events

(B) $K^+\Lambda(\gamma\pi^0)$ events

(C) $K^+\Lambda(\gamma)$ events

Figure 4.16 Excited-state hyperon mass as a function of topology. A) $K^+\Lambda(\pi^0)$ events. B) $K^+\Lambda(\gamma\pi^0)$ events. C) $K^+\Lambda(\gamma)$ events. In each case the missing particles are in parentheses.

Chapter 5

Physics Analysis

In this chapter I describe the steps towards obtaining a radiative branching ratio for excited-state hyperons in the $\Sigma(1385)/\Lambda(1405)$ region relative to the reasonably well-known $\Lambda\pi^0$ decay channel of the $\Sigma^0(1385)$. The task is to assess the contributions to the $M_X^2(K^+\Lambda X)$ distribution (figure 6.1), the missing mass squared for the reaction $\gamma p \rightarrow K^+\Lambda X$.

5.1 Background reactions

Not all of the counts in the π^0 peak in the $M_X^2(K^+\Lambda X)$ distribution (figure 6.1) come from the decay of the $\Sigma^0(1385)$. Two sources of “physics” backgrounds are K^* production and $\Sigma^+ \rightarrow p\pi^0$ decays.

5.1.1 $\Sigma^+\pi^-$ Events

Both the $\Sigma^0(1385)$ and the $\Lambda(1405)$ can decay to $\Sigma^+\pi^-$. The rates are dominated by the decay of the $\Lambda(1405)$ due to the larger branch (nominally 33% compared to 6%). The dominant decay branches for the $\Sigma^+(1189)$ are $n\pi^+$ (48.3%) and $p\pi^0$ (51.5%). The invariant mass of the π^- from the decay of the $\Lambda(1405)/\Sigma^0(1385)$ to $\Sigma^+\pi^-$ and the proton from the subsequent decay of the Σ^+ to $p\pi^0$ forms a wide bump

whose tail falls underneath the Λ peak in figure 4.13 and as shown in figure 5.1A. The missing mass off the π^- for Y^* events in the hyperon mass range 1.34-1.43 GeV shows a prominent peak at the $\Sigma^+(1189)$ mass (figure 5.1B). The shaded histogram shows the contribution to the peak due to the proton-pi minus pairs that satisfy the Λ cut.

The $\Sigma^0(1385)$ and $\Lambda(1405)$ can also decay to $\Sigma^-\pi^+$; however, the $\Sigma^-(1197)$ decays almost exclusively to $n\pi^-$ and so does not contribute to the analysis because a proton will always be required.

5.1.2 K^* production

The $J=1$ analogue of the kaon, the $K^*(892)$, is produced in the reactions

$$\gamma p \longrightarrow \begin{cases} K^{*+}\Lambda(1116); \\ K^{*0}\Sigma^+(1189). \end{cases} \quad (5.1)$$

The K^* subsequently decays to $K\pi \sim 100\%$ of the time. Consequently, not all the counts in the π^0 peak in figure 6.1 come from the decay $\Sigma^0(1385) \rightarrow \Lambda\pi^0$. The $\Lambda\pi^0/K^+\pi^0$ two-dimensional histogram (figure 5.2A) shows that the $\Sigma^0(1385)$ dominates the statistics, but the projection onto the $K^+\pi^0$ axis (figure 5.2B) shows a clear K^* peak.

Similarly, not all the Σ^+ counts in figure 5.1 come from the decay of excited-state hyperons. Figure 5.3 plots the $K^+\pi^-$ mass versus the hyperon mass for a wider range in hyperon mass. For the case where Σ^+ candidates are required (figure 5.3A), the K^{*0} clearly shows up as a horizontal band about $K^+\pi^-$ mass of 0.89 GeV. The cut for

(A) Missing mass from $Y^* \rightarrow \pi^- X$ vs. $p\pi^-$ mass

(B) Missing mass from $Y^* \rightarrow \pi^- X$

Figure 5.1 Overlap of $\Lambda(1116)$ and $\Sigma^+(1189)$. (A) Two dimensional histogram of the missing mass for the reaction $Y^* \rightarrow \pi^- (X)$ vs. the $p\pi^-$ invariant mass. (B) The mass of X from the reaction $Y^* \rightarrow \pi^- X$. The shaded histogram shows the contribution to the peak due to the $p\pi^-$ pairs that satisfy the Λ cut.

(A) $K^+\pi^0$ mass vs. $\Lambda\pi^0$ mass

(B) Projection onto $K^+\pi^0$ axis

Figure 5.2 K^* contamination. (A) Two-dimensional histogram of $K^+\pi^0$ mass versus $\Lambda\pi^0$ mass. The K^* shows up weakly as a horizontal band; the $\Sigma^0(1385)$ shows up as a vertical band. (B) projection onto the $K^+\pi^0$ mass axis. Here the $K^*(892)$ shows up clearly as a peak near 0.9 GeV.

identifying Σ^+ candidates was $1.16 < M_X(K^+\pi^-X) < 1.22$ GeV. For the case where Λ is identified in $p\pi^-$ mass for those events satisfying the Σ^+ cut (figure 5.3B), no evidence for a K^{*0} band exists. The plot of the $K^+\pi^-$ mass (figure 5.3C) for events satisfying the Λ cut shows no peak at the mass of the K^{*0} .

5.2 Acceptance calculations

The CLAS detector, although its acceptance is quite large, has several large gaps due to the six coils and the hole in the detector coverage in the downstream direction. The acceptance is further compounded by the decays in flight of the pions and especially the kaons. Therefore in order to measure the branching ratios the acceptance of the detector for different reaction types needs to be understood in some detail.

I studied the following reactions:

$$\begin{aligned}
 \Sigma^0(1385) &\rightarrow \begin{cases} \Lambda\gamma \\ \Lambda\pi^0 \\ \Sigma^+\pi^- \end{cases} \\
 \Lambda(1405) &\rightarrow \begin{cases} \Lambda\gamma \\ \Sigma^0\pi^0 \\ \Sigma^+\pi^- \end{cases} \\
 \gamma p &\rightarrow \Lambda K^{*+} \\
 &\quad \searrow \begin{cases} K^+\pi^0 \\ K^+\gamma \end{cases}
 \end{aligned}$$

Figure 5.3 Two-dimensional histogram of $K^+\pi^-$ mass vs Y^* mass for the case (A) where a Σ^+ is required and (B) where a Λ is required in addition to satisfying the Σ^+ cut; (C) is the $K^+\pi^-$ mass for the events satisfying the Λ and Σ^+ cuts.

5.2.1 Event Generation

Because of the raggedness of the true energy spectrum, a histogram (figure 2.6) of the spectrum from the data was used as the source for the Monte Carlo. Much of the spiky structure of the spectrum comes from variation in the size of the energy bite of the E-counter bins. Some of the E-counters were not functional during the G1 run, causing the counts to drop to zero in the spectrum.

The $\Sigma^0(1385)$ was generated with a relativistic Breit-Wigner shape [38]

$$\frac{d\sigma}{dM} \propto \frac{2MM_0\Gamma(M)}{(M^2 - M_0^2)^2 + M_0^2\Gamma^2(M)},$$

$$\Gamma(M) = \Gamma_0 \left(\frac{q}{q_0} \right)^{2L+1}, \quad (5.2)$$

where q (q_0) is the momentum of either of the decay particles in the rest frame of an excited-state hyperon with mass M (M_0), $M_0 = 1.3837$ GeV, $\Gamma_0 = 0.036$ GeV, and $L=0$. The choice of $L=0$ fit the data better than $L=1$. The rejection method [39] was used. A random number between the minimum (M_{min}) and maximum (M_{max}) mass values is generated and the Breit-Wigner distribution formula is evaluated at this value. Denote the result $f(M')$, with $M_{min} \leq M' \leq M_{max}$. A second random number between zero and the maximum value of the Breit-Wigner function (at the peak mass of the resonance) is generated and compared to $f(M')$. If $f(M')$ is greater than this second random number, then M' is returned; otherwise, the process is repeated.

The asymmetric shape of the $\Lambda(1405)$ (see, for example, figure 1.4) precludes using a Breit-Wigner form. A better fit to the shape is obtained by using the K-matrix

approach of Dalitz and Tuan[40]. For isospin zero reactions.

$$K = \begin{bmatrix} \alpha & \beta \\ \beta & \gamma \end{bmatrix}; \quad (5.3)$$

$$\alpha = \langle \bar{K}N | \hat{K} | \bar{K}N \rangle; \quad (5.4)$$

$$\beta = \langle \bar{K}N | \hat{K} | \Sigma\pi \rangle = \langle \Sigma\pi | \hat{K} | \bar{K}N \rangle \quad (5.5)$$

$$\gamma = \langle \Sigma\pi | \hat{K} | \Sigma\pi \rangle. \quad (5.6)$$

The K-matrix is related to the S-wave transition matrix T by

$$T^{-1} = K^{-1} - iQ, \quad (5.7)$$

where Q is a diagonal matrix containing the center-of-momentum momenta q and k for the $\Sigma\pi$ and $\bar{K}N$ systems, respectively. The $\Sigma\pi - \Sigma\pi$ matrix element T_{22} is give by

$$T_{22} = \frac{A}{1 - iqA}, \quad (5.8)$$

$$A = \gamma + \frac{ik\beta^2}{1 - ik\alpha} = \gamma - \frac{|k|\beta^2}{1 + |k|\alpha} \quad (5.9)$$

below the $\bar{K}N$ threshold where k is purely imaginary. The shape of the $\Sigma\pi$ mass spectrum is given by

$$\Sigma\pi \text{ mass distribution} \propto \frac{q}{f^2(|k|) + q^2}, \quad (5.10)$$

$$f(|k|) = \frac{1 + |k|\alpha}{\gamma + |k|(\gamma\alpha - \beta^2)} \quad (5.11)$$

below the $\bar{K}N$ threshold. Figure 5.4 shows the mass distributions used for the event generator.

Figure 5.4 Generated $\Sigma\pi$ mass distributions for the $\Lambda(1405)$.

Once the hyperon mass has been determined, the kaon is generated in the center-of-mass frame uniformly in ϕ but with a t -dependent distribution in θ_{cm} according to

$$P(t) \propto e^{2.0t}. \quad (5.12)$$

The decay products of the excited state hyperon are produced isotropically in the rest frame of the hyperon and boosted back to the lab frame.

A similar scheme is used for generating K^* events. The shape of the K^* mass distribution was determined by a Breit-Wigner form.

Gaussian distributions in x and y , both with sigmas of 0.6 cm, approximated the beam spot at the target. Events were generated uniformly along the length of the target from -9.0 cm to 9.0 cm along the beam axis.

5.2.2 Simulation of CLAS behavior

The generated events were fed to GSIM, the CLAS Geant-based Monte Carlo program. GSIM traces each particle through a model of the detector and creates banks of simulated hits on the various detector elements. The hits are written out in the same form (BOS banks) as the real data. The simulation contains decays-in-flight and models the effects of multiple scattering but does not take into account inefficiencies in the detector components.

5.2.3 GSIM post-processing and reconstruction

In the next step, the program gpp ("GSIM post-processor") knocks out bad drift chamber wires and scintillator paddles. It also smears the time-of-flight times and drift chamber hit positions by scaling default distributions. The factor for scaling the time-of-flight was 1.75. The factor for scaling the drift-chamber resolution was 2.0. The output of gpp is then processed by a1c and finally histograms are created by the analysis program.

5.2.4 Simulation results

As an initial guess, both the $\Sigma^0(1385)$ and the $\Lambda(1405)$ events were generated with cross sections that were assumed to be flat in E_γ . As shown in figure 5.5, this may be a poor approximation for the $\Lambda\pi^0$ channel, which is dominated by the $\Sigma^0(1385)$ (see figure 5.2A). The $\Lambda\pi^0$ channel was selected by requiring $0.007 < M_X^2(K^+\Lambda X) < 0.027 \text{ GeV}^2$. The cross section shows an enhancement peaking at $E_\gamma \approx 1.75 \text{ GeV}$. Assuming resonant production, these enhancements would correspond to a resonance with mass of $\sim 2 \text{ GeV}$. The acceptance was determined by the ratio of accepted events to thrown events in each bin of E_γ . The uncorrected yield was determined by the ratio of the raw events to the number of incident photons in each bin of E_γ . The density of the target liquid hydrogen was assumed to be constant. No background subtraction was performed.

Figure 5.5 $\Lambda\pi^0$ cross section. (A) $\Lambda\pi^0$ counts as a function of E_γ . (B) $\Lambda\pi^0$ distribution after Bremsstrahlung spectrum has been divided out. (C) Acceptance for $\Lambda\pi^0$ events as a function of E_γ . (D) Cross section for $\Lambda\pi^0$ events.

Figure 5.6 shows the cross section for the $\Sigma^0\pi^0$ channel, identified by requiring $M_X^2(K^+\Lambda X) > 0.027 \text{ GeV}^2$. The broad peak at $E_\gamma \approx 1.75 \text{ GeV}$ seen in the $\Lambda\pi^0$ channel is absent. The cross section is roughly flat.

The acceptance calculation for the $\Lambda\pi^0$ channel was iterated with an energy-dependence derived from the shape of the cross section in figure 5.5.

Figure 5.7 demonstrates that the exponential t -dependence in equation 5.12 provides a good match to the data for the $\Lambda\pi^0$ channel. In order to minimize the background due to K^{*+} production, the $K^+\pi^0$ invariant mass was required to be less than 0.85 GeV (approximately one width below the K^{*+} peak).

Figure 5.8 shows that the acceptance for the $\Lambda\pi^0$ channel is fairly flat over the angular range. For the $\Sigma^+\pi^-$ case, the acceptance distorts the angular distribution (figure 5.9): the acceptance is higher for $\cos\beta^* \rightarrow \pm 1$ than for $\cos\beta \sim 0$. Here $\cos\beta^*$ is the dot product between the normal to the production plane and the direction of the momentum of the ground state (decay) hyperon in the rest frame of the excited-state hyperon.

Figures 5.10 and 5.11 show the data corrected for the acceptance for the $\Lambda\pi^0$ and $\Sigma^+\pi^-$ channels. They demonstrate that the assumption of flat angular distributions for the decay products is a good approximation. The χ^2/ndf of a zero-slope linear fit for the $\Lambda\pi^0$ case was 1.75; the χ^2/ndf for the $\Sigma^+\pi^-$ case was 1.24.

Figure 5.12 compares the results of the $\Sigma^0(1385) \rightarrow \Lambda\pi^0$ and $\Lambda(1405) \rightarrow \Sigma^0\pi^0$ simulations. The $\Lambda(1405) \rightarrow \Sigma^0\pi^0$ counts show up above the π^0 peak because the

Figure 5.6 $\Sigma^0\pi^0$ cross section. (A) $\Sigma^0\pi^0$ counts as a function of E_γ . (B) $\Sigma^0\pi^0$ distribution after Bremsstrahlung spectrum has been divided out. (C) Acceptance for $\Sigma^0\pi^0$ events as a function of E_γ . (D) Cross section for $\Sigma^0\pi^0$ events.

(A) Acceptance in t

(B) Comparison of simulation with data

Figure 5.7 Dependence on t for the $\Lambda\pi^0$ channel. (A) Acceptance in t . (B) Comparison of acceptance corrected data (circles) with generated t -distribution (dashed histogram) using equation 5.12.

(A) Thrown angular distribution

(B) Accepted angular distribution

(C) Acceptance

Figure 5.8 $\Lambda\pi^0$ angular acceptance: A) Generated angular distribution. B) Accepted angular distribution C) Acceptance. The angle β^* is defined in the text.

Figure 5.9 $\Sigma^+\pi^-$ angular acceptance: A) Generated angular distribution. B) Accepted angular distribution. C) Acceptance.

(A) Raw angular distribution

(B) Acceptance-corrected angular distribution

Figure 5.10 $\Lambda\pi^0$ angular distribution: A) Raw data. (B) Acceptance-corrected data.

(A) Raw angular distribution

(B) Acceptance-corrected angular distribution

Figure 5.11 $\Sigma^+\pi^-$ angular distribution: A) Raw data. (B) Acceptance-corrected data.

Σ^0 decays to $\Lambda\gamma$, leaving two particles missing ($\gamma\pi^0$). This accounts for the structure above the π^0 peak in figure 6.1.

Figure 5.13 shows the $p\pi^-$ invariant mass for the $\Lambda(1405) \rightarrow \Sigma^+\pi^-$ case. The spectrum clearly overlaps with the $\Lambda(1116)$; the shape agrees with the hump next to the Λ peak in figure 4.13. The “overall” acceptance for these events is $(0.20 \pm 0.02) \times 10^{-3}$. The “overall” acceptance is the ratio of the number of events that survive the particle identification cuts described in chapter 4 to the number of events thrown with hyperon masses between 1.34 GeV and 1.43 GeV.

Figure 5.14 shows the effect of the acceptance and the cuts on the hyperon mass for $\Sigma^0(1385) \rightarrow \Lambda\pi^0$. In the region between 1.34-1.43 GeV, the acceptance is roughly flat; a straight-line fit to this region gives $(1.644 \pm 0.099) \times 10^{-3}$ ($\chi^2/ndf = 1.4$). Figures 5.15 and 5.16 show the acceptance for the $\Sigma^0(1385) \rightarrow \Lambda\gamma$ and the $\Lambda(1405) \rightarrow \Sigma^0\pi^0$ cases, respectively: the straight line fit results are $(1.69 \pm 0.12) \times 10^{-3}$ ($\chi^2/ndf = 1.1$) and $(1.481 \pm 0.077) \times 10^{-3}$ ($\chi^2/ndf = 0.81$), respectively.

Table 5.1 compares the “overall” acceptances for the $\Sigma^0(1385) \rightarrow \Lambda\pi^0$, $\Sigma^0(1385) \rightarrow \Lambda\gamma$, and $\Sigma^0(1385) \rightarrow \Sigma^+\pi^-$ channels for two different “models” for the $\Sigma(1385)$ production cross section: (1) resonant production peaking at $E_\gamma \approx 1.75$ GeV; (2) flat dependence on E_γ . One concern regarding this method is that events from the tails may migrate into the cut region or vice versa during simulation and reconstruction. Figures 5.14 and 5.15 suggests that if bin migration is occurring, roughly equal numbers of events leak out of the cut as leak in. This only becomes a problem if the

Figure 5.12 Comparison of simulation results for the channels $\Sigma^0(1385) \rightarrow \Lambda \pi^0$ (plain region) and $\Lambda(1405) \rightarrow \Sigma^0 \pi^0$ (hashed region). The missing mass squared for the reaction $Y^* \rightarrow \Lambda(X)$ is plotted.

(A) $p\pi^-$ invariant mass for $\Sigma^+\pi^-$ events

(B) Mass squared of missing particle

Figure 5.13 Results of $\Lambda(1405) \rightarrow \Sigma^+\pi^-$ simulation: A) $p\pi^-$ invariant mass. The shaded region shows the counts satisfying the Λ identification cut. B) Missing mass squared from the reaction $Y^* \rightarrow \Lambda X$.

Figure 5.14 Acceptance for $\Sigma^0(1385) \rightarrow \pi^0\Lambda$ events as a function of the $\Sigma^0(1385)$ mass. (A) Generated $\Sigma^0(1385)$ mass distribution. (B) Accepted mass distribution. (C) Acceptance.

(A) Generated $\Lambda\gamma$ mass distribution

(B) Accepted mass distribution

(C) Ratio of (B) to (A)

Figure 5.15 Acceptance for $\Sigma^0(1385) \rightarrow \gamma\Lambda$ events as a function of the $\Sigma^0(1385)$ mass. (A) Generated $\Sigma^0(1385)$ mass distribution. (B) Accepted mass distribution. (C) Acceptance.

Figure 5.16 Acceptance for $\Lambda(1405) \rightarrow \Sigma^0 \pi^0$ events as a function of the $\Sigma^0 \pi^0$ mass. (A) Generated $\Sigma^0 \pi^0$ mass distribution. (B) Accepted mass distribution. (C) Acceptance.

real events behave differently from the simulated events, so a certain amount of faith in the simulation is required. Another concern regarding this method is that it introduces event generator dependencies, as illustrated in table 5.1. The difference in the overall acceptance between iterations is a measure of the systematic error due to choice of model.

Model	$A(\Sigma^* \rightarrow \Lambda\pi^0)$	$A(\Sigma^* \rightarrow \gamma\Lambda)$	$A(\Sigma^* \rightarrow \Sigma^+\pi^-)$	$\frac{A(\Sigma^* \rightarrow \Lambda\pi^0)}{A(\Sigma^* \rightarrow \gamma\Lambda)}$	$\frac{A(\Sigma^* \rightarrow \Sigma^+\pi^-)}{A(\Sigma^* \rightarrow \Lambda\pi^0)}$
1	1.66 ± 0.08	1.78 ± 0.12	0.39 ± 0.07	0.93 ± 0.08	0.24 ± 0.04
2	1.98 ± 0.12	2.26 ± 0.12	0.40 ± 0.06	0.88 ± 0.07	0.20 ± 0.03

Table 5.1 Dependence of acceptance on models. All acceptance values are in units of 1.0×10^{-3} . Models 1 and 2 are described in the text in section 5.2.4.

Chapter 6

Results and Conclusions

6.1 Extracting yields

Figure 6.1 shows the $M_X^2(K^+\Lambda X)$ distribution for hyperon masses in the range 1.34-1.43 GeV. There is a small excess of counts at zero missing mass squared sitting on top of the tail of the π^0 peak from the decay $\Sigma^0(1385) \rightarrow \Lambda\pi^0$ and a roughly flat background. The counts above the π^0 peak come from the decay $\Lambda(1405) \rightarrow \Sigma^0\pi^0$. This chapter describes the method for calculating the radiative branching ratio from the counts in this distribution. The quantities that can be extracted directly from the data are $n(\gamma)$, the number of counts in the photon peak in $M_X^2(K^+\Lambda X)$; $n(\pi^0)$, the number of counts in the π^0 peak in $M_X^2(K^+\Lambda X)$; $n(\Sigma^0\pi^0)$, the number of counts above the π^0 due to the decay of the $\Lambda(1405)$; and $n(K^{*+} \rightarrow K^+\pi^0)$, the number of K^{*+} events surviving the cuts. I will be using lower case n for observed quantities and upper case N for acceptance corrected or derived quantities. Unless otherwise stated, all the particle identification and kinematic cuts described in chapter 4 will be required.

Figure 6.1 Missing mass squared from $Y^* \rightarrow \Lambda X$ for hyperon masses in the range 1.34-1.43 GeV. The curve is the result of the fit described in section 6.1.

To extract the number of photons, π^0 's, and counts due to $\Lambda(1405) \rightarrow \Sigma^0 \pi^0$, a three Gaussian fit plus a flat background is applied to the histogram in figure 6.1. Due to the slight asymmetric tail, two of the Gaussian were used to fit the shape of the π^0 peak. The number of π^0 candidates is 589.7 ± 34.5 . The third Gaussian was used to fit the γ peak. The data was fit in the range -0.04 GeV^2 to $+0.0275 \text{ GeV}^2$ with no fixed parameters. Two other fits of figure 6.1 were performed to assess the level of systematic uncertainty in the extraction of the yields. The third Gaussian was eliminated from these fits. For Fit 2, both Gaussians were used to fit the π^0 . This gives a minimum value for $n(\gamma)$ because the fitting routine attempts to include the counts near zero missing mass squared as part of the tail of the π^0 . In this case $n(\gamma)$ is the excess of counts above the tail. For Fit 3, one Gaussian was used for the π^0 and one for the γ peak. This gives a maximum estimate for $n(\gamma)$ because it assumes that the π^0 tail is minimal in the region of the γ peak. Fits 2 and 3 differ in the choices of initial guesses for the parameters. In all three cases, $n(\Sigma^0 \pi^0)$ was estimated from the number of counts N in the range $0.03\text{-}0.09 \text{ GeV}^2$ after subtracting the π^0 tail:

$$n(\Sigma^0 \pi^0) = \alpha(N - 26b), \quad (6.1)$$

where b is the flat background term obtained from the fit, 26 is the number of bins in the specified range, and $\alpha = 1.114 \pm 0.084$ is a correction factor derived from simulation to account for the fraction of the $\Sigma^0 \pi^0$ distribution that would fall outside the specified range. The value of $n(\pi^0)$ was corrected for the $\Sigma^0 \pi^0$ tail according to

$$\Delta n(\pi^0) = \left(\frac{\alpha - 1}{\alpha} \right) n(\Sigma^0 \pi^0). \quad (6.2)$$

The results of the three fits are tabulated in table 6.1.

Extracting the number of K^* counts is a bit tricky. As shown in figure 6.2A, relatively few K^* candidates survive the 1.34-1.43 GeV hyperon mass cut. The solid curve is a fit to the data using a fixed mass of 0.892 GeV and a fixed width of 0.05 GeV for the K^* and a Gaussian to approximate the background; the dashed curve is the result of a fit using just the Gaussian. The former fit gives an estimate on the number of K^* counts of 23.9 ± 17.7 with a χ^2/ndf of 0.999. The Gaussian-alone fit returns a slightly lower χ^2/ndf (0.935). The two fits of figure 6.2A are nearly equivalent. In order to assess the merit of adding the Breit-Wigner for the K^* peak, I performed an F-test[39]. I define

$$F = \left(\frac{\chi^2(m) - \chi^2(m+3)}{\chi^2(m+3)} \right) \frac{N - m - 3}{3}, \quad (6.3)$$

where $\chi^2(m)$ corresponds to the dashed curve in figure 6.2A, $\chi^2(m+3)$ corresponds to the solid curve in figure 6.2A, $N - m - 3 = 28$ and $m = 3$. I obtain 0.567, to be compared to $F \approx 3$, the value of F given the two degrees of freedom $\nu_1 = 3$ and $\nu_2 = 28$ for which the probability of exceeding F is 0.05. This implies that adding the

	Fit 1	Fit 2	Fit 3
$n(\gamma)$	12.2	8.3	14.4
$n(\pi^0)$	579.7	588.7	567.6
$n(\Sigma^0\pi^0)$	97.7	101.2	93.4
b	1.365	1.243	1.511
χ^2/ndf	0.926	1.119	1.185

Table 6.1 Results of the three fits of figure 6.1 described in the text.

Breit-Wigner to the fit is of dubious merit. Consequently, the K^{*+} contamination was also estimated in a second way. The $K^+\pi^0$ mass distribution for the 1.34-1.43 GeV hyperon mass range was fit with a single Gaussian in the range 0.72-1.1 GeV. The peak position and sigma of the Gaussian was fixed and fits of the $K^+\pi^0$ mass distributions for the set of hyperon mass intervals 1.34-1.5 GeV, 1.34-1.55 GeV, 1.34-1.6 GeV, 1.34-1.65 GeV, 1.34-1.7 GeV, 1.34-1.75 GeV, and 1.34-1.8 GeV were performed. The Gaussian shape was subtracted from each distribution and the result was fit with a Breit-Wigner. The trend was fit with a fourth-order polynomial and extrapolated to the 1.34-1.43 GeV mass interval, as shown in figure 6.2B. The curve was required to pass through the (0,0) point. The 1.34-1.43 GeV mass interval was not included in the fit of the trend. The number of K^{*+} counts obtained from the extrapolation technique was 19.9 ± 9.7 . There is a large systematic uncertainty in extracting $n(K^{*+} \rightarrow K^+\pi^0)$. Nevertheless, since the hyperon mass cut of 1.34-1.43 GeV fits within the kinematic range of the K^* and the evidence from figure 5.2 is clear that K^{*} 's are seen in the data, there is no reason to assume that the K^* contamination is zero; hence the result from the extrapolation will be used in the branching ratio calculation.

6.2 Calculating the branching ratio

Because of the large branch for the Σ^+ to decay to $p\pi^0$, the π^0 peak can contain contributions from both $\Sigma(1385)$ and $\Lambda(1405)$; furthermore, K^{*+} particles produced in the reaction $\gamma p \rightarrow K^{*+}\Lambda$ decay to $K^+\pi^0$ 1/3 of the time. Taking these contributions

Figure 6.2 Estimation of K^* contribution to the π^0 peak in $M_X^2(K^+\Lambda X)$. (A) Fit of the $K^+\pi^0$ mass distribution for events with Y^* mass in the range 1.34-1.43 GeV. (B) Plot of the K^* contribution for decreasing mass regions between 1.34-1.8 GeV and 1.34-1.5 GeV.

into account.

$$n(\pi^0) = N(\Sigma^* \rightarrow \Lambda \pi^0) + N(\Sigma^* \rightarrow \Sigma^+ \pi^-) + N(\Lambda^* \rightarrow \Sigma^+ \pi^-) + n(\Lambda K^{*+}) + N(\Sigma^+ K^{*0}), \quad (6.4)$$

where all terms are subject to the condition that there be a Λ in the event. The $\Sigma^+ K^{*0}$ contamination is negligible (see section 5.1.2), so $N(\Sigma^+ K^{*0})$ will be ignored.

Since the $\Sigma(1385)$ is isospin $I=1$, it cannot decay to $\Sigma^0 \pi^0$, whereas the $\Lambda(1405)$ does because it is $I=0$. This fact can be exploited to estimate the number of $\Lambda(1405)$ particles produced. The number $N(\Lambda(1405))$ is related to the number of counts $n(\Sigma^0 \pi^0)$ above the π^0 peak in 6.1 according to

$$N(\Lambda^*) = \frac{n(\Sigma^0 \pi^0)}{R(\Lambda^* \rightarrow \Sigma^0 \pi^0) A(\Lambda^* \rightarrow \Sigma^0 \pi^0)}, \quad (6.5)$$

where $R(\Lambda^* \rightarrow \Sigma^0 \pi^0)$ is the probability that the $\Lambda(1405)$ will decay to $\Sigma^0 \pi^0$ and $A(\Lambda^* \rightarrow \Sigma^0 \pi^0)$ is the probability that this decay channel will be seen after the various cuts have been applied. The Clebsch-Gordon tables indicate that each charged combination of $\Sigma \pi$ has equal probability: i.e.,

$$R(\Lambda^* \rightarrow \Sigma^0 \pi^0) = R(\Lambda^* \rightarrow \Sigma^+ \pi^-) = R(\Lambda^* \rightarrow \Sigma^- \pi^+) = \frac{1}{3} R(\Lambda^* \rightarrow (\Sigma \pi)^0) \approx \frac{1}{3}. \quad (6.6)$$

From this result an estimate of the number of counts in the π^0 peak in figure 6.1 coming from the reaction $\Lambda(1405) \rightarrow \Sigma^+ \pi^-$ can be calculated:

$$N(\Lambda^* \rightarrow \Sigma^+ \pi^-) = R(\Lambda^* \rightarrow \Sigma^+ \pi^-) A(\Lambda^* \rightarrow \Sigma^+ \pi^-) N(\Lambda^*) = \frac{A(\Lambda^* \rightarrow \Sigma^+ \pi^-)}{A(\Lambda^* \rightarrow \Sigma^0 \pi^0)} n(\Sigma^0 \pi^0) \quad (6.7)$$

We can estimate the number of $\Sigma(1385)$'s produced using the $\Lambda\pi^0$ channel:

$$N(\Sigma^*) = \frac{N(\Sigma^* \rightarrow \Lambda\pi^0)}{R(\Sigma^* \rightarrow \Lambda\pi^0) \mathcal{A}(\Sigma^* \rightarrow \Lambda\pi^0)}, \quad (6.8)$$

where $R(\Sigma^* \rightarrow \Lambda\pi^0)$ is the probability that the $\Sigma(1385)$ will decay to $\Lambda\pi^0$ and $\mathcal{A}(\Sigma^* \rightarrow \Lambda\pi^0)$ is the probability that this channel will be observed after the cuts. From this we can deduce how many $\Sigma(1385) \rightarrow \Sigma^+\pi^-$ counts would contribute to the π^0 peak:

$$\begin{aligned} N(\Sigma^* \rightarrow \Sigma^+\pi^-) &= R(\Sigma^* \rightarrow \Sigma^+\pi^-) \mathcal{A}(\Sigma^* \rightarrow \Sigma^+\pi^-) N(\Sigma^*) \\ &= \frac{R(\Sigma^* \rightarrow \Sigma^+\pi^-)}{R(\Sigma^* \rightarrow \Lambda\pi^0)} \frac{\mathcal{A}(\Sigma^* \rightarrow \Sigma^+\pi^-)}{\mathcal{A}(\Sigma^* \rightarrow \Lambda\pi^0)} N(\Sigma^* \rightarrow \Lambda\pi^0), \end{aligned} \quad (6.9)$$

where $R(\Sigma^* \rightarrow \Sigma^+\pi^-)$ is the probability that the $\Sigma(1385)$ will decay to $\Sigma^+\pi^-$ and $\mathcal{A}(\Sigma^* \rightarrow \Sigma^+\pi^-)$ is the probability that this channel will survive the cuts. To simplify the notation, we define

$$R \equiv R_{\Lambda\pi}^{\Sigma\pi} = \frac{R(\Sigma^* \rightarrow \Sigma^\pm\pi^\mp)}{R(\Sigma^* \rightarrow \Lambda\pi^0)} = 0.135 \pm 0.011, \quad (6.10)$$

using the average value quoted in [6]. The two charged combinations of $\Sigma\pi$ are equally probable.

Using equations 6.4 and 6.9, we solve for $N(\Sigma^* \rightarrow \Lambda\pi^0)$:

$$N(\Sigma^* \rightarrow \Lambda\pi^0) = \frac{n(\pi^0) - N(\Lambda^* \rightarrow \Sigma^+\pi^-) - N(\Lambda K^{*+})}{1 + \frac{R}{2} \frac{\mathcal{A}(\Sigma^* \rightarrow \Sigma^+\pi^-)}{\mathcal{A}(\Sigma^* \rightarrow \Lambda\pi^0)}}. \quad (6.11)$$

This is the contribution to the π^0 peak in figure 6.1 due to the decay $\Sigma^0(1385) \rightarrow \Lambda\pi^0$.

Both the $\Lambda(1405)$ and the $\Sigma(1385)$ can contribute to the photon peak in figure 6.1.

Furthermore, the K^{*+} decays radiatively 0.1% of the time. Taking these contributions

into account.

$$n(\gamma) = N(\Sigma^* \rightarrow \gamma\Lambda) + N(\Lambda^* \rightarrow \gamma\Lambda) + N(K^{*+} \rightarrow \gamma K^+). \quad (6.12)$$

We can use the previous experimental result for the decay $\Lambda(1405) \rightarrow \gamma\Lambda$ to estimate the $\Lambda(1405)$'s contribution:

$$\begin{aligned} N(\Lambda^* \rightarrow \gamma\Lambda) &= R(\Lambda^* \rightarrow \gamma\Lambda) A(\Lambda^* \rightarrow \gamma\Lambda) N(\Lambda^*) \\ &= 3R(\Lambda^* \rightarrow \gamma\Lambda) \frac{A(\Lambda^* \rightarrow \gamma\Lambda)}{A(\Lambda^* \rightarrow \Sigma^0 \pi^0)} n(\Sigma^0 \pi^0) \end{aligned} \quad (6.13)$$

where $R(\Lambda^* \rightarrow \gamma\Lambda) = 0.00054 \pm 0.00016$ based on a radiative width of 27 ± 8 keV[11] and a full width of 50 ± 2 MeV[6].

We can estimate how many K^{*+} radiative decays we should see by scaling the observed number of $K^{*+} \rightarrow K^+ \pi^0$ events. The estimated number of K^{*+} produced in the experiment is given by

$$N(K^{*+}) = \frac{n(K^{*+} \rightarrow K^+ \pi^0)}{R(K^{*+} \rightarrow K^+ \pi^0) A(K^{*+} \rightarrow K^+ \pi^0)}, \quad (6.14)$$

where $n(K^{*+} \rightarrow K^+ \pi^0)$ is the observed number of $K^{*+} \rightarrow K^+ \pi^0$ events, $R(K^{*+} \rightarrow K^+ \pi^0) = 1/3$ is the probability for the decay and $A(K^{*+} \rightarrow K^+ \pi^0)$ is the probability that this reaction will be observed and survive all the cuts. Then

$$\begin{aligned} N(K^{*+} \rightarrow K^+ \gamma) &= R(K^{*+} \rightarrow K^+ \gamma) A(K^{*+} \rightarrow K^+ \gamma) N(K^{*+}) \\ &= 3R(K^{*+} \rightarrow K^+ \gamma) \frac{A(K^{*+} \rightarrow K^+ \gamma)}{A(K^{*+} \rightarrow K^+ \pi^0)} n(K^{*+} \rightarrow K^+ \pi^0). \end{aligned} \quad (6.15)$$

Finally, the radiative branching ratio for the $\Sigma^0(1385)$ relative to the $\Lambda\pi^0$ channel is given by

$$\frac{\Gamma(\gamma\Lambda)}{\Gamma(\pi^0\Lambda)} = \left(\frac{n(\gamma) - N(\Lambda^* \rightarrow \gamma\Lambda) - N(K^{*-+} \rightarrow \gamma K^{*+})}{n(\pi^0) - N(\Lambda^* \rightarrow \Sigma^+\pi^-) - N(\Lambda K^{*-+})} \right) \times \left(\frac{A(\Sigma^* \rightarrow \Lambda\pi^0)}{A(\Sigma^* \rightarrow \gamma\Lambda)} \right) \left(1 + \frac{R}{2} \frac{A(\Sigma^* \rightarrow \Sigma^+\pi^-)}{A(\Sigma^* \rightarrow \Lambda\pi^0)} \right) \quad (6.16)$$

The various contributions to the number of π^0 and γ counts are tabulated in table 6.2. The values for the “overall” acceptances (see section 5.2.4) used in the calculations are enumerated in table 6.3. The result for the branching ratio is

$$R_{\Lambda\pi}^{\Lambda\gamma} \equiv \frac{\Gamma(\Sigma^0(1385) \rightarrow \gamma\Lambda)}{\Gamma(\Sigma^0(1385) \rightarrow \pi^0\Lambda)} = 0.0206 \pm 0.0084_{-0.0073}^{+0.0043}. \quad (6.17)$$

In order to compare this result with the model predictions, it is necessary to convert the branching ratio into a width in units of keV. The partial widths for all of the available channels should sum to $\Gamma_{TOT} = 36 \pm 5$ MeV:

$$\Gamma(\Lambda\pi^0) + \Gamma(\Sigma\pi) + \Gamma(\Lambda\gamma) + \Gamma(\Sigma^0\gamma) = \Gamma_{TOT}. \quad (6.18)$$

Rewriting equation 6.18 in terms of measured branching ratios and solving for $\Gamma(\Lambda\gamma)$, we obtain

$$\Gamma(\Lambda\gamma) = \frac{R_{\Lambda\pi}^{\Lambda\gamma} \Gamma_{TOT}}{1 + R_{\Lambda\pi}^{\Lambda\gamma} (1 + R_{\Lambda\gamma}^{\Sigma^0\gamma}) + R_{\Lambda\pi}^{\Sigma\pi}}. \quad (6.19)$$

Equation 6.19 includes the unmeasured branching ratio $R_{\Lambda\gamma}^{\Sigma^0\gamma}$. It is not possible to compute $\Gamma(\Lambda\gamma)$ without resorting to a model or assuming that $\Gamma(\Sigma^0\gamma)$ is negligible. Although the model predictions for $\Gamma(\Sigma^0\gamma)$ and $\Gamma(\Lambda\gamma)$ differ in magnitude, the ratio $\Gamma(\Sigma^0\gamma)/\Gamma(\Lambda\gamma)$ is approximately the same from model to model: the ratios are

	Yield	Statistical Uncertainty	Systematic Uncertainty
Estimated $\Sigma^0\pi^0$ counts	97.7	17.9	+3.5, -4.3
Raw " π^0 " counts	589.7	34.5	+9.4, -12.5
$\Sigma^0\pi^0$ tail	10.0	7.6	+0.4, -0.4
$\Lambda(1405) \rightarrow \Sigma^+\pi^-$	13.0	2.8	+0.5, -0.6
$\Sigma(1385) \rightarrow \Sigma^+\pi^-$	8.6	2.0	+0.2, -1.6
$K^* \rightarrow K^+\pi^0$	19.9	9.7	+4.0, -19.9
Corrected π^0 counts	538.2	36.8	+22.1, -13.1
Raw γ counts	12.17	4.7	+2.2, -3.9
$\Lambda(1405) \rightarrow \gamma\Lambda$	0.24	0.09	+0, -0.24
$K^{*+} \rightarrow \gamma K^+$	0.06	0.03	+0.01, -0.06
Corrected γ counts	11.87	4.7	+2.2, -3.9

Table 6.2 Breakdown of statistics for the $\Lambda\gamma$ and $\Lambda\pi$ channels.

Reaction	Acceptance (in units of 10^{-3})
$\Lambda(1405) \rightarrow \Sigma^0\pi^0$	1.505 ± 0.078
$\Lambda(1405) \rightarrow \Sigma^+\pi^-$	0.200 ± 0.020
$\Lambda(1405) \rightarrow \Lambda\gamma$	2.27 ± 0.24
$\Sigma(1385) \rightarrow \Lambda\gamma$	1.78 ± 0.12
$\Sigma(1385) \rightarrow \Lambda\pi^0$	1.664 ± 0.083
$\Sigma(1385) \rightarrow \Sigma^+\pi^-$	0.393 ± 0.065
$K^{*+} \rightarrow K^+\gamma$	0.56 ± 0.10
$K^{*+} \rightarrow K^+\pi^0$	0.57 ± 0.10

Table 6.3 Overall acceptances used in the calculation of the branching ratio. The uncertainties are statistical only.

0.081, 0.099, and 0.086 for the NRQM [9, 8], chiral bag [2], and relativized CQM [10] calculations, respectively. Using the average (0.0887) of these three ratios, I obtain

$$\Gamma(\Lambda\gamma) = 640 \pm 270^{+130}_{-220} \text{ keV}. \quad (6.20)$$

6.3 Requiring photons

So far the photons have been identified via the missing mass technique. If these correspond to true radiative events, some of these photons should be seen in the electromagnetic calorimeters (the limited coverage of the calorimeters precludes observing all of them). Figure 6.3 demonstrates that of the 22 γ candidates with missing mass squared between -0.007 GeV^2 and 0.0035 GeV^2 , only 7 point toward the active area of the forward electromagnetic calorimeters. Of these seven events, three contain one photon candidate detected by the electromagnetic calorimeter. Figure 6.4A plots the difference in θ_{LAB} between the predicted direction of the photon and the observed direction versus the difference in ϕ between the predicted and observed directions. In both dimensions the angular difference is small. Figure 6.4B plots the observed photon energy (in the lab frame) versus the predicted energy. Again, good agreement is found.

6.4 Discussion and conclusion

The radiative width (equation 6.20) I obtain is roughly twice as large as the largest theoretical prediction. The possibility that these predictions are too small needs to

(A) Predicted γ angular distribution

(B) Counts surviving fiducial cut

Figure 6.3 Predicted photon θ_{LAB} versus ϕ . (A) All counts with missing mass squared between -0.007 GeV^2 and 0.0035 GeV^2 . (B) Those counts satisfying a crude fiducial cut excluding the coils of the magnet and the hole in the forward direction ($< 8^\circ$). The horizontal dashed line indicates the maximum angle of coverage of the forward electromagnetic calorimeters.

Figure 6.4 “Golden” events for which a photon is identified in the calorimeter. (A) Scatter plot of the angular differences in θ_{LAB} and ϕ between the predicted direction and the observed direction. (B) Scatter plot of the observed γ energy versus the predicted γ energy (from the missing mass technique).

be considered. The MIT bag model calculation[41] and the NRQM calculation[7] for the partial width for the process $\Delta^+(\frac{3}{2}^+) \rightarrow p\gamma$ are 338 keV and 361 keV, respectively - both $\approx 50\%$ of the experimental result of 686 ± 37 keV[6] (using $M_\Delta = 1234.9 \pm 1.4$ MeV). The MIT bag model also under-predicts the Σ^0 width by $\approx 50\%$ [9]. The good agreement between the bag model calculations for the $\Lambda(1405) \rightarrow \Lambda\gamma$ width and the experimental result may be considered fortuitous considering the models' failure to obtain the observed $\Lambda(1405)$ mass. Clearly, there is some indication of something wrong or missing in the models.

In conclusion, I have presented evidence for the decay $\Sigma^0(1385) \rightarrow \gamma\Lambda$. The branching ratio relative to the $\Lambda\pi^0$ channel is $0.021 \pm 0.008^{+0.004}_{-0.007}$, corresponding to a partial width of $640 \pm 270^{+130}_{-220}$ keV. This is an exciting result because the radiative decay of the $\Sigma^0(1385)$ has never been observed before.

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Appendix A

Bremsstrahlung

The Bremsstrahlung cross section depends on the atomic number Z of the radiator material, the energy ν of the photon and the ratio ϵ of the Bremsstrahlung electron's energy E to the incident electron beam energy E_0 according to the formula [42]

$$\frac{d\sigma}{d\nu} = \frac{4Z^2 r_e^2 \alpha}{\nu} \left((1 + \epsilon^2) \left(\frac{\Phi_1(\xi)}{4} - \frac{\ln(Z)}{3} - f(Z) \right) - \frac{2\epsilon}{3} \left(\frac{\Phi_2(\xi)}{4} - \frac{\ln(Z)}{3} - f(Z) \right) \right), \quad (\text{A.1})$$

where $r_e^2 = (e^2/mc^2)^2$ and $f(Z)$ is a ‘‘Coulomb correction’’ whose dependence on Z has been empirically determined to be [43]

$$f(Z) \approx a^2((1 + a)^{-1} + 0.20206 - 0.0369a^2 + 0.0083a^4 - 0.002a^6), \quad a = Z/137. \quad (\text{A.2})$$

The terms $\Phi_1(\xi)$ and $\Phi_2(\xi)$ are screening functions that depend on the dimensionless parameter ξ defined by

$$\xi = \frac{100m_e c^2 h \nu}{E_0 E Z^{1/3}}. \quad (\text{A.3})$$

For the case of a gold radiator ($Z = 79$), ξ ranges between 0.0012 for $\nu = 0.2E_0$ and 0.09 for $\nu = 0.95E_0$. Since $\xi \ll 1$, the formula for the cross section simplifies to

$$\frac{d\sigma}{d\nu} \approx \frac{4Z^2 r_e^2 \alpha}{\nu} \left((1 + \epsilon^2 - \frac{2\epsilon}{3}) (\ln(183Z^{-1/3}) - f(Z)) + \frac{\epsilon}{9} \right), \quad (\text{A.4})$$

with the approximate $1/E_\gamma$ behavior alluded to in section 2.2.2.

Appendix B

Details of the Scintillator reconstruction

B.1 The SC calibration map

Subsystem: NMIP_ADC, nitems: 2

Item: left, length: 288, type: float, narray:22

Item: right, length: 288, type: float, narray:22

Subsystem: NMIP_ADCu, nitems: 2

Item: left, length: 288, type: float, narray:18

Item: right, length: 288, type: float, narray:18

Subsystem: T0_TDC, nitems: 2

Item: left, length: 288, type: float, narray:15

Item: right, length: 288, type: float, narray:15

Subsystem: T0_TDCu, nitems: 2

Item: left, length: 288, type: float, narray:14

Item: right, length: 288, type: float, narray:14

Subsystem: T1, nitems: 2

Item: left, length: 288, type: float, narray:14

Item: right, length: 288, type: float, narray:14

Subsystem: T1u, nitems: 2

Item: left, length: 288, type: float, narray:14

Item: right, length: 288, type: float, narray:14

Subsystem: T2, nitems: 2

Item: left, length: 288, type: float, narray:14

Item: right, length: 288, type: float, narray:14

Subsystem: T2u, nitems: 2

Item: left, length: 288, type: float, narray:14

Item: right, length: 288, type: float, narray:14

Subsystem: T_sigma, nitems: 2

Item: first_param, length: 288, type: float, narray:0

Item: second_param, length: 288, type: float, narray:0

Subsystem: WALK0u, nitems: 2

Item: left, length: 288, type: float, narray:7

Item: right, length: 288, type: float, narray:7

Subsystem: WALK1, nitems: 2

Item: left, length: 288, type: float, narray:9

Item: right, length: 288, type: float, narray:9

Subsystem: WALK1u, nitems: 2

Item: left, length: 288, type: float, narray:7

Item: right, length: 288, type: float, narray:7

```

Subsystem: WALK2,      nitems: 2

    Item: left,      length: 288,    type: float,    narray:9
    Item: right,     length: 288,    type: float,    narray:9

Subsystem: WALK2u,      nitems: 2

    Item: left,      length: 288,    type: float,    narray:7
    Item: right,     length: 288,    type: float,    narray:7

Subsystem: WALK_A0,     nitems: 2

    Item: left,      length: 288,    type: float,    narray:9
    Item: right,     length: 288,    type: float,    narray:9

Subsystem: atten_length, nitems: 2

    Item: left,      length: 288,    type: float,    narray:19
    Item: right,     length: 288,    type: float,    narray:19

Subsystem: atten_u,     nitems: 2

    Item: left,      length: 288,    type: float,    narray:19
    Item: right,     length: 288,    type: float,    narray:19

Subsystem: delta_T,     nitems: 4

    Item: dc2sc,     length: 24,     type: float,    narray:1
    Item: left_right, length: 288,    type: float,    narray:31
    Item: paddle2paddle, length: 288, type: float,    narray:72
    Item: panel2panel, length: 24,     type: float,    narray:8

Subsystem: delta_Tu,    nitems: 3

```



```

Item: left_right,      length: 288,    type: float,    narray:4
Item: paddle2paddle,   length: 288,    type: float,    narray:72
Item: panel2panel,     length: 24,     type: float,    narray:8
Subsystem: pedestals,  nitems: 2

Item: left,            length: 288,    type: float,    narray:34
Item: right,           length: 288,    type: float,    narray:34
Subsystem: pedu,       nitems: 2

Item: left,            length: 288,    type: float,    narray:25
Item: right,           length: 288,    type: float,    narray:25
Subsystem: pulser,     nitems: 1

Item: normalisation,   length: 1,      type: float,    narray:3
Subsystem: status,     nitems: 2

Item: left,            length: 288,    type: int,      narray:12
Item: right,           length: 288,    type: int,      narray:12
Subsystem: time_walk,  nitems: 1

Item: ref_adc,         length: 1,      type: float,    narray:1
Subsystem: veff,       nitems: 2

Item: left,            length: 288,    type: float,    narray:4
Item: right,           length: 288,    type: float,    narray:4
Subsystem: veffu,      nitems: 2

Item: left,            length: 288,    type: float,    narray:3

```

Item: right, length: 288, type: float, narray:3

B.2 The SC bank

The DDL definition of the SC bank is

```

!-----
!      BANKname BANKtype      ! Comments
TABLE SC      B16    ! create write display delete
! Scintillation counter event bank
!
!  ATtributes:
!  -----
!COL ATT-name FMT Min   Max  ! Comments
!
1  ID      I      1     48 ! the address of the hit detector element
2  TDCL    I      0  65536 ! tdc information (channels)
3  ADCL    I      0  65536 ! adc information (channels)
4  TDCR    I      0  65536 ! tdc information (channels)
5  ADCR    I      0  65536 ! adc information (channels)
!
!  RELations:
!  -----

```



```

!-----
!      BANKname BANKtype      ! Comments
TABLE SC1      ! create write display delete
! Scintillation counter hits bank
!
!  ATtributes:
!  -----
!COL ATT-name FMT Min   Max ! Comments
!
1  ID          I   1    48 ! the address of the hit detector element
2  time_l      F   0 100000 ! time for left paddle(ns)
3  energy_l    F   0  65536 ! energy in left paddle(MeV)
4  time_r      F   0  65536 ! time for right paddle(ns)
5  energy_r    F   0  65536 ! energy in right paddle(MeV)
6  dtime_l     F   0  65536 ! uncertainty in time for left paddle(ns)
7  denenergy_l F   0  65536 ! energy uncertainty for left paddle(MeV)
8  dtime_r     F   0  65536 ! uncertainty in time for right paddle(ns)
9  denenergy_r F   0  65536 ! energy uncertainty for right paddle(MeV)
!
!  RELations:
!  -----

```

```
!COL RELname  RELtype INTbank  ! Comments
```

```
!                               (COL)
```

```
!
```

```
END TABLE
```

```
!
```

The corresponding C data structure is

```
typedef struct {
    int id;           /* the address of the hit detector element */
    float time_l;     /* time for left paddle(ns) */
    float energy_l;   /* energy in left paddle(MeV) */
    float time_r;     /* time for right paddle(ns) */
    float energy_r;   /* energy in right paddle(MeV) */
    float dtime_l;    /* uncertainty in time for left paddle(ns) */
    float denergy_l;  /* uncertainty in energy in left paddle(MeV) */
    float dtime_r;    /* uncertainty in time for right paddle(ns) */
    float denergy_r;  /* uncertainty in energy in right paddle(MeV) */
} sc1_t;

typedef struct {
    bankHeader_t bank;

    sc1_t sc1[1];
} clasSC1_t;
```

B.4 The SCR bank

The DDL definition of the SCR bank is

```

!-----
!      BANKname BANKtype      ! Comments
TABLE  SCR                    ! create write display delete
! Sc Scintillator reconstruction hit bank
!
!  ATtributes:
!  -----
!COL ATT-name FMT Min    Max      ! Comments
1   id          I    0      48      ! paddle id
2   energy       F    0 100000.0    ! Energy (MeV)
3   time         F    0 100000.0    ! time(ns)
4   x            F -999.9  999.9 ! x position in sector coordinate system
5   y            F -999.9  999.9 ! y position in sector coordinate system
6   z            F -999.9  999.9 ! z position in sector coordinate system
7   dx           F -999.9  999.9 ! x error in sector coordinate system
8   dy           F -999.9  999.9 ! y error in sector coordinate system
9   dz           F -999.9  999.9 ! z error in sector coordinate system
10  status       I -999999  999999 ! status word defined in sc.h
11  denenergy     F    0 100000.0    ! uncertainty in Energy (MeV)

```

```

12  dtime      F   0 100000.0      ! uncertainty in time (ns)

!

!   RELations:

!   -----

!COL RELname  RELtype INTbank  ! Comments

!                               (COL)

!

END TABLE

!
```

The corresponding C data structure is

```

typedef struct {

    int id;      /* paddle id */

    float energy; /* Energy (MeV) */

    float time;  /* time(ns) */

    vector3_t pos; /* position in sector coordinate system */

    vector3_t err; /* error in sector coordinate system */

    int status; /* status word defined in sc.h */

    float denenergy; /* uncertainty in Energy (MeV) */

    float dtime; /* uncertainty in time (ns) */

} scr_t;

typedef struct {
```

```

    bankHeader_t bank;

    scr_t scr[1];

} clasSCR_t;

```

B.5 The SCRC bank

The DDL definition of the SCRC bank is

```

!-----
! Time-of-flight Cluster (final result) bank.
!-----

!      BANKname BANKtype      ! Comments
TABLE  SCRC                  ! create write display delete
!Sc Scintillator reconstruction hit bank

!

!  ATtributes:
!  -----

!COL ATT-name FMT Min      Max    ! Comments
1   id          I   0        48 ! cluster id
2   energy       F   0 100000.0 ! cluster Energy (MeV)
3   denergy      F   0 100000.0 ! error in cluster energy (ns)
4   time         F   0 100000.0 ! cluster (energy-weighted) time(ns)
5   dtime        F   0 100000.0 ! error in cluster time (ns)

```



```

6      x      F -999.9  999.9 ! x position in sector coordinate system
7      y      F -999.9  999.9 ! y position in sector coordinate system
8      z      F -999.9  999.9 ! z position in sector coordinate system
9      dx     F -999.9  999.9 ! x error in sector coordinate system
10     dy     F -999.9  999.9 ! y error in sector coordinate system
11     dz     F -999.9  999.9 ! z error in sector coordinate system
12     status  I -999999   999999 ! status word defined in sc.h
!
!      RELations:
!      -----
!COL RELname  RELtype INTbank ! Comments
!
!                                (COL)
!
!
!      END TABLE
!

```

The corresponding C data structure is

```

typedef struct {
    int id;      /* cluster id      */
    float energy; /* cluster Energy (MeV) */
    float denergy; /* error in cluster energy (ns) */
    float time;   /* cluster (energy-weighted) time(ns) */
}

```

```
float dtime;    /* error in cluster time (ns) */

vector3_t pos;  /* position in sector coordinate system */

vector3_t err;  /* {x,y,z} error in sector coordinate system */

int status;     /* status word defined in sc.h */

} scrc_t;

typedef struct {

    bankHeader_t bank;

    scrc_t scrc[1];

} clasSCRC_t;
```